# Overview

## Charm Production

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Charm-Production in $e^+e^-$ Annihilation Around 4 GeV

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Using the CLEO-c detector at the Cornell Electron Storage Ring, we have measured inclusive and exclusive cross sections for the production of $D^+$, $D^0$ and $D^*_s$ mesons in $e^+e^-$ annihilations at thirteen center-of-mass energies between 3.97 and 4.26 GeV. Exclusive cross sections are presented for final states consisting of two charmed mesons ($D\bar{D}$, $D^*\bar{D}$, $D^*_s\bar{D}_s^*$, $D^*_s\bar{D}_s^*$, and $D^*_s\bar{D}_s^*$) and for processes in which the charmed meson pair is accompanied by a pion.

1. Introduction

Hadron production in electron-positron annihilations just above $c\bar{c}$ threshold has been a subject of mystery and little intensive study for more than three decades since the discovery of charm. Recent developments, like the observation of the $Y(4260)$ reported by the BaBar collaboration [1] and subsequently confirmed by CLEO-c [2] and Belle [3], underscore our incomplete understanding and demonstrate the potential for discovery of new states, such as hybrids and glueballs. It is also clear that precise measurements of charm-meson properties are essential for higher-energy investigations of $b$-flavored particles and new states that might decay into $b$. They also offer unique opportunities to test the validity and guide the development of theoretical tools, like Lattice QCD, that are needed to interpret measurements of the CKM quark-mixing parameters [4]. Any comprehensive program of precise charm-decay measurements demands a detailed understanding of charm production.

Past studies of hadron production in the charm-threshold region have been dominated by measurements of the cross-section ratio $R = \sigma(e^+e^- \rightarrow \text{hadrons}, s)/\sigma(e^+e^- \rightarrow \mu^+\mu^-, s)$ that have been made over this energy range by many experiments [5]. Recent measurements with the Beijing Spectrometer (BES) [6] near charm threshold are especially noteworthy. There is a rich structure in this energy region, reflecting the production of $c\bar{c}$ resonances and the crossing of thresholds for specific charm-meson final states. Interesting features in the hadronic cross section between 3.9 and 4.2 GeV include a large enhancement at the threshold for $D^*\bar{D}^*$ production ($\sim 4$ GeV) and a fairly large plateau that begins at $D^*_s\bar{D}_s^*$-threshold. While there is considerable theoretical interest [7-10], there has been little experimental information about the composition of these enhancements.

In this paper we describe measurements of charm-meson production in $e^+e^-$ annihilations at thirteen center-of-mass energies between 3970 and 4260 MeV. These studies were carried out with the CLEO-c detector at the Cornell Electron Storage Ring (CESR) [11] in 2005-6. (Throughout this paper charge-conjugate modes are implied.) The principal objective of the CLEO-c energy scan was to determine the optimal running point for studies of $D^*_s$-meson decays. The same data sample has been used to confirm the direct production of $Y(4260)$ in $e^+e^-$ annihilations and to demonstrate $Y(4260)$ decays to final states in addition to $\pi^+\pi^-J/\psi$ [2]. Specific results presented in this paper include cross-section measurements for exclusive final states with $D^+, D^0$ and $D^+$ mesons and inclusive measurements of the total charm-production cross section and $R$.

2. Data Sample and Detector

The data sample for this analysis was collected with the CLEO-c detector. Both the fast-feedback analysis carried out as data were collected and the detailed analysis reported here are extensions of techniques developed for charm-meson studies at the $\psi(3770)$ [12].

An initial energy scan, conducted during August-October, 2005, consisted of twelve energy points between 3970 and 4260 MeV, with a total integrated luminosity of 60.0 pb$^{-1}$. The scan was designed to provide cross-section measurements at each energy for all accessible final states consisting of a pair of charmed mesons. At the highest energy point these include $DD$, $D^*\bar{D}$, $D^*\bar{D}^*$, $D^*_s\bar{D}_s^*$, $D^*_s\bar{D}_s^*$, and $D^*_s\bar{D}_s^*$, where the first three contain both charged and neutral states. A follow-up run beginning early in 2006 provided a larger sample of 178.9 pb$^{-1}$ at 4170 MeV that proved essential in understanding the composition of charm production throughout this energy region. The center-of-mass energies and integrated luminosities for the thirteen subsamples are listed in Table I. Integrated luminosity is determined by measuring the processes $e^+e^- \rightarrow e^+e^-$, $\mu^+\mu^-$, and $\gamma\gamma$, which are used since their cross sections are precisely determined by QED. Each of the three final states relies on different components of the detector, with different systematic effects. The three individual results are combined using a weighted average to obtain the luminosity needed for this analysis.

CLEO-c is a general-purpose magnetic spectrometer with most components inherited from the CLEO III detector [13], which was designed primarily to
study $B$ decays at the $\Upsilon(4S)$. Its cylindrical charged-particle tracking system covers 93% of the full $4\pi$ solid angle and consists of a six-layer all-stereo inner drift chamber and a 47-layer main drift chamber. These chambers are coaxial with a superconducting solenoid that provides a uniform 1.0-Tesla magnetic field throughout the volume occupied by all active detector components used for this analysis. Charged particles are required to satisfy criteria ensuring successful fits and vertices consistent with the $e^+e^-$ collision point. The resulting momentum resolution is $\sim 0.6\%$ at 1 GeV/c. Oppositely-charged and vertex-constrained pairs of tracks are identified as $K^0_S \rightarrow \pi^+\pi^-$ candidates if their invariant mass is within 4.5 standard deviations ($\sigma$) of the known mass ($\sim 12$ MeV/$c^2$).

The main drift chamber also provides $dE/dx$ measurements for charged-hadron identification, complemented by a Ring-Imageing Cherenkov (RICH) detector covering 80% of $4\pi$. The overall efficiency for pion or kaon identification is greater than 90%, and the misidentification probability is less than 5%.

An electromagnetic calorimeter consisting of 7784 CsI(Tl) crystals provides electron identification and neutral detection over 95% of $4\pi$, with photon-energy resolution of 2.2% at 1 GeV and 5% at 100 MeV. We select $\pi^0$ and $\eta$ candidates from pairs of photons with invariant masses within 3$\sigma$ of the known values $\pi^0$ ($\sim 6$ MeV/$c^2$ for $\pi^0$ and $\sim 12$ MeV/$c^2$ for $\eta$).

### 3. Event-Selection Procedures

The procedures and specific criteria for the selection of $D^+$, $D^0$ and $D^{+}_s$ mesons closely follow previous CLEO-c analyzes and are described in Refs. [12] and [14]. Candidates are identified based on their invariant masses and total energies, with selection criteria optimized on a mode-by-mode basis. We use only the cleanest final states for $D^0$ ($K^-\pi^+$) and $D^+$ ($K^-\pi^+\pi^+$) selection, since these provide sufficient statistics for precise cross-section determinations. For $D^+_s$ we optimize for efficiency by selecting eight decay modes: $\phi\pi^+$, $K^{*0}K^+$, $\eta\pi^+$, $\eta\rho^+$, $\eta'\pi^+$, $\eta'\rho^+$, $\phi\rho^+$, and $K^0_SK^+$. Accepted intermediate-particle decay modes (mass cuts) are $\phi \rightarrow K^+K^-$ ($\pm 10$ MeV), $K^{*0} \rightarrow K^-\pi^- ($$\pm 75$ MeV), $\eta' \rightarrow \eta\pi^+\pi^-$ ($\pm 10$ MeV), and $\rho^+ \rightarrow \pi^+\pi^0$ ($\pm 150$ MeV).

To determine the production yields and cross sections for the final states accessible at a particular center-of-mass energy, we classify events based on the energy of a $D_{(s)}$ candidate ($\Delta E \equiv E_{\text{beam}} - E_{D_{(s)}}$) and its momentum in the form of beam-constrained mass ($M_{bc} \equiv \sqrt{E_{\text{beam}}^2 - P^2_{D_{(s)}}}$). Figure 1 shows the expected behavior in a two-dimensional plot of $\Delta E$ vs. $M_{bc}$ for a Monte Carlo simulation of CLEO-c data at 4160 MeV with about ten times the statistics of our data sample at that energy. There is clear separation of events into the expected final states consisting of two charmed mesons. This separation was exploited during the scan run for a fast-feedback “cut-and-count” determination of event yields. It is also evident in plots of the momenta of charm-meson candidates selected by cutting on candidate mass that the composition of final states can be analyzed by fitting the momentum spectra of $D^0$, $D^+$ and $D^+_s$ candidates. Figure 2 illustrates this with the momentum spectra for $D^0 \rightarrow K^-\pi^+$ candidates within 15 MeV of the nominal mass both in the Monte Carlo sample of Fig. 1 and in 10.16 $pb^{-1}$ of CLEO-c data at 4160 MeV. While no background corrections have been applied to these distributions, the structure of distinct Doppler-

### Table I Center-of-mass energies and integrated luminosity totals for all data samples for the CLEO-c energy scan.

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<th>$E_{cm}$ (MeV)</th>
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<td>3970</td>
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<td>3990</td>
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<th>$\Delta E$ (GeV)</th>
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Figure 1: $\Delta E$ vs. $M_{bc}$ in a Monte Carlo simulation of expected final states at a center-of-mass energy of 4160 MeV, showing clear separation among the expected two-charm-meson final states.
Figure 2: Momentum spectra at 4160 MeV for \( D^0 \rightarrow K^-\pi^+ \) candidates with invariant masses within 15 MeV of the nominal value for Monte Carlo (top) and data (bottom). As described in the text, the concentrations of entries correspond to the dominant expected final states with two charmed mesons.

Figure 3: The mass spectrum of \( X \) in (a) \( e^+e^- \rightarrow D^0\pi^\pm X \) at 4170 MeV, (b) \( e^+e^- \rightarrow D^{\ast\pm}\pi^\mp X \) at 4170 MeV, and (c) \( e^+e^- \rightarrow D^{0\ast}\pi^\pm X \) at 4260 MeV. Peaks at the \( D^\ast \) mass in (a) and the \( D \) mass in (b) are evidence for the decay \( D^\ast D\pi \). The \( D \) peak in (c) confirms \( D^\ast D\pi \) and the \( D^\ast \) peak demonstrates that \( D^\ast D\pi \) is produced at 4260 MeV.

4. Evidence for Multi-Body Production

While the qualitative features of the measured charm-meson momentum spectra accorded with expectations (Fig. 2), initial attempts to fit the spectra did not produce acceptable results. It was quickly concluded that the two-body processes listed above are insufficient to account for all observed charm-meson production. Final states like \( D\bar{D}\pi \), \( D^{\ast}\bar{D}\pi \), in which the charm-meson pair is accompanied by one or more additional pions, emerged as the likely explanation. While not unexpected, these "multi-body" events have not previously been observed in the charm-threshold region, and there are no predictions of the cross sections for \( D^0 \) and \( D^+ \) production through multi-body final states.

To assess which multi-body final states (\( D\bar{D}\pi \), \( D^{\ast}\bar{D}\pi \), etc.) are measurably populated in our data we examine observables other than the charm-meson momenta, because ISR causes smearing of the peaks in the momentum spectra that can obscure the two-body kinematics. We applied \( D^{\ast}(\pm) \) momentum cuts to exclude two-body contributions and examined the distributions of missing mass against a \( D^{\ast}(\pm) \) and an accompanying charged or neutral pion, using charge correlations to suppress incorrect combinations. Figure 3 shows clear evidence for \( D^{\ast}\bar{D}\pi \) events at 4170 MeV, as well as indications of \( D^{\ast\ast}\bar{D}\pi \) in the sample of 13 pb\(^{-1}\) collected at 4260 MeV (Fig. 3c). These events cannot be attributed to two-body production with ISR, because radiative photons would destroy any peak in the missing-mass spectrum. The absence of a peak at the \( D \) mass in Fig. 3a indicates that there is no evidence for \( D D\pi \) production. Analysis of events with \( D_s \) reveals no evidence for multi-body production, consistent with expectations, since the \( D_s^+ D_s^- \pi^0 \) final state violates isospin conservation.

5. Momentum-Spectrum Fits and Cross-Section Results

Candidate momentum spectra for \( D^0 \), \( D^+ \) and \( D_{s}^{+} \) were selected by requiring the invariant mass to be within \( \pm 15 \) MeV of the nominal value. Backgrounds
are estimated with a sideband technique. Sideband regions are taken on both sides of the expected signal, and are significantly larger than the signal region to minimize statistical uncertainty in the background subtraction. Specific widths are set mode by mode based on the expectation of specific background processes.

Having identified the components of multi-body charm production, we determine yields for these channels and the two-body modes by fitting the sideband-subtracted $D^0, D^+$ and $D_s^+$ momentum spectra. Signal momentum distributions for specific channels are based on full GEANT [15] simulations using EvtGen [16] for the production and decay of charmed mesons. The EvtGen simulation incorporates all angular and time-dependent correlations by using individual amplitudes for each node in the decay chain. ISR is included in the simulation, which requires input of energy-dependent cross sections for each final state. We used simple parameterizations of these cross sections constructed by linearly interpolating between the preliminary measurements from our analysis. (In doing this we make the assumption that the energy dependence of the Born-level cross sections is adequately represented by the uncorrected cross sections.) For the multi-body $D^*D^0$ and $D^*D^0\pi$ final states we used a spin-averaged phase-space model within EvtGen.

Momentum fits for the large sample of data at 4170 MeV are shown in Fig. 4 for (a) $D^0 \to K^-\pi^+$, (b) $D^+ \to K^-\pi^+\pi^+$, and (c) $D_s^+ \to \phi\pi^+$ candidates. The lack of $D_s^+$ entries below $\sim 200$ MeV confirms the absence of multi-body $D_s$ production. Because of the relative simplicity of $D_s$ production demonstrated by the $D_s^+ \to \phi\pi^+$ fits and the limited statistics in the sample, we determined the final cross sections for $D_s^+D_s^-$, $D_s^+D_s^+$, and $D_s^+D_s^*$ by using a sideband subtraction technique to count signal events in a region of the $M_{bc}$ and $\Delta E$ plane. The cross sections are then determined from a weighted sum of the yields for the eight $D_s$ decay modes listed in Sect. 3 with weights minimizing the combined statistical and systematic uncertainties that were calculated from previously measured branching fractions and efficiencies determined by Monte Carlo. The cut-and-count analysis gives results that are consistent with momentum fits. There is good agreement among the separately-calculated cross sections for the different $D_s$ decay modes.

Each of the thirteen data subsamples has been analyzed with the techniques developed and refined on data at 4170 MeV. A complete set of fit results is provided in Ref. [17]. Figure 5 shows the $D^0, D^+$ and $D_s$ fits for data sample at 4260 MeV, which are of particular interest because the charm-production cross sections might provide insight to the nature of the $Y(4260)$ state. The fits at 4260 MeV behave similarly to those at lower energy, although a larger proportion of multi-body decays is apparent.

Cross sections for the two-body and multi-body final states are shown in Fig. 5a-c. The uncertainties on the data points are statistical and systematic combined in quadrature. Ref. [17] provides detailed descriptions of the systematic uncertainties of the cross-section determinations. Briefly, there are three sources of systematic uncertainty: determination of the efficiency of charm-meson selection, extraction of yields, and overall normalization. The total systematic uncertainty is not dominated by any one of these.

Track selection and particle identification closely follow previous CLEO-c analyses [12, 14]. The effi-
Figure 5: Sideband-subtracted momentum spectra for (a) $D^0 \rightarrow K^- \pi^+$, (b) $D^+ \rightarrow K^- \pi^+ \pi^+$, and (c) $D_s^+ \rightarrow \phi \pi^+$ (bottom) at 4260 MeV. Data are shown as points with errors and the total fit result is shown as the solid black line. The colored histograms represent fit components, mostly single $D$-production modes. For example, the primary $D^0$ in $D^{*0}\bar{D}^{*0}$, which peaks at 0.8 GeV/c, is shown in bright red. The secondary $D^0$ mesons from the primary $D^{*0}$ decaying via the emission of a $\pi^0$ form the broad peak at 0.7 GeV/c shown in light blue. The second broad peak, at 0.7 MeV/c, consists of $D^0$ mesons from the charged pion decay of the $D^{*+}$ in $D^{*+}\bar{D}^-$. The multi-body events are combined and result in the broad spectrum between 0 and 0.6 GeV/c shown in dark red for $D^{*}\bar{D}\pi$ and in black between 0 and 0.4 GeV/c for $D^{*}\bar{D}^*\pi$.

Figure 6: Exclusive cross sections for two-body and multi-body charm-meson final states, and total observed charm cross section with combined statistical and systematic uncertainties.

ciency for reconstructing charged tracks has been estimated by a missing-mass technique applied to events collected at the $\psi(2S)$ and $\psi(3770)$ resonances. There is good agreement between data and Monte Carlo, with an estimated relative uncertainty of ±0.7% per track. Pion and kaon identification have been studied with $D^0$ and $D^+$ decays in $\psi(3770)$ data, with estimated systematic uncertainties in the respective efficiencies of ±0.3% and ±1.3%.

The extraction of event yields by fitting the charmed-meson momentum spectra (non-$D_s$ modes) incurs systematic uncertainty primarily through the signal functions generated by Monte Carlo, which depend on details of ISR and, in the case of $D^{*}\bar{D}^*$, the helicity amplitudes and resulting $D$-meson angular distributions. As for the exclusive measurements, these details were studied with the large data sample at 4170 MeV, for which statistical uncertainties are small, and the resulting estimated relative systematic uncertainties are applied to all energy points. For the ISR calculation, the exclusive cross sections input to EvtGen were varied from their nominal shapes. While a qualitative constraint of consistency with our measured cross sections was imposed, some extreme variations are included in the final systematic uncertainty. Both the direct effect on the fitted yield of varying a specific mode and the indirect effect of varying other modes were computed, although the former dominates in quadrature.

The yields for $D_s$ final states are determined by direct counts after cutting on $M_{bc}$ and $\Delta E$. Systematic uncertainty arises in these measurements if the Monte Carlo simulation does not provide an accurate determination of the associated efficiency. This is probed by adjusting the cuts and recomputing the cross sections, again using the high-statistics sample at 4170 MeV. The systematic uncertainties assigned
based on these studies are ±3%, ±2.5% and ±5% for $D_s^+ D_s^−$, $D_s^+ D_s^{∗−}$, and $D_s^∗+ D_s^{∗−}$, respectively.

In converting the measured yields to cross sections we must correct for the branching fractions of the charm-meson decay modes. For the non-strange charmed mesons, only one mode is used and CLEO-c measurements [12] provide the branching fractions and uncertainties: ±3.1% for $D^0 \rightarrow K^- \pi^+$ and ±3.9% for $D^+ \rightarrow K^- \pi^+ \pi^+$. For $D_s$ modes we use CLEO-c measurements of the branching fractions for the eight decay modes included in the weighted sum [14]. The world-average value is used for the $D^{∗+} \rightarrow D^0 \pi^+$ branching fraction, with a systematic uncertainty of ±0.7% [3]. Finally, the cross-section normalization also depends on the absolute determination of the integrated luminosity for each data sample, with a systematic uncertainty of ±1%.

A mode-by-mode summary of the systematic uncertainties in the exclusive cross-section measurements is provided in Table II.

As a cross-check, for the two largest data samples (4170 MeV and 4260 MeV), the multi-body cross sections are also determined by fitting the distributions of missing mass against detected $D^0 \pi$, $D^+ \pi$ and $D^* \pi$ combinations. While these measurements are less precise, they show good agreement with the results of the momentum-spectrum fits.

### Table II: Total systematic errors in the exclusive cross-sections.

<table>
<thead>
<tr>
<th>Mode</th>
<th>Relative Error ($10^{-2}$)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Determined by Momentum Fits</td>
<td></td>
</tr>
<tr>
<td>$D\bar{D}$</td>
<td>4.5</td>
</tr>
<tr>
<td>$D\bar{D}^*$</td>
<td>3.4</td>
</tr>
<tr>
<td>$D^<em>\bar{D}^</em>$</td>
<td>4.7</td>
</tr>
<tr>
<td>$D^*\bar{D}\pi$</td>
<td>12.0</td>
</tr>
<tr>
<td>$D^*\bar{D}^{∗}\pi$</td>
<td>25.0</td>
</tr>
<tr>
<td>Determined by Counting</td>
<td></td>
</tr>
<tr>
<td>$D_s D_s$</td>
<td>5.6</td>
</tr>
<tr>
<td>$D_s D_s^+$</td>
<td>5.3</td>
</tr>
<tr>
<td>$D_s^* D_s^{∗}$</td>
<td>6.8</td>
</tr>
</tbody>
</table>

where the contributing cross sections are defined by $\sigma_D = N_D/\epsilon B \mathcal{L}$, where $\epsilon$ and $B$ are the efficiency and branching fraction for the decay mode used ($D^0 \rightarrow K^- \pi^+$, $D^+ \rightarrow K^- \pi^+ \pi^+$, and $D^{∗+} \rightarrow K^- K^+ \pi^+$), $\mathcal{L}$ is the integrated luminosity, and $N_D$ is the yield obtained by fitting the mass spectrum. In the case of $D^0$ and $D^+$, the invariant-mass distribution is fitted to a Gaussian signal and polynomial background. For $D_s$, the event-type requirements are maintained because of the relatively large background for the high-yield $K^- K^+ \pi^+$ decay mode. For our energy points below 4120 MeV, where $D_s$ production occurs only through $D_s^{∗+} D_s^{∗−}$, the yield is extracted by fitting $M_{bc}$ to a Gaussian signal and ARGUS background function [13]. For 4120 MeV and above, event types involving $D_s^{∗+}$ contribute. For all candidate events that pass the selection requirements for any of $D_s^{∗+} D_s^{∗−}$, $D_s^{∗+} D_s^{∗−}$, and $D_s^{∗+} D_s^{∗−}$ (the last only for 4260 MeV), a fit to the $D_s^{∗+}$ invariant mass is used to determine the yield.

The second cross-check is a determination of the total cross section made by counting multihadronic charm events. The contribution of $uds$ continuum production is estimated with measurements made at $E_{cm} = 3671$ MeV, below $c\bar{c}$ threshold, and extrapolated as $1/s$. Procedures for this measurement are identical to those used to determine the cross section for $e^+e^- \rightarrow \psi(3770) \rightarrow$ hadrons in CLEO-c data at $E_{cm} = 3770$ MeV [10].

Figure 6d shows the inclusive measurements (statistical and systematic uncertainties combined in quadrature) and the sum of the cross sections for the measured exclusive final states, without radiative corrections. The excellent agreement demonstrates that, to current precision, the measured exclusive two- and three-body final states saturate charm production in this region. Furthermore, charm is demonstrated to account for all production of multihadronic events above the extrapolated $uds$ cross section.

For the inclusive-charm cross-section measurements, the systematic uncertainties associated with the per-particle efficiencies for tracking and particle identification are identical to those of the exclusive measurements. The uncertainties in normalization (luminosity and branching fractions) are also identical. Systematic uncertainty in the yield extraction is dominated by the choice of fitting function. This is evaluated mode by mode and propagated into overall systematic uncertainties accounting for all correlations, with combined systematic uncertainties of ±4.3%, ±5.1%, and ±8.6% (±10.6%) for $D^0$, $D^+$, and $D_s^{∗+}$ below (above) 4120 MeV. For the hadron-counting inclusive cross sections, the systematic uncertainties are identical to those of Ref. [12].

For comparison with other experiments and theory it is necessary to obtain Born-level cross sections from the observed cross sections by correcting for ISR. We do this by calculating correction factors following the method of Kuraev and Fadin [20], which gives the
observed cross section at any $\sqrt{s}$:

$$\sigma_{\text{obs}}(s) = \int_0^1 dk \cdot f(k, s) \sigma_B(s_{\text{eff}}),$$  \hspace{1cm} (2)$$

where the Born cross section $\sigma_B$ is a function of the effective center-of-mass energy squared ($s_{\text{eff}} = s(1-(E_\gamma/E_{\text{beam}}))$, and $f(k, s)$ is the ISR kernel. The radiative-correction factor is also calculated following the alternative implementation of Bonneau and Martin [21]. A 4% difference between the two calculations is taken as the systematic uncertainty in the radiative correction.

Figure 7 shows that there is excellent agreement between our inclusive-charm measurement and $R$ measurements in this region made by BES [6] and Crystal Ball [22].

7. Summary and Conclusions

In summary, we have presented detailed information about charm production above $c\bar{c}$ threshold. Realizing the main objective of the CLEO-c scan run, we find the center-of-mass energy that maximizes the yield of $D_s$ to be 4170 MeV, where the cross section of $\sim 0.9$ nb is dominantly $D_s^+D_s^-$. This information has guided the planning of subsequent CLEO-c running, with initial results already presented on leptonic [26] and hadronic [14] $D_s$ decays. The total charm cross section between 3.97 GeV and 4.26 GeV has been measured both inclusively and for specific two-body and multi-body final states. Internal consistency is excellent and radiatively-corrected inclusive cross sections are consistent with previous experimental results. Figure 6 shows that the observed exclusive cross sections for $D\bar{D}$, $D^*\bar{D}$, $D^*\bar{D}^*$, $D_s^*\bar{D}_s^*$, $D_s^*\bar{D}_s^*$, $D^*\bar{D}^\pi$, and $D^*\bar{D}^*\pi$ exhibit structure that reflects the intricate behavior expected in the charm-threshold region. Figure 8 provides a comparison between our measured cross sections and the updated calculation of Eichten et al. [9, 24]. There is reasonable qualitative agreement for most of the two-charm-meson final states. The most notable exception is the cross section for $D^*D^*$ in the region between 4050 and 4200 MeV, where the measurement exceeds the prediction by as much as 2 nb. This corresponds to nearly a factor-of-two disagreement in the ratio of $D^*D^*$ to $D^*D$ production, accounting for about two thirds of the difference in the total charm cross section. This is a much larger effect than the absence of a multi-body component from the theoretical prediction.

It has been suggested by Dubynskiy and Voloshin [25] that the existence of a peak in the $D^*\bar{D}$ and $D_s\bar{D}_s$ channels at the $D^*D^*$ threshold, along with the observation that there is a minimum in $D\bar{D}$, can be interpreted as a possible new narrow resonance, but avail-
able data are insufficient for a definitive assessment.

The $D^*\bar{D}^*$ cross section exhibits a plateau just above its threshold. This contrasts with $D\bar{D}$, which we observe to peak at threshold, in agreement with recently presented preliminary results from Belle [23].

Acknowledgments

We gratefully acknowledge the effort of the CESR staff in providing us with excellent luminosity and running conditions. We thank E. Eichten and M. Voloshin for useful discussions. In addition, we thank the organizers of the conference.

References

Determination of the excited charmonia parameters

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The scan data for $R$ value measurement taken in 1999 with the BESII detector between 3.7 and 5.0 GeV are fitted to determine resonance parameters (mass, total width, electron width) of the high mass charmonium states, $\psi(3770)$, $\psi(4040)$, $\psi(4160)$ and $\psi(4415)$. Various effects, the phase angle of resonance, interference, energy-dependent width, and the initial state radiative correction, are considered, and the model dependent results are compared.

1. Introduction

Figure 1 shows the measured $R$ values between 3.0 and 4.8 GeV. Above the open charm thresholds where resonance structures show up, the measured $R$ values are used to determine the parameters of resonances with $J^{PC} = 1^{−−}$. For the high mass charmonium resonances, the $\psi(3770)$ was measured by MARK-I [1], DELCO [2], MARK-II [3] and BES [4, 5]; the $\psi(4040)$ and $\psi(4160)$ were measured by DASP [6]; and the $\psi(4415)$ was measured by DASP [6] and MARK-I [7]. There were also some other measurements of $R$ values as reported in Refs. [8–10], but no attempt was made to determine resonance parameters. The resonance parameters in the Particle Data Group (PDG)’s compilation remained unchanged for more than 20 years up to the 2004 edition [11]. The resonance parameters for the three high mass resonances were updated in PDG2006 [12], mainly based on Seth’s evaluation [13] using BESII [14, 15] and Crystal Ball [10] data.

![Figure 1: Measured R values between 3.0 and 4.8 GeV and the pQCD prediction](image)

Figure 1: Measured $R$ values between 3.0 and 4.8 GeV and the pQCD prediction

The $R$ value is measured at BES with the following expression [14, 15]

$$R_{\text{exp}} = \frac{N_{\text{obs}}^\text{had} - N_{\text{bg}}}{\sigma_\text{trg}^0 L e_\text{trg}^0 (1 + \delta_{\text{obs}})}.$$

where $N_{\text{obs}}^\text{had}$ is the number of observed hadronic events, $N_{\text{bg}}$ is the number of the residual background events, $L$ is the integrated luminosity, $(1 + \delta_{\text{obs}})$ is the effective correction factor of the initial state radiation (ISR), $e_\text{trg}^0$ is the hadronic detection efficiency determined by the Monte Carlo without bremsstrahlung simulation, $e_{\text{trg}}$ is the trigger efficiency, and $\sigma_\text{trg}^0$ is the theoretical Born cross section for $e^+ e^- \rightarrow \mu^+ \mu^-$. The measurement of $R$ values and determination of resonance parameters are intertwined: the factor $(1 + \delta_{\text{obs}})$ in Eq.(1) contains contributions from the resonances and depends on the resonance parameters. Therefore, the calculation of $(1 + \delta_{\text{obs}})$ and the measurement of $R$ need a number of iterations before stable and convergent results are obtained.

In this paper, the memo in BES about the global fit over the center-of-mass energy region from 3.7 to 5.0 GeV covering the high mass resonances $\psi(3770)$, $\psi(4040)$, $\psi(4160)$ and $\psi(4415)$, is abstracted. In the fitting, the energy-dependent full widths, the relative phases between the resonances are considered, and the results obtained by different models are compared.

2. Related models and formulae

2.1. Breit-Wigner amplitude

Resonance is an unstable particle. In quantum mechanics, the wave function of a non-stationary state with central angular frequency $\omega = M/\hbar$ and lifetime $\tau = \hbar/\Gamma$ is (in units of $\hbar = c = 1$)

$$\Psi(t) = \Psi(0) e^{-i\omega t} e^{-t/2\tau} = |\Psi(0)| e^{i\delta} e^{-it(M - i\Gamma/2)},$$

where $\theta(t)$ is the step function, $M$ is the nominal mass, $\Gamma$ the decay width, and $\delta$ is the initial phase angle at the moment of production, phase factor $e^{i\delta}$ is the intrinsic quantum characteristic of the particle. The amplitude as the function of energy $W$ is

$$\Phi(W) = \int \Psi(t) e^{iWt} dt = \frac{Ke^{i\delta}}{W - M - i\Gamma/2}. \quad (3)$$

where $K$ is a real number related to the properties of the resonance. Usually, the non-relativistic Breit-Wigner amplitude is written as

$$T(W) = \frac{\frac{1}{2} \Gamma e^{i\delta}}{W - M - i\Gamma/2} = \frac{\frac{1}{2} \sqrt{\Gamma^2 - \Gamma} e^{i\delta}}{W - M - i\Gamma/2}. \quad (4)$$
For a specialized process $e^+e^- \rightarrow \text{resonance} \rightarrow$ hadronic final state $f$, the Breit-Wigner amplitude is

$$T^f(W) = \sqrt{\frac{\Gamma^e\Gamma_f}{W - M - i\Gamma/2}} e^{i\delta} = \sqrt{\frac{\Gamma\Gamma_f}{W - M - i\Gamma/2}} e^{i\delta}, \quad (5)$$

where $\Gamma \equiv \Gamma^{tot}$ is the total width, $\Gamma^e \equiv B_e\Gamma$ is the electron width corresponding to the initial state $e^+e^-$, and $\Gamma_f \equiv B_f\Gamma$ is the hadronic width for the decay channel $f$. $B_e$ and $B_f$ are the branch ratios of the decay resonance $\rightarrow e^+e^-$ and $\rightarrow f$. The total width is

$$\Gamma^{tot} = \Gamma^e + \Gamma^\mu + \Gamma^\tau + \Gamma^{had} \quad (6)$$

where $\Gamma^e$ and $\Gamma^\mu$ are decay width for $\mu^+\mu^-$ and $\tau^+\tau^-$. Here the universality of the leptons for $\Gamma^e = \Gamma^\mu$ is used, and the kinematic suppression factors for $\Gamma^{tot}$, 0.48, 0.66, 0.72 and 0.78 for $\psi(3770)$, $\psi(4040)$, $\psi(4160)$ and $\psi(4415)$ respectively, are considered. In experiment, the value of the $\Gamma^e$ includes the contribution from the vacuum polarization effect.

In the relativistic case, the amplitude of the negative-energy state ($W \rightarrow -W$) must be included in direct way

$$T^f(W) = \left[ \frac{1}{W - M - i\Gamma/2} + \frac{1}{-W - M - i\Gamma/2} \right] e^{i\delta} = \frac{\sqrt{\Gamma\Gamma_f}(M - i\Gamma/2)}{\sqrt{W^2 - M^2 + i\Gamma M}} e^{i\delta} \quad (7)$$

For narrow resonances, $M \gg \Gamma$, the terms $i\Gamma/2$ in numerator and $\Gamma^2/4$ in denominator may be neglected, one can get the usually relativistic Breit-Wigner amplitude

$$T^f = \frac{M\sqrt{\Gamma\Gamma_f}}{W^2 - M^2 + i\Gamma M} e^{i\delta}. \quad (8)$$

The high mass charmonium (assigned as $r$) can decay into several two-body final states $f$. According to the Eichten model [16] and experimental data [17], the decay channels (including their conjugate states) are:

$$\psi(3770) \Rightarrow D\bar{D}; \quad \psi(4040) \Rightarrow D\bar{D}, D^*\bar{D}^*, DD^*, D_s\bar{D}_s; \quad (9)$$

$$\psi(4140) \Rightarrow D\bar{D}, D^*\bar{D}^*, DD^*, D_s\bar{D}_s, D_s\bar{D}_s^*; \quad \psi(4415) \Rightarrow D\bar{D}, D^*\bar{D}^*, DD^*, D_s\bar{D}_s, D_s\bar{D}_s^*, D_s^*\bar{D}_s^*, D\bar{D}_1, D\bar{D}_2. \quad (9)$$

The total squared inclusive amplitude of the resonances is the double summations: the inner is coherent sum for the same state $f$ decaying from different resonances $r$, and the outer is incoherent sum over all different decay channels $f$,

$$|T_{res}|^2 = \sum_f \sum_f |T_f(W)|^2. \quad (10)$$

The resonant cross section expressed in $R$ value is

$$R_{res} = \frac{12\pi}{s}(|T_\psi|^2 + |T_{res}|^2), \quad (11)$$

where the influence of the $\psi(2S)$ tail with the same parameters as in PGD2006 is considered.

### 2.2. Continuum backgrounds

The contribution from continuum background originating from light quark pairs ($u\bar{u}$, $d\bar{d}$ and $s\bar{s}$) can be calculated by pQCD above 2 GeV [12]. Since the fitted energy region is close to the open-cc threshold, their cross sections can be described by phenomenological models or empirical expressions. Here, following two forms are considered.

#### 2.2.1. Polynomial form

Assume that there are many continuum channels above the open-charm threshold, such as the multi-body states $DD'$, $DDM'$ and $DDMM'$ (where $M$ stands for the meson made of light and strange quarks, and $D$ for charmed meson), the continuum cross sections are expected to vary smoothly. For simplicity, the continuum backgrounds with a second order polynomial may be a reasonable choice,

$$R_{con} = C_0 + C_1(W - 2M_{D^\pm}) + C_2(W - 2M_{D^\pm})^2, \quad (12)$$

where $C_0$, $C_1$ and $C_2$ are free parameters, and $M_{D^\pm}$ is the mass of the lightest charmed meson $D^\pm$. Parameter $C_0 \sim R_{con}^{(s)}$ represents the contributions from light quarks, $C_1 > 0$ expresses the contribution of the charmed states that increases with energy, $C_2 < 0$ ensures the saturation of $R_{con}$ at energies well above the charm threshold.

#### 2.2.2. DASP-form

Assume that the continuum charm backgrounds are only the two-body states producing in the process $e^+e^- \rightarrow \gamma^* \rightarrow$ continuum mechanism $\rightarrow DD'$, DASP group used the following continuum background form [6]

$$\sigma_{con}(c) = \sigma(3.6) \frac{(3.6 \text{ GeV})^2}{s} + \sum_{k=1}^{6} A_k \beta_k^4 F^2 \frac{s}{s_0}, \quad (13)$$

where $F = (1 - s/s_0)^{-1}$ is the form factor, $s_0 = 3.1$ GeV, and $\sigma(3.6)$ is the background cross section at 3.6 GeV, index $k$ runs 1 to 6 including $DD$, $DD^*$, $D^*\bar{D}^*$, $D_s\bar{D}_s$, $D_s\bar{D}_s^*$ and $D_s^*\bar{D}_s^*$ states, $\beta_k$ is the velocity of the decaying particles, $A_k$ are the free parameters.
3. Energy-dependent width

The total width of the wide resonance is not a constant. The calculation for the energy-dependent width $\Gamma_f^I(s)$ concerns the strong interactions. In practice, some phenomenological methods are used to describe the behavior of $\Gamma_f^I(s)$.

In quantum mechanics, the decay of a resonance may proceed from barrier penetration, which predicts the hadronic width is dependent on the momentum $P_f$ and quantum number $L$ of the orbital angular momentum of the decaying final state $f$. The hadronic width $\Gamma_f^I(s)$ contains a phase space factor and a centrifugal barrier factor $B_L$. For the first level approximation, $B_L$ may be derived from a squared well potential model in the non-relativistic quantum mechanics [18],

$$\Gamma_f^I(s) = \hat{\Gamma}_r \sum_L Z_f^{2L+1} B_L.$$  \hspace{1cm} (14)

where $\hat{\Gamma}_r$ is the parameter to be determined by fitting experimental data, $Z_f \equiv \rho P_f$, $\rho$ is the radius of the interaction dimension which dimension is about $1 \sim 3$ fermi (which value is insensitive to the results), and $P_f$ is the decaying momentum of final state $f$. The decay partial wave is proportional to $P_f^{2L+1}$ due to the centrifugal barrier [19]. When the decay momentum $P_f$ is not large, the partial decay width decreases rapidly with the increase of $L$. This effect makes the unsymmetrical Breit-Weigner cross section and the lower partial waves are important. A relativistic correction factor $2M_r/(M_r + W)$ may be introduced in Eq.(14) [20]

$$\Gamma_f^I(s) = \hat{\Gamma}_r \frac{2M_r}{M_r + W} \sum_L Z_f^{2L+1} B_L.$$  \hspace{1cm} (15)

The first four energy-dependent partial wave functions $B_L$ are [18]

$$B_0 = 1, \quad B_1 = 1 + Z^2, \quad B_2 = 9 + 3Z^2 + Z^4,$$
$$B_3 = 225 + 45Z^2 + 6Z^4 + Z^6.$$  \hspace{1cm} (16)

When resonance decays in several hadronic channels, hadronic width is the summation of all its partial widths

$$\Gamma_{\text{had}}(s) = \sum_f \Gamma_f^I(s).$$  \hspace{1cm} (17)

3.1. Fitting schemes

The resonant parameters can be determined by fitting the scanned $R$-like values with the software package MINUIT [21] using a least squares method that minimize the objective function

$$\chi^2 = \sum_i \frac{[f_c \hat{\Gamma}_{\text{exp}}(W_i) - \hat{\Gamma}_{\text{the}}(W_i)]^2}{f_c \Delta \hat{\Gamma}_{\text{exp}}(W_i)^2} + \frac{(f_c - 1)^2}{\sigma_c^2},$$  \hspace{1cm} (18)

where $W_i$ is the energy of the measured point. The experimental and theoretical quantities are

$$\hat{R}_{\text{exp}} = \frac{N^{\text{obs}}_{\text{had}} - N_{\text{bg}}}{\sigma_{\mu\mu} N_{\text{elec}} \sigma_{\text{had}}}.$$  \hspace{1cm} (19)

and

$$\hat{R}_{\text{the}} = (1 + \delta_{\text{obs}}) R_{\text{the}}.$$  \hspace{1cm} (20)

If interferences between the continuum and resonant states are ignored (in fact, one can not treat them well for the limited knowledge of the strong interaction), $R_{\text{the}}$ is given by

$$R_{\text{the}} = R_{\text{con}} + R_{\text{res}}.$$  \hspace{1cm} (21)

The term $\Delta \hat{R}_{\text{exp}}(W_i)$ in Eq. (18) is the combined statistical and non-common systematic errors (except the error of ISR) and held a constant during the fit. The error common to all the points $\sigma_c$ (3%) is not included in $\Delta \hat{R}_{\text{exp}}(W_i)$. The scale factor $f_c$ reflects the influence of the common error. In each iteration, the resonant parameters used in the calculation of $(1 + \delta_{\text{obs}})$ and $R_{\text{the}}$ are updated to the new values.

In order to understand the model dependence of the results, some fitting schemes with different models are studied. Here, three of them are selected from the BES memo and given. In scheme A, the Breit-Wigner amplitude Eq.(8), continuum background Eq.(12) and energy-dependent width Eq.(15) are used, which is favored one by BES; in scheme B, the Eq.(12) is replaced by Eq.(13); in scheme C, the all functions are the same as scheme A, but the all of the phase angles $\delta_f$ fixed at 0.

4. Results

The values of the parameters of the high mass charmonium states determined by different schemes in this work and the values in PDG and reference [13] are listed in Table I. The differences among the different schemes may supply the some references for the estimation of the uncertainty of models. The figures of the updated $R$ values (the percentages of the error are the same as in paper [14, 15]) between 3.7 and 5.0 GeV and the fitting curves obtained by scheme A, B, and C are shown in Fig. 2. It is noticed that the mass of $\psi(4160)$ in scheme A and B is about 30 MeV higher than in scheme C and PDG, the differences are much larger than the given errors, the other parameters are consistent within the errors. One may see the interference curve with nonzero phase angles (scheme A and B) are different from the one with zero phase (scheme C). If the energy-dependent width Eq.(15) is replaced by another form derived with the effective interaction theory [22] and keep nonzero phase
angles in Eq. (8), the masses of the four charmonia are 3770.9±1.6, 4040.0±1.0, 4192.7±9.7 and 4412.3±8.7 MeV respectively. The comparisons between Fig. 2 show the affects of phase angles on the shape of the resonant structure is significant. The comparisons among scheme A, B and C show the mass difference of $\psi(4160)$ is due to the phase angles, other than by the interference with zero phase angle (the scheme C considered the interference other than phase angle). The comparisons between scheme A and B show the background forms are insensitive to the results. The fitted DASP parameters $A_2 - A_5 \approx 0$ in Eq. (13), which violate the experiments [17], and this may be understood as the inclusive data can not supply enough information to determine correct ratios among different channels. The comparisons of the measured $R$ values in different schemes, see Fig. 2, show that the affect of the ISR factors (in fact the values of resonance parameters) to the measured $R$ values is significant.

In the previous works, the resonance parameters were determined by fitting the fixed $R$ values where the ISR factors are constant in fitting. It is easy to understand that fitted resonance parameters are certainly different from the ones used in the calculations of ISR factor. In this work, the $R$ values and the resonance parameters are measured with the raw data in iterative way.

All of these comparisons show that the values of the resonant parameters determined by fitting are scheme dependent. The international high energy physics community should work out a set of common scheme to avoid the systematic bias among the experimental groups.

BES Collaboration favors the values of the resonances parameters obtained by scheme A [23], and other results can be the estimation of the model uncertainty.

References

Table I The comparisons among the results by different fitting schemes.

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Scheme</th>
<th>$\psi$(3770)</th>
<th>$\psi$(4040)</th>
<th>$\psi$(4160)</th>
<th>$\psi$(4415)</th>
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</thead>
<tbody>
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<td>$M$ (MeV/$c^2$)</td>
<td>PDG2004</td>
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<td>4159 ± 20</td>
<td>4415 ± 6</td>
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<td>4421 ± 4</td>
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<td>4425 ± 6</td>
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<tr>
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<td>BES(Seth)</td>
<td>—</td>
<td>4040 ± 1</td>
<td>4155 ± 5</td>
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<td>$\Gamma_{tot}$ (MeV)</td>
<td>A</td>
<td>3772.0 ± 1.9</td>
<td>4039.6 ± 4.3</td>
<td>4491.7 ± 6.5</td>
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<tr>
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<tr>
<td></td>
<td>C</td>
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<td>4048.4 ± 3.2</td>
<td>4156.2 ± 4.4</td>
<td>4405.2 ± 5.7</td>
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<td>$\Gamma_{ee}$ (keV)</td>
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<td>23.6 ± 2.7</td>
<td>52 ± 10</td>
<td>78 ± 20</td>
<td>43 ± 15</td>
</tr>
<tr>
<td></td>
<td>PDG2006</td>
<td>23.0 ± 2.7</td>
<td>80 ± 10</td>
<td>103 ± 8</td>
<td>62 ± 20</td>
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<tr>
<td></td>
<td>CB(Seth)</td>
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<td>85 ± 10</td>
<td>107 ± 10</td>
<td>119 ± 16</td>
</tr>
<tr>
<td></td>
<td>BES(Seth)</td>
<td>—</td>
<td>89 ± 6</td>
<td>107 ± 16</td>
<td>118 ± 35</td>
</tr>
<tr>
<td></td>
<td>A</td>
<td>30.4 ± 8.5</td>
<td>84.5 ± 12.3</td>
<td>71.8 ± 12.8</td>
<td>71.5 ± 19.0</td>
</tr>
<tr>
<td></td>
<td>B</td>
<td>32.3 ± 9.0</td>
<td>103.6 ± 13.4</td>
<td>61.6 ± 13.8</td>
<td>82.8 ± 26.8</td>
</tr>
<tr>
<td></td>
<td>C</td>
<td>30.8 ± 8.9</td>
<td>109.1 ± 14.9</td>
<td>74.4 ± 14.2</td>
<td>103.8 ± 26.0</td>
</tr>
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<td>$\delta$ (degree)</td>
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<td>—</td>
<td>—</td>
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</tr>
<tr>
<td></td>
<td>A</td>
<td>0</td>
<td>136 ± 46</td>
<td>293 ± 57</td>
<td>234 ± 88</td>
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<tr>
<td></td>
<td>B</td>
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<td>136 ± 42</td>
<td>302 ± 16</td>
<td>245 ± 90</td>
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Double $c\bar{c}$ production in $e^+e^-$ annihilations at high energy

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School of Physics, University of Sydney. NSW 2006, AUSTRALIA.

We review the current state of experimental knowledge on double $c\bar{c}$ production in $e^+e^-$ annihilation. The large cross-sections ($O(20\,fb)$) for $e^+e^-\rightarrow \gamma^*\rightarrow \psi^{(1)}(c\bar{c})_{\text{res}}$ processes have been confirmed by detailed tests and reproduced by a second group: they should now be considered well-established. The latest experimental results concern the case where the second $c\bar{c}$ system is above open-charm threshold: hidden-charm states continue to play a prominent role in the mass spectrum. Some “loose ends” in the field are also briefly discussed.

1. Introduction

Electron-positron annihilation to charmonium and an additional hidden-charm state, or a pair of open-charmed mesons — double $c\bar{c}$ production — has been known for only five years, but is now an established object of theoretical and experimental study. The field is still being driven by data. This presentation surveys current experimental knowledge: after reviewing the history (Section 2) and the now-established results on double-charmonium production $e^+e^-\rightarrow \psi^{(1)}(c\bar{c})_{\text{res}}$ (Section 3), we discuss the latest experimental work, which focusses on states above open-charm threshold (Section 4). Some “loose ends”, which might reward renewed attention, are also noted (Section 5). In closing (Section 6) we summarize both established and new results, and their relation to theory.

2. History

This field grew from studies of inclusive charmonium production $e^+e^-\rightarrow \psi(1)X$ [1,2]. An old CLEO analysis on a small sample [3] had presented evidence for direct decays $\Upsilon(4S)\rightarrow J/\psi X$, distinguished by $J/\psi$ momenta above the endpoint for $\Upsilon(4S)\rightarrow BB\rightarrow J/\psi X$. Using data from early B-factory running—29.4 fb$^{-1}$ on the $\Upsilon(4S)$ resonance, and 3.0 fb$^{-1}$ in the continuum 60 MeV below—Belle [4] excluded such production, setting a limit $B(\Upsilon(4S)\rightarrow J/\psi X)<1.9 \times 10^{-4}$ at 95% confidence. More importantly, they established picobarn cross-sections for $e^+e^-\rightarrow \psi(1)X$ processes in the continuum, and a peculiar momentum spectrum for the produced $\psi(1)$: in the $J/\psi$ sample, the cross-section fell to zero well below the momentum endpoint (see Fig. 3 of Ref. [1]).

The $e^+e^-$ initial state being known, a simple rescaling of the $J/\psi$ momentum gives the mass of the system recoiling against it in the final state:

$$M_{\text{recoil}}(\psi) = \sqrt{(\sqrt{s} - E_0^2)^2 - (p_\psi^*\omega)^2}$$  \hspace{1cm} (1)

An upper bound on momentum thus corresponds to a lower bound on the mass of the recoiling system. Studied in this way [4], the surprising conclusion (Fig. 1) was that the interaction $e^+e^-\rightarrow J/\psi X$ does not proceed if the mass of the recoiling system $X$ is below $c\bar{c}$ threshold, while immediately above threshold there is significant two-body production, with $X=\eta_c, \chi_{c0}, \eta_c'$. Although familiar now, it is worth remembering how completely unexpected this result was at the time.

It was then accepted [5] that charmonium production at $\sqrt{s}\approx 10.6$ GeV was dominated by the $e^+e^-\rightarrow \psi gg$ process, with a momentum spectrum extending to the kinematic endpoint and thus a $M_{\text{recoil}}$ spectrum to low masses; $e^+e^-\rightarrow \psi c\bar{c}$ was an $O(10\%)$ correction. An additional contribution from $e^+e^-\rightarrow \psi g$, where the charmonium develops from a colour-octet $c\bar{c}$ pair, was expected at momentum endpoint (low $M_{\text{recoil}}$). The contradiction with data led to reexamination of both theoretical and experimental methods.

In the recoil mass technique, the system $X$ is not reconstructed and identification is thus indirect. Alternative explanations of the data have sprung from this limitation, while improvements to the method have focused on partial reconstruction and constraints. The original analysis [3] already used a mass-vertex constraint in $J/\psi$ reconstruction, improving $M_{\text{recoil}}$ resolution by a factor of two. Contributions from QED processes are more troublesome. Initial-state radiation (ISR) leads to a high-$M_{\text{recoil}}$ tail (typically model-
dependent) on the lineshape of any peak; more than four tracks are required to suppress contributions from other (low-multiplicity) QED interactions.

The leading alternative interpretation of the data relied on such a process: \(e^+e^- \rightarrow \gamma^* \gamma^* \rightarrow \psi X\) [6]. While \(e^+e^- \rightarrow \gamma^* \rightarrow \psi X\) production requires the second state to be even under charge conjugation, \(\xi^+\xi^- = +1\), the two-virtual-photon process allows \(\xi^+\xi^- = \pm 1\) and in particular permits \(e^+e^- \rightarrow \gamma^* \gamma^* \rightarrow J/\psi J/\psi\). As only the \(J/\psi \eta_c\) signal was statistically significant in the 2002 analysis [1], it was attractive to ascribe the peak (in part) to events with a second \(J/\psi\). A more ambitious (and essentially mirror-image) proposal [7] was that \(e^+e^- \rightarrow J/\psi \eta_c\) did in fact dominate charmonium production, with the gluons sometimes coupling to a glueball state close to the \(\eta_c\) mass.

These and other interpretations were effectively ruled out by an updated Belle analysis [8], using 155 fb\(^{-1}\) of data and techniques designed to discriminate between the theoretical options. Three \(e^+e^- \rightarrow \psi \eta_c\) events were fully reconstructed, to be compared with the 2.6 ± 0.8 expected from the inclusive yield. The recoil mass scale was also calibrated using \(e^+e^- \rightarrow \gamma_{ISR} \psi'\) events and found to have a bias of less than 3 MeV, ruling out confusion between \(\eta_c\) and \(J/\psi\). Attempts to include \(e^+e^- \rightarrow \psi \eta_c(\psi')\) components in the recoil mass fit resulted in negative yields, and restrictive upper limits on their contribution (see Fig. 2, upper plot). An eventual confirmation by BaBar [9] found comparable results (lower plot).

Belle also performed an angular analysis of \(J/\psi \eta_c(\psi')\) and \(J/\psi \chi_{c0}\) events. The distinctive forward peak in \(J/\psi\) production angle (\(\cos \theta_{\text{prod}} \rightarrow 1\)) predicted by the two-virtual-photon model for \(J/\psi J/\psi\) (Fig. 2 of Ref. [6]) was not seen in data: results for all three states were consistent with shapes \(1 + \alpha \cos^2 \theta\), and equal coefficients for production- and helicity-angle distributions (see Table I) as expected for a single virtual photon. The \(J/\psi \eta_c(\psi')\) fits were consistent with \(\alpha = +1\) (P-wave production) as required by \(\eta_c\) quantum numbers, and strongly disfavoured the \(-0.87\) expectation for a spin-zero glueball [7]. One- rather than two-virtual photon production (or the glueball explanation) is thus favoured by all experimental tests.

Table I Coefficients from fits of the function \(1 + \alpha \cos^2 \theta\) to \(J/\psi\) production (\(\theta_{\text{prod}}\)) and helicity angle (\(\theta_{\text{hel}}\)) data at Belle [4], for \(e^+e^- \rightarrow \psi \eta_c(\psi')\). See the text for results under the constraint \(\alpha_{\text{hel}} = \alpha_{\text{prod}}\), and the expectation.

<table>
<thead>
<tr>
<th>((\psi')_{\text{res}})</th>
<th>(\alpha_{\text{prod}})</th>
<th>(\alpha_{\text{hel}})</th>
<th>(\alpha_{\text{hel}} = \alpha_{\text{prod}})</th>
<th>Expectation</th>
</tr>
</thead>
<tbody>
<tr>
<td>(\eta_c)</td>
<td>(1.4_{-0.8}^{+1.1})</td>
<td>(+0.5_{-0.5}^{+0.7})</td>
<td>(+0.93_{-0.47}^{+0.57})</td>
<td>+1 (P)</td>
</tr>
<tr>
<td>(\chi_{c0})</td>
<td>(-1.7 \pm 0.5)</td>
<td>(-0.7_{-0.5}^{+0.7})</td>
<td>(-1.01_{-0.33}^{+0.38})</td>
<td>-1 (S)</td>
</tr>
<tr>
<td>(\eta_c)</td>
<td>(+1.9_{-1.2}^{+2.0})</td>
<td>(+0.3_{-0.7}^{+1.0})</td>
<td>(+0.8_{-0.64}^{+0.86})</td>
<td>+1 (P)</td>
</tr>
</tbody>
</table>

Figure 2: Baseline \(e^+e^- \rightarrow J/\psi (\psi')_{\text{res}}\) results from Belle (upper plot), and the confirmation from BaBar (lower plot). In the upper plot the dashed line shows the upper limit contribution from final states where the fitted yield is insignificant or negative: \(J/\psi (J/\psi; \chi_{c1,c2}; \psi')\).

3. Baseline results

Belle [8] and BaBar [9] measurements of \(e^+e^- \rightarrow \psi(\psi')_{\text{res}}\) cross-sections are summarised in Table I, with competing explanations excluded, and reasonable agreement between the two experiments, these results are no longer in serious dispute. Remarkably, there seems to be no suppression of radially-excited states: cross-sections for \(\psi \eta_c, \psi' \eta_c, \psi' \eta_c, \psi' \eta_c\) are all comparable. This presumably contains some hint as to the production mechanism.

Low-order perturbative calculations, as embodied in nonrelativistic quantum mechanics (NRQCD, [10][11][12]) underestimate the cross-sections by an order of magnitude or more; the discrepancy is reduced, but not removed, when relativistic corrections are taken into account [10]. Neither are other features of the data well-explained: NRQCD predicts \(\alpha \simeq +0.25\) for \(e^+e^- \rightarrow J/\psi \chi_{c0}\) [10], disfavoured by the Belle analysis (Table I), which prefers pure S-wave production.

A calculation in the light-cone formalism [13], however, appears to match at least the \(J/\psi \eta_c\) cross-section. A variety of theoretical approaches are currently being pursued, and it is no longer easy to characterise the issues at stake: on comparison of NRQCD and light-cone estimates, see for example the very ex-
Table II Double-charmonium production cross-sections from Belle and BaBar, with theoretical pre- and post-dictions. Due to background-suppression criteria, the experiments report effective cross-sections for the case where the unreconstructed state (c)_{res} decays to at least 2 ("\(> 0\)") or 4 ("\(> 2\)") charged tracks, and thus underestimate the cross-section.

<table>
<thead>
<tr>
<th>(\sigma(e^+e^- \rightarrow \psi(nS)(c)_{res}) ) [fb]</th>
<th>(\eta_c(1S))</th>
<th>(\chi_{c0})</th>
<th>(\eta_c(2S))</th>
</tr>
</thead>
<tbody>
<tr>
<td>(\psi(2S) \times B_{&gt;0})</td>
<td>Belle [5]</td>
<td>16.3 ± 4.6 ± 3.9</td>
<td>12.5 ± 3.8 ± 3.1</td>
</tr>
<tr>
<td>(\psi(1S) \times B_{&gt;2})</td>
<td>Belle [6]</td>
<td>25.6 ± 2.8 ± 3.4</td>
<td>6.4 ± 1.7 ± 1.0</td>
</tr>
<tr>
<td></td>
<td>BABAR [9]</td>
<td>17.6 ± 2.8^{+1.5}_{-2.1}</td>
<td>10.3 ± 2.5^{+1.4}_{-1.8}</td>
</tr>
<tr>
<td>(\psi(1S))</td>
<td>Braaten and Lee [10]</td>
<td>3.78 ± 1.26</td>
<td>2.40 ± 1.02</td>
</tr>
<tr>
<td>... with relativistic corrections [11]</td>
<td></td>
<td>7.4^{+10.9}_{-4.1}</td>
<td>–</td>
</tr>
<tr>
<td></td>
<td>Liu, He, and Chao [12]</td>
<td>5.5</td>
<td>6.9</td>
</tr>
<tr>
<td></td>
<td>Zhang, Gao, and Chao [12]</td>
<td>14.1</td>
<td>–</td>
</tr>
</tbody>
</table>

We note that a theoretical analysis has appeared since the workshop [15], claiming that the discrepancy in the \(J/\psi\eta_c\) cross-section between NRQCD and experiment is now resolved.

4. The new cutting edge: states above open-charm threshold

Active experimental work has now shifted to the case where the system recoiling against the \(\psi\) is above open-charm threshold. Even allowing for the various \(e^+e^- \rightarrow D^{(*)}\overline{D}^{(*)}\) continuum components in fitting the inclusive \(M_{\text{recoil}}(\psi)\) spectrum, Belle found a 5.0σ peak at \((3936 ± 14)\) MeV (Fig. 1 of Ref. [10]). The background under this peak (called \(X(3940)\)) being too large for further detailed study, Belle explicitly reconstructed a \(D\)-meson and then considered cases where the remaining system was close to the \(D \) or \(D^*\) in mass: \(M_{\text{recoil}}(\psi D) \approx m_{D^{(*)}}\). Constraining such cases to match the \(D^{(*)}\) mass also improved the \(M_{\text{recoil}}(\psi)\) resolution. This allowed the reconstruction of a clear \(X(3940) \rightarrow D\overline{D}\) peak, and an upper limit on the same structure in \(D\overline{D}\) (Fig. 3 of [10]).

Belle has released an updated analysis [17] based on systematic use of this \(D^{(*)}\)-tagging technique. Fig. 3 shows the \(M_{\text{recoil}}(\psi D^{(*)})\) spectrum for 693 fb\(^{-1}\) of data, after reconstruction and mass-constraint of the \(J/\psi\), and then a \(D^0\), \(D^+\), or \(D^{*+}\) meson. A simultaneous fit with the \(D^{(*)}\)-mass sidebands is performed: clear and significant peaks are seen, corresponding to processes \(e^+e^- \rightarrow \psi D\overline{D}\), \(\psi D\overline{D}\), and \(\psi D\overline{D}\).

Monte Carlo study shows that we indeed expect these processes to be reconstructed in this way, with recoil mass resolution of about 30 MeV: smaller than the difference in \(D\) and \(D^*\) masses. Disjoint samples \(|M_{\text{recoil}}(\psi D^{(*)}) - m_{\text{tag}}| < 70\) MeV where the unreconstructed system is tagged as a \(D\) or a \(D^*\) are thus selected (ISR leads to a 10% \(\psi D\overline{D} \rightarrow \psi D D^*\) cross-feed), and \(M_{\text{recoil}}(\psi D^{(*)})\) is then constrained to the mass of the tagged meson. This improves the resolution on \(M(D^{(*)}\overline{D}^{(*)})\) by a factor of 3–10: results for the three samples are shown in Fig. 4. Peaks above the background are seen near threshold in each sample.

Figure 3: The mass of the system recoiling against fully-reconstructed \(J/\psi\) and \(D\) (upper plot) or \(D^{(*)}\) (lower plot) mesons at Belle [17]. The shaded histogram shows the distribution in \(D^{(*)}\)-mass sidebands: the fitted curve is discussed in the text.

Combinatorial backgrounds are taken into account via simultaneous fits to the data in the reconstructed \(D^{(*)}\)-mass signal and sideband regions. The excess over background is fitted with the sum of a threshold function to represent non-resonant \(e^+e^- \rightarrow\)
Figure 4: \(M(D^{(*)}\bar{D}^{(*)})\) spectra for \(e^+e^- \rightarrow J/\psi D^{(*)}\bar{D}^{(*)}\) events at Belle [17], where the first D-meson is explicitly reconstructed, and the associated D-meson is recovered via recoil mass selections and constraints: (upper) \(D\bar{D}\), (middle) \(D\bar{D}'\), and (lower) \(D^{*}\bar{D}'\) samples are shown. Distributions in the reconstructed \(D^{(*)}\)-mass sidebands are shown in yellow, and reflections (middle) \(D\bar{D} \rightarrow D\bar{D}'\) and (lower) \(D\bar{D}' \rightarrow D^{*}\bar{D}'\) in green. The solid curves show the fits described in the text, and the dashed curves their combinatorial background and reflection components. The second panel in the upper and middle plots shows the data and fit with these components subtracted.

<table>
<thead>
<tr>
<th>Final state</th>
<th>(D\bar{D})</th>
<th>(D\bar{D}')</th>
<th>(D^{*}\bar{D}')</th>
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<tr>
<td>Reconstructed</td>
<td>D</td>
<td>D</td>
<td>D*</td>
</tr>
<tr>
<td>Resonant term</td>
<td>—</td>
<td>(X(3940))</td>
<td>(X(4160))</td>
</tr>
<tr>
<td>Mass (MeV)</td>
<td>3878 ± 48</td>
<td>3942^{+7}_{-6} ± 6</td>
<td>4150^{+25}_{-20} ± 15</td>
</tr>
<tr>
<td>Width (MeV)</td>
<td>347^{+316}_{-143}</td>
<td>37^{+126}_{-15} ± 8</td>
<td>139^{+111}_{-61} ± 21</td>
</tr>
<tr>
<td>Significance</td>
<td>4.4(\sigma)</td>
<td>6.0(\sigma)</td>
<td>5.5(\sigma)</td>
</tr>
<tr>
<td>Fit behaviour</td>
<td>unstable</td>
<td>stable</td>
<td>stable</td>
</tr>
</tbody>
</table>

\(\eta_{(e^+e^- \rightarrow \psi X)} \times B_{D^{(*)}\bar{D}^{(*)}}\) (fb) \(= 13.9^{+6.4}_{-4.1} ± 2.2\) \(24.7^{+12.8}_{-8.8} ± 5.0\)

\(\psi D^{(*)}\bar{D}^{(*)}\) production, and an S-wave relativistic Breit-Wigner. In all cases the threshold term is insignificant, and a significant Breit-Wigner peak is seen. In \(M(D\bar{D})\) the peak is broad, and the fit unstable against variations of its conditions; in \(M(D\bar{D}')\) and \(M(D^{*}\bar{D}')\) the fit is stable and the resonant peak significant at over 5\(\sigma\). Various cross-checks are performed, including Monte Carlo and data tests of the background shape in \(D^{(*)}\)-mass sidebands and signal regions; study of charged- and neutral-D subsamples; and fitting of events with reconstructed \(D^*\), and associated \(\bar{D}\): the latter obtains results consistent with the \(J/\psi D\bar{D}^*\) analysis of Fig. 3 (middle), but with lower efficiency and significance. Parameters for the resonant enhancements are summarised in Table III.

The \(X(4160)\) enhancement has not previously been reported. The \(X(3940)\) mass and yield results are consistent with those of the earlier analysis [16], while the width is larger than the published value of \((15.1 ± 10.1)\) MeV: a likelihood function non-parabolic in the width parameter had been noted in that case, with a 52 MeV upper limit at 90\% confidence. The corresponding limit in the new analysis is \(\Gamma < 76\) MeV.

Cross-sections for \(J/\psi X(3940)\) and \(J/\psi X(4160)\) production (Table III, last row) are in the 20 femtobarn class, as for all the significant \(e^+e^- \rightarrow \psi X\) measurements in Fig. 4 (right) are seen. In

5. Sidelines

The concentration on experimentally fruitful problems — establishing and measuring quasi-two-body processes \(e^+e^- \rightarrow \psi(\not{c}\not{\bar{c}})_{\text{res}}\) — has led to a relative neglect of inclusive \(e^+e^- \rightarrow \psi(\not{c}\not{\bar{c}})X\) studies. The 2002 Belle analysis [1] established the fraction

\[
\frac{\sigma(e^+e^- \rightarrow \psi \not{c}\not{\bar{c}})}{\sigma(e^+e^- \rightarrow \psi X)} = 0.59^{±0.15}_{−0.3} ± 0.12
\]

(2)
using the following method: reconstruction and mass-constraint of $J/\psi$, and then an associated $D^{(*)}$ meson; rejection of contamination from $e^+e^- \rightarrow B\bar{B}$ events using momentum requirements; and a two-dimensional fit to obtain $\psi D^{(*)}X$ yields (see Fig. 2 of Ref. [4]). The cross-section thus obtained is model-dependent, relying on simulated $c\bar{c} \rightarrow D^{(*)}X$ fragmentation by PYTHIA to establish efficiencies. There has been no published update of this remarkable result.

(Model-independent but since-unpublished Belle results were presented to the Quarkonium Working Group in 2002, relying instead on counting of $D^0$, $D^+$, and all ground-state charmed hadrons under minimal cuts. Resolution was similar, with a lower limit on the $J/\psi c\bar{c}$ fraction of 0.48 at 95% confidence.)

A complementary problem is the nature of $e^+e^- \rightarrow J/\psi X$ production by processes other than $J/\psi c\bar{c}$. Such production does seem to occur: there is a continuous component in inclusive recoil mass spectra (Fig. 5) in excess of background, below the open-charm threshold. The mystery is why this component — due to the $e^+e^- \rightarrow \psi g$ process? — should obey the $c\bar{c}$ threshold. A problem for experimental study here is the lack of theoretical guidance: predictions for the various $e^+e^- \rightarrow \psi X$ processes that so preoccupied the original Belle [1] and BaBar [2] analyses have been discredited, but new predictive studies have not taken their place. Such work is overdue.

6. Summary

The period of fundamental doubt about double charmonium production is now over. Questions concerning the experimental method have been addressed, and the potential for confusion by $e^+e^- \rightarrow \gamma^*\gamma^* \rightarrow J/\psi X$ or other more exotic processes has been excluded [3]. And the fear that some untraceable mistake had been made is effectively dispelled by BaBar’s confirmation [9] of the Belle results [4].

Those results establish that $e^+e^- \rightarrow \psi c\bar{c}$ dominates charmonium production in the continuum at $\sqrt{s} \approx 10.58$ GeV, while $\psi^{(1)}(c\bar{c})_{res}$ cross-sections are at the 20 fb level, with no suppression of radially excited states. Recent work shows that prominent resonant contributions continue above open-charm threshold, with similar cross-sections; this process is proving fruitful in the search for new hidden-charm states.

The ball now lies in the court of theory. Interpretive work on the new states $X(3940)$ and $X(4160)$ is already underway, but a predictive account of $e^+e^- \rightarrow \psi^{(1)}X$ amplitudes is still lacking. An advantage of the NRQCD approach is its pretension to universal application: at the Tevatron, for example, and the LHC. We await an accurate account of double-charmonium production in $e^+e^-$ annihilation that can also embrace quarkonium production at other facilities.

References

Charm (and Beauty) Production at the Tevatron

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We present recent results on heavy flavour production at Tevatron Run II for typically \( 1 \text{ fb}^{-1} \) of analysed \( pp \) data at \( \sqrt{s} = 1.96 \text{ TeV} \). This includes results on single and correlated open charm and bottom cross sections, charm pair production kinematics, \( J/\psi \), \( \psi(2S) \) and \( \chi_{cJ} \) cross sections and polarisation measurements in \( J/\psi \), \( \psi(2S) \), \( \Upsilon(1S) \), and \( \Upsilon(2S) \).

I. INTRODUCTION

The large \( b\bar{b} \) and \( c\bar{c} \) cross section at \( \sqrt{s} = 1.96 \text{ TeV} \) \( pp \) makes the Tevatron a unique place for the study of the production and decay of heavy flavour. Already in Run I, heavy flavour production measurements challenged theory, finding heavy flavour production cross sections significantly higher than predicted by Next-to-Leading-Order (NLO) QCD calculations (see for example [1] and [2]). The discrepancy between data and theory was particularly dramatic in the quarkonium production, where the “Colour Singlet Model” leading-order QCD calculation underestimates the measured cross section by more than an order of magnitude (see e.g. [3]).

Since the first hadroproduction measurements in Run I, a number of theoretical advances have been made. Fixed-Order Next to Leading Logarithm (FONLL) calculations [3], describe the open charm and \( b \) production cross sections well. Competing models have been put forward that describe the observed quarkonium production rates and \( p_T \) spectra well, but disagree on their results for quarkonium polarisation.

There has also been dramatic experimental progress. In Run II, which started in 2001, the Tevatron collides protons and antiprotons at a centre of mass energy of 1.96 TeV with a bunch crossing every 396 ns at each interaction point. Some of the bunches are by design empty, so while the detectors have to be able to cope with peak rates of 2.5 MHz, the average collision rate is \( \sim 1.7 \text{ MHz} \). Since the start of data taking, the Tevatron has delivered more than \( 3 \text{ fb}^{-1} \) of integrated luminosity at each interaction point, and is now reaching peak luminosities of typically \( 2 \cdot 10^{32} \text{ cm}^{-2} \text{s}^{-1} \), with the best runs exceeding \( 2.8 \cdot 10^{32} \text{ cm}^{-2} \text{s}^{-1} \). Two general purpose detectors take data at the Tevatron, CDF and DØ. Both collaborations have analysed approximately \( 1 \text{ fb}^{-1} \) of nearly \( 3 \text{ fb}^{-1} \) each has on tape.

II. THE TEVATRON, CDF AND DØ

A. The Tevatron Run II

The Tevatron in Run II collides protons and antiprotons at a centre of mass energy of 1.96 TeV with a bunch crossing every 396 ns at each interaction point. Some of the bunches are by design empty, so while the detectors have to be able to cope with peak rates of 2.5 MHz, the average collision rate is \( \sim 1.7 \text{ MHz} \). Since the start of data taking, the Tevatron has delivered more than \( 3 \text{ fb}^{-1} \) of integrated luminosity at each interaction point, and is now reaching peak luminosities of typically \( 2 \cdot 10^{32} \text{ cm}^{-2} \text{s}^{-1} \), with the best runs exceeding \( 2.8 \cdot 10^{32} \text{ cm}^{-2} \text{s}^{-1} \). Two general purpose detectors take data at the Tevatron, CDF and DØ. Both collaborations have analysed approximately \( 1 \text{ fb}^{-1} \) of nearly \( 3 \text{ fb}^{-1} \) each has on tape.

B. The CDF and DØ Detectors

Both the CDF and the DØ collaboration have a strong heavy flavour physics programme, and in the last upgrade many features have been added to the detectors to facilitate this programme. These features include precise vertexing, extended \( \mu \) coverage, and sophisticated read-out electronics and triggers. A description of the DØ detector can be found at [5] and of the CDF detector at [6].

The same detector characteristics that make DØ and CDF powerful B physics experiments, also provide the basis for a wide-ranging charm physics programme. The feature that makes DØ stand out as a heavy flavour experiment particularly is its large \( \mu \) coverage [5], up to \( |\eta| < 2 \). CDF’s most distinctive detector element for B and charm physics is its fully hadronic track trigger [6].

C. Triggers

As we will see below, heavy flavour production at hadron colliders is a vibrant field which is lead by experiment rather than theory. A particular challenge for theory is the quarkonium polarisation, for which we present new results from the Tevatron in this paper.

All numbers, unless accompanied by a reference to a journal publication, are preliminary.
discarded instantly. The trigger is crucial, and has to operate in a challenging environment. There are two basic strategies to trigger on heavy flavour events.

- Trigger on leptons ($\mu$, $e$). This provides a clean signature in an hadronic environment, where most tracks are pions. Both experiments use this strategy. DØ benefits particularly from its large $\mu$ coverage, giving large number of semileptonic and leptonic B and charm events, as well as charmonium and bottomonium decays, e.g. 180k $J/\psi$ in the first 0.1 fb$^{-1}$, and 28k $B \rightarrow J/\psi K$ in 1.6 fb$^{-1}$ (after selection cuts).

- Trigger on the long lifetimes of B and D mesons. This is very challenging in a hadronic environment because it requires very fast track and impact parameter reconstruction in very busy events, in time for the trigger decision. Both experiments use this strategy. But only CDF’s displaced track trigger has enough bandwidth to trigger on fully hadronic decays alone, while DØ have to require an additional lepton in the event to keep the trigger rates manageable. This high-bandwidth displaced track trigger gives CDF unique access to fully hadronic heavy flavour decays, such as 13M $D^0 \rightarrow K^- \pi^+$ events (here, and in similar expressions in the rest of the note, the charge-conjugate decay is always implied) , 0.3M $D^+ \rightarrow \phi(K^+K^-)\pi^+$ event as well as 53k $B_s \rightarrow D_s^+\pi^-$ in $\sim 1.6$ fb$^{-1}$ (numbers after selection cuts).

III. CROSS SECTIONS

A. D and B production cross sections

Amongst the earliest Tevatron Run II results were heavy flavour cross section measurements. Figure 11 shows CDF’s measurement of the prompt differential charm production cross section versus $p_T$ for $|\eta| < 1$, for $D^0$, $D^+$, $D^{*+}$ and $D_s$ mesons. Using only 5.8 pb$^{-1}$, less than 0.5% of the presently analysed dataset, the statistical error, given by the inner error bars, is already smaller than the systematic uncertainty; the combined statistical and systematic uncertainty is given by the outer error bars. The uncertainty on the FONLL calculation 11 is given as the grey band, with the central value shown as a solid line. While the measured cross sections appear to be consistently higher than the FONLL prediction, the results are compatible within the (correlated) errors.

The analyses use the reconstructed impact parameter of the D to distinguish prompt D mesons (with zero impact parameter) from those originating from B decays (with large impact parameters).

The same technique of using either impact parameters or decay length has been used to measure inclusive B cross sections in $B \rightarrow D^0\mu^+\nu X$, $B \rightarrow D^{*+}\mu^+\nu X$ and $B \rightarrow J/\psi X$ decays, where long decay lengths or large impact parameters identify $D^0\mu$, $D^{*+}\mu$ and $J/\psi$ originating from B mesons. The inclusive differential cross sections measured in the $D^0\mu$, $D^{*+}\mu$ and the $J/\psi$ channel are shown in Figures 2 and 3. The results are in good agreement with FONLL calculations. The preliminary result for the integrated inclusive B cross section for $b$-hadrons with $p_T > 6$ GeV and $|y| < 0.6$ is:

$$\sigma(p\bar{p} \rightarrow H_b) = 1.34 \mu b$$

$$\pm 0.08 \mu b(stat)_{+0.13}^{+0.13} \mu b(sys) \pm 0.07 \mu b(BR)$$

for the analysis using $H_b \rightarrow D^0(K^-\pi^+)\mu X$, and

$$\sigma(p\bar{p} \rightarrow H_b) = 1.47 \mu b$$

$$\pm 0.18 \mu b(stat)_{+0.17}^{+0.17} \mu b(sys) \pm 0.11 \mu b(BR)$$

for the analysis using $H_b \rightarrow D^{*+}(D^0(K^-\pi^+)\pi^+)\mu X$, where $H_b$ stands for a generic $b$ hadron. The last uncertainty is due to the uncertainty in the branching fractions of the specific final states of the $D^0$ and $D^*$ being investigated. The result is in good agreement with the FONLL value of 1.39$\pm$0.49 $\mu$b 17.

CDF performed a measurement of the exclusive $B^+ \rightarrow J/\psi K^+$ measurement, using 0.74 fb$^{-1}$, finding

$$\sigma(p_T > 6 GeV, |y| < 1) = (2.65 \pm 0.12(stat) \pm 0.21(sys)) \mu b$$

12. The differential cross sections for the exclusive measurement can be seen in Fig 2 together with the $B \rightarrow J/\psi X$ inclusive results. All measurements disagree significantly with NLO calculations. The agreement with FONLL is however very good.

B. Correlated $b\bar{b}$ and $c\bar{c}$ cross sections

For correlated $b\bar{b}$ and $c\bar{c}$ cross sections, i.e. cross sections where both the quark and the antiquark are within a certain, central rapidity range, higher order terms are expected to be smaller and consequently NLO calculations are expected to describe the data
FIG. 1: Differential charm cross section measured in fully hadronic charm decays using only 5.8 fb\(^{-1}\) at CDF [10]. The inner error bars represent the statistical uncertainty, the outer error bars represent the total uncertainty of the measurement, and the black line and grey band the central value and uncertainty of the FONLL calculation by [11]. No calculation for the \(D_s\) cross section was available at the time.


FIG. 3: Comparing inclusive \(B \rightarrow D \mu \nu X\) with inclusive \(B \rightarrow J/\psi X\) [6] and the FONLL calculation by [17].

momentum \(p_T > 3\) GeV, pseudorapidity \(|\eta| < 0.7\) and invariant mass \(m_{\mu\mu} \in [5, 80]\) GeV. This corresponds to \(b\bar{b}\) pairs with \(p_T \geq 2\) GeV and a rapidity \(|y| \leq 1.3\). Prompt \(\mu\) are separated from \(\mu\) created in charm decays, and those from \(\mu\) created in B decays, using the impact parameters of each muon in the \(\mu^+\mu^-\) pair. The 1-D projection of the impact parameter distribution, and fit, is shown in Fig. 4.

The preliminary result for the correlated \(b\bar{b}\) production cross section, where each \(b\) quark decays to a \(\mu\), is

\[\sigma_{b\rightarrow\mu,\bar{b}\rightarrow\mu} = (1549 \pm 133) \text{ pb}\]

and for \(c\bar{c}\) production where each charm quark decays to a \(\mu\) is

\[\sigma_{c\rightarrow\mu,\bar{c}\rightarrow\mu} = (624 \pm 104) \text{ pb}\]
TABLE II: Measured correlated $b\bar{b}$ and $c\bar{c}$ cross sections in 0.74 fb$^{-1}$ at CDF, divided by NLO predictions, for different assumptions for the Peterson fragmentation parameter $\epsilon$.

<table>
<thead>
<tr>
<th></th>
<th>$b\bar{b}$ (0.006)</th>
<th>$b\bar{b}$ (0.002)</th>
<th>$c\bar{c}$ (0.006)</th>
<th>$c\bar{c}$ (0.002)</th>
</tr>
</thead>
<tbody>
<tr>
<td>measured</td>
<td>$\sigma_{b\bar{b}}$</td>
<td>$\sigma_{b\bar{b}}$</td>
<td>$\sigma_{c\bar{c}}$</td>
<td>$\sigma_{c\bar{c}}$</td>
</tr>
<tr>
<td>NLO</td>
<td>$\sigma_{b\bar{b}}$</td>
<td>$\sigma_{b\bar{b}}$</td>
<td>$\sigma_{c\bar{c}}$</td>
<td>$\sigma_{c\bar{c}}$</td>
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<tr>
<td>measured</td>
<td>$\sigma_{b\bar{b}}$</td>
<td>$\sigma_{b\bar{b}}$</td>
<td>$\sigma_{c\bar{c}}$</td>
<td>$\sigma_{c\bar{c}}$</td>
</tr>
<tr>
<td>NLO</td>
<td>$\sigma_{b\bar{b}}$</td>
<td>$\sigma_{b\bar{b}}$</td>
<td>$\sigma_{c\bar{c}}$</td>
<td>$\sigma_{c\bar{c}}$</td>
</tr>
</tbody>
</table>

Translating the $b \rightarrow \mu$, $\bar{b} \rightarrow \mu$ cross section to an inclusive $b\bar{b}$ cross section, independent of the final state:

$$\sigma_{bb}(p_T \geq 6 \text{ GeV}, |y| \leq 1) =$$

$$(1618 \pm 148 \pm [\sim 400 \text{ fragmentation}]) \text{ nb}$$

where the dominant error comes from the uncertainty in the fragmentation function, i.e. the fraction of $b$ quarks that decay to $\mu$. Different values for the Peterson fragmentation parameter $\epsilon$ result in significantly different results, especially for the correlated $c\bar{c}$ cross section. This can be seen in Table II where the ratio of the measured cross sections and the NLO predictions is given for two values of the Peterson fragmentation parameter $\epsilon$. The value traditionally used, $\epsilon = 0.006$, is obtained from fits to $e^+, e^- \rightarrow \eta$ data. However, [20] point out that these calculations, made on the basis of LO parton level cross sections evaluated with the leading-log approximation of the parton shower event generator, cannot be consistently used with exact NLO calculations, and find that a more suitable value of the Peterson fragmentation parameter would be $\epsilon = 0.002$. For $b\bar{b}$ the results for both values of $\epsilon$ are now in good agreement with NLO predictions, in contrast to previous measurements. For $c\bar{c}$

IV. CHARM PAIR CROSS SECTION

Different production mechanisms for heavy flavour quark-antiquark pairs lead to different kinematic distributions. The leading production mechanism, and their kinematic signature, are depicted in Fig 5. The lowest order diagrams ("Flavour Creation") favours back-to-back production of the quark-antiquark pair, while "Gluon Splitting" favours collinear production. Measuring the angular distribution in charm pair production therefore gives clues about the $c\bar{c}$ production mechanism.

Figure 6 shows the $D_0^\ast$, $D_0^\ast$ pair cross section and the $D^\ast$, $D^\ast$ pair cross section as a function of the angle between the two charm mesons, for the kinematic range $|y^D_0^\ast| < 1$, $p_T^D_0^\ast \in [5.5, 20]$ GeV and $|y^D_0^\ast| < 1$, $p_T^D_0^\ast \in [5.5, 20]$ GeV. Using a $D^\ast$ in the reconstruction ensures clean data samples. It can be seen that in both cases, collinear production is approximately of the same size as back-to-back production. The figure also shows the result from the Pythia event generator, configured to run with leading order matrix elements plus parton shower (Tune A). The total pair production cross sections agree well between the Pythia simulation and the data. However, the simulation over-estimates back-to-back production and under-estimates collinear production.
nism, illustrated in Fig 7(a), was first used to describe quarkonium production \cite{28,29}, but dramatically fails to describe the observed data at the Tevatron, underestimating the $J/\psi$ and $\psi'$ production in Run I by more than an order of magnitude (see e.g. \cite{30}). The discussion about a solution to this problem is dominated by two approaches:

- The “Colour Octet” mechanism proposed by \cite{31} is an effective field theory model based on Non-Relativistic QCD (NRQCD). It combines results from \cite{32} and elements of the Colour Evaporation model (CEM). The CEM was originally proposed in \cite{33,34}; more recent discussions can be found in \cite{35,36}. In the CEM the $J/\psi$ is essentially formed in a coloured state, so no “bleaching gluon” is needed. Then the colour “evaporates” in the emission of soft gluons.

- \cite{27} perform a calculation based on full, relativistic QCD, adding higher-order terms to the colour singlet term of the type shown in Fig 7(b), where the “bleaching gluon” is absorbed by a spectator parton. Although each additional term is small, there is such a large number of such diagrams that the sum of them does indeed make a large enough contribution to account for the order-of-magnitude difference between the observed cross sections and those predicted by the colour-singlet model.

Both NRQCD colour-octet and higher-order perturbative QCD describe the observed $J/\psi$ and $\psi'$ cross sections and $p_T$ spectra well (see for example \cite{27,31,37}), where the NRQCD approach has a number of adjustable hadronisation parameters that allow a certain level of tuning. However, NRQCD makes a firm prediction that the $J/\psi$ should be transversely polarised \cite{38} while perturbative QCD predicts a longitudinal polarisation of the $J/\psi$ \cite{27}.

The $k_T$ factorisation approach \cite{39,40,41,42,43,44,45,46,47}, can be combined with the colour singlet and the colour octet mechanism. In contrast to the usual collinear approach, $k_T$ factorisation takes the non-vanishing transverse momentum of the interacting gluons into account when calculating the hadronic matrix element. Especially when combined with the colour-octet mechanism, it describes the production cross sections well, although there is a number of adjustable parameters that are not yet fixed by experimental data \cite{45,46}. In contrast to the usual NRQCD colour-octet mechanism, $k_T$ factorisation predicts a longitudinal polarisation of the quarkonium that increases with $p_T$.

For a more complete and detailed review of the theory of quarkonium production at hadron colliders, including comparison to data, see \cite{48} and references therein.

V. CHARMONIUM, BOTTOMIUM

Producing a colour-neutral $J^P = 1^-$ state directly by gluon-gluon fusion is not possible. The simplest solution to the problem is to produce a colour-charged $c\bar{c}$ (or $b\bar{b}$) pair with gluon fusion and “bleach” it by radiating off a gluon. The “colour-singlet” mechanism,
A. The \(\psi(2S)\) Cross Section

The differential cross section of prompt \(\psi(2S)\) from \(\sim 1\) fb\(^{-1}\) at CDF Run II and, for comparison the differential cross section of prompt \(J/\psi\) from 39.7 pb at CDF Run II is shown in Fig 8. The preliminary result for the integrated cross section of \(p\bar{p} \rightarrow \psi(2S)X\) at \(\sqrt{s} = 1.96\) TeV, with a subsequent decay \(\psi(2S) \rightarrow \mu^+\mu^-\), in the kinematic range \(|y_{\psi(2S)}| < 0.6, p_T > 2\) GeV, is

\[
\begin{align*}
\sigma(p\bar{p} \rightarrow \psi(2S)X, |y_{\psi(2S)}| < 0.6, p_T > 2\text{ GeV}) \\
x \times Br(\psi(2S) \rightarrow \mu^+\mu^-) \\
= (2.60 \pm 0.05(\text{stat})^{+0.19}_{-0.18}(\text{syst})) \text{ nb}
\end{align*}
\]

B. Measurement of charmonium polarisation

The \(J/\psi\) can have three polarisation states, two transverse and one longitudinal. In \(J/\psi \rightarrow \mu^+\mu^-\) decays, the transverse and longitudinal polarisation states can be disentangled by measuring the angle \(\theta^*\) between the \(J/\psi\) and the \(\mu^+\) in the the \(J/\psi\) restframe. It is useful to define the parameter \(\alpha\) in terms of the cross section of transversely polarised \(J/\psi\), \(\sigma_T\), and the cross section of longitudinally polarised \(J/\psi\), \(\sigma_L\), as:

\[
\alpha \equiv \frac{\sigma_T - 2\sigma_L}{\sigma_T + 2\sigma_L}
\]

If the one longitudinal and the two transverse polarisation states are all equally populated, one would measure \(\alpha = 0\); for longitudinal polarisation, \(\alpha < 0\), and for transverse polarisation, \(\alpha > 0\). The parameter \(\alpha\) can be extracted from the distribution of events as a function of \(\cos\theta^*\):

\[
\frac{dN}{d(\cos\theta^*)} \propto 1 + \alpha \cos^2\theta^*
\]

In 0.8 fb\(^{-1}\), CDF finds 0.8M prompt \(J/\psi \rightarrow \mu^+\mu^-\) decays. The mass distribution of the \(\mu^+\mu^-\) pairs is shown in Fig 10. \(J/\psi\) originating from B decays are rejected using an impact parameter cut. The \(\cos\theta^*\) distributions are analysed in six bins of different \(p_T\), as shown in Fig 10. The result in terms of the parameter \(\alpha\) as a function of \(p_T\) is shown in Fig 10 (a). The plot shows \(p_T\)-dependent, longitudinal polarisation of the prompt \(J/\psi\). This contradicts NRQCD which predicts increasingly transverse polarisation with higher momenta; \(k_T\) factorisation on the other hand appears to over-estimate the longitudinal polarisation. The
FIG. 11: The polarisation parameter $\alpha$ vs momentum for prompt $J/\psi$ (a) and prompt $\psi(2S)$ (b) at CDF. Negative $\alpha$ correspond to longitudinal polarisation, positive $\alpha$ to transverse polarisation. The NRQCD calculation $^{38, 50, 51}$ is superimposed in turquoise, the $k_T$ factorisation calculation $^{46}$ in magenta.

FIG. 12: The polarisation parameter $\alpha$ vs momentum for prompt $J/\psi$ at CDF Run I and Run II, where the Run II data were re-analysed to match the binning in Run I.

The polarisation has also been measured in the theoretically cleaner (no feed down from higher states) $\psi'$ channel; the result is shown in Fig 11(b). The event numbers are much lower than for the $J/\psi$ but the results indicate a trend for the longitudinal polarisation fraction to increase with higher $p_T$, inconsistent with the NRQCD calculation. These results have been published in $^{49}$.

FIG. 13: The $\mu^+\mu^-$ mass distribution at DØ after selection cuts, on the left for all events, on the right only for those events selected by the di-muon trigger.

VI. MEASUREMENT OF THE $\Upsilon$ POLARISATION

The theoretical methods describing the production and polarisation of $\Upsilon(1S)$ and $\Upsilon(2S)$ in $p-\bar{p}$ collisions are equivalent to those for charmonium presented in the previous section.

The DØ experiment measures the polarisation of the $\Upsilon(1S)$ and the $\Upsilon(2S)$ as a function of the $\Upsilon$ transverse momentum in 1.3 fb$^{-1}$ of data in Tevatron Run II. Data are selected in the di-muon channel. To achieve a more reliable estimate of the trigger efficiency in the cross section calculation, only events selected by the di-muon trigger are used in the analysis. The invariant mass of the $\mu^+\mu^-$ pairs near the $\Upsilon$ mass is shown in Fig 14 for all data passing the selection cut (a), and those that also were selected by the di-muon trigger (b). Figure 14 shows the di-muon mass spectrum for one bin in $|\cos \theta^*|$ and $p_T$, with the fit to the data superimposed, and the Gaussians describing the individual contribution of the $\Upsilon(1S)$, $\Upsilon(2S)$.
and $\Upsilon(3S)$. Because of the relatively small number of $\Upsilon(3S)$ events, to ensure a stable fit, the width, relative position and the fraction of events in the $\Upsilon(3S)$ peak were taken from MC simulation and fixed in the fit. In another fit, the position of the $\Upsilon(3S)$ peak is allowed to float. The difference between the two approaches is taken as a systematic error. The $\Upsilon(1S)$ polarisation vs $p_T$ is shown in Fig.15. The data represent the result for an admixture of directly produced $\Upsilon(1S)$ and $\Upsilon(1S)$ from other sources, in particular from $\Upsilon(2S)$, $\Upsilon(3S)$ and $\chi_b(2P)$. For comparison, the CDF Run I result [53] is shown as the green triangles with error bars.

**VII. RELATIVE PRODUCTION CROSS SECTION $\chi_{c2}$ AND $\chi_{c1}$**

The measurement of the $\chi_{cJ}$ ($J = 1, 2$) cross section is an interesting measurement in its own right, as well as important input to $J/\psi$ production measurements to which it provides an important source of feed-down. Experimentally, this measurement has always been plagued by the poor mass resolution of the reconstructed $\chi_{cJ}$ which is due to the soft photon in the decay chain $\chi_{cJ} \rightarrow J/\psi\gamma$. The high luminosity at the Tevatron now provides sufficient statistics to restrict the analysis only to events where the photon undergoes conversions to an $e^+e^-$ pair. The momenta of the electrons can be measured precisely resulting in a far superior energy resolution of the photon than would be possible by measuring the photon’s energy in the calorimeter. This leads to an excellent resolution of the reconstructed $\chi_{cJ}$ mass, allowing a separation of the $J = 0, 1, 2$ states. CDF reconstructs $\chi_{cJ}$ in in the channel $\chi_{cJ} \rightarrow J/\psi(\mu\mu)\gamma(e^+e^-)$ within the kinematic range $p_{T(J/\psi)} \in [4, 20] \text{ GeV}$. To select photon
conversions, the $e^+e^-$ pair is a required to form a well-reconstructed vertex a large distance (> 12 cm) from the beam, well inside the instrumented region of CDF. The spectrum of the ($\mu^+\mu^-$) invariant mass is shown in Fig 17 showing two well-separated peaks at the $\chi_{c1}$ mass and the $\chi_{c2}$ mass. No significant evidence for $\chi_{c0}$ production can be seen. After all selection cut total number of $\chi_{c1} \rightarrow J/\psi(\mu\mu) \gamma(e^+e^-)$ events found in 1.1 $fb^{-1}$ is $\sim 7k$.

The distance between the beamline and the $\mu^+\mu^-$ vertex is used to separate the prompt contribution from $\chi_{c1}$ originating from B decays. The measured ratio for prompt production is:

$$\frac{\sigma(\chi_{c2})}{\sigma(\chi_{c1})} = 0.70 \pm 0.04({\text{stat}}) \pm 0.03({\text{sys}}) \pm 0.06({\text{BF}})$$

for $p_T(J/\psi) \in [4, 20]$ GeV. The last contribution to the error is due to the uncertainty in the branching fraction of $\chi_{cJ} \rightarrow J/\psi \gamma$. This result is at odds with the expectation from the colour octet model, which, by counting spin-states, predicts a ratio of 5/3.

VIII. CONCLUSIONS

The large charm and bottom production cross section at $\sqrt{s} = 1.96$ TeV proton-antiproton collisions, combined with the capabilities of the DØ and the CDF detector provide the opportunity for new measurements with unprecedented statistics and precision.

In this paper we presented Run II results on single and correlated open charm and bottom production as well as quarkonium production and polarisation. The same theoretical framework that managed to describe successfully the unexpectedly large charmonium production cross observed at Tevatron Run I now fails to account for the significant longitudinal polarisation of charmonium and bottomium observed in Run II, nor the $\chi_{c2}$ and $\chi_{c1}$ cross section ratio. Several alternative models are being developed, but at the time of writing this paper, none has provided a detailed quantitative post-diction of the charmonium polarisation vs $p_T$ that matches the observed data.

So heavy flavour production is, refreshingly, a field that is clearly led by experiment. In response to the data from the Tevatron, some of which have been presented here, we can look forward to new calculations and models, offering new descriptions of the mechanism of heavy flavour, and in particular quarkonium, production. We will also see further new measurements making use of the large amount of data yet to be analysed. Only about 1/3 of the Tevatron data taken so far have been used for the results presented here, and with the machine going stronger than ever, we can expect to at least double the integrated Tevatron luminosity before the end of Run II.


[26] R. D. Field (CDF) (2002), rick Field’s webpage with up-to-date information on Pythia Tune Set A.


Observations by the PHENIX and STAR collaborations suggest that a strongly coupled quark-gluon plasma is produced in heavy-ion collisions at the Relativistic Heavy Ion Collider (RHIC). After a brief introduction to heavy-ion physics, measurements of heavy-quark production in heavy-ion collisions and the modification of heavy-quark spectra by the QGP are presented. Measurements of the total charm cross-section in several different collision systems confirm that $c\bar{c}$ pairs are produced through parton hard-scattering in the initial stages of the collisions. Non-photonic $e^\pm$ (proxies for heavy quarks) are suppressed by a factor of $\sim 5$ in central $Au + Au$ collisions relative to $p+p$ collisions. This is larger than most current theoretical predictions and has led to a re-examination of heavy-quark energy loss in the medium. The relative contributions of $c$ and $b$ decays to the non-photonic $e^\pm$ spectrum have been predicted by perturbative QCD calculations; STAR measurements of azimuthal correlation functions of non-photonic $e^\pm$ and hadrons agree with these predictions.

### 1. Introduction

The main purpose of the PHENIX and STAR experiments at the Relativistic Heavy Ion Collider is to study the properties of strongly interacting matter at high temperatures. Figure 1 shows a schematic theoretical phase diagram of nuclear matter for various temperatures and baryon chemical potentials. At normal baryon chemical potentials and low temperatures, strongly interacting matter exists as nuclei or a gas of interacting hadrons. At higher temperatures, however, the quarks may become deconfined due to asymptotic freedom; the degrees of freedom of the system are not hadrons, but individual quarks and gluons. This state of matter, called the quark-gluon plasma (QGP), is thought to have existed in the first few microseconds following the Big Bang, and may also have been produced in high-energy nucleus-nucleus collisions such as those at RHIC at Brookhaven National Laboratory and the SPS at CERN.

Lattice QCD calculations predict a sudden increase in the number of degrees of freedom of strongly interacting matter near a critical temperature $T_c \sim 170$ MeV. Whether or not there is a true phase transition and its exact nature has yet to be determined. It has been proposed that a critical point may exist in the phase diagram, and that for higher baryon densities, a first-order phase transition may be observed. A search for this critical point may be conducted in the next few years at RHIC by reducing the collision CM energy to lower values ($\sim 5$GeV) than the collider’s present operating range (22 GeV to 200 GeV).

The collision of two nuclei introduces a large amount of energy into a region of space approximately the size of a nucleus for a short period of time. Lattice calculations predict a critical energy density of $\approx 700$ MeV/fm$^3$ needed for QGP formation. The highest-energy RHIC collisions reach an energy density of at least 4.9 GeV/fm$^3$, seven times the critical energy density. The QCD vacuum is "melted" and a quark-gluon plasma is produced. It is believed that the QGP quickly reaches thermal equilibrium, expands, and cools for a few fm/$c$ until the transition temperature is reached. At this point, the partons become confined into hadrons. The hadron gas expands and the hadrons scatter inelastically until chemical freeze-out. The hadron gas continues to expand for a few more fm/$c$ with elastic hadron-hadron interactions until thermal freeze-out, after which hadronic interactions are negligible. The resulting shower of particles can be detected in detector systems at RHIC.

Many phenomena in heavy-ion physics depend upon the degree of overlap between the two colliding nuclei, called the centrality of the collision. If the distance between the centers of the nuclei (impact parameter) is small, the overlap between the nuclei is large. Such collisions are called "central" events. A peripheral event has a small overlap and the impact parameter approaches the sum of the radii of the two nuclei. $N_{\text{part}}$ is the number of nucleons that partici-
part in collisions and \(N_{\text{binary}}\) is the number of binary collisions between those participating nucleons. \(N_{\text{part}}\) and \(N_{\text{binary}}\) are large for central events and small for peripheral events. The centrality of an event is estimated using the multiplicity of charged tracks at mid-rapidity. Central collisions are characterized by a higher charged-particle multiplicity than peripheral events. \(N_{\text{part}}\) and \(N_{\text{binary}}\) are estimated using the Glauber model of nucleus-nucleus interactions [5, 6].

The ratios of various particle yields have been fit with statistical thermal models, which indicate that at the time of chemical freeze-out, the system is thermalized and has a temperature \(T = 170 - 180\) MeV for RHIC collision energies [7]. A non-central collision will have a spatially asymmetric overlap region, roughly elliptically shaped. The spatial asymmetry translates into a momentum-space asymmetry in the spectra of final state particles, with hadrons emitted preferentially near the reaction plane (the plane containing the beam axis and the impact parameter vector). Measurements of this phenomenon (called elliptic flow) are well described by ideal hydrodynamical models, indicating that the state of matter produced in the collision has a very low viscosity and is a nearly perfect liquid [8, 9, 10]. All RHIC experiments have observed the production of back-to-back di-jets in \(p+p\), \(d+Au\), and peripheral \(Au+Au\) collisions. In central \(Au+Au\) collisions, however, di-jets are not observed: the away-side jets appear to have been quenched by the medium [11, 12].

\[
R_{AA} \equiv \frac{d^2N(A+A)}{dy \, dp_T} \frac{\langle N_{\text{binary}} \rangle}{d^2N(p+p)} \langle dN_{\text{binary}}/dy \rangle \]

If no medium is produced, then a nucleus-nucleus collision can be viewed as an incoherent superposition of nucleon-nucleon collisions and \(R_{AA}\) will be unity. Deviations from unity indicate the effects of nuclear matter and the quark-gluon plasma on particle yields. Central collisions produce a larger medium than peripheral collisions, which should result in a greater suppression of high-\(p_T\) particles.

Figure 2 shows measurements of \(R_{AA}\) for direct photons, \(\pi^0\), and \(\eta\) by PHENIX. \(R_{AA}\) is unity for direct photons [13], indicating that there is no suppression of direct photons in nucleus-nucleus collisions. This is expected since photons do not interact strongly and should not be affected by the presence of a QGP. However, the hadrons [16, 17, 18] are suppressed by about a factor of 5 at high \(p_T\). Also shown is a theoretical prediction of light-flavor-hadron suppression from the GLV model of parton-QGP interactions [19, 20]. \(dN_g/df\) is the gluon density, a parameter in the GLV model related to the opacity of the medium. The observed light-flavor-hadron suppression was used to constrain the value of this parameter, giving \(dN_g/df \approx 1100\). \(R_{AA}\) has also been measured for heavy-flavor decay products and compared to models; this will be discussed in subsequent sections. These and other measurements indicate that the matter produced in RHIC collisions is a strongly coupled quark-gluon plasma (sQGP) [21].

2. Experiments

The experiments described in these proceedings are conducted at the Relativistic Heavy Ion Collider (RHIC) at Brookhaven National Laboratory on Long Island, NY [22]. Ions are accelerated from the tandem Van de Graaff accelerators or the proton sources, through the AGS Booster and the AGS, and are injected into RHIC. Electrons are stripped off at several locations along the way. RHIC consists of two
synchrotron rings, 3.8 km in circumference. The two counter-circulating beams of ions intersect in each of the six interaction regions. The PHENIX and STAR detectors sit at two of these interaction regions. At RHIC, physicists can create and maintain beams of ions ranging from protons (both polarized and unpolarized) to the heaviest nuclei. RHIC can be used to collide protons with a center-of-mass collision energy up to $\sqrt{s} = 500$ GeV and a luminosity of $\sim 10^{32}$ cm$^{-2}$s$^{-1}$. RHIC can also collide ions ranging in mass from deuterons to gold nuclei with CM energy per nucleon pair $\sqrt{s_{NN}} \leq 200$ GeV and a luminosity of $\sim 10^{26}$ cm$^{-2}$s$^{-1}$. 

Figure 3: The STAR detector. Measurements described below use the Time Projection Chamber (TPC), Time-of-Flight detector (TOF), and the electromagnetic calorimeter (EMC).

Figure 4 shows a diagram of the STAR (Solenoidal Tracker At RHIC) detector. The barrel and forward time-projection chambers (TPC and FTPC) record particle trajectories inside a room-temperature solenoidal 0.5-T magnet. The TPC covers pseudorapidity $|\eta| < 1.8$, while the FTPCs cover $2.5 < |\eta| < 4$. After passing through the TPC, particles enter the barrel electromagnetic calorimeter (BEMC) or the endcap electromagnetic calorimeter (EEMC, not shown). The STAR calorimeters together cover pseudorapidity $-1 < \eta < 2$. The Shower Maximum Detectors (SMDs), located approximately 5 radiation lengths inside each EMC tower module, provide additional particle identification based on the shape of the electromagnetic shower produced in the calorimeters. The SMDs allow shower shapes to be measured to high precision ($\Delta \eta = 0.007$, $\Delta \phi = 0.007$ rad). A silicon vertex tracker (SVT) covers $|\eta| < 1$ between the beampipe and the TPC, providing accurate particle tracking near the collision vertex. A prototype Time-of-Flight detector (TOF) has been installed outside the TPC covering $-1 < \eta < 0$ and $\Delta \phi = 0.04$ rad. The TOF provides a precise measurement of particle velocity. Plans call for the TOF to be extended to full azimuthal coverage over pseudorapidity $|\eta| < 1$. A full description of the STAR detector is given in [23].

Figure 4 shows a diagram of the PHENIX (Pioneering High-Energy Nuclear Interaction Experiment) detector. The two central spectrometer arms sit in an axial magnetic field and cover $|\eta| < 0.35$ and $\pi/2$ each in azimuth. Particle tracking is provided by three sets of pad chambers and (in the east arm) the time expansion chamber. The Ring-Imaging Cherenkov detectors and the Time-of-Flight detector provide particle identification. Beyond these detector systems sit lead-glass and lead-scintillator calorimeters. The two forward muon spectrometer arms sit in radial magnetic fields. They consist of drift chambers for precision tracking and muon identifiers. The muon identifiers are made up of alternating layers of steel absorber plates and streamer-tube tracking layers. A full description of the PHENIX detector, including the azimuthal and pseudorapidity coverage of each detector subsystem, is given in [24].
3. Heavy Flavors

Due to the high luminosity at RHIC, even particles with comparatively low production cross-sections, such as $c$ and $b$ quarks, can be used to probe the strongly interacting matter produced. In nucleus-nucleus collisions, the dominant production mechanism for heavy quarks is gluon-gluon fusion in the initial hard scattering of nucleons. Thermal production of heavy quark pairs in the medium is believed to be negligible after the initial stages of the collision and the number of heavy quarks is essentially "frozen." Therefore, heavy quarks are probes sensitive to all stages in the evolution of the QGP, from its initial formation to hadronization and freeze-out.

The presence of a QGP is expected to lead to the dissociation and suppression of heavy quarkonia through Debye screening of color charges. Some theoretical calculations predict sequential dissociation of heavy quarkonia, with the more weakly bound resonances dissociating at lower temperatures. If this is true, measurements of heavy quarkonium suppression could lead to a measurement of the initial temperature of the medium. PHENIX measurements of $R_{AA}$ for $J/\psi$ show less suppression than was expected based on SPS data. This may be due to the regeneration of quarkonia after the initial stages of the collision.

4. Measurements of Open Heavy Flavors

The RHIC experiments have studied several different heavy-flavor decay channels, including the hadronic decays of $D^0$ mesons to pions and kaons, and the semileptonic decays of heavy-flavor hadrons to muons and electrons.

4.1. Direct Reconstruction of $D^0$ Decays

The STAR collaboration has found the yields of $D^0$ and $D^0$ mesons by reconstructing the $D^0(\bar{D}^0) \rightarrow K^+ + \pi^-$ decays, which have a branching ratio of 3.83%. The pions and kaons are identified by their energy loss in the Time Projection Chamber. STAR cannot reconstruct the full decay topology since $c\tau(D^0) = 124 \mu m$. The TPC does not have sufficient track projection resolution to distinguish $D^0$ decay products from tracks coming directly from the primary collision vertex. The $D^0$ invariant mass spectrum was obtained by pairing each kaon with oppositely charged pions from the same event. The combinatorial background was estimated through event mixing techniques and subtracted. Figure 5 shows the $K\pi$ invariant mass spectrum for $m < 1$ in Au + Au collisions at $\sqrt{s_{NN}} = 200$ GeV. A clear $D^0$ peak is visible. Because of the small branching ratio for this decay and the lack of a dedicated trigger, this analysis is limited by statistics to $p_T < 3$ GeV/c.

4.2. Decays to Muons

The STAR collaboration has identified low-$p_T$ muons (0.17 GeV/c < $p_T$ < 0.25 GeV/c) through measurements of energy loss in the Time Projection Chamber and $m^2 = (p/(3\gamma))^2$ in the Time-of-Flight detector and the TPC. Figure 5 shows the $m^2$ distribution of particles after energy-loss selection. The muon and pion peaks are clearly visible. In addition to muons from heavy flavor decays, the muon peak contains a large number of muons from pion and kaon decays; it was necessary to remove these "background" muons. HIJING was used to simulate the DCA (distance of closest approach) of muon tracks with the primary collision vertex (see Figure 5b). The DCA distribution for muons from charm decays has a maximum much closer to zero than the distribution for muons from pion and kaon decays. The observed DCA distribution was fit with a linear combination of these two simulated DCA distributions to obtain the contribution of charm-decay muons to the total muon yield.

4.3. Decays to Single Electrons

Both the PHENIX and STAR collaborations have also studied the spectrum of single electrons, i.e. $e^\pm$.
Figure 5: STAR heavy-flavor identification methods. (a) $D^0$ meson reconstruction in 200-GeV Au + Au collisions. The $K\pi$ invariant mass distribution is shown after the subtraction of the combinatorial background (estimated using event-mixing). The residual background is approximated by a linear function and subtracted. (b) $m^2$ distribution measured by TOF detector in 200-GeV Au + Au collisions. The particles have passed the energy-loss cut used to identify low-$p_T$ muons. (c) DCA distributions used to identify charm-decay contribution to the inclusive muon spectrum (see explanation in text).

produced with (anti)neutrinos in weak decays. The spectrum of single $e^\pm$ is expected to be dominated by heavy-flavor semileptonic decays (e.g. $D^0 \rightarrow e^+ + K^- + \nu_e$). The main sources of background to the single $e^\pm$ signal are $e^+e^-$ pairs from $\pi^0$ and $\eta$ Dalitz decays and photon conversions [37]. For this reason, single $e^\pm$ are called non-photonic $e^\pm$, while the background $e^\pm$ are called photonic $e^\pm$. The decays of vector mesons (e.g. $\rho$, $\phi$, $\phi$, and $\omega$) make small contributions to the photonic $e^\pm$ background. Background $e^\pm$ from photon conversion is less significant at PHENIX due to the reduced amount of material relative to STAR. $K_{e3}$ decays ($K^\pm \rightarrow \pi^0 + e^\pm + \nu_e(\bar{\nu}_e)$ and $K_L^0 \rightarrow \pi^+ + e^\pm + \nu_e(\bar{\nu}_e)$) make small contributions to the single $e^\pm$ signal.

The STAR collaboration identifies $e^\pm$ using two different methods. For $p_T < 3.5$ GeV/c, $e^\pm$ are identified through measurements of energy loss and $m^2$ with the Time Projection Chamber and Time-of-Flight detector (similar to the method of muon identification described above) [38]. For $p_T > 1.5$ GeV/c, $e^\pm$ are identified through measurements of TPC energy loss, the energy $E$ deposited in the electromagnetic calorimeter, and the shape of the electromagnetic shower measured in the shower maximum detector [41]. The use of a trigger allows this measurement to extend up to $p_T \approx 8$ GeV/c. An energy-loss cut of approximately 3.5 keV/cm $< dE/dx < 5$ keV/cm is used to identify $e^\pm$ (the exact bounds are varied slightly depending on track momentum and event charged-track multiplicity). Compared to hadrons, $e^\pm$ produce larger showers and deposit more of their energy in the EMC. The ratio $p/E$ has a maximum near 1 for $e^\pm$ (see Figure 6 (12)) and additional $e^\pm$ identification is provided by cuts on the shower size in the SMD. The remaining hadron contamination is $\approx 2\%$ at $p_T \sim 2$ GeV/c and $\approx 20\%$ at $p_T \sim 8$ GeV/c. STAR identifies photonic (background) $e^\pm$ through invariant-mass reconstruction of $e^+e^-$ pairs. The invariant-mass distribution of $e^+e^-$ pairs from photon conversions and $\pi^0$ and $\eta$ Dalitz decays has a maximum near 0 mass. When a cut of $M_{inv} (e^+e^-) < 150$ MeV/c$^2$ is used, photonic electrons are identified with an efficiency around 70%.

The PHENIX collaboration identifies $e^\pm$ using the Ring-Imaging Cherenkov detector, measurements of the shower shape in the electromagnetic calorimeter, and a cut on the energy-to-momentum ratio [43, 44]. Photonic $e^\pm$ are identified using two different methods. In the ”cocktail subtraction” method, the spectra of $e^\pm$ from various sources of background are simulated. Measured yields of $\pi^0$, $\eta$, direct photons, and other sources of background are used as input for the simulation generator. In the ”converter subtraction” method, a photon converter (a thin brass sheet of 1.67% $X_0$) is inserted around the beam pipe. $\Delta N_e$, the increase in the $e^\pm$ yield due to the converter, is
measured. GEANT simulations are used to determine $R_\gamma$, the fractional increase in the photonic $e^\pm$ yield caused by the converter ($R_\gamma \approx 2.3$ in the 2006 $p+p$ run). Knowledge of $\Delta N_e$ and $R_\gamma$ allows the photonic $e^\pm$ yield to be determined and removed from the inclusive $e^\pm$ yield. Where $N_{\gamma e}^{NC}$ is the photonic $e^\pm$ yield with no converter present,

$$N_{\gamma e}^{NC} = \frac{\Delta N_e}{R_\gamma - 1}. \quad (3)$$

The non-photonic $e^\pm$ yields measured using the cocktail and converter subtraction methods are consistent with each other.

5. Total Charm Cross-Section

STAR determines the total charm cross-section, $\sigma_{cc}$, for each collision system through a combined fit of the $D^0$, muon, and non-photonic $e^\pm$ measurements described above (e.g. Figure 7a for 200-GeV Au + Au collisions) [39]. PHENIX determines the total charm cross-section from the measurement of the non-photonic $e^\pm$ yield (e.g. Figure 7a for 200-GeV $p+p$ collisions) [43]. Figure 8a shows measurements of the scaled total charm cross-section by STAR and PHENIX for $p+p$ [43], $d+Au$ [36], and $Au+Au$ [41, 44] collisions (in different centrality bins) at $\sqrt{s_{NN}} = 200$ GeV. A preliminary STAR measurement for $Cu+Cu$ using only $D^0$ reconstruction is also shown [43]. The cross-section is divided by $\langle N_{\text{binary}} \rangle$, the average number of binary collisions for the given collision system. Within each experiment, the charm cross-section scales with the number of binary collisions. This is a confirmation that charm is indeed produced through initial parton hard-scattering and that charm production through other mechanisms (such as thermal production in the QGP) is not significant.

Note that the cross-sections measured by STAR are higher than those measured by PHENIX by a factor of $\sim 2$. This is still under investigation.

Due to the large quark masses, the hard-scattering processes that produce heavy flavor can be calculated using perturbative QCD [46]. The most advanced perturbative calculation scheme is the Fixed-Order plus Next-to-Leading-Log-resummed approximation, or FONLL. A FONLL prediction for the charm cross-section [47, 48] is shown as prediction (1) in Figure 8a. The PHENIX data are consistent with this prediction; the STAR data are greater than the FONLL prediction by a factor of $\sim 5$ and sit well above the upper uncertainty bound. However, the charm cross-section predictions are sensitive to the number of active flavors, the choice of scale, and the parton densities in the collision system. A new calculation by R. Vogt [48] (prediction (II) in Figure 8b) indicates that the uncertainties on the perturbative calculation may be larger than previously thought. The total charm cross-section measured by STAR is consistent with the new perturbative calculation.

Figure 8b shows the ratios of the measured non-photonic $e^\pm$ yields [36, 41, 43] to the FONLL prediction as functions of $p_T$ for $p+p$ collisions at $\sqrt{s} = 200$ GeV. The dashed horizontal lines indicate the ratio of the measured cross-sections to the FONLL cross-section. Figure 8b indicates that STAR’s disagreement with PHENIX (by a factor of $\sim 2$) and FONLL (by a factor of $\sim 5$) exists even at high $p_T$. However, the FONLL prediction does describe the shape of the STAR and PHENIX non-photonic $e^\pm$ spectra well.

6. Medium Modification of Non-photonic $e^\pm$ Spectra

The cross-section discrepancy between PHENIX and STAR cancels in the nuclear modification factor $R_{AA}$, which is the scaled ratio of particle yields (see definition above). Figure 9a shows $R_{AA}$ of non-photonic $e^\pm$ for Au + Au collisions at $\sqrt{s_{NN}} = 200$ GeV for both PHENIX [44] and STAR [41, 49]. $R_{AA}$ is plotted as a function of $\langle N_{\text{part}} \rangle$, the average number of nucleons participating in collisions in a given centrality bin. The most central collisions have the highest values of $\langle N_{\text{part}} \rangle$. The PHENIX and STAR $R_{AA}$ are consistent with each other across the range of $\langle N_{\text{part}} \rangle$ shown. Plotted in shaded bands are the values of $R_{AA}$ measured by PHENIX for $\pi^0$ [16] and by
Figure 7: Measurements used in calculation of total charm cross-section. (a) STAR $D^0$, muon, and $e^\pm$ spectra for 200-GeV Au + Au collisions in various centrality bins. (b) PHENIX non-photonic $e^\pm$ spectrum for 200-GeV $p + p$ collisions. Also shown are pQCD predictions (FONLL) for the $c^-$ and $b$-decay contributions to the spectrum (see discussion below and Figure 10).

Figure 8: Comparisons of RHIC measurements to FONLL predictions. (a) Total charm cross-section (scaled by $1/\langle N_{\text{binary}} \rangle$) in different collision systems at $\sqrt{s_{NN}} = 200$ GeV. Theoretical predictions are also shown. (b) Ratios of non-photonic $e^\pm$ yields in central 200-GeV Au + Au collisions for STAR and PHENIX to FONLL prediction in central 200-GeV Au + Au collisions. STAR $e^\pm$ have been selected using measurements of energy-loss ($dE/dx$) in the TPC or $m_2$ in the TOF. STAR for charged light-flavor hadrons [50]. The measured suppression of non-photonic $e^\pm$ is similar to the suppression observed for light-flavor hadrons. This was unexpected, as heavy quarks were expected to lose less energy in the medium than light quarks. As a result, non-photonic $e^\pm$ (proxies for heavy quarks) would be suppressed less than light-flavor hadrons.

Figure 9b shows a comparison of the non-photonic $e^\pm$ suppression measured by STAR and PHENIX to several theoretical models of heavy-quark interactions with a quark-gluon plasma. The data are from central Au + Au collisions at $\sqrt{s_{NN}} = 200$ GeV. In $d + Au$ collisions at this energy a ~ 50% enhancement of non-photonic $e^\pm$ is observed (data [41] not
Figure 9: Nuclear modification factor ($R_{AA}$) for non-photonic $e^\pm$ in Au + Au collisions at 200 GeV. (a) $R_{AA}$ vs. $\langle N_{\text{part}} \rangle$. Also shown are light-flavor hadron $R_{AA}$ data. (b) $R_{AA}$ vs. $p_T$ for central collisions. The data are compared to theoretical predictions.

The enhancement is explained by the Cronin effect (multiple scattering in normal nuclear matter). Therefore, the strong non-photonic-$e^\pm$ suppression observed in central Au + Au collisions does not appear to be due to such cold nuclear matter effects. The calculations that produce curves (I) and (II) include only energy loss through gluon radiation; these calculations under-predict the suppression of non-photonic $e^\pm$, especially at high $p_T$. The model of Adil and Vitev [55] (curve (VI)) uses the fragmentation of heavy quarks and the subsequent dissociation of $D$- and $B$-meson states in the medium to explain the suppression pattern of non-photonic $e^\pm$. This model does describe the observed suppression well.

In addition to energy loss through gluon radiation, partons can lose energy through collisions with other partons in the medium. Calculations [13] indicate that, for heavy quarks, collisional energy loss is as important as radiative energy loss. Curve (III) is generated using the DGLV model (which produced curve (I)) including both radiative and collisional energy loss [53]. Curve (IV) is generated using the model of van Hees et al. [54], which includes heavy-quark energy loss through elastic collisions in the medium and the formation of resonant $D$- and $B$-meson states through quark coalescence. Curves (III) and (IV) seem to describe the data better than curves (I) and (II), which include only radiative energy loss, but still tend to under-predict the observed suppression of non-photonic $e^\pm$, especially at high $p_T$. Curve (V) (which includes both radiative and collisional energy loss) is consistent with the measured suppression. It is therefore important to disentangle the $c$ and $b$ contributions and determine their relative strengths; this will be discussed in the next section.

7. Electron-Hadron Azimuthal Correlations

Figure 10 shows FONLL calculations of the $D$- and $D$-meson-decay contributions to the non-photonic $e^\pm$ spectrum. Curve (V) is generated with the assumption that only $c$-quark decays contribute to the non-photonic $e^\pm$ spectrum and that the $b$-quark contribution is insignificant for the $p_T$ range shown [53]. Curve (V) (which includes both radiative and collisional energy loss) is consistent with the measured suppression. It is therefore important to disentangle the $c$ and $b$ contributions and determine their relative strengths; this will be discussed in the next section.
tion scales. The charm contribution is dominant at low $p_T$. The bottom contribution becomes larger than the charm contribution around $p_T \sim 5 \text{ GeV}/c$.

Figure 10: FONLL predictions for the contributions of $D$- and $B$-meson decays to the mid-rapidity non-photonic $e^\pm$ cross-section at $\sqrt{s} = 200 \text{ GeV}$.

To disentangle the charm and bottom contributions, the STAR collaboration has studied azimuthal correlations between non-photonic $e^\pm$ and hadrons. Figure 11 shows the difference in azimuthal angle ($\Delta \phi$) between non-photonic $e^\pm$ and hadrons in 200-GeV proton-proton collisions for one $p_T$ bin [56]. Due to decay kinematics, there will be an azimuthal correlation between the leptons and hadrons produced in heavy-flavor semi-leptonic decays. Due to the larger $B$-meson mass, a $B$-meson can give more kinetic energy to its decay products than a $D$-meson, resulting in a broader near-side ($\Delta \phi \sim 0$) correlation peak. The expected charm and bottom contributions to this distribution are simulated with PYTHIA. Varying the quark fragmentation functions inside PYTHIA does not significantly alter the shapes of the simulated $\Delta \phi$ distributions. The measured distribution is then fit with a linear combination of the simulated charm and bottom distributions to find their relative strengths. $B/(B+D)$, the fraction of the total non-photonic $e^\pm$ cross-section due to $B$-meson decays, is plotted in Figure 12. The measurement of $B/(B+D)$ through $e^-h$ correlations [56] is consistent with FONLL predictions. This is an indication that $b$-quark suppression should be taken into account in describing the suppression of non-photonic $e^\pm$ at moderate $p_T$ (see the discussion in the previous section).

8. Conclusions

Measurements of jet quenching and elliptic flow suggest that the state of matter created in a heavy-ion collision at RHIC energies is a strongly coupled quark-gluon plasma (sQGP) [21] and is a near-perfect liquid [10]. PHENIX and STAR have calculated the total charm cross-section $\sigma_{c\bar{c}}$ for $p+p$, $d+Au$, $Cu+Cu$, and $Au+Au$ collisions at $\sqrt{s_{NN}} = 200 \text{ GeV}$ [36, 41, 43, 44, 45]. The cross-section scales with the number of binary nucleon-nucleon collisions, an indication that charm quarks are indeed produced through initial parton hard-scattering. PHENIX and STAR have measured the nuclear modification factor $R_{AA}$ for light-flavor hadrons [18] and for non-photonic $e^\pm$, which come predominantly from heavy-flavor semileptonic decays [41, 43, 44]. The suppression of non-photonic $e^\pm$ in central Au + Au collisions is greater than expected. It has been difficult for theoretical models [51, 52, 53, 54, 55] to describe the suppression of light-flavor hadrons and non-photonic $e^\pm$ simultaneously. Perturbative QCD calculations (FONLL) [47] predict that the $b$-decay contribution to the non-photonic $e^\pm$ becomes comparable to the
c-decay contribution at $p_T \sim 5$ GeV/c. STAR measures the difference in azimuth angle $\phi$ for pairs of non-photonic $e^\pm$ and hadrons [56]. This distribution is fit with PYTHIA simulations of the expected $c$- and $b$-decay contributions to determine their relative strengths; these data are consistent with the FONLL predictions.

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References

**$D^0 - D^0$ Mixing at BABAR**

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The BABAR and Belle collaborations have recently found evidence for mixing within the $D$ meson system. We present some of the mixing search techniques used by BABAR and their status as of the beginning of the summer 2007. These have culminated in a measurement in the $K\pi$ decay final state of the $D$ that is inconsistent with the no-mixing hypothesis with a significance of 3.9 standard deviations.

1. Introduction

Mixing among the lightest neutral mesons of each flavor has traditionally provided important information on the electroweak interactions, the CKM matrix, and the possible virtual constituents that can lead to mixing. Among the long-lived mesons, the $D$ meson system exhibits the smallest mixing phenomena. The B-factories have now accumulated sufficient luminosity to observe mixing in the $D$ system and we can expect to see more detailed results as more luminosity is accumulated and additional channels sensitive to mixing are analyzed. The B-factories produce about 1.3 million Charm events per fb$^{-1}$ of integrated luminosity accumulated. The BABAR integrated luminosity of about 384 fb$^{-1}$ used for the evidence for mixing result we will present corresponds to about 500 million charm events produced. The present BABAR integrated luminosity is approximately 500 fb$^{-1}$. BABAR is a high acceptance general-purpose detector providing excellent tracking, vertexing, particle ID, and neutrals detection. All of these capabilities are crucial for making the difficult mixing measurement.

2. Mixing Measureables for the $D$ System

The propagation eigenstates, including the electroweak interactions for the $D$ mesons are given by:

$$|D_{1,2}\rangle = p|D^0\rangle \pm q|\bar{D}^0\rangle , \quad |p|^2 + |q|^2 = 1. \quad (1)$$

Propagation parameters that determine the time-evolution for the two states are given by:

$$\Gamma = \frac{1}{2}(\Gamma_1 + \Gamma_2), \quad \Delta M = M_1 - M_2, \quad \Delta \Gamma = (\Gamma_1 - \Gamma_2); \quad (2)$$

with the observable oscillations determined by the scaled parameters

$$x = \frac{\Delta M}{\Gamma}, \quad y = \frac{\Delta \Gamma}{2\Gamma}. \quad (3)$$

In the case of CP conservation the two $D$ eigenstates are the CP even and odd combinations. We will choose $D_1$ to be the CP even state. The sign choice for the mass and width difference varies among papers, we use the choice above.

Assuming CP conservation, small mixing parameters, and an initial state tagged as a $D^0$, we can write the time dependence to first order in $x$ and $y$:

$$D(t) = \left(D^0 + \bar{D}^0(-y - ix)\frac{\Gamma}{2}\right) e^{-(\Gamma/2 + m_M)t}. \quad (4)$$

Projected this onto a final state $f$ gives to first order the amplitude for finding $f$:

$$\left(A_f + \bar{A}_f(-y - ix)\frac{\Gamma}{2}\right) e^{-(\Gamma/2 + m_M)t}. \quad (5)$$

This leads to a number of ways to measure the effect of mixing, for example:

1. Wrong sign semileptonic decays. Here $A_f$ is zero and we measure directly the quantity, after integrating over decay times:

$$R_M = (x^2 + y^2)/2. \quad (6)$$

Limits using this measurement, however, are not yet sensitive enough to get down to the $10^{-4}$ level for $R_M$. Using 334 fb$^{-1}$ of data, electron decays only, and a double tag technique, BABAR measures $R_M = 0.4 \times 10^{-4}$, with a 68% confidence interval ($-5.6, 7.4) \times 10^{-4}$\cite{1}.

2. Cabibbo favored, right sign (RS) hadronic decays (for example $K^-\pi^+$). These are used to measure the average lifetime, with the correction from the term involving $x$ and $y$ usually ignored (provides a correction of $O(10^{-3})$).

3. Singly suppressed decays (for example $K^+\pi^-$ or $\pi^+\pi^-$). In this case tagging the initial state isn’t necessary. For CP even final states: $A_f = \bar{A}_f$. This provides the most direct way to measure $y$. With tagging we can also check for CP violation, by looking at the value of $y$ for each tag type. BABAR will be updating this measurement with the full statistics later this year. The initial BABAR measurement was based on 91 fb$^{-1}$ and gave the result $y = 0.8\%$, with statistical and systematic errors each about 0.4% \cite{2}, measurement \cite{3}.
4) Doubly suppressed and mixed, wrong sign (WS) decays (for example $K^+\pi^-$). Mixing leads to an exponential term multiplied by both a linear and a quadratic term in $t$. The quadratic term has a universal form depending on $R_M$. For any point in the decay phase space the decay rate is given by

$$\left( |A_f|^2 + |A_j| |\bar{A}_j| y \Gamma_t + |\bar{A}_j|^2 R_M \left( \Gamma_t^2 \right)^{\frac{3}{2}} \right) e^{-\Gamma_t t}.$$  \hspace{1cm} (7)

Here $y' = y \cos \delta - x \sin \delta$, where $\delta$ is a strong phase difference between the Cabibbo favored and doubly suppressed amplitudes. For the $K^+\pi^-\pi^0$ decay there is just the one phase and the ratio of $|A_f|^2$ to $|A_j|^2$ is defined to be $R_D$. For multibody decays the strong phase varies over the phase space and the term proportional to $t$ will involve a sum with different phases if we add all events in a given channel.

$\text{BaBar}$ has analyzed the decay channel $K^+\pi^-\pi^0$, with a mass cut that selects mostly $K^+\rho^- \to \pi^0$ decays, the largest channel for the Cabibbo allowed amplitude arising from mixing. Based on 230 fb$^{-1}$, $\text{BaBar}$ measures $[4]$

$$\alpha y' = (- 1.2^{+0.8}_{-0.8} \pm 0.2)\%$$

$$R_M = (0.023^{+0.18}_{-0.13} \pm 0.004)\%.$$  \hspace{1cm} (8)

The parameter $\alpha$ allows for the phase variation over the region summed over. A fit to the full Dalitz plot would allow more events to be used in the mixing study. This, however, requires a model for all the resonant and smooth components that contribute to the given channel, which may introduce uncertainties. $\text{BaBar}$ is working on such a fit, which will be based on approximately 1500 signal events.

Another important 3-body channel is the $K_S\pi^+\pi^-$ decay channel. Analysis of this channel was pioneered by CLEO $[5]$. It contains: CP-even, CP-odd, and mixed-CP resonances. Now one must correctly model the relative amounts of CP-odd and CP-even contributions (including smooth components) to get the correct lifetime difference. This channel also provides the possibility to directly measure $x$. $\text{BaBar}$ is working on this channel; Belle has published their results $[6]$. In the Standard Model $y$ and $x$ are mainly due to long-distance effects. They may be comparable in value but this depends on physics that is difficult to model. Long-distance effects control how complete the SU(3) cancellation is, which would make both parameters vanish in the symmetry limit. The exact values therefore depends on SU(3) violations in matrix elements and phase space. Also, the sign of $x/y$ provides an important measurement. One might expect the $x$ and $y$ parameters to be in the range $O(10^{-8}$ to $10^{-2})$. Thus the present data are consistent with the Standard Model. Searches for CP violation are important goals of the B-factories, since observation at a non-negligible level would signify new physics.

We will turn now to the strongest evidence for $D$-mixing from $\text{BaBar}$, using the $K\pi$ final state. This result has recently been published $[7]$.  

### 3. Analysis of the $K\pi$ channel

We study the right-sign (RS), Cabibbo-favored (CF) decay $D^0 \to K^+\pi^-$ $[8]$ and the wrong-sign (WS) decay $D^0 \to K^+\pi^-$. The latter can be produced via the doubly Cabibbo-suppressed (DCS) decay $D^0 \to K^+\pi^+$ or via mixing followed by a CF decay $D^0 \to \bar{D^0} \to K^+\pi^-$. The DCS decay has a small rate $R_D$ of order $\tan^2 \theta_C \approx 0.3\%$ relative to the CF decay with $\theta_C$ the Cabibbo angle. We tag the $D^0$ at production using the decay $D^{++} \to \pi^+D^0$ where the $\pi^+$ is referred to as the “slow pion”. In RS decays the $\pi^+$ and kaon have opposite charges, while in WS decays the charges are the same. The time dependence of the WS decay rate is used to separate the contributions of DCS decays from $D^0\bar{D^0}$ mixing.

We study both CP-conserving and CP-violating cases. For the CP-conserving case, we fit for the parameters $R_D$, $x^{\pm}$, and $y'$. To search for CP violation, we apply Eq. (7) to the $D^0$ and $\bar{D^0}$ samples separately, fitting for the parameters $\{R_{D}\pm, x^{\pm}, y'^{\pm}\}$ for $D^0$ ($+$) decays and $\bar{D^0}$ ($-$) decays.

We select $D^0$ candidates by pairing oppositely-charged tracks with a $K^{+}\pi^{\mp}$ invariant mass $m_{K\pi}$ between 1.81 and 1.92 GeV/$c^2$. We require the $\pi^+_s$ to have a momentum in the laboratory frame greater than 0.1 GeV/$c$ and in the $e^+e^-$ center-of-mass (CM) frame below 0.45 GeV/$c$.

To obtain the proper decay time $t$ and its error $\sigma_t$ for each $D^0$ candidate, we refit the $K^+$ and $\pi^+$ tracks, constraining them to originate from a common vertex. We also require the $D^0$ and $\pi^+_s$ to originate from a common vertex, constrained by the position and size of the $e^+e^-$ interaction region. The vertical RMS size of each beam is typically 6 $\mu$m. We require the $\chi^2$ probability of the vertex-constrained combined fit $P(x^2)$ to be at least 0.1%, and the $m_{K\pi}\pi^+_s-m_{K\pi}$ mass difference $\Delta m$ to satisfy $0.14 < \Delta m < 0.16$ GeV/$c^2$.

To remove $D^0$ candidates from $B$-meson decays and to reduce combinatorial backgrounds, we require each $D^0$ to have a momentum in the CM frame greater than 2.5 GeV/$c$. We require $-2 < t < 4$ ps and $\sigma_t < 0.5$ ps, (the most probable value of $\sigma_t$ for signal events is 0.16 ps). For $D^{++}$ candidates sharing one or more tracks with other $D^{++}$ candidates, we retain only the candidate with the highest $P(x^2)$. After applying all criteria, we keep approximately 1.229,000 RS and 64,000 WS $D^0$ and $\bar{D^0}$ candidates.

The mixing parameters are determined in an unbinned, extended maximum-likelihood fit to the RS and WS data samples over the four observables $m_{K\pi}$, $\Delta m$, $t$, and $\sigma_t$. The fit is performed in several stages.
First, RS and WS signal and background shape parameters are determined from a fit to $m_{K\pi}$ and $\Delta m$, and are not varied in subsequent fits. Next, the $D^0$ proper-time resolution function and lifetime are determined in a fit to the RS data using $m_{K\pi}$ and $\Delta m$ to separate the signal and background components. We fit to the WS data sample using three different models. The first model assumes both CP conservation and the absence of mixing. The second model allows for mixing, but assumes no CP violation. The third model allows for both mixing and CP violation.

The RS and WS [$m_{K\pi}, \Delta m$] distributions are described by four components: signal, random $\pi^+_s$ misreconstructed $D^0$ and combinatorial background. The signal component has a characteristic peak in both $m_{K\pi}$ and $\Delta m$. The random $\pi^+_s$ component models reconstructed $D^0$ decays combined with a random slow pion and has the same shape in $m_{K\pi}$ as signal events, but does not peak in $\Delta m$. Misreconstructed $D^0$ events have one or more of the $D^0$ decay products either not reconstructed or reconstructed with the wrong particle hypothesis. They peak in $\Delta m$, but not in $m_{K\pi}$. For RS events, most of these are semileptonic $D^0$ decays. For WS events, the main contribution is RS $D^0 \to K^-\pi^+$ decays where the $K^-$ and the $\pi^+$ are misidentified as $\pi^-$ and $K^+$, respectively. Combinatorial background events are those not described by the above components; they do not exhibit any peaking structure in $m_{K\pi}$ or $\Delta m$.

The functional forms of the probability density functions (PDFs) for the signal and background components are chosen based on studies of Monte Carlo (MC) samples. However, all parameters are determined from two-dimensional likelihood fits to data over the full $m_{K\pi}$ and $\Delta m$ region.

We fit the RS and WS data samples simultaneously with shape parameters describing the signal and random $\pi^+_s$ components shared between the two data samples. We find $1,141,500 \pm 1,200$ RS signal events and $4,030 \pm 90$ WS signal events. The dominant background component is the random $\pi^+_s$ background. Projections of the WS data and fit are shown in Fig. 1.

The measured proper-time distribution for the RS signal is described by an exponential function convolved with a resolution function whose parameters are determined by the fit to the data. The resolution function is the sum of three Gaussians with widths proportional to the estimated event-by-event proper-time uncertainty $\sigma_t$. The random $\pi^+_s$ background is described by the same proper-time distribution as signal events, since the slow pion has little weight in the vertex fit. The proper-time distribution of the combinatorial background is described by a sum of two Gaussians, one of which has a power-law tail to account for a small long-lived component. The combinatorial background and real $D^0$ decays have different $\sigma_t$ distributions, as determined from data using a background-subtraction technique based on the fit to $m_{K\pi}$ and $\Delta m$.

The fit to the RS proper-time distribution is performed over all events in the full $m_{K\pi}$ and $\Delta m$ region. The PDFs for signal and background in $m_{K\pi}$ and $\Delta m$ are used in the proper-time fit with all parameters fixed to their previously determined values. The fitted $D^0$ lifetime is found to be consistent with the world-average lifetime [9].

The measured proper-time distribution for the WS signal is modeled by Eq. (7) convolved with the resolution function determined in the RS proper-time fit. The random $\pi^+_s$ and misreconstructed $D^0$ backgrounds are described by the RS signal proper-time distribution since they are real $D^0$ decays. The proper-time distribution for WS data is shown in Fig. 2. The fit results with and without mixing are shown as the overlaid curves.

The fit with mixing provides a substantially better description of the data than the fit with no mixing. The significance of the mixing signal is evaluated based on the change in negative log likelihood with respect to the minimum. Figure 3 shows confidence-level (CL) contours calculated from the change in negative log likelihood ($-2\Delta \ln L$) in two dimensions ($x^2$ and $y'$) with systematic uncertainties included. The likelihood maximum is at the unphysical value of $x^2 = -2.2 \times 10^{-4}$ and $y' = 9.7 \times 10^{-3}$. The value of $-2\Delta \ln L$ at the most likely point in the physically allowed region ($x^2 = 0$ and $y' = 6.4 \times 10^{-3}$) is 0.7 units. The value of $-2\Delta \ln L$ for no-mixing is 23.9 units. Including the systematic uncertainties, this corresponds to a significance equivalent to 3.9 standard deviations ($1 - CL = 1 \times 10^{-3}$) and thus constitutes evidence for mixing. The fitted values of the mixing parameters and $R_D$ are listed in Table I. The correlation coefficient between the $x^2$ and $y'$ parameters is $-0.94$.

Allowing for the possibility of CP violation, we calculate the values of $R_D = \sqrt{R_D^+ R_D^-}$ and $A_D = (R_D^+ - R_D^-)/(R_D^+ + R_D^-)$ listed in Table I, from the fitted $R_D^\pm$ values. The best fit points ($x^{2\pm}, y'^{\pm}$) shown in Table I are more than three standard deviations away.

![Figure 1](image-url)  
Figure 1: a) $m_{K\pi}$ for wrong-sign (WS) candidates with $0.1445 < \Delta m < 0.1465\text{ GeV}/c^2$, and b) $\Delta m$ for WS candidates with $1.843 < m_{K\pi} < 1.883\text{ GeV}/c^2$. The fitted PDFs are overlaid.
and without mixing. The solid curve shows the difference between fits with (

\begin{equation}
\chi^2 \quad -0.22 \pm 0.30 \pm 0.21 \\
y' \quad 9.7 \pm 4.4 \pm 3.1
\end{equation}

) CL contours for 1

\begin{equation}
R_D \quad 3.03 \pm 0.16 \pm 0.10 \\
\sigma_D^2 \quad 3.03 \pm 0.16 \pm 0.10 \\
\sigma_D^2+ \quad -0.24 \pm 0.43 \pm 0.30 \\
\sigma_D^2- \quad -0.20 \pm 0.41 \pm 0.29 \\
y' \quad 9.6 \pm 6.1 \pm 4.3
\end{equation}

Figure 3. All cross checks indicate that the close agreement between the separate

D^0 and D\bar{D} fits to slices in measured proper time is coincidental.

As a cross-check of the mixing signal, we perform independent \(\{m_{K\pi}, \Delta m\}\) fits with no shared parameters for intervals in proper time selected to have ap-

proximately equal numbers of RS candidates. The fitted WS branching fractions are shown in Fig. 4 and are seen to increase with time. The slope is consistent with the measured mixing parameters and inconsistent with the no-mixing hypothesis.

\begin{equation}
R_{WS} \quad 0.45 \\
\chi^2 \quad 24.0
\end{equation}

We validated the fitting procedure on simulated data samples using both MC samples with the full detector simulation and large parametrized MC samples. In all cases we found the fit to be unbiased. As a further cross-check, we performed a fit to the RS data proper-time distribution allowing for mixing in the signal component; the fitted values of the mixing parameters are consistent with no mixing.

In evaluating systematic uncertainties in \(R_D\) and the mixing parameters we considered variations in the fit model and in the selection criteria. We also considered alternative forms of the \(m_{K\pi}, \Delta m\), proper time, and \(\sigma_t\) PDFs. We varied the \(t\) and \(\sigma_t\) requirements. In addition, we considered variations that keep or reject all \(D^+\) candidates sharing tracks with other candidates.
For each source of systematic error, we compute the significance

\[ s_i^2 = 2t\left[ \ln L(x'^2_i, y'_i) - \ln L(x'^2_i, y'_i) \right]/2.3 \]

where \((x'^2_i, y'_i)\) are the parameters obtained from the standard fit, \((x'^2_i, y'_i)\) the parameters from the fit including the \(i^{th}\) systematic variation, and \(L\) the likelihood of the standard fit. The factor 2.3 is the 68% confidence level for 2 degrees of freedom. To estimate the significance of our results in \((x'^2, y')\), we reduce \(-2\Delta \ln L\) by a factor of \(1 + \Sigma s_i^2 = 1.3\) to account for systematic errors. The largest contribution to this factor, 0.06, is due to uncertainty in modeling the long decay time component from other \(D\) decays in the signal region. The second largest component, 0.05, is due to the presence of a non-zero mean in the proper time signal resolution PDF. The mean value is determined in the RS proper time fit to be 3.6 fs and is due to small misalignments in the detector. The error of \(15 \times 10^{-3}\) on \(A_D\) is primarily due to uncertainties in modeling the differences between \(K^+\) and \(K^-\) absorption in the detector.

In conclusion we summarize the \(\text{BaBar}\) evidence for \(D^0-\bar{D}^0\) mixing. Our result is inconsistent with the no-mixing hypothesis at a significance of 3.9 standard deviations. We measure \(y' = [9.7 \pm 4.4 \text{ (stat.)} \pm 3.1 \text{ (syst.)}] \times 10^{-3}\), while \(x'^2\) is consistent with zero.

The use of charge-conjugate modes is implied unless otherwise noted.

\section*{References}

[8] The use of charge-conjugate modes is implied unless otherwise noted.
**D^0 Mixing at Belle**

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We report the recent two results of D^0-D^0 mixing studies at Belle in D^0 → K^+K^-/π^+π^- and D^0 → K^0_Sπ^+π^- decays. The former measures the relative difference of the lifetimes y_{CP}, giving the evidence of D^0-D^0 mixing; the latter measures the D^0 mixing parameters x and y.

1. Introduction

Mixing phenomenon, i.e. the oscillation of a neutral meson into its corresponding anti-meson as a function of time, has been observed in the K^0, B^0, and most recently B^0_s systems. This process is also possible in the D-meson system, but has not previously been observed.

Mixing in heavy flavor systems such as that of B^0 and B^0_s is governed by the short-distance box diagram. However, in the D^0 system this diagram is both GIM-suppressed and doubly-Cabibbo-suppressed relative to the amplitude dominating the decay width, and thus the short-distance rate is very small. Consequently, D^0-D^0 mixing is expected to be dominated by long-distance processes that are difficult to calculate; theoretical estimates for the mixing parameters x = (m_1-m_2)/Γ and y = (Γ_1-Γ_2)/2Γ range over two-three orders of magnitude [1]. Here, m_1, m_2 (Γ_1, Γ_2) are the masses (decay widths) of the mass eigenstates |D_1,2⟩ = p|D^0⟩ ± q|D^0⟩, and Γ = (Γ_1+Γ_2)/2. The parameters p and q are complex coefficients satisfying |p|^2 + |q|^2 = 1.

The general experimental method identifies the flavor of the neutral D meson when produced by reconstructing the decay D^*+ → D^0π^+ or D^*^- → D^0^-π^-; the charge of the accompanying pion identifies the D flavor. Because the energy release in D^* decays is only ~ 6 MeV, the background is largely suppressed. The D^0 decay time (t) is calculated via (l/p) × m_{π^0}, where l is the distance between the D^* decay vertices and p is the D^0 momentum. The D^* vertex position is taken to be the intersection of the D^0 momentum with the beampot profile. To reject D^*(s) decays originating from B decays, one requires p_{D^*} > 2.5 GeV, which is the kinematic endpoint.

2. CP-eigenstates K^+K^- and π^+π^-

We have studied the decays to CP eigenstates D^0 → K^+K^- and D^0 → π^+π^-; treating the decay-time distributions as exponential, we measured the quantity

$$y_{CP} = \frac{\tau_{K^-K^+}}{\tau_{K^+K^-}} - 1,$$

where τ_{K^-K^+} and τ_{K^+K^-} are the lifetimes of D^0 → K^-π^+ and D^0 → K^+π^- (or D^0 → π^-π^+) decays. It can be shown that y_{CP} = y cos φ - x sin φ [3], where A_M parameterizes CPV in mixing and φ is a weak phase. If CP is conserved, A_M = φ = 0 and y_{CP} = y. This method has been used by numerous experiments to constrain y_{CP} [4]. Our measurement, based on 540 fb^-1 data, yields a nonzero value of y_{CP} with > 3σ significance [5]. We also searched for CPV by measuring the quantity

$$A_Γ = \frac{τ(D^0 → K^-K^+) - τ(D^0 → K^+K^-)}{τ(D^0 → K^-K^+) + τ(D^0 → K^+K^-)};$$


this observable equals A_Γ = 1/2A_My cos φ - x sin φ [3].

We reconstruct D^0+ → D^0π^+π^- and D^0 → K^+K^+, K^-π^+, and π^+π^- decays. Candidate D^0 mesons are selected using two kinematic observables: the invariant mass of the D^0 decay products, M, and the energy release in the D^0 decay, Q = (M_D^0 - M - m_{π^+})c^2. According to Monte Carlo (MC) simulated distributions of t, M and Q, background events fall into four categories: (1) combinatorial, with zero apparent lifetime; (2) true D^0 mesons combined with random slow pions (this has the same apparent lifetime as the signal); (3) D^0 decays to three or more particles, and (4) other charm hadron decays. The apparent lifetime of the latter two categories is 10-30% larger than τ_{D^0}.

For the lifetime measurements, we select the events satisfying |ΔM|/σ_M < 2.3, |Q - 5.9 MeV| < 0.80 MeV and σ_t < 370 fs, where ΔM = M - m_{D^0}, and σ_t is the decay time uncertainties calculated event-by-event. The invariant mass resolution σ_M varies from 5.5-6.8 MeV/c^2, depending on the decay channel. The selection criteria are chosen to minimize the expected statistical error on y_{CP} using the MC. We find 111 × 10^3 K^+K^-π^0, 1.22 × 10^6 K^-π^+π^0, and 49 × 10^3 π^+π^- signal events, with purities of 98%, 99% and 92% respectively.

The relative lifetime difference y_{CP} is determined by performing a simultaneous binned maximum likelihood fit to the D^0 → K^+K^-π^0, D^0 → K^-π^+π^0, D^0 → π^+π^- decay time distributions. Each distribution is assumed to be a sum of signal and background contributions, with the signal contribution being a convolution of an exponential and a detector resolution.
The resolution function $R(t - t')$ is constructed from the normalized distribution of the decay time uncertainties $\sigma_t$. The $\sigma_t$ of a reconstructed event ideally represents an uncertainty with a Gaussian probability density: in this case, bin $i$ in the $\sigma_t$ distribution is taken to correspond to a Gaussian resolution term $\sigma_{i,k}^\text{pull}$ and fractions $w_k$, constrained to the same mean. Therefore, we choose the parameterization

$$R(t - t') = \sum_{i=1}^{n} f_i \sum_{k=1}^{3} w_k G(t - t'; \sigma_{ik}, t_0),$$

with $\sigma_{ik} = s_k \sigma_{i,k}^\text{pull} \sigma_t$, where the $s_k$ are three scale factors introduced to account for differences between the simulated and real $\sigma_{i,k}^\text{pull}$, and $t_0$ allows for a (common) offset of the Gaussian terms from zero.

The background $B(t)$ is parameterized assuming two lifetime components: an exponential and a $\delta$ function, each convolved with corresponding resolution functions as parameterized by Eq. (4). Separate $B(t)$ parameters for each final state are determined by fits to the $t$ distributions of events in $M$ sidebands. The MC is used to select the sideband region that best reproduces the timing distribution of background events in the signal region.

Fitting the $K^-\pi^+, K^+K^-$, and $\pi^+\pi^-$ decay time distributions (Figs. 1(a)-(c)) shows a statistically significant difference between the $K^-\pi^+$ and $h^+h^-$ lifetimes. The effect is visible in Fig. 1d, which plots the ratio of event yields $N_{h^+h^-}/N_{K^-\pi^+}$ as a function of decay time. The fitted lifetime of $D^0$ meson in the $K^-\pi^+$ final states is $408.7 \pm 0.6$ fs, which is consistent with the PDG value [6] (and actually has greater statistical precision). We measure

$$y_{\text{CP}} = (1.31 \pm 0.32 \pm 0.25)\%,$$

which deviates from zero by $3.2\sigma$. The systematic error is dominated by uncertainty in the background decay time distribution, variation of selection criteria, and the assumption that $t_0$ is equal for all three final states. The analysis also measures

$$A^\tau = (0.01 \pm 0.30 \pm 0.15)\%,$$

which is consistent with zero (no $CPV$). The sources of systematic error for $A^\tau$ are similar to those for $y_{\text{CP}}$. 

3. Dalitz Plot Analysis of $D^0 \rightarrow K_S^0 \pi^+\pi^-$

The time dependence of the Dalitz plot for $D^0 \rightarrow K_S^0 \pi^+\pi^-$ decays is sensitive to mixing parameters $x$ and $y$ without ambiguity due to strong phases. For a particular point in the Dalitz plot $(m_2^+, m_2^-)$, where $m_+ \equiv m(K_S^0 \pi^+)$ and $m_- \equiv m(K_S^0 \pi^-)$, the overall decay amplitude is

$$A_{D^0}(m_2^+, m_2^-) = \frac{e_{1}(t) + e_{2}(t)}{2} + \frac{(2p)}{A_{D^0}(m_2^+, m_2^-)} \frac{e_{1}(t) - e_{2}(t)}{2} \left( \frac{e_{(1,2)}(t)}{A_{D^0}(m_2^+, m_2^-)} \right),$$

where $e_{(1,2)}(t) = e^{-(m_1^2 + \Gamma_1)^2 t}$. The first term represents the (time-dependent) amplitude for $D^0 \rightarrow K_S^0 \pi^+\pi^-$, and the second term represents the amplitude for $D^0 \rightarrow (\overline{D})^0 \rightarrow K_S^0 \pi^+\pi^-$. Taking the modulus squared of Eq. (7) gives the decay rate or, equivalently, the density of points $\rho(m_2^+, m_2^-; t)$. The result contains terms proportional to $\cosh(y \Gamma t)$, $\cos(x \Gamma t)$, and $\sin(x \Gamma t)$, and thus fitting the time-dependence of $\rho(m_2^+, m_2^-; t)$ determines $x$ and $y$. This method was developed by CLEO [7].

To use Eq. (7) requires choosing a model for the decay amplitudes $A_{D^0}(m_2^+, m_2^-)$. This is usually taken to be the “isobar model” [8], and thus, in addition to $x$ and $y$, one also fits for the magnitudes and phases of various intermediate states. Specifically, $A_{D^0}(m_2^+, m_2^-) = \sum_j a_j e^{i \delta_j} A_j$, where $\delta_j$ is a
strong phase, $A_j$ is the product of a relativistic Breit-Wigner function and Blatt-Weiskopf form factors, and the parameter $j$ runs over all intermediate states. This sum includes possible scalar resonances and, typically, a constant non-resonant term. For no direct CPV, $A_{\Gamma 0}(m_{\pi^+}^2, m_{\pi^+}^2) = A_{D0}(m_{\pi^+}^2, m_{\pi^+}^2)$; otherwise, one must consider separate decay parameters $(a_j, \delta_j)$ for $D^0$ decays and $(\bar{a}_j, \delta_j)$ for $\bar{D}^0$ decays.

We have fit a large $D^0 \rightarrow K_s^0 \pi^+ \pi^-$ sample selected from 540 fb$^{-1}$ of data [9]. The analysis proceeds in two steps. First, signal and background yields are determined from a two-dimensional fit to variables $M(K \pi) \pi K_{\pi}$ and $Q = M(K \pi K_{\pi}) - M(K \pi) - m_{\pi^+}$. Within a signal region $|M(K \pi K_{\pi}) - m_{D0}| < 15 \text{ MeV/c}^2$ and $|Q - 5.9 \text{ MeV}| < 1.0 \text{ MeV}$ (corresponding to $3 \sigma$ in resolution), there are 534,000 signal candidates with 95% purity. These events are fit for $x$ and $y$; the (unbinned ML) fit variables are $m_{\pi^+}^2, m_{\pi^+}^2$, and the decay time $t$. Most of the background is combinatoric, i.e., the $D^0$ candidate results from a random combination of tracks. The decay-time distribution of this background is modeled as the sum of a delta function and an exponential function convolved with a Gaussian resolution function, and all parameters are determined from fitting events in the sideband $30 \text{ MeV/c}^2 < |M(K \pi) - m_{D0}| < 55 \text{ MeV/c}^2$.

The results from two separate fits are listed in Table I. In the first fit CP conservation is assumed, i.e., $q/p = 1$ and $A_{\Gamma 0}(m_{\pi^+}^2, m_{\pi^+}^2) = A_{D0}(m_{\pi^+}^2, m_{\pi^+}^2)$, the free parameters are $x, y, \tau_{D0}$, some timing resolution function parameters, and decay model parameters $(a_j, \delta_j)$. The results for the latter are listed in Table II. The results for $x$ and $y$ indicate that $x$ is positive, about $2 \sigma$ from zero. Projections of the fit are shown in Fig. 2. The fit also yields $\tau_{D0} = (409.9 \pm 1.0) \text{ fs}$, which is consistent with the PDG value [6] (and actually has greater statistical precision).

Table I: Fit results and 95% C.L. intervals for $x$ and $y$, from analysis of $D^0 \rightarrow K_s^0 \pi^+ \pi^-$ decays. The errors are statistical, experimental systematic, and decay-model systematic, respectively.

<table>
<thead>
<tr>
<th>Fit Param.</th>
<th>Result</th>
<th>95% C.L. inter.</th>
</tr>
</thead>
<tbody>
<tr>
<td>No $x$ (%)</td>
<td>0.80 ± 0.29</td>
<td>(0.0, 1.6)</td>
</tr>
<tr>
<td>CPV $y$ (%)</td>
<td>0.33 ± 0.24</td>
<td>(0.0, 0.96)</td>
</tr>
<tr>
<td>$CPV \ y$ (%)</td>
<td>0.31 ± 0.13</td>
<td>(0.0, 1.6)</td>
</tr>
<tr>
<td>$q/p$ (%)</td>
<td>0.25 ± 0.07</td>
<td>(0.0, 1.0)</td>
</tr>
<tr>
<td>$\phi$ (°)</td>
<td>−14.6 ± 1.2</td>
<td></td>
</tr>
</tbody>
</table>
possible sign change, the effect upon $x$ and $y$ is small, and the results for $|q/p|$ and $\phi$ are consistent with no CPV. The sets of Dalitz parameters $(a_r, \delta_r)$ and $(\bar{a}_r, \bar{\delta}_r)$ are consistent with each other, indicating no direct CPV. Taking $a_j = \bar{a}_j$ and $\delta_j = \bar{\delta}_j$ (i.e., no direct CPV) and repeating the fit gives $|q/p| = 0.95^{+0.22}_{-0.20}$ and $\phi = (-2^{+10}_{-11})^\circ$.

The dominant systematic errors are from the time dependence of the Dalitz plot background, and the effect of the $p_D^{-}$ momentum cut used to reject $D^*$'s originating from $B$ decays. The default fit includes $\pi\pi$ scalar resonances $\sigma_1$ and $\sigma_2$; when evaluating systematic errors, the fit is repeated without any $\pi\pi$ scalar resonances using K-matrix formalism [10]. The influence upon $x$ and $y$ is small and included as a systematic error.

The 95% C.L. contour for $(x, y)$ is plotted in Fig. 3. The contour is obtained from the locus of points where $-2 \ln L$ rises by 5.99 units from the minimum value; the distance of the points from the origin is subsequently rescaled to include systematic uncertainty. We note that for the CPV-allowed case, the reflections of the contours through the origin are also allowed regions.

![Figure 3: 95% C.L. contours for $(x, y)$: dotted (solid) is statistical (statistical plus systematic) contour for no CPV; dashed-dotted (dashed) is statistical (statistical plus systematic) contour allowing for CPV. The point is the best-fit value for no CPV.](image)

**References**

[2] Charge-conjugate modes are included unless noted otherwise.
Measurement of the Strong Phase in $D^0 \to K^+\pi^-$ Using Quantum Correlations

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We exploit the quantum coherence between pair-produced $D^0$ and $\bar{D}^0$ in $\psi(3770)$ decays to study charm mixing and to make a first measurement of the relative strong phase $\delta$ between $D^0 \to K^+\pi^-$ and $\bar{D}^0 \to K^+\pi^-$. Using 281 pb$^{-1}$ of $e^+e^-$ collision data collected with the CLEO-c detector at $E_{cm}$ = 3.77 GeV, as well as branching fraction input from other experiments, we make a preliminary determination of $\cos \delta = 1.03 \pm 0.19 \pm 0.08$, where the uncertainties are statistical and systematic, respectively. By further including other external mixing parameter measurements, we obtain an alternate measurement of $\cos \delta = 0.93\pm0.32\pm0.04$, where the systematic uncertainty from assuming $x \sin \delta = 0$ has not been included.

1. Introduction

Recent measurements of $D^0$-$\bar{D}^0$ mixing parameters [1, 2, 3, 4] highlight the need for information on the relative phase between the Cabibbo favored decay $D^0 \to K^-\pi^+$ and the doubly Cabibbo suppressed decay $\bar{D}^0 \to K^-\pi^-$. Here, we present a measurement that takes advantage of the correlated production of $D^0$ and $\bar{D}^0$ mesons in $e^+e^-$ collisions. If there are no accompanying particles, the $D^0\bar{D}^0$ pair is in a quantum-coherent state $C = -1$. Because the initial state (the virtual photon) has $J^{PC} = 1^{-+}$, there follows a set of selection rules for the decays of the $D^0$ and $\bar{D}^0$ [5, 6, 7, 8, 9, 10, 11, 12, 13, 14, 15]. For example, both $D^0$ and $D^0$ cannot decay to $CP$ eigenstates with the same eigenvalue. On the other hand, decays to $CP$ eigenstates of opposite eigenvalue are enhanced by a factor of two. More generally, final states that can be reached by both $D^0$ and $D^0$ (such as $K^-\pi^+$) are subject to similar interference effects. As a result, the effective $D^0$ branching fractions in this $D^0\bar{D}^0$ system differ from those measured in isolated $D^0$ mesons. Moreover, using time-independent rate measurements, it becomes possible to probe $D^0$-$\bar{D}^0$ mixing as well as the relative strong phases between $D^0$ and $\bar{D}^0$ decay amplitudes to any given final state.

In the Standard Model, $D^0$-$\bar{D}^0$ mixing is suppressed both by the GIM mechanism and by CKM matrix elements, although sizeable mixing could arise from both by the GIM mechanism and by CKM matrix $D$ mixing as well as the relative strong phases between $D^0$ and $\bar{D}^0$ measurements, it becomes possible to probe $D^0$-$\bar{D}^0$ mixing as well as to make a first measurement of the relative strong phase $\delta$ between $D^0 \to K^+\pi^-$ and $\bar{D}^0 \to K^+\pi^-$. Using 281 pb$^{-1}$ of $e^+e^-$ collision data collected with the CLEO-c detector at $E_{cm}$ = 3.77 GeV, as well as branching fraction input from other experiments, we make a preliminary determination of $\cos \delta = 1.03 \pm 0.19 \pm 0.08$, where the uncertainties are statistical and systematic, respectively. By further including other external mixing parameter measurements, we obtain an alternate measurement of $\cos \delta = 0.93\pm0.32\pm0.04$, where the systematic uncertainty from assuming $x \sin \delta = 0$ has not been included.

1. Introduction

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In the Standard Model, $D^0$-$\bar{D}^0$ mixing is suppressed both by the GIM mechanism and by CKM matrix elements, although sizeable mixing could arise from new physics [14]. Charm mixing is conventionally described by two small dimensionless parameters:

$$x = \frac{M_2 - M_1}{\Gamma_2 + \Gamma_1} \quad (1)$$
$$y = \frac{\Gamma_2 - \Gamma_1}{\Gamma_2 + \Gamma_1} \quad (2)$$

where $M_{1,2}$ and $\Gamma_{1,2}$ are the masses and widths, respectively, of the neutral $D$ meson $CP$ eigenstates, $D_1$ ($CP$-odd) and $D_2$ ($CP$-even), which are defined as follows:

$$|D_1| = \frac{|D^0| + |\bar{D}^0|}{\sqrt{2}} \quad (3)$$
$$|D_2| = \frac{|D^0| - |\bar{D}^0|}{\sqrt{2}} \quad (4)$$

assuming $CP$ conservation. The mixing probability is then denoted by $R_M = (x^2 + y^2)/2$, and the width of the $D^0$ and $\bar{D}^0$ flavor eigenstates is $\Gamma = (\Gamma_1 + \Gamma_2)/2$.

Many previous searches for charm mixing have focused on $D^0$ decay times. Direct measurements of $y$ come from comparing lifetimes in $D^0 \to K^+K^-$ and $\pi^+\pi^-$ decay to that in $D^0 \to K^-\pi^+\pi^-$. An indirect measure of $y$ is provided by the “wrong-sign” process $D^0 \to K^+\pi^-$, where interference between the doubly-Cabibbo-suppressed (DCS) amplitude and the mixing amplitude manifests itself in the apparent $D^0$ lifetime. These analyses are sensitive to $y' \equiv y \cos \delta - x \sin \delta$, where $-\delta$ is the phase of the amplitude ratio $\langle K^+\pi^-|D^0\rangle/\langle K^-\pi^+|\bar{D}^0\rangle$. Below, we also denote the magnitude of this ratio by $r$, which is measured to be approximately 0.06. Because $\delta$ has not previously been measured, the separate determinations of $y$ and $y'$ above have not been directly comparable.

In this note, we present an implementation of the method described in Ref. [10] for measuring $y$ and $\cos \delta$ using quantum correlations at the $\psi(3770)$ resonance. Our experimental technique is an extension of the double tagging method previously used to determine absolute hadronic $D$-meson branching fractions at CLEO-c [17]. This method combines yields of fully-reconstructed single tags (ST), which are individually reconstructed $D^0$ or $\bar{D}^0$ candidates, with yields of double tags (DT), which are events where both $D^0$ and $\bar{D}^0$ are reconstructed, to give absolute branching fractions without needing to know the luminosity or $D^0\bar{D}^0$ production cross section. Given a set of input yields, efficiencies, and background estimates, a least-squares fitter [18] extracts the number of $D^0\bar{D}^0$ pairs produced $(N)$ and the branching fractions $(B)$ of the reconstructed $D^0$ final states, while accounting for all statistical and systematic uncertain-
ties and their correlations. We employ a modified version of this fitter that also determines $y$, $x^2$, $r^2$, and $r \cos \delta$ using the following categories of reconstructed final states: $D \to K^+ \pi^\mp$, $CP$-even ($S_\pm$) and $CP$-odd ($S_\mp$) eigenstates, and semileptonic decays ($e^\pm$). For optimal precision on $\delta$, we also incorporate measurements of branching fractions and mixing parameters from other CLEO-c analyses or from external sources. $CP$ violation in $D$ and $K$ decays are negligible second order effects that we ignore.

2. Formalism

To first order in $x$ and $y$, the $C$-odd width $\Gamma_{D^0 \bar{D}^0}(i,j)$ for $D^0 \bar{D}^0$ decay to final state $i/j$ follows from the anti-symmetric amplitude $M_{ij}$:

$$\Gamma_{D^0 \bar{D}^0}(i,j) \propto |M_{ij}|^2 = |\langle i|D_2\rangle\langle j|D_1\rangle - \langle i|D_1\rangle\langle j|D_2\rangle|^2,$$

where $A_i \equiv \langle i|D^0\rangle$, $\bar{A}_i \equiv \langle i|\bar{D}^0\rangle$. The total width, $\Gamma_{D^0 \bar{D}^0}$, is the same as for uncorrelated decay, as are ST rates. However, unlike the case of uncorrelated $D^0 \bar{D}^0$, we can consider the $C$-odd $D^0 \bar{D}^0$ system as a $D_1D_2$ pair. If only flavored final states are considered, as in Ref. [17], then the effects of quantum correlations are negligible. In this analysis, we also include $CP$ eigenstates, which brings additional sensitivity to $y$ and $\delta$, as demonstrated below.

Quantum-correlated semileptonic rates probe $y$ because the decay width does not depend on the $CP$ eigenvalue of the parent $D$ meson, as this weak decay is only sensitive to flavor content. However, the total width of the parent meson does depend on its $CP$ eigenvalue: $\Gamma_{1,2} = \Gamma(1 \mp y)$, so the semilepton branching fraction for $D_1$ or $D_2$ is modified by $1 \pm y$. If we reconstruct a semilepton decay in the same event as a $D_2 \to S_+\pi^-$ decay, the semilepton $D$ must be a $D_1$. Therefore, the effective quantum-correlated $D^0 \bar{D}^0$ branching fractions ($\mathcal{F}_{\text{corr}}$) for $CP$-tagged semileptonic final states depend on $y$:

$$\mathcal{F}_{S_{\pm}\pi^\mp} \approx 2B_{S_{\pm}}B_{\ell}(1 \pm y).$$

Combined with estimates of $B_{\ell}$ and $B_{S_{\pm}}$ from ST yields, external sources, and flavor-tagged semileptonic yields, this equation allows $y$ to be determined.

Similarly, if we reconstruct a $D \to K^- \pi^+$ decay in the same event as a $D_2 \to S_+\pi^-$, then we know the $K^- \pi^+$ was produced from a $D_1$. The effective branching fraction for this DT process is therefore

$$\mathcal{F}_{K^-\pi^+/K^+\pi^-} \approx |\langle S_+|D_2\rangle\langle K^-\pi^+|D_1\rangle|^2 = A_{S_+}^2|A_{K^-\pi^+} + \bar{A}_{K^-\pi^+}|^2 = A_{S_+}^2A_{K^-\pi^+}^2|1 + r e^{-i\delta}|^2 \approx B_{S_+}B_{K^+\pi}(1 + R_{WS} + 2r \cos \delta + y),$$

where $R_{WS} \equiv \Gamma(D^0 \to K^- \pi^+)/\Gamma(D^0 \to K^- \pi^+) = r^2 + ry^\prime + R_M$, and we have used $B_{S_{\pm}} \propto A_{S_{\pm}}^2(1 \mp y)$ and $B_{K^\mp} \propto A_{K^\mp}^2(1 + ry\cos\delta + rx\sin\delta)$. In an analogous fashion, we find $\mathcal{F}_{S_{\mp}\pi^-/K^+\pi^+} \approx B_{S_{\mp}}B_{K^+\pi}(1 + R_{WS} - 2r \cos \delta - y)$. When combined with knowledge of $\mathcal{B}_{S_{\pm}}, y$, and $r$, the asymmetry between these two DT yields gives $\cos \delta$. In the absence of quantum correlations, the effective branching fractions above would be $B_{S_{\pm}}B_{K^+\pi}(1 + R_{WS})$.

More concretely, we evaluate Eq. 5 with the above definitions of $r$ and $\delta$ to produce the expressions in Table I. In doing so, we use the fact that inclusive ST rates are given by the incoherent branching fractions since each event contains one $D^0$ and one $\bar{D}^0$. Comparison of $\mathcal{F}_{\text{corr}}$ with the uncorrelated effective branching fractions, $\mathcal{F}_{\text{unc}}$, also given in Table I allows us to extract $r^2$, $r \cos \delta$, $y$, and $x^2$. Information on $B_{\ell}$ is obtained from ST yields at the $\psi(3770)$ and from external measurements using incoherently-produced $D^0$ mesons. These two estimates of $B_{\ell}$ are averaged by the fitter to obtain $\mathcal{F}_{\text{unc}}$.

<table>
<thead>
<tr>
<th>Mode</th>
<th>C-odd</th>
<th>Uncorr.</th>
</tr>
</thead>
<tbody>
<tr>
<td>$K^-\pi^+$</td>
<td>1 + $R_{WS}$</td>
<td>1 + $R_{WS}$</td>
</tr>
<tr>
<td>$S_+$</td>
<td>2</td>
<td>2</td>
</tr>
<tr>
<td>$S_-$</td>
<td>2</td>
<td>2</td>
</tr>
<tr>
<td>$K^-\pi^+/K^+\pi^-$</td>
<td>$R_M$</td>
<td>$R_{WS}$</td>
</tr>
<tr>
<td>$K^-\pi^+/K^+\pi^-(1 + R_{WS})^2 - 4r \cos \delta(\cos \delta + y) + 1$</td>
<td>$1 + R_{WS}^2$</td>
<td></td>
</tr>
<tr>
<td>$K^-\pi^+/S_+$</td>
<td>$1 + R_{WS} + 2r \cos \delta + y$</td>
<td>$1 + R_{WS}$</td>
</tr>
<tr>
<td>$K^-\pi^+/S_-$</td>
<td>$1 + R_{WS} - 2r \cos \delta - y$</td>
<td>$1 + R_{WS}$</td>
</tr>
<tr>
<td>$K^-\pi^+/e^-$</td>
<td>$1 - ry \cos \delta - rx \sin \delta$</td>
<td>1</td>
</tr>
<tr>
<td>$S_+/S_+$</td>
<td>0</td>
<td>1</td>
</tr>
<tr>
<td>$S_-/S_-$</td>
<td>0</td>
<td>1</td>
</tr>
<tr>
<td>$S_+/S_-</td>
<td>4</td>
<td>2</td>
</tr>
<tr>
<td>$S_+/-e^-$</td>
<td>1 + $y$</td>
<td>1</td>
</tr>
<tr>
<td>$S_-/-e^-$</td>
<td>1 - $y$</td>
<td>1</td>
</tr>
</tbody>
</table>

3. Fit Inputs

We analyze 281 pb$^{-1}$ of $e^+e^-$ collision data produced by the Cornell Electron Storage Ring (CESR) at $E_{cm} = 3.77$ GeV and collected with the CLEO-c detector, which is described in detail elsewhere [19]. We reconstruct the $D^0$ and $\bar{D}^0$ final states listed in Table II with $\pi^0/\eta \to \gamma \gamma$, $\omega \to \pi^+\pi^-\pi^0$, and $K^0_S \to \pi^+\pi^-$. Signal and background efficiencies, as well as crossfeed probabilities among signal modes, are
determined from simulated events that are processed in a fashion similar to data.

Table II  D final states reconstructed in this analysis.

<table>
<thead>
<tr>
<th>Type</th>
<th>Final States</th>
</tr>
</thead>
<tbody>
<tr>
<td>Flavored</td>
<td>$K^-\pi^+, K^+\pi^-$</td>
</tr>
<tr>
<td>$S_+</td>
<td>K^+K^-, \pi^+\pi^-, K_S^0\pi^0, K_L^0\pi^0</td>
</tr>
<tr>
<td>$S_-$</td>
<td>$K_S^0\eta, K_S^0\eta, K_S^0\omega$</td>
</tr>
<tr>
<td>$e^+\nu$</td>
<td>Inclusive $Xe^+\nu, Xe^-\bar{\nu}$</td>
</tr>
</tbody>
</table>

Hadronic final states without $K_L^0$ mesons are fully reconstructed via two kinematic variables: the beam-constrained candidate mass, $M_{c} \equiv \sqrt{E_{b}^{2} - p_{D}^{2}}$, where $p_{D}$ is the $D^{0}$ candidate momentum and $E_b$ is the beam energy, and $\Delta E \equiv E_{D} - E_{b}$, where $E_{D}$ is the sum of the $D^{0}$ candidate daughter energies. We extract ST and DT yields from $M$ distributions using unbinned maximum likelihood fits (ST) or by counting candidates in signal and sideband regions (DT).

Because most $K_L^0$ mesons and neutrinos produced at CLEO-c are not detected, we only reconstruct modes with these particles in DTs, by demanding that the other $D$ in the event be fully reconstructed. Ref. [20] describes the missing mass technique used to identify $K_L^0\pi^0$ candidates. For semileptonic decays, we use inclusive, partial reconstruction to maximize efficiency, demanding that only the electron be identified with a multivariate discriminant [21] that combines measurements from the tracking chambers, the electromagnetic calorimeter, and the ring-imaging Čerenkov counter.

Table III gives yields and efficiencies for 8 ST modes and 58 DT modes, where the DT modes have been grouped into categories. Fifteen of the DT modes are forbidden by $CP$ conservation and are not included in the standard fit. In general, crossfeed among signal modes and backgrounds from other $D$ decays are smaller than 1%. Modes with $K_S^0\pi^0\pi^0$ have approximately 3% background, and yields for $K^{+}\pi^{-}/K^{+}\pi^{-}$ and $S_{+}/S_{-}$ are consistent with being entirely from background.

External inputs to the fit include measurements of $R_M$, $R_WS$; $B_{K^+\pi^{-}}$, and $B_{S_{+}}$, as well as an independent $B_{K^0_S\pi^0}$ from CLEO-c, as shown in Table IV. The external $R_WS$ is required to constrain $r^2$, and thus, to determine $\cos \delta$ from $r \cos \delta$. We also use the external mixing parameter measurements shown in Table V. The fit incorporates the full covariance matrix for these inputs, accounting for statistical overlap with the yields in this analysis. Covariance matrices for the fits in Ref. [27] have been provided by the CLEO, Belle, and BABAR collaborations.

Table IV Averages of external measurements used in the standard fit. Charge-averaged $D^0$ branching fractions are denoted by final state.

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Average</th>
</tr>
</thead>
<tbody>
<tr>
<td>$R_{WS}$</td>
<td>0.00409 ± 0.00022 [22]</td>
</tr>
<tr>
<td>$R_{M}$</td>
<td>0.00017 ± 0.00039 [23]</td>
</tr>
<tr>
<td>$K^-\pi^+$</td>
<td>0.0381 ± 0.0009 [24]</td>
</tr>
<tr>
<td>$K^+\pi^-/K^-\pi^+$</td>
<td>0.1010 ± 0.0016 [25]</td>
</tr>
<tr>
<td>$\pi^-\pi^+/K^-\pi^+$</td>
<td>0.0359 ± 0.0005 [25]</td>
</tr>
<tr>
<td>$K_S^0\pi^0$</td>
<td>0.0097 ± 0.0003 [20]</td>
</tr>
<tr>
<td>$K_S^0\pi^0$</td>
<td>0.0115 ± 0.0012 [24]</td>
</tr>
<tr>
<td>$K_S^0\eta$</td>
<td>0.00380 ± 0.00060 [24]</td>
</tr>
<tr>
<td>$K_S^0\omega$</td>
<td>0.0130 ± 0.0030 [24]</td>
</tr>
</tbody>
</table>

Table V Averages of external measurements used in the standard and extended fits.

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Average</th>
</tr>
</thead>
<tbody>
<tr>
<td>$y$</td>
<td>0.00062 ± 0.00021 [2, 25, 26]</td>
</tr>
<tr>
<td>$x$</td>
<td>0.00811 ± 0.00334 [26]</td>
</tr>
<tr>
<td>$r^2$</td>
<td>0.00339 ± 0.00012 [27]</td>
</tr>
<tr>
<td>$y'$</td>
<td>0.0034 ± 0.0030 [27]</td>
</tr>
<tr>
<td>$x'^2$</td>
<td>0.00006 ± 0.00018 [27]</td>
</tr>
</tbody>
</table>

Systematic uncertainties associated with efficiencies for reconstructing tracks, $K_S^0$ decays, $\pi^0$ decays, and
for hadron identification are assigned as described in Ref. [29]. Other sources of efficiency uncertainty include: $\Delta E$ requirements (0.5–5.5%), $\eta$ reconstruction (4.0%), electron identification (1.0%), modeling of particle multiplicity and detector noise (0.1–1.3%), simulation of initial and final state radiation (0.5–1.2%), and modeling of resonant substructure in $K_S^0\pi^0$ (0.7%). We also include additive uncertainties of 0.0–0.9% to account for variations of yields with fit function.

These systematic uncertainties are included directly in the covariance matrix given to the fitter, which propagates them to the fit parameters. The other fit inputs determined in this analysis are ST and DT yields and efficiencies, crossfeed probabilities, background branching fractions and efficiencies, and statistical uncertainties on all of these measurements. Quantum correlations between signal and background modes are accounted for using assumed values of amplitudes ratios and strong phases that are systematically varied and found to have negligible effect. We validated our analysis technique in a simulated $C$-odd $D^0\bar{D}^0$ sample 15 times the size of our data sample.

4. Preliminary Fit Results

Our standard fit excludes the 15 same-CP DT modes and includes the measurements in Table IV but not Table V. In this fit, there is not enough information to reliably determine $x\sin\delta$, so we fix it to zero, and the associated systematic uncertainty is negligible. We obtain a first measurement of $\cos\delta = 1.03 \pm 0.19 \pm 0.08$, consistent with being at the boundary of the physical region. The fit results for $y$, $r^2$, $x^2$, and branching fractions are consistent with previous measurements.

The likelihood curve for $\cos\delta$, computed as $\mathcal{L} = e^{-(x^2 - x_{\text{min}}^2)/2}$ and shown in Figure 1 is slightly non-Gaussian. For values of $|\cos\delta| < 1$, we also show $\mathcal{L}$ as a function of $|\sin\delta|$. We integrate these curves within the physical region to obtain 95% confidence level limits of $\cos\delta > 0.54$ and $|\sin\delta| < 0.72$.

We also perform an extended fit that includes the previous measurements of $y$ and $y'$ in Table V in addition to all the inputs to the standard fit above. In this fit, we find $\cos\delta = 0.93 \pm 0.32 \pm 0.04$. The systematic uncertainty does not include the contribution from assuming $x\sin\delta = 0$, which is still under study. From the corresponding likelihood functions shown in Figure 1, we determine 95% confidence level limits of $\cos\delta > 0.38$ and $|\sin\delta| < 0.84$.

The $\cos\delta$ uncertainty in the extended fit is larger than in the standard fit because of a non-linear effect. Most of the information on $r^2$ (and therefore on $r$) is provided by $R_{\text{WS}}$. Because $R_{\text{WS}}$ also depends on $y \cdot r \cos\delta$, the sign of the correlation between $r^2$ and $r \cos\delta$ is given by the sign of $y$. In the standard fit, $y$ attains a more negative central value than in the extended fit, where $y$ is constrained to the precise external measurements. Hence, the uncertainty on $\cos\delta$ becomes inflated in the extended fit.

By observing the change in $1/\sigma^2_{\cos\delta}$ as each fit input is removed, we identify the major contributors of information about $\cos\delta$ to be the $K\pi/S_\pm$ DT yields and the ST yields. We also find that no single input or group of inputs exerts a pull larger than three standard deviations on $\cos\delta$ or $y$. Moreover, removing all external inputs gives branching fractions consistent with those in Table V.

We also allow for a $C$-even $D^0\bar{D}^0$ admixture in the initial state, which is expected to be $O(10^{-8})$ [30], by including the 15 $S_\pm/S_\pm$ DT yields in the fit. These modes limit the $C$-even component, which can modify the other yields as described in Ref. [10]. In both the standard and extended fits, we find a $C$-even fraction consistent with zero with an uncertainty of 2.4%, and neither the fitted $\cos\delta$ values nor their uncertainties are shifted noticeably from the results quoted above.

5. Summary

Using 281 pb$^{-1}$ of $e^+e^-$ collisions produced at the $\psi(3770)$, we make a preliminary first determination of the strong phase $\delta$, with $\cos\delta = 1.03 \pm 0.19 \pm 0.08$. By
further including external mixing parameter measurements in our analysis, we obtain an alternate measurement of $\cos \delta = 0.93 \pm 0.32 \pm 0.04$, where the systematic uncertainty from assuming $x \sin \delta = 0$ has not been included. Knowledge of $\delta$ allows independent measurements of $y$ and $y'$ to be combined, thereby improving our overall knowledge of charm mixing parameters.

Acknowledgments

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References

HFAG Charm Mixing Averages

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Recently the first evidence for charm mixing has been reported by several experiments. To provide averages of these mixing results and other charm results, a new subgroup of the Heavy Flavor Averaging Group has been formed. We here report on the method and results of averaging the charm mixing results.

1. Introduction

Almost since the discovery of charm mesons, mixing of $D^{0} \rightarrow \bar{D}^{0}$ mesons have been sought in analogy to the well-known $K^{0} \rightarrow \bar{K}^{0}$ mixing. Due to very effective GIM suppression, the expected mixing rate in the charm system is much smaller than for kaons. Only very recently, the BaBar [1] and Belle [2] collaborations have reported the first evidence of charm mixing\(^1\). These results have renewed the interest from the theory community as the observed mixing rate could be caused by physics beyond the standard model or at least provide additional constraints on new physics.

None of the mixing measurement have a significance above four standard deviations, but several have similar precision for the mixing parameters. By combining the measurements we therefore obtain more precise values for the mixing parameters and exclude the non-mixing hypothesis with larger confidence. Combining the different mixing measurements is not completely straightforward, since not all measurements are sensitive to the same charm mixing parameters.

The Heavy Flavor Averaging Group (HFAG) in 2006 created a subgroup with the responsibility of providing averages of charm physics measurements. One of the high priority tasks of this group is to combine the charm mixing measurements into world-average values for the fundamental mixing parameters. The first average assuming CP conservation was shown at FPCP [3]. Besides those results, we here report the first results of combining mixing measurements where we allow for CP violation.

2. Averaging Method

Mixing is present in the $D^{0} \rightarrow \bar{D}^{0}$ system if the mass eigenstates, $|D_{1}\rangle$ and $|D_{2}\rangle$, differ from the flavor eigenstates, $|D^{0}\rangle$ and $|\bar{D}^{0}\rangle$. Generally one can write $|D_{1,2}\rangle = p|D^{0}\rangle \pm q|\bar{D}^{0}\rangle$. The variables of fundamental interest are the mass difference, $\Delta M = M_{1} - M_{2}$ and decay width difference, $\Delta \Gamma = \Gamma_{1} - \Gamma_{2}$ between the two mass eigenstates. Traditionally, in charm mixing one uses the dimensionless variables, $x = \Delta M / \Gamma$ and $y = \Delta \Gamma / 2 \Gamma$, where $\Gamma$ is the average decay width. CP violation in mixing or in the interference between mixing and decay would manifest itself as $|q/p| \neq 1$ and $\phi = \arg(q/p) \neq 0$, respectively\(^2\). In addition CP violation could show up in the decay itself giving rise to decay mode dependent parameters.

Most measurements do not directly measure $(x, y)$. For instance in mixing measurements using $D^{0} \rightarrow K^{+}\pi^{-}$ decays there is a unknown strong phase, $\delta_{K\pi}$, so the results obtained are for $x' = x \cos \delta_{K\pi} + y \sin \delta_{K\pi}$ and $y' = -x \sin \delta_{K\pi} + y \cos \delta_{K\pi}$. In the averaging procedure, we first combine measurements of the same parameters to obtain the more precise observables. Most measurements are performed using likelihood fits and the combination is therefore performed by multiplying likelihood functions from each measurement and finding the new maximum. By combining likelihoods, correlations between observables and possible non-Gaussian tails are taken into account. For measurements which are not using likelihoods, we construct a likelihood using symmetrized Gaussian uncertainties. To combine different types of measurements, the different combined likelihoods are recalculated as a function of $(x, y, \delta_{K\pi})$ minimizing over any other variables. $\delta_{K\pi}$ is included since there is both a direct measurement [4] and by combining the $D^{0} \rightarrow K^{+}\pi^{-}$ measurement with the other measurement of $x$ and $y$, one can also get a precise measurement of $\delta_{K\pi}$. When plotting confidence contours for $(x, y)$ we minimize the likelihood over $\delta_{K\pi}$.

The combining of likelihood functions is currently only done for the CP conserving case. In principle it can be done also for the CP violating case by simply having two more variables, $|q/p|$ and $\phi$, in the final likelihood function. Unfortunately not all likelihoods are currently available for the measurements which allow for CP violation. A simple combination

\(^1\)Shortly after the CHARM2007 workshop additional results with evidence for charm mixing has been reported by the BaBar and CDF collaborations. In these proceedings we will summarize the status at the time of the workshop.

\(^2\)The phase $\phi$ is for the moment assumed to be independent of decay mode.
is therefore performed by forming a $\chi^2$ of all measurements expressed in terms of the fundamental mixing parameters. The $\chi^2$ assumes Gaussian errors, but correlations between observables in each individual measurement is taken account by using the full covariance matrix for each result.

3. CP Conserving Averages

The following averages were performed by adding log likelihoods from fits where CP conservation was assumed.

3.1. Lifetime Ratio Average

One can observe charm mixing by finding a difference in the lifetime measured in decays to CP eigen states such as $D^0 \rightarrow K^+ K^-$ and $D^0 \rightarrow \pi^+ \pi^-$ and the mixed-CP decay $D^0 \rightarrow \pi^+ K^-$. We combine six results [2, 3, 6, 8, 9] from such analyzes. All of these measure $y_{CP} = \tau_{K\pi}/\tau_{hh} - 1$. In the limit of CP conservation one has $y_{CP} = y$. The average of the six measurements is $y_{CP} = (1.12 \pm 0.32) \times 10^{-2}$. This is $3.5\sigma$ from the no-mixing hypothesis. As can be seen from Figure 1, this average is mainly driven by the recent Belle measurement.

3.2. Mixing Rate Average

Wrong-signed semileptonic decays provide a clean way of searching for charm mixing, but the measure-

ments are only sensitive to the integrated mixing rate $R_M = (x^2 + y^2)/2$. Four measurements [10, 11, 12, 13] are combined and give an average of $R_M = (1.7 \pm 3.9) \times 10^{-4}$. In addition to the semileptonic decays, $R_M$ can also measured in the analysis of fully hadronic decays. The semileptonic result is therefore combined from two hadronic analyzes [14, 15] and in addition an analysis of tagged decays at the $\psi(3770)$ [4]. The combination is illustrated in Figure 2 and gives an average value of $R_M = (2.1 \pm 1.1)^{-4}$. In the transformation to a likelihood in $(x, y)$, we ignore the non-physical region of $R_M < 0$.

![Figure 2: The mixing rate from measurements using semileptonic $D^0$ decays are averaged with results from multi-body hadronic charm decays.](image)

3.3. $(x, y)$ Average

One can measure $x$ and $y$ directly using a time-dependent Dalitz plot analysis of $D^0 \rightarrow K_S^0 \pi^+ \pi^-$ decays. Two measurements [16, 17] have been published and these have been averaged by HFAG and gives $x = (8.1 \pm 3.3) \times 10^{-3}$ and $y = (3.1 \pm 2.8) \times 10^{-3}$. Combining this average with the averages above for $R_M$ and $y_{CP}$ using likelihoods mapped as a function of $(x, y)$ we obtain $x = (9.2 \pm 3.4) \times 10^{-3}$ and $y = (7.0 \pm 2.2) \times 10^{-3}$. Contours of the combined likelihood function at the levels corresponding to 1 to $3\sigma$ confidence levels are shown in Figure 3. Note that the confidence levels shown correspond to two-dimensional coverage probabilities of 68.27%, 95.45%, etc., and therefore $2\Delta \ln L = 2.30, 6.18$, etc.
3.4. Averages for $D^0 \rightarrow K^+\pi^-$ Decays

As mentioned above, one can measure $x'$ and $y'$ using the doubly-Cabibbo suppressed (DCS) decay $D^0 \rightarrow K^+\pi^-$. The likelihood functions are available for two measurements [1,18] of this type. These are combined and gives the averages $x'^2 = (-0.1 \pm 2.0) \times 10^{-4}$ and $y' = (5.5^{+2.8}_{-3.7}) \times 10^{-3}$. The corresponding likelihood contours are shown in Figure 4.

3.5. World Average

The combined likelihood for $D^0 \rightarrow K^+\pi^-$ decays can be expressed as a function of $(x, y, \delta_{K\pi})$ ignoring the part with $x'^2 < 0$. This likelihood can be combined with the likelihood from the combination of the other mixing results in Section 3.3 which do not depend on $\delta_{K\pi}$. An additional constraint comes from a CLEO-c measurement [2] of $\cos \delta_{K\pi} = 1.09 \pm 0.66$, where a small dependence on $x$ and $y$ is ignored in the combination. Figure 6 shows the likelihood contours in $(x, y)$ after minimizing over $\delta_{K\pi}$. The region around the central value is almost unchanged with respect to the result without the $D^0 \rightarrow K^+\pi^-$ decays (Figure 5). This is also reflected in the over all average for $x$ and $y$ which are

$$x = (8.7^{+3.0}_{-3.4}) \times 10^{-3},$$

$$y = (6.6 \pm 2.1) \times 10^{-3}.$$

The $D^0 \rightarrow K^+\pi^-$ measurements do not contribute much to the central value, because of the poorly known phase $\delta_{K\pi}$. However they do help exclude the no-mixing hypothesis and cause the dip seen in the contours close to $(x, y) = (0, 0)$. At $(x, y) = (0, 0)$ we obtain $2\Delta \ln L = 37$ with respect to the minimum. This corresponds to a significance of the combined mixing signal of 5.7$\sigma$.

4. CP Violating Averages

Measurements of charm mixing can be done without assuming CP conservation by fitting $D^0$ and $\bar{D}^0$ mesons as separate samples. Most of the measurements above have done that and we therefore can combine those to also provide constraints on the CP violating parameters. When allowing for CP violation,
the measured parameters are related slightly differently to the mixing parameters. For the lifetime ratio measurements, one has
\[
2y_{CP} = (|q/p| + |p/q|) \cos \phi - (|q/p| - |p/q|) x \sin \phi,
\]
\[
2A_{\Gamma} = (|q/p| - |p/q|) y \cos \phi - (|q/p| + |p/q|) x \sin \phi,
\]
where \(A_{\Gamma}\) is the measured relative lifetime difference for \(D^0 \rightarrow h^+h^-\) and \(D^0 \rightarrow h^+h^-\). For \(D^0 \rightarrow K^+\pi^-\) decays, the \(x'\) and \(y'\) measured for \(D^0\) and \(D^0\) are related as follows
\[
x'^\pm = \left( \frac{1 \pm A_M}{1 \mp A_M} \right)^{1/4} (x' \cos \phi \pm y' \sin \phi),
\]
\[
y'^\pm = \left( \frac{1 \pm A_M}{1 \mp A_M} \right)^{1/4} (y' \cos \phi \mp x' \sin \phi),
\]
where \(A_M = \frac{|q/p|^2 - |p/q|^2}{|q/p|^2 + |p/q|^2}\). For \(D^0 \rightarrow K_S^0 \pi^+\pi^-\) decays the measurement directly gives \(x, y, |q/p|\) and \(\phi\), while for the \(R_M\) analysis the results are not separated and therefore just measure \(R_M = (x^2 + y^2)/2\).

The measurement of \(\delta_{K\pi}\) from CLEO-c is not done separately for \(D^0\) and \(D^0\) mesons and is not included in the combined result allowing for CP violation.

In total 22 measurements are combined in a \(\chi^2\)-fit to extract seven parameters, the four mixing and CP violation parameters, \(x, y, |q/p|\) and \(\phi\), and three characterizing \(D^0 \rightarrow K^+\pi^-\), namely \(\delta_{K\pi}\), the DCS rate \(R_D\), and the direct decay rate asymmetry \(A_D\). The fit gives \(\chi^2 = 14.4\) and the following mixing parameters
\[
x = (8.4^{+3.2}_{-3.4}) \times 10^{-3},
\]
\[
y = (6.9 \pm 2.1) \times 10^{-3},
\]
\[
|q/p| = 0.88^{+0.23}_{-0.20},
\]
\[
\phi = (-0.09 \pm 0.17) \text{ rad}.
\]

The mixing parameters are almost unchanged with respect to the CP conserving average. This is also seen from the confidence levels shown in Figure 7. One can also draw the 1 to 5\(\sigma\) confidence level contour for \(\phi\) versus \(|q/p|\) using \(\Delta \chi^2\). This is shown in Figure 8. The no-CP violation hypothesis is seen to lie well within the 1\(\sigma\) contour.

The combined results for the \(D^0 \rightarrow K^+\pi^-\) parameters are
\[
\delta_{K\pi} = 0.33^{+0.26}_{-0.29} \text{ rad},
\]
\[
R_D = (3.35 \pm 0.11) \times 10^{-3},
\]
\[
A_D = (-0.8 \pm 3.1).
\]

There is little change in \(\delta_{K\pi}\) with respect to the CP conserving average and no evidence for direct CP violation as \(A_D\) is consistent with zero.

5. Summary

Evidence of charm mixing has been reported from several experiments in the last year. A new subgroup of HFAG has performed an average of these and other existing charm mixing results. The combined result has a signal significance in excess of 5 standard deviations and gives the mixing parameters
\[
x = (8.4^{+3.2}_{-3.4}) \times 10^{-3},
\]
\[
y = (6.9 \pm 2.1) \times 10^{-3}.
\]
CP violation parameters have also been combined and gives

\[ |q/p| = 0.88^{+0.23}_{-0.20}, \]
\[ \phi = (-0.09^{+0.17}_{-0.19}) \text{rad}. \]

This is fully consistent with no CP violation being present in charm mixing. HFAG intends to periodically update these averages as new results become available in order to provide the most precise mixing parameters to the community.

References


Charm Mixing - Theory

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We discuss Standard Model (SM) and New Physics (NP) descriptions of $D^0$ mixing. The SM part of the discussion addresses both quark-level and hadron-level contributions. The NP part describes our recent works on the rate difference $\Delta \Gamma_D$ and the mass difference $\Delta M_D$. In particular, we describe how the recent experimental determination of $\Delta M_D$ is found to place tightened restrictions on parameter spaces for 17 of 21 NP models considered in a recent paper by Hewett, Pakvasa, Petrov and myself.

I. INTRODUCTION

Given the forthcoming operation of the LHC, perhaps the dominant role of experimental flavor studies in particle physics will be supplanted by discoveries in the so-called new physics. Even if flavor physics faces an unsure future, all would acknowledge its remarkable recent progress via the observation of rare phenomena such as CP-violation in $B$-mesons or particle-antiparticle mixing for $B_s$ and $D^0$ mesons. If new physics is indeed observed, the continued exploration of rare observables could well be an asset in deciphering exotic LHC events.

My purpose in this talk is to describe two recent theoretical contributions to $D^0$ mixing [1,2,22]. Both Ref. [1] and Ref. [2] should be considered in the context of the recent HFAG values [3].

$$x_D \equiv \frac{\Delta M_D}{\Gamma_D} = (8.4^{+3.2}_{-2.4}) \cdot 10^{-3}$$

$$y_D \equiv \frac{\Delta \Gamma_D}{2\Gamma_D} = (6.9 \pm 2.1) \cdot 10^{-3} .$$

In light of the Physical Review Letters criteria of ‘observation’ (> 5σ) or ‘evidence’ (3σ-to-5σ), we see that the above 2.4σ determination for $x_D$ amounts to a ‘measurement’ (< 3σ). As such, we all await improvements in sensitivity for charm mixing.

The observed signal is seen to occur at about the 1% level. Whether or not this is the magnitude expected for the SM signal is a topic I will discuss shortly. At any rate, I wish to also consider the possibility of a NP component in $D^0$ mixing amplitude,

$$M_{\text{mix}} = M_{\text{SM}} + M_{\text{NP}} .$$

The relative phase between $M_{\text{SM}}$ and $M_{\text{NP}}$ is not known. Thus, in our detailed study of various NP contributions to $x_D$ in Ref. [2] we most often compared the NP predictions to $\pm 1\sigma, \pm 2\sigma$ windows relative to the central $x_D$ value of Eq. [1].

### A. Operator Product Expansion (OPE) and Renormalization Group

An important technical aspect of Refs. [1,2] is the process of relating an amplitude at some NP scale $\mu = M$ to one at, say, the charm scale $\mu = m_c$. This takes the form

$$\langle f | \mathcal{H}_{\text{NP}} | i \rangle = G \sum_{i=1} C_i(\mu) \langle f | Q_i | i \rangle(\mu) ,$$

where the prefactor $G$ has the dimension of inverse-squared mass, the $C_i$ are dimensionless Wilson coefficients, and the $Q_i$ are the effective operators. At the leading order of dimension six, it turns out that there are eight four-quark operators,

$$Q_1 = (\bar{\pi} L \gamma_\mu c L) (\bar{\pi} L \gamma^\mu c L) ,$$

$$Q_2 = (\bar{\pi} L \gamma_\mu c L) (\bar{\pi} R \gamma^\mu c R) ,$$

$$Q_3 = (\bar{\pi} L c R) (\bar{\pi} R c L) ,$$

$$Q_4 = (\bar{\pi} R c L) (\bar{\pi} L c R) ,$$

$$Q_5 = (\bar{\pi} R \sigma^{\mu \nu} c L) (\bar{\pi} R \sigma_{\mu \nu} c L) ,$$

$$Q_6 = (\bar{\pi} R \gamma_\mu c R) (\bar{\pi} R \gamma^\mu c R) ,$$

$$Q_7 = (\bar{\pi} L c R) (\bar{\pi} L c R) ,$$

$$Q_8 = (\bar{\pi} L \sigma^{\mu \nu} c R) (\bar{\pi} L \sigma_{\mu \nu} c R) .$$

Any given NP contribution will often involve several of these, but in all events never more than these eight. The evolution is determined by solving the RG equations obeyed by the Wilson coefficients,

$$\frac{d}{d \log \mu} \tilde{C}(\mu) = \tilde{\gamma} T \tilde{C}(\mu) ,$$

where $\tilde{\gamma}$ is the $8 \times 8$ anomalous dimension matrix [4]. The output of this calculation is a set of RG factors $r_i(\mu, M)$ which are expressed in terms of ratios of QCD fine structure constants evaluated at different scales, e.g. as with

$$r_1(\mu, M) = \left( \frac{\alpha_s(M)}{\alpha_s(m_i)} \right)^{2/7} \left( \frac{\alpha_s(m_i)}{\alpha_s(m_b)} \right)^{6/25} \left( \frac{\alpha_s(m_b)}{\alpha_s(\mu)} \right)^{6/25} .$$

### B. Operator Matrix Elements

One needs ultimately to evaluate the $D^0$-to-$\bar{D}^0$ matrix elements of the eight operators $\{ Q_i \}$. In general, eight non-perturbative parameters would need to be...
evaluated by some means such as a lattice determination. As a practical matter, the method used in Refs. \[1, 2\] is to introduce a ‘modified vacuum saturation’ (MVS), where all such matrix elements are written in terms of the known matrix elements of (V-A) × (V-A) and (S-P) × (S+P) matrix elements $B_D$ and $F_D^{(S)}$ \[3\],

$$
(Q_1) = \frac{2}{3} f_D^2 M_D^2 B_D , \\
(Q_2) = -\frac{1}{2} f_D^2 M_D^2 B_D - \frac{1}{N_c} f_D^2 M_D^2 B_D^{(S)} , \\
(Q_3) = \frac{1}{4N_c} f_D^2 M_D^2 B_D + \frac{1}{2} f_D^2 M_D^2 B_D^{(S)} , \\
(Q_4) = -\frac{2N_c-1}{4N_c} f_D^2 M_D^2 B_D^{(S)} , \\
(Q_5) = \frac{3}{N_c} f_D^2 M_D^2 B_D^{(S)} , \\
(Q_6) = (Q_1) , \\
(Q_7) = (Q_4) , \\
(Q_8) = (Q_5) ,
$$

where the number of colors is $N_c = 3$ and, as in Ref. \[3\], we define

$$
F_D^{(S)} \equiv B_D^{(S)} \cdot \frac{M_D^2}{(m_c + m_u)^2} .
$$

With the above theoretical machinery in hand, we are now ready to consider SM and NP contributions to $D^0$ mixing.

### II. STANDARD MODEL ANALYSIS

One can use quarks or hadrons as the basic degrees of freedom in carrying out the SM analysis of $D^0$ mixing. In principle, these should give the same result. However, as we shall see, rather different features appear in each description.

#### A. Quark-level Analysis

At leading order in the SM, the OPE for $D^0$ mixing consists of two dimension-six four-quark operators \[2\]. The next order contains fifteen dimension-nine six-quark operators. For each increasing order in the OPE, there are still more local quark and gluon operators and the problem of determining operator matrix elements becomes ever more severe. For this reason, the dimension six sector has received by far the most attention.

The dimension six amplitude is depicted in Fig. 1. Since the $b$-quark is essentially decoupled due to the tiny $V_{ub}$ value, only the light $d, s$ quarks propagate in the loop. The Cabibbo dependence of this diagram, $\sin^2 \theta_c$, itself seems to suggest that the experimental signal (near the 0.01 level) is easily understood. But not so fast! For convenience, let us set $m_d = 0$. Then the only mass ratio that appears in the problem is

$$
z \equiv (m_s/m_c)^2 \simeq 0.006 .
$$

Table II examines one of the loop-functions for $\Delta \Gamma_D$ and shows the results of carrying out an expansion in powers of $z$. We see that the contributions of the individual intermediate states in the mixing diagram are not intrinsically small – in fact, they begin to contribute at $O(z^0)$. However, flavor cancellations remove all contributions through $O(z^2)$ for $\Delta \Gamma_D$, so the net result is $O(z^3)$. Charm mixing clearly experiences a remarkable GIM suppression!

We understand the reason for this. $D^0$ mixing vanishes in the limit of exact SU(3) flavor symmetry. It is nonzero only because flavor SU(3) is broken, and indeed, $D^0$ mixing occurs at second order in SU(3) breaking \[8\]. A factor of $z$ will accompany each order of SU(3) breaking and the rate difference $y_D$ will experience an additional factor of $z$ due to helicity suppression.

\begin{table}[h]
\centering
\begin{tabular}{|c|c|c|c|}
\hline
Intermediate State & $O(z^0)$ & $O(z^1)$ & $O(z^2)$ \\
\hline
$s\bar{s}$ & 1/2 & $-3z$ & $3z^2$ \\
$d\bar{d}$ & 1/2 & 0 & 0 \\
$s\bar{d} + d\bar{s}$ & -1 & $3z$ & $-3z^2$ \\
\hline
Total & 0 & 0 & 0 \\
\hline
\end{tabular}
\caption{Flavor cancellations in $\Delta \Gamma_D$.}
\end{table}

Of course, this is just the leading order (LO) result in QCD, and we should consider the next-to-leading order result as well,

$$
x_D = x_D^{(LO)} + x_D^{(NLO)} , \\
y_D = y_D^{(LO)} + y_D^{(NLO)} .
$$

This has been done in Ref. \[6\] and the results are summarized in Table II, which reveals that $y_D$ is given by $y_D^{NLO}$ to a reasonable approximation (due to the removal of helicity suppression by virtual gluons) whereas $x_D$ is greatly affected by destructive interference between $x_D^{LO}$ and $x_D^{NLO}$. The net effect is to render $y_D$ and $x_D$ of similar small magnitudes, at
least through this order of analysis, as compared to the experiment signal.

It is not inconceivable that the quark-level prediction of $x_D$ and $y_D$ just described might be considerably affected by a higher order in the OPE \cite{2} which suffers less $z$ suppression. Simple dimensional analysis \cite{10} suggests the magnitudes $x_D \sim y_D \sim 10^{-3}$ might be achievable, although order-of-magnitude cancellations or enhancements are possible.

B. Hadron-level Analysis

Most of the work involving the hadron degree-of-freedom has been done on $y_D$. One starts with the following general expression for $\Delta \Gamma_D$,

$$\Delta \Gamma_D = \frac{1}{M_D} \Im I$$ (11)

$$I \equiv \langle \bar{D}^0 \mid i \int d^4x T \{ \mathcal{H}_w^{[\Delta C=1]}(x) \mathcal{H}_w^{[\Delta C=1]}(0) \} \mid D^0 \rangle .$$

To utilize this relation, one inserts intermediate states between the $[\Delta C]=1$ weak hamiltonian densities $\mathcal{H}_w^{[\Delta C=1]}$. Although this can be done using either quark or hadron degrees of freedom, let us consider the latter here. Clearly, some knowledge of the matrix elements $\langle \eta | \mathcal{H}_w^{[\Delta C=1]} | D^0 \rangle$ is required.

One approach is to model $[\Delta C]=1$ decays theoretically and fit various model parameters to charm decay data. Some time ago, $\Delta \Gamma_D$ was determined in this manner and the result $y_D \approx 10^{-3}$ was found \cite{11}. This value is smaller than the recent BaBar and Belle central values.

Alternatively, one can arrange for charm decay data to play a somewhat different role. The earliest work in this regard focused on the $P^+P^- = \pi^+\pi^-, K^+K^-, K^-\pi^+, K^+\pi^-\pi^+$ states \cite{12,13}. In the flavor SU(3) limit, this subset of states gives zero contribution due to cancellations. But SU(3) breaking had already been known to be significant in individual charm decays. Since the study of charm decays in the 1980’s lacked an abundance of data, these references could only conclude that $y_D$ might be large.

A modern version of this approach now exists, although the analysis takes an unexpected direction \cite{3}. Since SU(3) breaking occurs at second order in $D^0$ mixing, let us hypothesize that the contribution of the $P^+P^-$ sector is in fact negligible due to flavor cancellations. Likewise for all other sectors whose decays are kinematically allowed. However, this cannot be true for four-pseudoscalars because decay into four-kaon states is kinematically forbidden. In Ref. \cite{8} it is estimated that these ‘kinematically-challenged’ sectors can provide enough $SU(3)$ violation to induce $y_D \sim 10^{-2}$. I personally find such an argument to be an important advance in our understanding of the subject. At the same time, it is unfortunately more persuasive than compelling due to the uncontrollable uncertainties inherent in this line of reasoning.

To summarize, we have just described how the observed $D^0$ mixing signal could well arise from SM physics, but the associated numerical prediction is seen to be lacking in precision. This conceivably leaves room for some NP mechanism to co-contribute or even dominate the SM signal. In the following we consider in turn NP analyses of the width difference $y_D$ and the mass difference $x_D$. 

III. NP AND THE WIDTH DIFFERENCE

At first glance, it would appear unlikely that NP could affect $y_D$ because the particles contributing to the loop amplitude of Fig. 1 must be on-shell. Since NP particles will be heavier than the charm mass, ‘there can be no NP contribution to $y_D$’. Or so goes the argument.

However, as explained in Ref. \cite{1}, NP effects in $\mathcal{H}_w^{[\Delta C=1]}$ can generally contribute to $y_D$. In the loop amplitude of Fig. 1 the NP contribution (empirically small for $\Delta C = -1$ processes) arises from either of the two vertices. We represent the NP $\Delta C = -1$ hamiltonian as (indices $i, j, k, \ell$ represent color),

$$\mathcal{H}_{NP}^{\Delta C=-1} = \sum_{q,q'} D_{qq'} \left[ C_1^{\ast}(\mu) O_1 + C_2^{\ast}(\mu) O_2 \right] ,$$

$$O_1 = \bar{u}_i \Gamma_1 q_j \bar{\tau}_j \Gamma_2 c_i ,$$

$$O_2 = \bar{u}_i \Gamma_1 q'_j \bar{\tau}_j \Gamma_2 c_j ,$$ (12)

where $D_{qq'}$ and the spin matrices $\Gamma_{1,2}$ encode the NP model. $\bar{C}_{1,2}(\mu)$ are Wilson coefficients evaluated at energy scale $\mu$ and the flavor sums on $q, q'$ extend over the $d, s$ quarks.

This leads to a prediction for the NP contribution to $y_D$. For a generic NP interaction, one finds (with the number of colors $N_c = 3$)

$$y_D = -\frac{4\sqrt{2} G_F}{M_D \Gamma_D} \sum_{q,q'} V_{cq}^{\ast} V_{aq} D_{qq'} (K_1 \delta_{i\delta_j} \delta_{\ell})$$

$$+ K_2 \delta_{i\delta_j} \delta_{\ell}^5 \sum_{a=1}^5 I_a(x, x') \langle D^0 | O^{i\delta_j \ell}_a | D^0 \rangle ,$$ (13)

where $\{K_a\}$ are combinations of Wilson coefficients,

$$K_1 = (C_1 \bar{C}_1 N_1 + (C_1 \bar{C}_2 + \bar{C}_1 C_2)) ,$$

$$K_2 = C_2 \bar{C}_2 ,$$ (14)
and the $\{O^{ijkl}\}$ are four-quark operators written down in Ref. [1]. Numerical results for some NP models are displayed in Table III.

<table>
<thead>
<tr>
<th>Model</th>
<th>$\gamma_D$</th>
<th>Comment</th>
</tr>
</thead>
<tbody>
<tr>
<td>RPV-SUSY</td>
<td>$6 \cdot 10^{-6}$</td>
<td>Squark Exchange</td>
</tr>
<tr>
<td></td>
<td>$-4 \cdot 10^{26}$</td>
<td>Slepton Exchange</td>
</tr>
<tr>
<td>Left-right</td>
<td>$-5 \cdot 10^{-6}$</td>
<td>‘Manifest’</td>
</tr>
<tr>
<td></td>
<td>$-8.8 \cdot 10^{-5}$</td>
<td>‘Nonmanifest’</td>
</tr>
<tr>
<td>Multi-Higgs</td>
<td>$2 \cdot 10^{-10}$</td>
<td>Charged Higgs</td>
</tr>
<tr>
<td>Extra Quarks</td>
<td>$10^{-8}$</td>
<td>Not Little Higgs</td>
</tr>
</tbody>
</table>

One sees that the entries, aside from R-parity violating SUSY, produce small contributions. We emphasize, however, that Eq. (12) and Eq. (13) represent general formulae for the contribution of all NP models of $|\Delta C| = 1$ interactions, encompassing those not included in Table III.

### IV. NP AND THE MASS DIFFERENCE

As the operation of the LHC looms near, the number of potentially viable NP models has never been greater. In this section, I will give an overview of Ref. [2], whose hallmark is the study of many (21 in all) NP models. Perhaps the best way to start is to consider the different ways that ‘extras’ can be added to the SM:

- Extra gauge bosons (LR models, etc)
- Extra scalars (multi-Higgs models, etc)
- Extra fermions (little Higgs models, etc)
- Extra dimensions (universal extra dims., etc)
- Extra global symmetries (SUSY, etc)

Although this approach does not provide a totally clean partition of NP models (e.g. obviously SUSY contains extra particles appearing in other categories), it proved useful to the authors of Ref. [2].

The broad menu of NP models which were analyzed is listed in Table IV. The extensive content of this list (e.g. there are four different SUSY realizations and three involving large extra dimension) indicates how rich the field of NP models has become. Of course, the subject of NP is by now fairly mature (in preparing this talk, I realized that my first paper on charged Higgs bosons [14] was written nearly 30 years ago!) and thus many models have been well exposed to the scrutiny of experiment. This would seem to imply that parameter spaces for the various models have shrunk so much that a measurement like $D^0$ mixing would have little impact. In fact, in giving this talk in several venues I challenged each audience to predict how many of the 21 models considered here were constrained by the $D^0$ mixing values or equivalently how many evaded constraint. Before answering this question, we consider a specific NP example in some detail.

Suppose a vector-like quark of charge $Q = +2/3$ is added to the SM. Recall that a vector-like quark is one whose electric charge is either $Q = +2/3$ or $Q = -1/3$ and which is an $SU(2)_L$ singlet. Both choices of charge are actually well motivated, as such fermions appear explicitly in several NP models. For example, weak isosinglets with $Q = -1/3$ appear in $E_6$ GUTs [16, 17], with one for each of the three generations ($D$, $S$, and $B$). Weak isosinglets with $Q = +2/3$ occur in Little Higgs theories [18, 19] in which the Standard Model Higgs boson is a pseudo-Goldstone boson, and the heavy iso-singlet $T$ quark cancels the quadratic divergences generated by the top quark in the mass of the Higgs boson. We restrict our attention here to the $Q = +2/3$ case. Since the electroweak quantum number assignments are different than those for the SM fermions, flavor changing neutral current interactions will be generated in the left-handed up-quark sector. Thus, there will also be FCNC couplings with the $Z^0$ boson [17]. These couplings contain a mixing parameter $\lambda_{uc}$ which is con-

<table>
<thead>
<tr>
<th>Model</th>
</tr>
</thead>
<tbody>
<tr>
<td>Fourth Generation</td>
</tr>
<tr>
<td>$Q = -1/3$ Singlet Quark</td>
</tr>
<tr>
<td>$Q = +2/3$ Singlet Quark</td>
</tr>
<tr>
<td>Little Higgs</td>
</tr>
<tr>
<td>Generic $Z'$</td>
</tr>
<tr>
<td>Family Symmetries</td>
</tr>
<tr>
<td>Left-Right Symmetric</td>
</tr>
<tr>
<td>Alternate Left-Right Symmetric</td>
</tr>
<tr>
<td>Vector Leptoquark Bosons</td>
</tr>
<tr>
<td>Flavor Conserving Two-Higgs-Doublet</td>
</tr>
<tr>
<td>Flavor Changing Neutral Higgs</td>
</tr>
<tr>
<td>FC Neutral Higgs (Cheng-Sher ansatz)</td>
</tr>
<tr>
<td>Scalar Leptoquark Bosons</td>
</tr>
<tr>
<td>Higgsless</td>
</tr>
<tr>
<td>Universal Extra Dimensions</td>
</tr>
<tr>
<td>Split Fermion</td>
</tr>
<tr>
<td>Warped Geometries</td>
</tr>
<tr>
<td>Minimal Supersymmetric Standard</td>
</tr>
<tr>
<td>Supersymmetric Alignment</td>
</tr>
<tr>
<td>Supersymmetry with RPV</td>
</tr>
<tr>
<td>Split Supersymmetry</td>
</tr>
</tbody>
</table>
strained by the unitarity condition
\[ \lambda_{uc} \equiv - (V_{ud}^* V_{cd} + V_{us}^* V_{cs} + V_{ub}^* V_{cb}) \, . \] (15)

A tree-level contribution to \( \Delta M_D \) is thus generated from \( Z^0 \)-exchange (see Fig. 2). It is straightforward to calculate that

\[ x_D^{(2/3)} = \frac{G_F \lambda_{uc}^2}{\sqrt{2} M_D \Gamma_D} r_1(m_c, M_Z) \langle \bar{D} | Q_1 | D^0 \rangle \]

\[ = \frac{2 G_F f_D^2 M_D}{3 \sqrt{2} \Gamma_D} B_D (\lambda_{uc})^2 r_1(m_c, M_Z) \quad (16) \]

where we have made use of Eq. (1) and Eq. (7). The result is displayed in the graph of Fig. 3 which contrasts a \( \pm 1 \sigma \) window (dashed lines) about the HFAG central value with the NP prediction \( x_D^{(2/3)} \) (solid line), as a function of the mixing parameter \( \lambda_{uc} \). The bound on \( \lambda_{uc} \) from \( D^0 \) mixing turns out to be roughly two orders of magnitude better than that from the CKM unitarity constraint.

Upon performing analogous analyses for the other NP models, we arrive at a set of constraints (mainly in the form of graphs) like the one depicted in Fig. 3. It is not possible here to summarize the results for all 21 models of Table IV. One must refer to Ref. [2] for that. However, we can answer the question raised earlier about how many models avoid being meaningfully constrained by the \( D^0 \) mixing data. The answer is just 4 of the 21 models, which came as a surprise to many. These 4 models are Split SUSY, Universal Extra Dimensions, LR Symmetric and Flavor-conserving Higgs Doublet.

It is of interest to briefly consider a few of these, in order to understand how the \( D^0 \) mixing constraints can be evaded:

1. Split SUSY [20]: This is a relatively new variant of SUSY (2003-4) in which SUSY breaks at a very large scale \( M_S \gg 1000 \) TeV. All scalars except the Higgs have mass \( M \sim M_S \) whereas fermions have the usual weak-scale mass. It is known that large \( D^0 \) mixing in SUSY will involve squark amplitudes. But since squark masses in Split SUSY are huge, the mixing becomes suppressed.

2. Universal Extra Dimensions [21]: UED is a variant of the idea that TeV-\( 1 \)-sized extra dimensions exist. There are no branes appearing in this approach, so all SM fields reside in the bulk and just one extra dimension is usually considered. Each SM field will have an infinity of KK excitations. It turns out that GIM cancellations suppress all but a few b-quark KK terms, but these are CKM inhibited.

So we see that suppressions can arise from more than one source and that the suppressing mechanism will depend on the specific model.

V. CONCLUSIONS

At long last, signals for \( x_D \) and \( y_D \) have been observed. These experimental findings, although greatly welcome, whet our appetite for even more precise determinations. Hopefully these will be forthcoming, so we can put aside any lingering concerns that all the excitement has been the result of statistical fluctuations.

The SM analysis, as is so often the case, is not without its difficulties. At the quark level, theoretical analysis in the dimension six sector through NLO gives \( x_D \sim y_D \sim 10^{-6} \). These values are tiny compared to the reported experimental signals. It is evident that the triple sum over the operator dimension \( d_{10} \), the QCD coupling \( \alpha_s \) and the mass expansion parameter \( z \) of Eq. (3) is slowly convergent. This approach remains inconclusive at best.

A more promising avenue is to study \( y_D \) with the hadronic degree of freedom. This yields a plausible, and quite possibly correct, explanation for reaching the \( y_D \sim 0.01 \) level. Again, however, the effect of strong interaction uncertainties mars predictive power.

Finally, the work of Refs. [1, 2] has explored which NP models can yield sizable values for \( x_D \), \( y_D \) and
which cannot. Charm mixing data has been found to infer useful constraints on NP parameters spaces, and as should be clear to all, provides a most welcome addition to the High Energy Physics community.

Acknowledgments

The author’s contribution to the work described above was supported in part by the U.S. National Science Foundation under Grant PHY–0555304.

[22] These two papers contain extensive references of contributions not possible to cite here.
CP violation in charm

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CP-violating asymmetries in charm provide a unique probe of physics beyond the Standard Model. I review several topics relevant to searches for CP-violation in charmed meson and baryon transitions.

1. Introduction

Charm transitions play a unique role in the modern investigations of flavor physics. They provide valuable supporting measurements for studies of CP-violation in $B$-decays, such as formfactors and decays constants, as well as outstanding opportunities for indirect searches for physics beyond the Standard Model (SM). It must be noted that in many dynamical models of new physics the effects of new particles observed in s, c, and b transitions are correlated. Therefore, such combined studies could yield the most stringent constraints on parameters of those models. For example, loop-dominated processes such as $D^0 - \overline{D^0}$ mixing or flavor-changing neutral current (FCNC) decays are influenced by the dynamical effects of down-type particles, whereas up-type particles are responsible for FCNC in the beauty and strange systems. Finally, from the practical point of view, charm physics experiments provide outstanding opportunities for studies of New Physics because of the availability of large statistical samples of data.

CP-violation can be introduced in Quantum Field Theory in a variety of ways [1]. One way, CP-violation can be introduced explicitly through dimension-4 operators (the so-called “hard” CP-breaking). This is how CP-invariance is broken in the Standard Model via quark Yukawa interactions,

$$\mathcal{L}_Y = \xi_{ik} \bar{\psi}_i \psi_k \phi + \text{h.c.} \quad (1)$$

The complex Yukawa couplings $\xi_{ik}$ lead to complex-valued Cabibbo-Kobayashi-Maskawa (CKM) quark mixing matrix providing the natural source of CP-violation for the case of the Standard Model with three (or more) generations. Another way could be via operators of dimensions less than four (the “soft” CP-breaking), which is popular in supersymmetric models. Yet another way is to break CP-invariance spontaneously. This method, which is somewhat aesthetically appealing, introduces CP-violating ground state with CP-conserved Lagrangian. It is realized in a class of left-right-symmetric models or multi-Higgs models. All these mechanisms can be probed in charm transitions. In fact, observation of CP-violation in the current round of charm experiments is arguably one of the cleanest signals of physics beyond the Standard Model (BSM).

It can be easily seen why manifestation of new physics interactions in the charm system is associated with the observation of (large) CP-violation. This is due to the fact that all quarks that build up the hadronic states in weak decays of charm mesons belong to the first two generations. Since $2 \times 2$ Cabibbo quark mixing matrix is real, no CP-violation is possible in the dominant tree-level diagrams which describe the decay amplitudes. CP-violating amplitudes can be introduced in the Standard Model by including penguin or box operators induced by virtual $b$-quarks. However, their contributions are strongly suppressed by the small combination of CKM matrix elements $V_{cb} V_{ub}^*$. It is thus widely believed that the observation of (large) CP violation in charm decays or mixing would be an unambiguous sign for new physics. This fact makes charm decays a valuable tool in searching for new physics, since the statistics available in charm physics experiment is usually quite large.

As with other flavor physics, CP-violating contributions in charm can be generally classified by three different categories:

(I) CP violation in the $\Delta C = 1$ decay amplitudes. This type of CP violation occurs when the absolute value of the decay amplitude for $D$ to decay to a final state $f$ ($A_f$) is different from the one of corresponding CP-conjugated amplitude (“direct CP-violation”). This can happen if the decay amplitude can be broken into at least two parts associated with different weak and strong phases,

$$A_f = |A_1| e^{i\delta_1} e^{i\phi_1} + |A_2| e^{i\delta_2} e^{i\phi_2}, \quad (2)$$

where $\phi_i$ represent weak phases ($\phi_i \to -\phi_i$ under CP-transformation), and $\delta_i$ represents strong phases ($\delta_i \to \delta_i$ under CP-transformation). This ensures that CP-conjugated amplitude, $\overline{A_f}$, would differ from $A_f$.

(II) CP violation in $D^0 - \overline{D^0}$ mixing matrix. Introduction of $\Delta C = 2$ transitions, either via SM or NP one-loop or tree-level NP amplitudes leads to non-diagonal entries in the $D^0 - \overline{D^0}$ mass matrix,

$$\left[ M - i \frac{\Gamma}{2} \right]_{ij} = \begin{pmatrix} A & p^2 \ A \end{pmatrix} \begin{pmatrix} q^2 & -p^2 \ A \end{pmatrix} \quad (3)$$
This type of CP violation is manifest when $R_m^\pm = |p/q|^2 = (2M_{12} - i\Gamma_{12})/(2M_{12} - i\Gamma_{12}^\star) \neq 1$.

(III) CP violation in the interference of decays with and without mixing. This type of CP violation is possible for a subset of final states to which both $D^0$ and $\bar{D}^0$ can decay.

For a given final state $f$, CP violating contributions can be summarized in the parameter

$$\lambda_f = \frac{q}{p} \frac{A_f}{\bar{A}_f} = R_m e^{i(\phi + \delta)} \frac{|A_f|}{|\bar{A}_f|},$$

where $A_f$ and $\bar{A}_f$ are the amplitudes for $D^0 \rightarrow f$ and $\bar{D}^0 \rightarrow f$ transitions respectively and $\delta$ is the strong phase difference between $A_f$ and $\bar{A}_f$. Here $\phi$ represents the convention-independent weak phase difference between the ratio of decay amplitudes and the mixing matrix.

The non-diagonal entries in the mixing matrix of Eq. (3) lead to mass eigenstates of neutral $D$-mesons that are different from the weak eigenstates,

$$|D_{\pm} \rangle = \frac{1}{\sqrt{2}} \left[ |D^0 \rangle \pm |\bar{D}^0 \rangle \right],$$

where the complex parameters $p$ and $q$ are obtained from diagonalizing the $D^0 - \bar{D}^0$ mass matrix with $|p|^2 + |q|^2 = 1$. If CP-violation in mixing is neglected, $p$ becomes equal to $q$, so $|D_{1,2} \rangle$ become CP eigenstates, $CP|D_{\pm} \rangle = \pm |D_{\pm} \rangle$.

The mass and width splittings between these eigenstates are given by

$$x = \frac{m_2 - m_1}{\Gamma}, \quad y = \frac{\Gamma_2 - \Gamma_1}{2\Gamma}.$$

It is known experimentally that $D^0 - \bar{D}^0$ mixing proceeds extremely slowly, which in the Standard Model is usually attributed to the absence of superheavy quarks destroying GIM cancellations $[2,3,4]$. As we shall see later, this fact additionally complicates searches for CP-violation in charmed mesons.

2. CP-violation in mesons

CP violation can be searched for by a variety of methods. In general, one can separate two ways. One way employs “static” observables, such as electric dipole moment of a baryon. Another way, more applicable to charm physics, employs “dynamical” observables, i.e. decay probabilities and asymmetries. Here we shall concentrate on this methods of searching for CP-violation.

a. CP-violation in transitions, forbidden by CP-invariance. This method is based on the idea that if both initial and final states are prepared as CP-eigenstates, the transition from the initial to final state would be forbidden if their CP-eigenvalues do not match. If CP is broken then transition probability would be proportional to CP-breaking parameter.

While neither of $D$-mesons constitute a CP-eigenstates, a linear combination of neutral $D$-mesons of Eq. (6) is. Thus such measurements can be performed at threshold charm factories, such as CLEO-c or BES-III, using quantum coherence of the initial state.

An example of this type of signal is a decay $(D^0\bar{D}^0 \rightarrow f_1f_2) at \psi(3770)$ with $f_1$ and $f_2$ being the different final CP-eigenstates of the same CP-parity.

These types of signals are very easy to detect experimentally. The corresponding CP-violating decay rate for the final states $f_1$ and $f_2$ is

$$\Gamma_{f_1f_2} = \frac{1}{2R_m^2} \left[ (2 + x^2 - y^2) |\lambda_{f_1} - \lambda_{f_2}|^2 \right.$$

$$\left. + (x^2 + y^2) |1 - \lambda_{f_1}\lambda_{f_2}|^2 \right] \Gamma_{f_1f_2}. \quad (8)$$

The result of Eq. (8) represents a slight generalization of the formula given in Ref. [5]. It is clear that both terms in the numerator of Eq. (8) receive contributions from CP-violation of the type I and III, while both terms in the numerator of Eq. (8) receive contributions from CP-violation of the type II.

Moreover, for a large set of the final states the first term would be additionally suppressed by SU(3)$_F$ symmetry, as for instance, $\lambda_{\pi\pi} = \lambda_{KK}$ in the SU(3)$_F$ symmetry limit. This expression is of the second order in CP-violating parameters (it is easy to see that in the approximation where only CP violation in the mixing matrix is retained, $\Gamma_{f_1f_2} \propto |1 - R_m^2| \propto A_m^2$).

The existing experimental constraints $[6]$ demonstrate that CP-violating parameters are quite small in the charm sector, regardless of whether they are produced by the Standard Model mechanisms or by some new physics contributions. Since the above measurements involve CP-violating decay rates, these observables are of second order in the small CP-violating parameters, a challenging measurement.

b. CP-violation in decay asymmetries.

Most of the experimental techniques that are sensitive to CP violation make use of decay asymmetries, which are similar to the ones employed in B-physics $[1]$.

$$a_f = \frac{\Gamma(D \rightarrow f) - \Gamma(\bar{D} \rightarrow \bar{f})}{\Gamma(D \rightarrow f) + \Gamma(\bar{D} \rightarrow \bar{f})}. \quad (9)$$

One can also introduce a related asymmetry,

$$a_\tau = \frac{\Gamma(D \rightarrow \tau\nu) - \Gamma(\bar{D} \rightarrow \bar{\tau}\bar{\nu})}{\Gamma(D \rightarrow \tau\nu) + \Gamma(\bar{D} \rightarrow \bar{\tau}\bar{\nu})}. \quad (10)$$
For charged D-decays the only contribution to the asymmetry of Eq. (10) comes from the multi-component structure of the \( \Delta C = 1 \) decay amplitude of Eq. (12). In this case,

\[
\alpha_f = \frac{2Im(A_1A_2^*) \sin \delta}{|A_1|^2 + |A_2|^2 + 2ReA_1A_2^* \cos \delta} = 2r_f \sin \phi \sin \delta, \tag{11}
\]

where \( \delta = \delta_1 - \delta_2 \) is the CP-conserving phase difference and \( \phi \) is the CP-violating one. \( r_f = |A_2/A_1| \) is the ratio of amplitudes. Both \( r_f \) and \( \delta \) are extremely difficult to compute reliably in D-decays. However, the task can be significantly simplified if one only concentrates on detection of New Physics in CP-violating asymmetries in the current round of experiments [3], i.e. at the \( \mathcal{O}(1\%) \) level. This is the level at which \( \alpha_f \) is currently probed experimentally, as summarized in Table 1. As follows from Eq. (11), in this case one should expect \( r_f \sim 0.01 \).

It is easy to see that the Standard Model asymmetries are safely below this estimate. First, Cabibbo-favored \( (A_f \sim \lambda^0) \) and doubly Cabibbo-suppressed \( (A_f \sim \lambda^1) \) decay modes proceed via amplitudes that share the same weak phase, so no CP-asymmetry is generated [1]. Moreover, presence of NP amplitudes does not significantly change this conclusion [3]. On the other hand, singly-Cabibbo-suppressed decays \( (A_f \sim \lambda^1) \) readily have two-component structure, receiving contributions from both tree and penguin amplitudes. In this case the same conclusion follows from the consideration of the charm CKM unitarity,

\[
V_{ud}V_{us}^* + V_{ub}V_{cs}^* + V_{ub}V_{cb}^* = 0. \tag{12}
\]

In the Wolfenstein parameterization of CKM, the first two terms in this equation are of the order \( \mathcal{O}(\lambda^2) \) (where \( \lambda \approx 0.22 \)), while the last one is \( \mathcal{O}(\lambda^3) \). Thus, CP-violating asymmetry is expected to be at most \( \alpha_f \sim 10^{-3} \) in the Standard Model. Model-dependent estimates of this asymmetry exist and are consistent with this estimate [1].

Asymmetries of Eq. (9) can also be introduced for the neutral D-mesons. In this case a much richer structure becomes available due to interplay of CP-violating contributions to decay and mixing amplitudes [5, 11],

\[
\alpha_f = a_f^d + a_f^m + a_f^i, \quad a_f^d = 2r_f \sin \phi \sin \delta, \quad a_f^m = -R_f \frac{y'}{2} (m_m - m_m^{-1}) \cos \phi, \quad a_f^i = R_f \frac{x'}{2} (m_m + m_m^{-1}) \sin \phi,
\]

where \( a_f^d, a_f^m, \) and \( a_f^i \) represent CP-violating contributions from decay, mixing and interference between decay and mixing amplitudes respectively. For the final states that are also CP-eigenstates, \( f = \bar{f} \) and \( y' = y \).

As can be seen from Eq. (13), the CP-violating asymmetries in neutral D-decays depend on \( D^0 - \bar{D}^0 \) mixing parameters \( x \) and \( y \). Presently, experimental information about the \( D^0 - \bar{D}^0 \) mixing parameters \( x \) and \( y \) comes from the time-dependent analyses that can roughly be divided into two categories. First, more traditional studies look at the time dependence of \( D \to f \) decays, where \( f \) is the final state that can be used to tag the flavor of the decayed meson. The most popular is the non-leptonic doubly Cabibbo suppressed decay \( D^0 \to K^+ \pi^- \). Time-dependent studies allow one to separate the DCS from the mixing contribution \( D^0 \to \bar{D}^0 \to K^+ \pi^- \),

\[
\Gamma[D^0 \to K^+ \pi^-] = e^{-\Gamma t} |A_{K^{-\pi^+}}|^2 
\left[ R + \sqrt{RR_m}(y \cos \phi - x' \sin \phi) \Gamma t \right] 
+ \frac{R_m}{4} (y^2 + x'^2)(\Gamma t)^2, \tag{14}
\]

where \( R \) is the ratio of DCS and Cabibbo favored (CF) decay rates. Since \( x \) and \( y \) are small, the best constraint comes from the linear terms in \( t \) that are also linear in \( x \) and \( y \). A direct extraction of \( x \) and \( y \) from Eq. (14) is not possible due to unknown relative strong phase \( \delta_D \) of DCS and CF amplitudes [12],

\[
as x' = x \cos \delta_D + y \sin \delta_D, \quad as y' = y \cos \delta_D - x \sin \delta_D.
\]

This phase can be measured independently [13]. The corresponding formula can also be written [11] for \( D^0 \) decay with \( x' \to -x' \) and \( R_m \to R_m^{-1} \).

Second, \( D^0 \) mixing can be measured by comparing the lifetimes extracted from the analysis of \( D \) decays into the CP-even and CP-odd final states. This study is also sensitive to a linear function of \( y \) via

\[
\tau(D \to K^+ \pi^-) \tau(D \to \bar{K}^0 \pi^+) - 1 = y \cos \phi - x \sin \phi \left( \frac{R_m^2 - 1}{2} \right). \tag{15}
\]

Time-integrated studies of the semileptonic transitions are sensitive to the quadratic form \( x^2 + y^2 \) and

\begin{table}
\begin{tabular}{|c|c|}
\hline
Decay mode & CP asymmetry \\
\hline
\( D^+ \to K_S \pi^+ \) & \(-0.016 \pm 0.017 \) \\
\( D^+ \to K_S \pi^+ \) & \(+0.017 \pm 0.062 \) \\
\( D^+ \to K^+ \pi^- \pi^+ \) & \(+0.007 \pm 0.008 \) \\
\( D^+ \to \pi^+ \pi^- \pi^+ \) & \(-0.017 \pm 0.042 \) \\
\( D^+ \to K_{S} K^+ \pi^+ \) & \(-0.042 \pm 0.068 \) \\
\hline
\end{tabular}
\end{table}
Table II Current experimental constraints on CP-violating asymmetries in neutral D-decays.

<table>
<thead>
<tr>
<th>Decay mode</th>
<th>CP asymmetry</th>
</tr>
</thead>
<tbody>
<tr>
<td>(D^0 \to K^+ K^-)</td>
<td>(+0.0136 \pm 0.012)</td>
</tr>
<tr>
<td>(D^0 \to K_S K_S)</td>
<td>(-0.23 \pm 0.19)</td>
</tr>
<tr>
<td>(D^0 \to \pi^+ \pi^-)</td>
<td>(+0.0127 \pm 0.0125)</td>
</tr>
<tr>
<td>(D^0 \to \pi^0 \pi^0)</td>
<td>(+0.001 \pm 0.048)</td>
</tr>
<tr>
<td>(D^0 \to K^- \pi^+ \pi^-)</td>
<td>(-0.031 \pm 0.086)</td>
</tr>
<tr>
<td>(D^0 \to K^+ K^- \pi^+ \pi^-)</td>
<td>(-0.082 \pm 0.073)</td>
</tr>
</tbody>
</table>

at the moment are not competitive with the analyses discussed above.

Three experimental collaborations (BaBar, Belle and CDF) have recently announced evidence for observation of \(D^0 - \bar{D}^0\) mixing using the analyses described above. The results reported by these collaborations were combined by the Heavy Flavor Averaging Group (HFAG) to yield [2]

\[
x = (8.4^{+3.2}_{-3.4}) \times 10^{-3},
\]

\[
y = (6.9 \pm 2.1) \times 10^{-3}.
\]

Once again, it can be seen that the results depend on hadronic parameters, such as the strong phase \(\delta_D\). While the observed values of \(x\) and \(y\), which are believed to be dominated by the Standard Model contributions (for recent analyses of NP contributions see [3]) and happen to be quite large, the SM CP-violating phases are still quite small. Thus, one can talk about almost background-free search for CP-violation induced by BSM interactions. Current experimental constraints on CP-violating asymmetries in neutral D-decays are summarized in Table II. As one can see, most measurements are the percent sensitivity. One should note that the rate asymmetries of Eq. (9) for neutral D-mesons require tagging of the initial state with the consequent reduction of the available dataset.

One question that can be asked is what models of New Physics can be probed via CP-violating observables in D-decays in the near future. A decomposition of Eq. (13) allows to address this question by studying parameters that enter Eq. (13). In particular, one needs to study the amplitude ratio \(r_f\) to see the feasibility of constraining a given NP model via charge asymmetries in D-decays. A general conclusion of the recent study [2] is that \(O(1\%)\) asymmetries are possible for SUSY models where new contributions come from QCD penguin operators and especially from chromomagnetic dipole operators, while tree-level direct CP violation in various known models is constrained to be much smaller than \(10^{-2}\) (see Table III. Clearly, neutral D decays could exhibit contributions from indirect or direct CP violation (or both). One can experimentally distinguish between these possibilities by selecting particular combinations of final states. For instance, combined analysis of \(D \to K \pi\) and \(D \to K K\) can yield interesting constraints on CP-violating parameters, which are universal [11].

\[
\Delta Y_{KK} = \frac{\Gamma'(D^0 \to K^+ K^-) - \Gamma'(\bar{D}^0 \to K^+ K^-)}{\Gamma'(D^0 \to K^+ K^-) + \Gamma'(\bar{D}^0 \to K^+ K^-)} = a_{KK}^m + a_{K^+ K^-},
\]

where \(\Gamma'(D^0 \to K^+ K^-)\) and \(\Gamma'(\bar{D}^0 \to K^+ K^-)\) are the modified decay rate parameters [11].

\[
\Gamma'(D^0 \to K^+ K^-) = \Gamma_D (1 + \eta_f^C P R_m (y \cos \phi - x \sin \phi)) ,
\]

\[
\Gamma'(\bar{D}^0 \to K^+ K^-) = \Gamma_D (1 + \eta_f^C P R_m (y \cos \phi + x \sin \phi)) .
\]

Table III. Tree-level NP contributions to \(r_f\) [2].

<table>
<thead>
<tr>
<th>Model</th>
<th>(r_f)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Extra quarks in vector-like rep</td>
<td>(&lt;10^{-3})</td>
</tr>
<tr>
<td>R-parity violating SUSY</td>
<td>(&lt;1.5 \times 10^{-4})</td>
</tr>
<tr>
<td>Two-Higgs doublet</td>
<td>(&lt;4 \times 10^{-4})</td>
</tr>
</tbody>
</table>

It is possible to use a method that both does not require flavor or CP-tagging of the initial state and results in the observable that is first order in CP violating parameters [10]. Let’s concentrate on the decays of D mesons to final states that are common for \(D^0\) and \(\bar{D}^0\). If the initial state is not tagged the quantities that one can easily measure are the sums

\[
\Sigma_i = \Gamma_i(t) + \bar{\Gamma}_i(t)
\]

for \(i = f\) and \(\bar{f}\). A CP-odd observable which can be formed out of \(\Sigma_i\) is the asymmetry

\[
A_{CP}^{ij} (f, t) = \frac{\Sigma_f - \Sigma_\bar{f}}{\Sigma_f + \Sigma_\bar{f}} = \frac{N(t)}{D(t)} .
\]

We shall consider both time-dependent and time-integrated versions of the asymmetry [20]. Note that this asymmetry does not require quantum coherence of the initial state and therefore is accessible in any D-physics experiment. It is expected that the numerator
and denominator of Eq. \( \text{20} \) would have the form,
\[
N(t) = \Sigma_f - \Sigma_T = e^{-T} \left[ A + BT + CT^2 \right],
\]
\[
D(t) = 2e^{-T} \left( |A_f|^2 + |A_T|^2 \right), \tag{21}
\]
where we neglected direct CP violation in \( D(t) \). Integrating the numerator and denominator of Eq. \( \text{20} \) over time yields
\[
A_{\text{CP}}^U(f) = \frac{1}{D} [A + B + 2C], \tag{22}
\]
where \( D = \Gamma \int_0^\infty dt \ D(t) \).

Both time-dependent and time-integrated asymmetries depend on the same parameters \( A, B, \) and \( C \). The result is
\[
A = |A_f|^2 - |A_T|^2 - |A_f|^2, \\
B = -2y\sqrt{R} \left[ \sin \phi \sin \delta \left( |A_f|^2 + |A_T|^2 \right) \right. \\
- \cos \phi \cos \delta \left( |A_f|^2 - |A_T|^2 \right), \tag{23}
\]
\[
C = y^2 / 2A.
\]

We neglect small corrections of the order of \( \mathcal{O}(m_{\pi}, r_f, \ldots) \) and higher. It follows that Eq. \( \text{24} \) receives contributions from both direct and indirect CP-violating amplitudes. Those contributions have different time dependence and can be separated either by time-dependent analysis of Eq. \( \text{20} \) or by the “designer” choice of the final state. Note that this asymmetry is manifestly first order in CP-violating parameters.

In Eq. \( \text{23} \), non-zero value of the coefficient \( A \) is an indication of direct CP violation. This term is important for singly Cabibbo suppressed (SCS) decays. The coefficient \( B \) gives a combination of a contribution of CP violation in the interference of the decays with and without mixing (first term) and direct CP violation (second term). Those contributions can be separated by considering DCS decays, such as \( D \to K^{(*)}\pi \) or \( D \to K^{(*)}\rho \), where direct CP violation is not expected to enter. The coefficient \( C \) represents a contribution of CP-violation in the decay amplitudes after mixing. It is negligibly small in the SM and all models of new physics constrained by the experimental data. Note that the effect of CP-violation in the mixing matrix on \( A, B, \) and \( C \) is always subleading.

Eq. \( \text{23} \) is completely general and is true for both DCS and SCS transitions. Neglecting direct CP violation we obtain a much simpler expression,
\[
A = 0, \quad C = 0, \\
B = -2y\sin \delta \sin \phi \sqrt{R} \left( |A_f|^2 + |A_T|^2 \right). \tag{24}
\]

For an experimentally interesting DCS decay \( D^0 \to K^+\pi^- \) this asymmetry is zero in the flavor \( SU(3)_F \) symmetry limit, where \( \delta = 0 \). Since \( SU(3)_F \) is badly broken in \( D \)-decays, large values of \( \sin \delta \) are possible. At any rate, regardless of the theoretical estimates, this strong phase could be measured at CLEO-c. It is also easy to obtain the time-integrated asymmetry for \( K^\pi \). Neglecting small subleading terms of \( \mathcal{O}(\lambda^4) \) in both numerator and denominator we obtain
\[
A_{\text{CP}}^U(K\pi) = -y \sin \delta \sin \phi \sqrt{R}. \tag{25}
\]

It is important to note that both time-dependent and time-integrated asymmetries of Eqs. \( \text{24} \) and \( \text{25} \) are independent of predictions of hadronic parameters, as both \( \delta \) and \( R \) are experimentally determined quantities and could be used for model-independent extraction of CP-violating phase \( \phi \). Assuming \( R \approx 0.4\% \) and \( \delta \approx 40^\circ \) \( \text{12} \) and \( y \approx 1\% \) one obtains \( |A_{\text{CP}}^U(K\pi)| \approx (0.04\%) \sin \phi \). Thus, one possible challenge of the analysis of the asymmetry Eq. \( \text{23} \), is that it involves a difference of two large rates, \( \Sigma_{K^+\pi^-} \) and \( \Sigma_{K^-\pi^+} \), which should be measured with the sufficient precision to be sensitive to \( A_{\text{CP}}^U \), a problem tackled in determinations of tagged asymmetries in \( D \to K\pi \) transitions.

Alternatively, one can study SCS modes, where \( R \approx 1 \), so the resulting asymmetry could be \( \mathcal{O}(1\%) \sin \phi \). However, the final states must be chosen such that \( A_{\text{CP}}^U \) is not trivially zero. For example, decays of \( D \) into the final states that are CP-eigenstates would result in zero asymmetry (as \( \Gamma_f = \Gamma_T \) for those final states) while decays to final states like \( K^+K^- \) or \( \rho^+\pi^- \) would not. It is also likely that this asymmetry is larger than the estimate given above due to contributions from direct CP-violation (see eq. \( \text{23} \)). The final state \( f \) can also be a multiparticle state. In that case, more untagged CP-violating observables could be constructed, for instance involving asymmetries of the Dalitz plots, such as the ones proposed for B-decays \( \text{18} \).

As any rate asymmetry, Eq. \( \text{20} \) requires either a “symmetric” production of \( D^0 \) and \( \bar{D}^0 \), a condition which is automatically satisfied by all \( p\bar{p} \) and \( e^+e^- \) colliders, or a correction for \( D^0/\bar{D}^0 \) production asymmetry.

### 3. CP-violation in baryons

Charmed baryons provide another system for searches for CP-violation in charm. The fact that baryons are spin-1/2 particles allows us to form CP-violating asymmetries that are different from the ones in the meson systems.

Taking \( \Lambda_c \) as an example, a charmed baryon decay
amplitude can be parameterized as
\[ A(\Lambda_c \to B\pi) = \pi_B(p, s) \left[ A_S + A_P \gamma_5 \right] u_{\Lambda_c}(p_{\Lambda_c}, s_{\Lambda_c}), \] (26)
where \( B \) is a charmless baryon, and \( A_S \) and \( A_P \) parameterize \( s \)- and \( p \)-wave decay amplitudes respectively. They can be combined in an “asymmetry parameter” \( \alpha_{\Lambda_c} \) as
\[ \alpha_{\Lambda_c} = \frac{2Re(A_S^* A_P)}{|A_S|^2 + |A_P|^2}. \] (27)
This parameter can be directly measured experimentally using angular distribution of decay products in \( \Lambda_c \) decay,
\[ \frac{dW}{d\theta} = \frac{1}{2} (1 + P\alpha_{\Lambda_c} \cos \theta). \] (28)
Here \( P \) is polarization of the initial-state baryon. If this analysis can be done for \( \bar{\Lambda}_c \) decay as well, then a CP-violating asymmetry can be formed,
\[ A_f = \frac{\alpha_{\Lambda_c} + \alpha_{\bar{\Lambda}_c}}{\alpha_{\Lambda_c} - \alpha_{\bar{\Lambda}_c}}, \] (29)
which follows from the fact that \( \alpha_{\Lambda_c} \to -\alpha_{\bar{\Lambda}_c} \) under CP-transformation (if CP is conserved).

There were some experimental studies of this observable. In particular, FOCUS collaboration reported [19]
\[ A_{\Lambda\pi} = -0.07 \pm 0.19 \pm 0.24. \] (30)
New studies of CP-asymmetries in charmed baryon decays are urged, which could be performed at LHCb or even in one of the new experiments associated with Project-X at FNAL [20].

4. Conclusions

In summary, charm physics, and in particular studies of CP-violation, could provide new and unique opportunities for indirect searches for New Physics. Large statistical samples of charm data allow unique sensitive measurements of charm mixing and CP-violating parameters. While unambiguous theoretical predictions of CP-violating asymmetries in charm transitions are hard, observation of CP-violation at the level of \( O(1\%) \) would indicate new physics contribution to charm decays.

Acknowledgments

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References


Experimental Prospects for \( CP \) and \( T \) Violation Studies in Charm*  

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We present the current status of experimental results and prospects for the determination of \( CP \) and \( T \) violation in the charm sector. Such measurements have acquired renewed interest in recent years in view of theoretical work, which has highlighted the possibility to probe experimental signatures from New Physics beyond the Standard Model, since the effect of \( CP \) violation due to Standard Model processes is expected to be highly suppressed in \( D \) decays. The current limits of experimental sensitivities for these studies are reaching the interesting theoretical regimes. We include new measurements from the Belle, \babar, and CLEO-c collaborations.

1. Introduction

The amount of \( CP \) violation (CPV) currently discovered in nature is not sufficient to explain the universe as we see it. Looking in the charm sector is a natural extension of this task. Three are the kinds of CPV we deal with: CPV in the \( D^0 - \bar{D}^0 \) mixing matrix, which is expected to be insignificant in the charm sector, CPV in the decay amplitudes, and CPV in the interference between mixing and decay, which should be very small as well. The second one is also known as direct CPV and will be covered in this paper.

The expression for the \( CP \) asymmetry resulting from a process \( f \) and its \( CP \) conjugate \( f \) is given by:

\[
A_{CP} = \frac{\Gamma(f) - \Gamma(\bar{f})}{\Gamma(f) + \Gamma(\bar{f})} = \frac{2\Im(A_1 A_2^*) \sin(\delta_1 - \delta_2)}{|A_1|^2 + |A_2|^2 + 2\Re(A_1 A_2^*) \cos(\delta_1 - \delta_2)}
\]

\( A_1 \) and \( A_2 \) are two components of the decay amplitude and \( \delta_1 - \delta_2 \) the corresponding strong phase difference. It follows that two amplitudes with different strong as well as weak phases are needed to have CPV. In the realm of the SM, usually this means a tree and a penguin amplitude. The kinds of processes described in the following are categorized as Cabibbo favored (CF, \( c \to s\bar{d}u \)), suppressed (CS, \( c \to s\bar{s}u, c \to d\bar{d}u \)), and doubly suppressed (DCS, \( c \to d\bar{s}u \)), according to the kind of vertices that intervene in the charm quark decay.

In contrast to the beauty sector, the Standard Model (SM) charm sector is largely \( CP \) conserving, as it involves 4 quarks and the \( 2 \times 2 \) Cabibbo mixing matrix is real. In singly Cabibbo suppressed decays diluted weak phases can produce asymmetries of the order \( 10^{-3} - 10^{-4} \), while no weak phases, hence no CPV, exist in CF and DCS decays, except for some minimal asymmetry in the \( D^+ \to K^- \pi^+ \pi^- \) mode. It is interesting to notice that it is possible in principle to distinguish direct and indirect CPV, either combining direct \( CP \) asymmetries with time-dependent measurements both for \( CP \) eigenstates, or just using time integrated measurements for CF \( CP \) eigenstate modes (assuming negligible CPV in CF modes) as \( K_S \pi^0 \) [1].

New Physics (NP) can contain \( CP \) violating couplings that could show up at the percent level [1] [2] [3] [4]. Several extensions of the SM predict such asymmetries, including models with leptoquarks, a fourth generation of fermions, right-handed weak currents, or extra Higgs doublets. Precision measurements and theory are required to detect NP. The charm sector is in a unique position to test physics beyond the SM. In particular it can test models where CPV is generated in the up-like quark sector. Flavor models where the CKM mixing is generated in the up sector generally predict large D-mixing and sizable CPV in charm, but smaller effects in the beauty sector. Furthermore, SCS \( D \) decays are now more sensitive to gluonic penguin amplitudes than are charmless B decays [1]. In summary, finding CPV in CF and DCS decays or finding CPV above 0.1% in SCS decays would indicate NP.

2. Current Experimental Results

There are several ways direct \( CP \) and \( T \) violation can be looked for: by measuring asymmetries in time integrated partial widths or in final state distributions of Dalitz plots or by measuring \( T \) odd correlations with 4-body \( D \) decays. As an example of charged \( D \) decays, Fig. [1] shows the reconstructed mass distributions in \( D^+ \to K^- K^+ \pi^+ \) and \( D^+ \to \pi^+ \pi^+ \pi^- \) candidates in the \babar\ detector, with fairly large datasets (80 \( fb^{-1} \) [2]). The \( CP \) asymmetry results of these and many other analyses are listed in Table [1].

As mentioned earlier, in the SM we can expect direct CPV at the \( 10^{-3} \) level in SCS decays, but no CPV in CF modes. In the \( K^+ K^- \) and \( \pi^+ \pi^- \) modes it is puzzling that the ratio of their branching ratios is so different from 1 (~2.8). This entails the presence of

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large final state interactions (FSI) and/or large penguin contributions, which could be fertile ground for NP to manifest itself. Phenomenological calculations set SM limits at or well below the $10^{-3}$ level. The CDF collaboration has to date the best CP measurements for these modes. The asymmetries are normalized to the CF $K\pi$ mode. Difficulties for this analysis are due to the track charge asymmetry which is calibrated with $K_S$ control samples and to the partially reconstructed $D$ background for the $K^+K^-$ mode. Total systematics are slightly above 0.5%.

A high precision analysis at BABAR reports $A_{CP} = 0.00 \pm 0.34 \pm 0.13$ for the $K^+K^-$ mode and $A_{CP} = -0.24 \pm 0.52 \pm 0.22$ for the $\pi^+\pi^-$ mode. To keep the systematics so low, the keys are to calibrate charge and tagging asymmetries using data, namely the CF $K\pi$ mode, and to account for forward-backward asymmetries due to QED effects, which can produce detection asymmetries in a detector as BABAR, due to the boost of the center of mass system with respect to the laboratory.

Charm factories benefit with respect to the beauty modes from a pure $D\bar{D}$ final state with low multiplicity, hence high tagging efficiency. This makes them competitive with the high statistics at BABAR and Belle. Single tag efficiencies range from 25 to 65%, values unimaginable at the B-factories. Most of the new CP violation results from CLEO-c, with 281 pb$^{-1}$ of data, are for CF modes, with the exception of $D^+ \rightarrow K^+K^-\pi^+$ (see Tables I and II). The uncertainties are of the order of 1% in most cases, For modes with charged kaons, the kaon systematics are the largest ones.

In case of indirect CPV and final CP eigenstates, the time integrated and time dependent CP asymmetries are universal and equal to each other. In contrast, for direct CPV, the time-integrated asymmetries in principle are not expected to be universal. Hence parts of phase-space in a multi-body decay might have different asymmetries (which may even cancel each other out when integrated over the whole phase-space). In addition, NP might not show up in the decay rates asymmetries but instead in the phase difference between amplitudes. 3-body decays permit the measurement of such phase differences. The Dalitz plot technique allows increased sensitivity to CP asymmetry by probing the decay amplitude rather than the decay rate and access to both CP eigenstates and non CP eigenstates with relatively high statistics. The CLEO-c collaboration has measured the CP asymmetry in the $\pi^+\pi^-\pi^0$ mode (integrated over the sum of all amplitudes in the Dalitz plot) and has also performed a full fledged Dalitz plot analysis of the $D^0 \rightarrow K_S\pi^+\pi^-$ decay. The BABAR and Belle collaborations can exploit their larger datasets for similar measurements.

T violation measurements can be performed exploiting T-odd correlations between the momenta of the decay products of 4-body $D$ decays as $KK\pi\pi$, while assuming CPT conservation:

$$C_T = p_{K^+} \cdot (p_{\pi^+} \times p_{\pi^-}).$$

Under time reversal $C_T$ changes sign, but its being different from 0 is not sufficient to establish T-violation as final state interactions can fake this asymmetry. To overcome this problem the analogous quantity from the CP conjugate decay can be defined as:

$$\overline{C_T} = p_{K^-} \cdot (p_{\pi^-} \times p_{\pi^+}).$$

Finding $\overline{C_T} \neq -C_T$ establishes T violation. T-odd asymmetries can be built as:

$$A_T = \frac{\Gamma(C_T > 0) - \Gamma(C_T < 0)}{\Gamma(C_T > 0) + \Gamma(C_T < 0)}$$

$$\overline{A_T} = \frac{\Gamma(\overline{C_T} > 0) - \Gamma(\overline{C_T} < 0)}{\Gamma(\overline{C_T} > 0) + \Gamma(\overline{C_T} < 0)}$$

and the T violation asymmetry as:

$$A_{T-viol} = \frac{1}{2} (A_T - \overline{A_T})$$

and if this is different from 0, T violation is established, even in the presence of strong phases.
As for the available measurements the only ones to date are from the FOCUS collaboration (see Tables I and III). The CLEO, BaBar, and Belle experiments should better this analysis with their larger and cleaner data samples.

Most of the measurements of CP and T violation in neutral D decays to date are shown in Table I. No evidence of direct CPV has been found. The best limits are of the order of one to few percent statistical errors with systematics of similar magnitude; few measurements have errors below the 1% level. Most of these are old measurements, except for the new ones from the CLEO-c collaboration and the ones “byproducts” of the mixing analyses of the BaBar and Belle collaborations.

Table I: ACP measurements to date using neutral D decays. The last row reports a measurement of $A_T$ by the FOCUS collaboration.

<table>
<thead>
<tr>
<th>Experiment(year)</th>
<th>Decay mode</th>
<th>$A_{CP} %$</th>
</tr>
</thead>
<tbody>
<tr>
<td>CDF(2005)</td>
<td>$D^0 \rightarrow K^+K^-$</td>
<td>$2.0 \pm 1.2 \pm 0.6$</td>
</tr>
<tr>
<td>CLEO(2002)</td>
<td>$D^0 \rightarrow K^+K^-$</td>
<td>$0.0 \pm 2.2 \pm 0.8$</td>
</tr>
<tr>
<td>FOCUS(2000)</td>
<td>$D^0 \rightarrow K^+K^-$</td>
<td>$-0.1 \pm 2.2 \pm 1.5$</td>
</tr>
<tr>
<td>CDF(2005)</td>
<td>$D^0 \rightarrow \pi^+\pi^-$</td>
<td>$1.0 \pm 1.3 \pm 0.6$</td>
</tr>
<tr>
<td>CLEO(2002)</td>
<td>$D^0 \rightarrow \pi^+\pi^-$</td>
<td>$1.9 \pm 3.2 \pm 0.8$</td>
</tr>
<tr>
<td>FOCUS(2000)</td>
<td>$D^0 \rightarrow \pi^+\pi^-$</td>
<td>$4.8 \pm 3.9 \pm 2.5$</td>
</tr>
<tr>
<td>CLEO(2001)</td>
<td>$D^0 \rightarrow K^0_S\pi^0_S$</td>
<td>$-23 \pm 19$</td>
</tr>
<tr>
<td>CLEO(2001)</td>
<td>$D^0 \rightarrow \pi^0\pi^0\phi$</td>
<td>$0.1 \pm 4.8$</td>
</tr>
<tr>
<td>CLEO(2001)</td>
<td>$D^0 \rightarrow K^0_S\pi^0$</td>
<td>$0.1 \pm 1.3$</td>
</tr>
<tr>
<td>CLEO(1995)</td>
<td>$D^0 \rightarrow K^0_S\phi$</td>
<td>$2.8 \pm 9.4$</td>
</tr>
<tr>
<td>CLEO(2005)</td>
<td>$D^0 \rightarrow \pi^+\pi^-\pi^0$</td>
<td>$1^{+0}_{-5}$</td>
</tr>
<tr>
<td>CLEO(2004)</td>
<td>$D^0 \rightarrow K^0_S\pi^+\pi^-$</td>
<td>$-0.9 \pm 2.1^{+0.5}_{-0.5}$</td>
</tr>
<tr>
<td>Belle(2005)</td>
<td>$D^0 \rightarrow K^+\pi^+\pi^-\pi^-$</td>
<td>$-1.8 \pm 4.4$</td>
</tr>
<tr>
<td>FOCUS(2005)</td>
<td>$D^0 \rightarrow K^+\pi^-\pi^+\pi^-$</td>
<td>$-8.2 \pm 5.6 \pm 4.7$</td>
</tr>
<tr>
<td>CLEO(2007)</td>
<td>$D^0 \rightarrow K^-\pi^+$</td>
<td>$0.4 \pm 0.5 \pm 0.9$</td>
</tr>
<tr>
<td>CLEO(2007)</td>
<td>$D^0 \rightarrow K^0_S\pi^+$</td>
<td>$0.2 \pm 0.4 \pm 0.8$</td>
</tr>
<tr>
<td>Belle(2005)</td>
<td>$D^0 \rightarrow K^-\pi^+\pi^-\pi^+$</td>
<td>$0.7 \pm 0.5 \pm 0.9$</td>
</tr>
<tr>
<td>Belle(2005)</td>
<td>$D^0 \rightarrow K^-\pi^+\pi^-\pi^+$</td>
<td>$-0.6 \pm 5.3$</td>
</tr>
<tr>
<td>BaBar(2007)</td>
<td>$D^0 \rightarrow K^-\pi^+\pi^-\pi^+$</td>
<td>$2.1 \pm 5.2 \pm 1.5$</td>
</tr>
<tr>
<td>BaBar(2007)</td>
<td>$D^0 \rightarrow K^-\pi^+\pi^-\pi^+$</td>
<td>$2.3 \pm 4.7$</td>
</tr>
<tr>
<td>FOCUS(2005) $A_T$</td>
<td>$D^0 \rightarrow K^+\pi^-\pi^-\pi^+$</td>
<td>$1.0 \pm 5.7 \pm 3.7$</td>
</tr>
</tbody>
</table>

Table II: ACP measurements to date using charged D decays. The last two rows report measurements of $A_T$ by the FOCUS collaboration.

<table>
<thead>
<tr>
<th>Experiment(year)</th>
<th>Decay mode</th>
<th>$A_{CP} %$</th>
</tr>
</thead>
<tbody>
<tr>
<td>BaBar(2005)</td>
<td>$D^+ \rightarrow K^+K^-\pi^+$</td>
<td>$1.4 \pm 1.0 \pm 0.8$</td>
</tr>
<tr>
<td>BaBar(2005)</td>
<td>$D^+ \rightarrow \phi\pi^+$</td>
<td>$0.2 \pm 1.5 \pm 0.6$</td>
</tr>
<tr>
<td>BaBar(2005)</td>
<td>$D^+ \rightarrow K^0_S\pi^+$</td>
<td>$0.9 \pm 1.7 \pm 0.7$</td>
</tr>
<tr>
<td>CLEO(2007)</td>
<td>$D^+ \rightarrow K^+\pi^-\pi^+$</td>
<td>$-0.1 \pm 1.5 \pm 0.8$</td>
</tr>
<tr>
<td>FOCUS(2000)</td>
<td>$D^+ \rightarrow K^+\pi^-\pi^+$</td>
<td>$0.6 \pm 1.1 \pm 0.5$</td>
</tr>
<tr>
<td>E791(1997)</td>
<td>$D^+ \rightarrow K^+\pi^-\pi^+$</td>
<td>$-1.4 \pm 2.9$</td>
</tr>
<tr>
<td>E791(1997)</td>
<td>$D^+ \rightarrow \phi\pi^+$</td>
<td>$-2.8 \pm 3.6$</td>
</tr>
<tr>
<td>E791(1997)</td>
<td>$D^+ \rightarrow K^0_S\pi^+$</td>
<td>$-1.0 \pm 5.0$</td>
</tr>
<tr>
<td>FOCUS(2002)</td>
<td>$D^+ \rightarrow K^0_S\pi^+$</td>
<td>$-1.6 \pm 1.5 \pm 0.9$</td>
</tr>
<tr>
<td>CLEO(2007)</td>
<td>$D^+ \rightarrow K^0_S\pi^+$</td>
<td>$-0.6 \pm 1.0 \pm 0.3$</td>
</tr>
<tr>
<td>CLEO(2007)</td>
<td>$D^+ \rightarrow K^0_S\pi^+\pi^0$</td>
<td>$0.3 \pm 0.9 \pm 0.3$</td>
</tr>
<tr>
<td>CLEO(2007)</td>
<td>$D^+ \rightarrow K^0_S\pi^+\pi^0\pi^0$</td>
<td>$0.1 \pm 1.1 \pm 0.6$</td>
</tr>
<tr>
<td>CLEO(2007)</td>
<td>$D^+ \rightarrow K^0_S\pi^+\pi^+\pi^-$</td>
<td>$-0.5 \pm 0.4 \pm 0.9$</td>
</tr>
<tr>
<td>CLEO(2007)</td>
<td>$D^+ \rightarrow K^0_S\pi^+\pi^+\pi^0$</td>
<td>$1.0 \pm 0.9 \pm 0.9$</td>
</tr>
<tr>
<td>CLEO(2007)</td>
<td>$D^+ \rightarrow K^+\eta$</td>
<td>$-20 \pm 18$</td>
</tr>
<tr>
<td>CLEO(2007)</td>
<td>$D^+ \rightarrow K^-\eta^*$</td>
<td>$-17 \pm 37$</td>
</tr>
<tr>
<td>CLEO(2007)</td>
<td>$D^+ \rightarrow K^0_S\pi^+$</td>
<td>$27 \pm 11$</td>
</tr>
<tr>
<td>CLEO(2007)</td>
<td>$D^+ \rightarrow K^+\pi^0$</td>
<td>$2 \pm 29$</td>
</tr>
<tr>
<td>E791(1997)</td>
<td>$D^+ \rightarrow \pi^+\pi^-\pi^+$</td>
<td>$-1.7 \pm 4.2$</td>
</tr>
<tr>
<td>FOCUS(2005) $A_T$</td>
<td>$D^+ \rightarrow K^0_S\pi^+\pi^+\pi^-$</td>
<td>$2.3 \pm 6.2 \pm 2.2$</td>
</tr>
<tr>
<td>FOCUS(2005) $A_T$</td>
<td>$D^+ \rightarrow K^0_S\pi^+\pi^+\pi^+$</td>
<td>$-3.6 \pm 6.7 \pm 2.3$</td>
</tr>
</tbody>
</table>

3. Future Prospects

The future prospects for these measurements are very promising. For the $KK$ and $\pi\pi$ modes both B-factories and CDF are expected to reach very interesting sensitivities of the order of few per thousand. The issue at the Tevatron will be whether the trigger can cope with the increase of luminosity. The $D^+ \rightarrow KK\pi$ mode should hit interesting limits as well, if systematics can be held under control. Very promising are also measurements from Dalitz plot analyses using SCS modes, where we have the added puzzle that it is not known where (if anywhere) CPV can show up in the Dalitz plane. Furthermore, the asymmetry could be large, but confined to only a part of the phase-space.

For the T-correlation analyses, as aforementioned there are large datasets of 4-body $D$ decays available. The BaBar and Belle collaborations could achieve statistical uncertainties at or below 0.5% if systematics can be kept as low.

If the present machines cannot fully probe the extent of the CP asymmetries allowed by NP or given to us by nature, new experiments at BESIII and at the super B-factories at KEK and/or Frascati, or LHC-b
Table III Average CP asymmetry measurements by mode.

<table>
<thead>
<tr>
<th>Decay mode</th>
<th>$A_{CP} %$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$D^0 \rightarrow K^+ K^-$</td>
<td>$+1.4 \pm 1.2$</td>
</tr>
<tr>
<td>$D^0 \rightarrow K^0_S K^0_S$</td>
<td>$-2.3 \pm 1.9$</td>
</tr>
<tr>
<td>$D^0 \rightarrow \pi^+ \pi^-$</td>
<td>$+1.3 \pm 1.3$</td>
</tr>
<tr>
<td>$D^0 \rightarrow \pi^0 \pi^0$</td>
<td>$+0.1 \pm 4.8$</td>
</tr>
<tr>
<td>$D^0 \rightarrow \pi^+ \pi^- \pi^0$</td>
<td>$+1 \pm 9$</td>
</tr>
<tr>
<td>$D^0 \rightarrow K^0_S \pi^0$</td>
<td>$+0.1 \pm 1.3$</td>
</tr>
<tr>
<td>$D^0 \rightarrow K^- \pi^+$</td>
<td>$-0.4 \pm 1.0$</td>
</tr>
<tr>
<td>$D^0 \rightarrow K^- \pi^+ \pi^0$</td>
<td>$+0.2 \pm 0.9$</td>
</tr>
<tr>
<td>$D^0 \rightarrow K^+ \pi^-$</td>
<td>$-0.8 \pm 3.1$</td>
</tr>
<tr>
<td>$D^0 \rightarrow K^+ \pi^- \pi^0$</td>
<td>$-0.1 \pm 5.2$</td>
</tr>
<tr>
<td>$D^0 \rightarrow K^+ S^0 + \pi^0$</td>
<td>$-0.9 \pm 4.2$</td>
</tr>
<tr>
<td>$D^0 \rightarrow K^+ \pi^- \pi^+ \pi^-$</td>
<td>$-1.8 \pm 4.4$</td>
</tr>
<tr>
<td>$D^0 \rightarrow K^+ S^0 + \pi^0 \pi^{-}$</td>
<td>$-8.2 \pm 7.3$</td>
</tr>
<tr>
<td>$D^+ \rightarrow K^0_S \pi^+$</td>
<td>$-0.9 \pm 0.9$</td>
</tr>
<tr>
<td>$D^+ \rightarrow K^0_S \pi^+ \pi^-$</td>
<td>$0.3 \pm 0.9$</td>
</tr>
<tr>
<td>$D^+ \rightarrow K^0_S \pi^+ \pi^- \pi^0$</td>
<td>$0.1 \pm 1.3$</td>
</tr>
<tr>
<td>$D^+ \rightarrow K^- \pi^+ \pi^+$</td>
<td>$0.5 \pm 1.0$</td>
</tr>
<tr>
<td>$D^+ \rightarrow K^- \pi^+ + \pi^- \pi^0$</td>
<td>$1.0 \pm 1.3$</td>
</tr>
<tr>
<td>$D^+ \rightarrow K^+ S^0 + K^+$</td>
<td>$+7.1 \pm 6.2$</td>
</tr>
<tr>
<td>$D^+ \rightarrow K^+ S^0 + K^- \pi^-$</td>
<td>$+0.6 \pm 0.8$</td>
</tr>
<tr>
<td>$D^+ \rightarrow \pi^+ \pi^- \pi^+$</td>
<td>$-1.7 \pm 4.2$</td>
</tr>
<tr>
<td>$D^+ \rightarrow K^0_S \pi^+ \pi^-$</td>
<td>$-4.2 \pm 6.8$</td>
</tr>
</tbody>
</table>

4. Conclusions

Charm physics provides unique opportunities for indirect searches for new physics. The theoretical calculations of the $D^0$ mixing parameters have large uncertainties, hence physics beyond the standard model will be hard to rule out from $D^0$ mixing measurements alone. The observation of large CPV would instead be a clear and robust signal of new physics. There have been some exciting new results this year from the CLEO-c, Belle, and BABAR collaborations. The total uncertainties are at the 1% level in several modes, but still far from observation. Experiments are just now entering the interesting domain. The future ahead is very promising with good sensitivities achievable by current experiments and high precision measurements expected with future and planned efforts.

Acknowledgments

I would like to thank the conference organizers for their help and gracious hospitality in Ithaca, my charming mentors, Brian Meadows and Mike Sokoloff, and students, Kalanand Mishra and Rolf Andreassen, from the University of Cincinnati, as well as all my other colleagues who have been, are, and will be working on this exciting subject. This work was supported in part by the National Science Foundation Grant Number PHY-0457336 and by the Department of Energy Contract Number DE-AC03-76SF00515.

References

Rare D Decays

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We discuss several recent measurements of rare charmed hadron decays. Focus is placed on radiative and annihilation topologies highlighting their sensitivity to new physics and pointing out the strengths and weaknesses of different channels. We compare the different measurement techniques employed at fixed target and $e^+e^-$ dedicated charm experiments, B-factories, and the Tevatron experiments. Comparisons are also made to similar topologies in the beauty, strange, and top systems where appropriate.

I. INTRODUCTION

Many extensions of the standard model (SM) predict anomalous effects in rare decays of beauty, charmed, and strange hadrons that could significantly alter their decay rate with respect to SM expectations. In $B$ meson decays, the experimental sensitivity has reached the SM expected rates for many rare processes. In contrast, GIM suppression \cite{1} in $D$ meson decays is significantly stronger and the SM branching fractions, in the case of radiative $D$ meson decays, are expected to be as low as $10^{-9}$ \cite{2, 3}. This leaves a large window of opportunity still available to search for new physics in charm decays.

Annihilation topologies of charged mesons can be used to probe new charged current phenomena that would appear at tree level such as models with charged Higgs bosons \cite{4}. Here, the advantage is the SM decay rate can be precisely calculated and the rates are experimentally accessible. Given apriori knowledge of the decay constants and CKM elements, measurements of these processes can place strict bounds on new phenomena.

As a third generation particle, sizable corrections are expected to $B^+$ annihilation in SUSY models with high tan $\beta$ \cite{5}. Sensitivity to new physics in these decays are currently limited by statistics but will eventually be limited by errors in $V_{ub}$ and $f_B$. As a second generation particle, the corrections are expected to be less evident in $D^{*0}$ annihilation \cite{4}. However, statistics are now available to make precision measurements of both $D_s \rightarrow \tau \nu$ and $D_s \rightarrow \mu \nu$. The ratio of these channels can then provide an extremely clean test for models that do not preserve lepton universality.

Radiative meson decay and annihilation of neutral mesons are sensitive to tree level neutral current phenomena or almost any new particle that can interact at higher order through penguin or box diagrams. The SM rate is absent at tree level and thus always suppressed. The precision to which the SM rate can be calculated varies drastically depending on generation and topology. For radiative beauty transitions such as $K_L \rightarrow \pi^0 \nu \bar{\nu}$, precision calculations exist and the SM rates are expected to be accessible in the next generation of kaon experiments.

For radiative charmed hadron decays such as $c \rightarrow u l^+ l^-$, the SM rate is extremely difficult to estimate. However, given the present level of experimental sensitivity, the errors in imperfect cancellation through the GIM mechanism can be ignored and we can essentially treat these decays as forbidden. Thus at the current level of sensitivity, any signal in the charm sector would unambiguously signal new physics. This relation between current experimental sensitivity and SM expectations is also true for annihilation of neutral $B$ and $D$ mesons. In this situation, any improvement of experimental limits allows us to place further constraints on new phenomena.

II. EXPERIMENTAL ENVIRONMENTS

Results are available from a extremely diverse set of experiments. The cleanest environment is $e^+e^-$ at charm threshold such as CLEO-c. Here, beam constraints are a powerful tool in background reduction and CESR has now delivered enough luminosity at particular resonances to allow for competitive studies of rare decays.

Some of the largest charm samples are available at the B-factory experiments Belle and BaBar where the direct charm production cross section is similar to the $\Upsilon(4S)$ production cross section and all species of charmed hadrons are available in the same data set. Since the final state is dominated pions, the excellent particle ID capabilities of these experiments greatly reduces the combinatorial background in $D$ and $\Lambda_c$ decays where either multiple kaons or protons are present. While not at threshold, the isolation of direct charm production still allows for strong background reduction through global event variables such as the total and missing energy in the event.

Results are available from many fixed target experiments conducted in the last decade at Fermilab with the best limits on rare decays coming from FOCUS \cite{7} that set the bar for the current experiments. The large
boost and excellent vertexing capabilities of these experiments led to large high purity samples of all charm species. While these data sets have now been surpassed by other experiments, they still remind us of opportunities that will become available at LHCb or possibly future dedicated fixed target charm experiments at Fermilab [8] that will have similar analysis strategies but much larger data sets.

Run II of the Fermilab Tevatron has brought the study of rare charm to the energy frontier experiments DØand CDF. Here again all species are available and the enormous production cross sections more than compensate for the lower luminosity. However for rare decays, the large backgrounds lead to stringent limitations on the channels available for study and to date, only final states containing dimuons have been considered.

III. ANNIHILATION

A. Charged Meson Annihilation

New results are available this summer from the Belle Collaboration on the decay $D_s \to \mu \nu$ [9]. Belle reports

$$B(D_s \to \mu \nu) = (6.44 \pm 0.76 \pm 0.52) \times 10^{-3}.$$  

Combining this measurement with PDG’06 [10] and BaBar [11] and CLEO-c [12] measurements from 2007 indicate an experimental sensitivity on the order of 8% in this branching fraction and indicate that the ratio of experimental measurement to theoretical prediction for $D_s \to \tau \nu/D_s \to \mu \nu$ can now be determined to roughly 15%. This can be compared to the experiment to theory ratio in $B^+ \to \tau \nu$ that is measured to a precision of about 44%, the recently observed Belle measurement of $B \to D^+\tau \nu$ [13] that has a precision of about 30%, or the recent measurement of $t \to b \tau \nu$ production cross section with $t \to b \tau \nu$ [14] that also has a precision of about 30%. So while the $c \to \tau$ transition is not expected to have contributions as large as those in the top and $b$ systems, it makes up for it with both experimental and theoretical precision.

B. Neutral Meson Annihilation

The best limits on $D^0$ annihilation have recently been reported by the CDF [15] and BaBar [16] collaborations. For normalization purposes, both analyses first reconstruct a large sample of $D^+\tau \nu$ tagged $D^0 \to \pi^+\pi^-$ decays. CDF reconstructs about 1.4k $D^0 \to \pi^+\pi^-$ decays in a 65 pb$^{-1}$ data sample while BaBar reconstructs greater than 7k $D^0 \to \pi^+\pi^-$ decays in a 122 fb$^{-1}$ data sample. The CDF analysis focuses on the dimuon final state while BaBar reconstructs both $\mu^+\mu^-$ and $e^+e^-$. The possible peaking background from double misidentification of $D^0 \to \pi^+\pi^-$ as $\mu^+\mu^-$ is studied using large samples of $D^*\tau \nu$ tagged $D^0 \to K\pi$ decays. CDF sets a 90% CL upper limit of

$$\mathcal{B}(D^0 \to \mu^+\mu^-) < 2.5 \times 10^{-6},$$

while BaBar sets 90% CL upper limits of

$$\mathcal{B}(D^0 \to \mu^+\mu^-) < 1.3 \times 10^{-6},$$

$$\mathcal{B}(D^0 \to e^+e^-) < 1.2 \times 10^{-6}.$$  

The final dilepton invariant mass distributions are shown in Fig. 1.

![CDF Run II Preliminary](image)

FIG. 1: Dilepton invariant mass distributions from CDF in the dimuon channel (above) and BaBar in the dimuon and dielectron channels in the $D^0 \to l^+l^-$ analyses.

IV. RADIATIVE DECAY

The first radiative charm decay to be observed is the decay $D_s \to \phi \gamma$ [17] where Belle measures

$$\mathcal{B}(D_s \to \phi \gamma) = (2.6^{+0.70}_{-0.61} \pm 0.15) \times 10^{-5}.$$
This is a beautiful measurement where many of the peaking backgrounds such as $\phi\pi^0$ and $\phi\eta$ could not be constrained using previous information and thus were concurrently measured along with the $\phi\gamma$ final state.

This result is also an excellent example of the inherent problems caused by long distance effects in the charm system. In the above channel, one can not distinguish between the quark level $c\bar{u} \rightarrow s\bar{s}\gamma$ transition and long distance rescattering of intermediate $D^0 \rightarrow \phi\rho$ or $\phi\omega$ transitions into the $\phi\gamma$ final state. Since the rate of these final state interactions can not be calculated with acceptable precision, no limits can be placed on new phenomena using the above channel [18].

This situation can be solved by moving from two-body to three-body radiative decays where the extra kinematic information in the final state allows for a separation of long distance and short distance components [2, 3]. For instance in the decay $D^+ \rightarrow \pi^+\mu^+\mu^-$ the long distance rescattering of $\phi \rightarrow \mu^+\mu^-$ can be extracted from the dimuon invariant mass spectra. Since the short distance component is expected to be three orders of magnitude below the long distance component, any excess in the dimuon mass spectra away from the $\phi$ resonance would clearly indicate new physics.

The best limits on the $c \rightarrow ul^+l^-$ transition come from CLEO-c [19], BaBar [20], and DØ [21]. The CLEO-c analysis is based on a data sample of 281 pb$^{-1}$ recorded at the $\psi(3770)$ resonance. The excellent calorimetry at CLEO-c leads to a focus on the $\pi^+e^-e^-$ final state. The BaBar analysis is based on 281 fb$^{-1}$. The combination of powerful hadron and lepton ID systems allow BaBar to search for both dimuon and dielectron final states of $D^+$, $D_s$, and $\Lambda_c$. The DØ analysis is based on a 1.3 fb$^{-1}$ data sample. The excellent dimon trigger system leads to a focus on the dimuon final state. Since the background reduction techniques rely heavily on secondary vertices reconstructed away from the interaction point, focus is placed on the $D^+$ meson rather than the $D_s$ or $\Lambda_c$ due to their shorter lifetimes.

As a first step, all three collaborations attempt to establish the long distance component $D^+ \rightarrow \phi\pi^+ \rightarrow l^+l^-\pi^+$ by requiring the dilepton invariant mass be consistent with a $\phi$. The results are shown in Fig. 2. CLEO-c finds two events with an expected background of 0.04 events. BaBar sees 19 signal events over a background of about 30 events. DØ sees 115 signal events over a background of roughly 850 events. The differences in the environments are clearly seen in these yield and background comparisons. The three collaborations measure

$\mathcal{B}(D^+ \rightarrow \phi\pi^+ \rightarrow e^+e^-\pi^+) =$

$$2.7^{+3.6}_{-1.8} \pm 0.2 \times 10^{-6} \text{ (CLEO)},$$

$$2.7^{+3.6}_{-1.8} \pm 0.2 \times 10^{-6} \text{ (BaBar)},$$

FIG. 2: Results of the search for $D^+ \rightarrow \pi\phi \rightarrow \pi l^+l^-$. The top figure is the beam constrained mass versus beam energy difference from CLEO in the dielectron channel. The middle figure is the $\pi e^+e^-$ invariant mass from BaBar. The lower figure is the $\pi\mu^+\mu^-$ invariant mass distribution from DØ.

With the long distance contribution established, each analysis proceeds to search for the short distance $c \rightarrow ul^+l^-$ transition by looking for an excess of events away from the $\phi$ resonance. CLEO-c takes advantage of beam constrained variables and detector hermiticity to specifically veto the dominant background of two semileptonic $D$ decays and arrives at a 90% CL upper limit of

$$\mathcal{B}(D^+ \rightarrow \pi^+e^+e^-) < 7.4 \times 10^{-6} \text{ (CLEO)}. $$
The BaBar analysis requires high momentum $D$ candidates consistent with direct $c\bar{c}$ production to remove backgrounds from semileptonic $B$ decay and then also relies on hermiticity to remove backgrounds from two semileptonic charm decays. Using $\Lambda_c$ decays easily distinguished using particle ID, they set the best 90% CL upper limit in the dielectron channel of

$$B(\Lambda_c \rightarrow p e^+ e^-) < 3.6 \times 10^{-6} \text{ (BaBar)}.$$  

The missing energy resolution of the DØdetector does not allow them to veto semileptonic events where the neutrinos typically carry away a few GeV of energy and the long lived backgrounds from semileptonic charm and $b$ hadron decay are essentially irreducible. However the much more dominant background is from light quark and Drell-Yann production that can be removed using flight length significance, vertex quality, and topological requirements and attempts are made to optimize the analysis for both direct $D$ meson production and $D$ mesons produced in $B$ meson decay. Background reduction based on these variables allow DØto set the best 90% upper limit in the dimuon channel of

$$B(D^+ \rightarrow \pi^+ e^+ e^-) < 3.9 \times 10^{-6} \text{ (D/O)}.$$  

The results are shown in Fig. 3. Since many scenarios of new phenomena predict different rates of excess in the dimuon and dielectron channels, its encouraging that together BaBar and DØcan cover both channels.

In conclusion, the last round of results in rare charm decays is producing precision measurements of $D_s$ annihilation branching fractions. The combination of statistical power and results in both the $\tau\nu$ and $\mu\nu$ channel may help add to knowledge recently gained from measurements of $B^+ \rightarrow \tau\nu$, $B \rightarrow D^*\tau\nu$ and $t \rightarrow b\tau\nu$.

The last round of results has also pushed limits on neutral annihilation and radiative decay from the $10^{-5}$ level to the $10^{-6}$ level with much of the data currently on tape yet to be analyzed. A complete analysis of the full B factory and Tevatron data sets as well as data at a super B factory and LHCb should push these results to the $10^{-7}$ level and hopefully yield an anomalous excess.

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I would like to thank the organizers for an excellent conference on a beautiful campus. I would also like to thank Paoti Chang, Jim Olsen, and Brian Peterson for help with the $B$ factory results. I would also like to acknowledge the papers of Burdman, Golowich, Hewett, and Pakvasa as well as those of Fajfer, Prelovsek, and Singer that played an important role in motivating these experimental studies. It only takes one person per experiment to continue a healthy rare charm decay program.
[14] V. Abazov et al., (D0Collaboration), D0note 5451-CONF.
Recent studies of Charmonium Decays at CLEO

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Department of Physics and Astronomy, University of Rochester, Rochester, New York 14627, USA

Recent results on Charmonium decays are reviewed which includes two-, three- and four-body decays of $\chi_{cJ}$ states, observations of $Y(4260)$ through $\pi\pi J/\psi$ transitions, precise measurements of $M(D^0)$, $M(\eta)$ as well as $B(\eta \rightarrow X)$.

1. Introduction

Decays of a bound state of a quark and its antiquark, quarkonium, provide an excellent laboratory for studying QCD. Particularly, heavy quarkonia such as charmonium states are less relativistic, thus play a special role in probing strong interactions.

CLEO recently has accumulated data taken at the $\psi(2S)$ resonance, providing a total of 27M $\psi(2S)$ decays. With the combination of this large statistical sample and the excellent CLEO detector, we will explore an unprecedented world of charmonia. While many analyses are currently being carried out, in this note we present recent results on multi-body $\chi_{cJ}$ decays which employed the pre-existing 3M $\psi(2S)$ sample.

We also present recent studies on decays of one of the exotic states, $Y(4260)$, as well as precision measurement on $M(D^0)$ that has an implication on properties of X(3872).

Finally, based on the full sample of $\psi(2S)$ data, we have results on properties of one of the light mesons, $\eta$.

2. Factory of $\chi_{c}(1^3P_J)$ states

$\chi_{c}(1^3P_J)$ states, which have one unit of orbital angular momentum and total spin of $J=0$, 1, or 2, cannot be produced directly from $e^+e^-$ collisions. They can be reached from $\psi(2S)$ through radiative (electronic dipole) transitions. Since $B(\psi(2S) \rightarrow \gamma \chi_{cJ}) = (9.3 \pm 0.4, 8.8 \pm 0.4, and 8.1 \pm 0.4) \times 10^{-2}$ for $J=0, 1,$ and 2 respectively [1], 27M $\psi(2S)$ decays of the new data provides 2-2M decays of each spin state of $\chi_{cJ}$ which should give us a greater understanding of the decay mechanisms of the $\chi_{cJ}$ mesons.

In this section, we present recent results of studies of $\chi_{cJ}$ decays based on 3M $\psi(2S)$ decays which should serve as the foundation for the future precision measurements by employing the full data sample of 27M of $\psi(2S)$ decays.

2.1. Two-body decay

We present results on $\chi_{cJ}$ decay into combinations of $\eta$ and $\eta'$. Figure 1 shows invariant masses of combinations of $\eta$ and $\eta'$. No $\chi_{c1}$ is seen as expected from conservation of spin-parity.

![Figure 1: Invariant masses of $\eta\eta$ (a), $\eta'\eta$ (b), and $\eta'\eta'$ (c)](image)

We measured $B(\chi_{c0} \rightarrow \eta\eta)$ to be $(0.31 \pm 0.05 \pm 0.04 \pm 0.02)\%$ [2] where the first uncertainty is statistical, the second is systematic, and the third is systematic due to the uncertainty in $B(\psi(2S) \rightarrow \chi_{cJ})$. This is slightly higher, but consistent with, the two previously published measurements. The BES Collaboration measured this branching ratio to be $(0.194 \pm 0.085 \pm 0.059)\%$ [3] and the E-835 Collaboration [4] had $(0.198 \pm 0.068 \pm 0.037)\%$. We also measured $B(\chi_{c0} \rightarrow \eta'\eta')$ to be $(0.17 \pm 0.04 \pm 0.02 \pm 0.01)\%$ for the first time. We set upper limits for $B(\chi_{c0} \rightarrow \eta'\eta') < 0.05\%$, $B(\chi_{c2} \rightarrow \eta\eta) < 0.047\%$, $B(\chi_{c2} \rightarrow \eta'\eta') < 0.023\%$, and $B(\chi_{c2} \rightarrow \eta'\eta') < 0.031\%$ at 90% confidence level.

Our result can be compared to predictions based on the model of Qiang Zhao [5]. He translates these decay rates into a QCD parameter, $r$, which is the ratio of doubly- to singly-OZI suppressed decay diagrams. In his model, our results indicate that the singly-OZI suppressed diagram dominates in these decays.
2.2. Three-body decay

We have also looked at three-body decays of $\chi_{cJ}$ states (one neutral and 2 charged hadrons) [6]. They are $\pi^+\pi^-\eta$, $K^+K^-\eta$, $p\bar{p}\eta$, $\pi^+\pi^-\eta'$, $K^+K^-\pi^0$, $p\bar{p}\pi^0$, $\pi^+K^-K_0^S$, and $K^0\bar{p}\Lambda$. Measured branching fractions are summarized in Table I. Again, our results are consistent with the results from BES Collaboration [7], with better precision.

In three of the above modes we looked for, $\pi^+\pi^-\eta$, $K^+K^-\pi^0$, and $\pi^+K^-K_0^S$, we observed significant signals of $\chi_{c1}$ decays which are shown in Figure 2.

![Figure 2: Invariant masses of $\pi^+\pi^-\eta$ (top-left), $K^+K^-\pi^0$ (top-right), and $K_0^S\bar{K}^+\pi^+$ (bottom).](image)

We performed Dalitz plot analyses based on these 3 modes in which we neglected any possible interference effects between resonances and polarization of $\chi_{c1}$. We estimated there could be $\sim 20(15)\%$ variations in fit fractions for the $\pi\pi\eta$ ($KK\pi$) mode due to such a simplified model. Figure 3 shows the Dalitz plot for $\pi^+\pi^-\eta$ and Table II shows its resultant fit fractions for each source. It is interesting to note that our data demand a relatively large yield of a $\sigma$ pole.

As for the $KK\pi$ mode, we performed simultaneous fits between $\chi_{c1} \rightarrow K^+K^-\pi^0$ and $\chi_{c1} \rightarrow K_0\bar{K}\pi$ by taking advantage of isospin symmetry. Dalitz plots for these modes are shown in Figures 4 and 5. Table III shows their resultant fit fractions.

![Figure 3: Dalitz plots for $\chi_{c1} \rightarrow \eta\pi^+\pi^-$.](image)

![Figure 4: Dalitz plots for $\chi_{c1} \rightarrow K^+K^-\pi^0$.](image)

2.3. Four-body decay

We present a preliminary result on four-body decay of $\chi_{cJ}$ states in which we reconstructed $h^+h^-\pi^0\pi^0$, where $h = \pi$, $K$, $p$; $K^+K^-\eta\pi^0$; and $K^+\pi^+K_0^S\pi^0$. Results of this kind of study, many-body decays of $\chi_{cJ}$ states, should help to build a comprehensive understanding about the P-wave dynamics.

Clean signals were seen in all modes except $\chi_{c1} \rightarrow p\bar{p}\pi^0\pi^0$ for the first time as can be seen in Figure 6. Many resonant substructures were also seen for which we only considered significant ones ($\pi^+\pi^-\pi^0\pi^0$ and $K^0\pi^+\bar{K}^0\pi^0$). The results are summarized in Table V.
Table I Branching fractions in units of $10^{-3}$. Uncertainties are statistical, systematic due to detector effects plus analysis methods, and a separate systematic due to uncertainties in the $\psi(2S)$ branching fractions. Limits are at the 90% confidence level.

<table>
<thead>
<tr>
<th>Mode</th>
<th>$\chi_{c0}$</th>
<th>$\chi_{c1}$</th>
<th>$\chi_{c2}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\pi^+\pi^-\eta$</td>
<td>$&lt; 0.21$</td>
<td>$5.0 \pm 0.3 \pm 0.4 \pm 0.3$</td>
<td>$0.49 \pm 0.12 \pm 0.05 \pm 0.03$</td>
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<tr>
<td>$K^+K^-\eta$</td>
<td>$&lt; 0.24$</td>
<td>$0.34 \pm 0.10 \pm 0.03 \pm 0.02$</td>
<td>$&lt; 0.33$</td>
</tr>
<tr>
<td>$pp\eta$</td>
<td>$0.39 \pm 0.11 \pm 0.04 \pm 0.02$</td>
<td>$&lt; 0.16$</td>
<td>$0.19 \pm 0.07 \pm 0.02 \pm 0.01$</td>
</tr>
<tr>
<td>$\pi^+\pi^-\eta'$</td>
<td>$&lt; 0.38$</td>
<td>$2.4 \pm 0.4 \pm 0.2 \pm 0.2$</td>
<td>$0.51 \pm 0.18 \pm 0.05 \pm 0.03$</td>
</tr>
<tr>
<td>$K^+K^-\pi^0$</td>
<td>$&lt; 0.06$</td>
<td>$1.95 \pm 0.16 \pm 0.18 \pm 0.14$</td>
<td>$0.31 \pm 0.07 \pm 0.03 \pm 0.02$</td>
</tr>
<tr>
<td>$pp\pi^0$</td>
<td>$0.59 \pm 0.10 \pm 0.07 \pm 0.03$</td>
<td>$0.12 \pm 0.05 \pm 0.01 \pm 0.01$</td>
<td>$0.44 \pm 0.08 \pm 0.04 \pm 0.03$</td>
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<tr>
<td>$\pi^+K^-\pi^0$</td>
<td>$&lt; 0.10$</td>
<td>$8.1 \pm 0.6 \pm 0.6 \pm 0.5$</td>
<td>$1.3 \pm 0.2 \pm 0.1 \pm 0.1$</td>
</tr>
<tr>
<td>$K^+\bar{p}\Lambda$</td>
<td>$1.07 \pm 0.17 \pm 0.10 \pm 0.06$</td>
<td>$0.33 \pm 0.09 \pm 0.03 \pm 0.02$</td>
<td>$0.85 \pm 0.14 \pm 0.08 \pm 0.06$</td>
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</table>

Table II Fit results for $\chi_{c1} \rightarrow \eta\pi^+\pi^-$ Dalitz plot analysis. The uncertainties are statistical and systematic. Allowing for interference among the resonances changes the fit fractions by as much as 20% in absolute terms as discussed in the text.

<table>
<thead>
<tr>
<th>Mode</th>
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<td>$a_0(980)^+\pi^-$</td>
<td>75.1 $\pm$ 3.5 $\pm$ 4.3</td>
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<tr>
<td>$f_2(1270)\eta$</td>
<td>14.4 $\pm$ 3.1 $\pm$ 1.9</td>
</tr>
<tr>
<td>$\sigma\eta$</td>
<td>10.5 $\pm$ 2.4 $\pm$ 1.2</td>
</tr>
</tbody>
</table>

Table III Results of the combined fits to the $\chi_{c1} \rightarrow K^+K^-\pi^0$ and $\chi_{c1} \rightarrow \pi KK^0$ Dalitz plots. Allowing for interference among the resonances changes the fit fractions by as much as 15% in absolute terms as discussed in the text.

<table>
<thead>
<tr>
<th>Mode</th>
<th>Fit Fraction (%)</th>
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<tr>
<td>$K^+(9892)K^-$</td>
<td>31.4 $\pm$ 2.2 $\pm$ 1.7</td>
</tr>
<tr>
<td>$K_0^*(1430)K$</td>
<td>30.4 $\pm$ 3.5 $\pm$ 3.7</td>
</tr>
<tr>
<td>$K_0^*(1430)K^0$</td>
<td>23.1 $\pm$ 3.4 $\pm$ 7.1</td>
</tr>
<tr>
<td>$a_0(980)/\pi$</td>
<td>15.1 $\pm$ 2.7 $\pm$ 1.5</td>
</tr>
</tbody>
</table>

Figure 5: Dalitz plots for $\chi_{c1} \rightarrow \pi KK^0$.

The measured branching fraction of $\chi_{cJ} \rightarrow \rho^+\pi^-\pi^0$ is consistent with that of $\chi_{cJ} \rightarrow \rho^0\pi^+\pi^-$ as expected from isospin symmetry. Similar isospin symmetry is also seen in Table IV where the partial width of $\chi_c \rightarrow K^{*0}K^{0}\pi^0$ and that of $\chi_c \rightarrow K^{*\pm}K^{\mp}\bar{K}^0$ are expected to be equal. Table IV also shows another good agreement with the isospin expectation of $B(\chi_c \rightarrow K^{*0}K^{0}\pi^0)/B(\chi_c \rightarrow K^{*0}K^{\mp}\pi^\pm) = 0.5$ and $B(\chi_c \rightarrow K^{*0}K^{0}\pi^0)/B(\chi_c \rightarrow K^{*\pm}K^{\mp}\bar{K}^0) = 0.5$.

Figure 6: Preliminary result on 4-body decays of $\chi_{cJ}$. Invariant masses of various combinations of hadrons are shown.
Table IV Branching fractions and combined error measurements for the isospin related $K^* \pi$ intermediate modes are listed.

<table>
<thead>
<tr>
<th>Mode</th>
<th>$\chi_{c0}$ B.F. (%)</th>
<th>$\chi_{c1}$ B.F. (%)</th>
<th>$\chi_{c2}$ B.F. (%)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$K^{\pm} K^0 \pi^0$</td>
<td>0.56±0.15</td>
<td>0.38±0.11</td>
<td>0.59±0.14</td>
</tr>
<tr>
<td>$K^{*0} K^{\pm} \pi^0$</td>
<td>-</td>
<td>-</td>
<td>0.90±0.25</td>
</tr>
<tr>
<td>$K^{+} K^0 \pi^0$</td>
<td>0.74±0.18</td>
<td>-</td>
<td>0.57±0.13</td>
</tr>
<tr>
<td>$K^0 K^{+} \pi^0$</td>
<td>0.96±0.25</td>
<td>-</td>
<td>0.90±0.25</td>
</tr>
</tbody>
</table>

The observation of $Y(4260) \rightarrow \pi^0 \pi^0 J/\psi$ is inconsistent with the $\chi_{cJ} \eta^0$ molecular model [11]. Our observation of $\pi^0 \pi^0 J/\psi$ rate being about half of $\pi^+ \pi^- J/\psi$ rate disagrees with the prediction of the baryonium model [12]. Evidence for the $K^+ K^- J/\psi$ signal is not compatible with these two models either. Table VI also shows that $Y(4160)$ behaves very differently compared to other charmonium states above $D\bar{D}$ threshold such as $\psi(4040)$ and $\psi(4160)$ for which we set upper limits in terms of cross section ($\sigma(e^+e^- \rightarrow X)$) and branching fractions.

## 3. Charmonium-like states above DD

### 3.1. Y(4260)

$Y(4260)$ was first discovered by the BaBar Collaboration via the reaction of $e^+e^- \rightarrow \gamma Y(4260) \rightarrow \gamma \pi^+ \pi^- J/\psi$, $J/\psi \rightarrow \ell^+ \ell^-$. Through the same production mechanism of initial state radiation, we also confirmed their observation based on data taken around the $Y(nS)$ resonances, where $n$ is 1, 2, 3, and, 4 [9]. The invariant mass of $\pi^+ \pi^- J/\psi$ based on such ISR production is shown in Figure 7.

![Figure 7: Distribution of invariant mass of $\pi^+ \pi^- J/\psi$ produced in ISR.](3850605-001)

The observation of $Y(4260) \rightarrow \pi^0 \pi^0 J/\psi$ is inconsistent with the $\chi_{cJ} \eta^0$ molecular model [11]. Our observation of $\pi^0 \pi^0 J/\psi$ rate being about half of $\pi^+ \pi^- J/\psi$ rate disagrees with the prediction of the baryonium model [12]. Evidence for the $K^+ K^- J/\psi$ signal is not compatible with these two models either. Table VI also shows that $Y(4160)$ behaves very differently compared to other charmonium states above $D\bar{D}$ threshold such as $\psi(4040)$ and $\psi(4160)$ for which we set upper limits in terms of cross section ($\sigma(e^+e^- \rightarrow X)$) and branching fractions.

## 3.2. X(3872) and Mass of neutral D meson

Since $X(3872)$ was discovered by Belle Collaboration [13] and subsequently confirmed by other experiments [14,15,16], many theoretical models have been proposed. Perhaps the most provocative theoretical suggestion is that $X(3872)$ is a loosely bound state of $D^0$ and $\bar{D}^{\ast 0}$ mesons [17].

### 4. Properties of $\eta$

It has been almost half a century since the $\eta$ meson was discovered [19]. Since then, many measurements have been made by many experiments. Still, almost all exclusive branching fractions are determined as relatives to other $\eta$ decays.

Based on the $27M \psi(2S)$ sample, we measured almost all the major modes (99% of generic decays of $\eta$) which allowed us to determine the major branching fractions [20]. We obtained the $\eta$ sample through...
Table V Branching fractions (B.F.) with statistical and systematic uncertainties are shown. The symbol "×" indicates product of B.F.’s. The third error in each case is due to the $\psi(2S) \rightarrow \gamma \chi_c$ branching fractions. Upper limits shown are at 90% C.L and include all the systematic errors. The measurements of the three-hadron final states are inclusive branching fractions, and do not represent the amplitudes for the three-body non-resonant decays.

<table>
<thead>
<tr>
<th>Mode</th>
<th>$\chi_{c0}$ B.F. (%)</th>
<th>$\chi_{c1}$ B.F. (%)</th>
<th>$\chi_{c2}$ B.F. (%)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\pi^+ \pi^- \pi^0 \pi^0$</td>
<td>3.54±0.10±0.43±0.18</td>
<td>1.28±0.06±0.16±0.08</td>
<td>1.87±0.07±0.23±0.13</td>
</tr>
<tr>
<td>$\rho^+ \pi^- \pi^0$</td>
<td>1.48±0.13±0.18±0.08</td>
<td>0.78±0.09±0.09±0.05</td>
<td>1.12±0.08±0.14±0.08</td>
</tr>
<tr>
<td>$\rho^- \pi^+ \pi^0$</td>
<td>1.56±0.13±0.19±0.08</td>
<td>0.78±0.09±0.09±0.05</td>
<td>1.11±0.09±0.13±0.08</td>
</tr>
<tr>
<td>$K^+ K^- \pi^0 \pi^0$</td>
<td>0.59±0.05±0.08±0.03</td>
<td>0.12±0.02±0.02±0.01</td>
<td>0.21±0.03±0.03±0.01</td>
</tr>
<tr>
<td>$p\bar{p}\pi^0\pi^0$</td>
<td>0.11±0.02±0.02±0.01</td>
<td>&lt; 0.05</td>
<td>0.08±0.02±0.01±0.01</td>
</tr>
<tr>
<td>$K^+ K^- \eta \pi^0$</td>
<td>0.32±0.05±0.05±0.02</td>
<td>0.12±0.03±0.02±0.01</td>
<td>0.13±0.04±0.02±0.01</td>
</tr>
<tr>
<td>$K^+ K^- K^0 \pi^0$</td>
<td>2.64±0.15±0.31±0.14</td>
<td>0.92±0.09±0.11±0.06</td>
<td>1.41±0.10±0.16±0.10</td>
</tr>
<tr>
<td>$K^0 K^0 \pi^0 \times K^0 \rightarrow K^\pm \pi^\mp$</td>
<td>0.37±0.09±0.04±0.02</td>
<td>0.25±0.06±0.03±0.02</td>
<td>0.39±0.07±0.05±0.03</td>
</tr>
<tr>
<td>$K^0 \pi^0 \pi^0 \times K^0 \rightarrow K^0 \pi^0 \pi^0$</td>
<td>0.30±0.07±0.04±0.02</td>
<td></td>
<td></td>
</tr>
<tr>
<td>$K^\pm K^\pm \pi^0 \times K^\pm \rightarrow K^\pm \pi^0$</td>
<td>0.49±0.10±0.06±0.03</td>
<td></td>
<td></td>
</tr>
<tr>
<td>$K^\pm \pi^\mp \pi^\mp \times K^0 \rightarrow K^\pm \pi^0$</td>
<td>0.32±0.07±0.04±0.02</td>
<td></td>
<td></td>
</tr>
<tr>
<td>$\rho^\pm K^0 \pi^0$</td>
<td>1.28±0.16±0.15±0.07</td>
<td>0.54±0.11±0.06±0.03</td>
<td>0.42±0.11±0.05±0.03</td>
</tr>
</tbody>
</table>

Table VI For each mode $e^+ e^- \rightarrow X$, for three center-of-mass regions: the detection efficiency, $\epsilon$; the number of signal [background] events in data, $N_s$ [$N_b$]; the cross-section $\sigma(e^+ e^- \rightarrow X)$; and the branching fraction of $\psi(4040)$ or $\psi(4160)$ to $X, B$. Upper limits are at 90% CL. "–" indicates that the channel is kinematically or experimentally inaccessible.

<table>
<thead>
<tr>
<th>Channel</th>
<th>$\sqrt{s} = 3970 - 4060$ MeV</th>
<th>$\sqrt{s} = 4120 - 4200$ MeV</th>
<th>$\sqrt{s} = 4260$ MeV</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\pi^+ \pi^- J/\psi$</td>
<td>$\epsilon$ $N_s$ $N_b$ $\sigma$ $B$</td>
<td>$\epsilon$ $N_s$ $N_b$ $\sigma$ $B$</td>
<td>$\epsilon$ $N_s$ $N_b$ $\sigma$ $B$</td>
</tr>
<tr>
<td>($%$) ($\text{pb}$) ($10^{-3}$)</td>
<td>($%$) ($\text{pb}$) ($10^{-3}$)</td>
<td>($%$) ($\text{pb}$) ($10^{-3}$)</td>
<td></td>
</tr>
<tr>
<td>$\pi^+ \pi^- J/\psi$</td>
<td>37 12 3.7</td>
<td>9$^{+6}_{-3}$±2 &lt; 4</td>
<td>38 13 3.7</td>
</tr>
<tr>
<td>$\pi^0 \pi^0 J/\psi$</td>
<td>20 1 1.9</td>
<td>&lt; 8 &lt; 2</td>
<td>21 5 0.9</td>
</tr>
<tr>
<td>$K^+ K^- J/\psi$</td>
<td>7 1 0.07</td>
<td>&lt; 20 &lt; 5</td>
<td>21 3 0.07</td>
</tr>
<tr>
<td>$\eta J/\psi$</td>
<td>19 12 9.5</td>
<td>&lt; 29 &lt; 7</td>
<td>16 15 8.8</td>
</tr>
<tr>
<td>$\pi^0 J/\psi$</td>
<td>23 2 &lt; 10 &lt; 2</td>
<td>22 1 &lt; 6 &lt; 1</td>
<td>22 1 &lt; 12</td>
</tr>
<tr>
<td>$\eta J/\psi$</td>
<td>11 4 2.5</td>
<td>&lt; 23 &lt; 5</td>
<td>11 0 1.5</td>
</tr>
<tr>
<td>$\pi^+ \pi^- \pi^0 J/\psi$</td>
<td>21 1 &lt; 8 &lt; 2</td>
<td>21 0 &lt; 4 &lt; 1</td>
<td>22 0 &lt; 7</td>
</tr>
<tr>
<td>$\eta J/\psi$</td>
<td>6 1</td>
<td>&lt; 4</td>
<td>19 0</td>
</tr>
<tr>
<td>$\eta J/\psi$</td>
<td>12 0</td>
<td>&lt; 15 &lt; 4</td>
<td>15 0</td>
</tr>
<tr>
<td>$\omega J/\psi$</td>
<td>9 11 11.5</td>
<td>&lt; 234</td>
<td></td>
</tr>
<tr>
<td>$\gamma J/\psi$</td>
<td>26 9 8.1</td>
<td>&lt; 50 &lt; 11</td>
<td>26 11 8.7</td>
</tr>
<tr>
<td>$\gamma J/\psi$</td>
<td>25 6 8.0</td>
<td>&lt; 76 &lt; 17</td>
<td>26 10 8.6</td>
</tr>
<tr>
<td>$\pi^+ \pi^- \pi^0 \chi_{c1}$</td>
<td>6 0</td>
<td>&lt; 47 &lt; 11</td>
<td>8 0</td>
</tr>
<tr>
<td>$\pi^+ \pi^- \pi^0 \chi_{c2}$</td>
<td>4 0</td>
<td>&lt; 141 &lt; 32</td>
<td>8 0</td>
</tr>
<tr>
<td>$\pi^+ \pi^- \phi$</td>
<td>17 26 3.0</td>
<td>&lt; 12 &lt; 3</td>
<td>17 17 6.0</td>
</tr>
</tbody>
</table>

...two-body decay of $\psi(2S)$, $\psi(2S) \rightarrow J/\psi$ where $J/\psi$ subsequently decays to two leptons ($e^+ e^-$ or $\mu^+ \mu^-$). That is, we have sample of about 0.1M $\eta$ decays with a di-lepton tag on $J/\psi$.

We first constrained the invariant mass of di-leptons to be the known mass of $J/\psi$. We then combined the fitted $J/\psi$ with $\eta$ decay products and constrained further to be the mass of $\psi(2S)$. In this analysis, the $\eta$ decay modes we considered were $\eta \rightarrow \gamma \gamma$, $3\pi^0$, $\pi^+ \pi^- \pi^0$, $\pi^+ \pi^- \gamma$ and $e^+ e^- \gamma$. According to Ref. [1], the sum of these 5 rates amounts to 99.88% of the total $\eta$ decays. We then took ratios between efficiency-corrected yields separately for each of $J/\psi \rightarrow e^+ e^-$ and $\mu^+ \mu^-$ cases in which all lepton related systematic.
uncertainties were canceled. The resulting ratios of $\eta$ branching fractions are summarized in Table VII.

Figure 11 shows a graphical version of comparison in terms of ratios of branching fractions to the single most precise other measurements (top of Figure 11).

By assuming that the 5 exclusive channels we considered in this analysis cover the all $\eta$ decay modes, we were also able to extract absolute branching fractions of these 5 modes. Other possible $\eta$ decay modes are either now allowed and/or found to be less than 0.2% of generic decays $\eta$ [11]. We included 0.3% as a possible systematic uncertainty in the absolute branching fraction measurements. The results are summarized in Table VIII and also shown graphically in Figure 11 in terms of ratio of our branching fractions to the PDG 2006 global fit (bottom of Figure 11). Several of the relative and absolute branching fractions obtained in this analysis are either the most precise to date or first measurements.

Further more, we also measured the mass of $\eta$ meson [22]. This was motivated by two recent precision measurements that were inconsistent with each other. In 2002, the NA48 Collaboration reported $M_{\eta} = 547.843 \pm 0.030 \pm 0.041$ MeV [23], while in 2005, GEM Collaboration reported $M_{\eta} = 547.311 \pm 0.028 \pm 0.032$ MeV [24] which was 8 standard deviations below NA48’s result.

We used the same $\eta$ sample described previously in this Section but used only 4 decay modes, $\eta \rightarrow \gamma\gamma$, $3\pi^0$, $\pi^+\pi^-\pi^0$, and $\pi^+\pi^-\gamma$ while, again, Constraining masses of $J/\psi$ and $\psi(2S)$. Our result, the average of the 4 $\eta$ decay modes, is $M_{\eta} = 547.785 \pm 0.017 \pm 0.057$ MeV which has comparable precision to both NA48 and GEM results, but is consistent with the former and 6.5 standard deviations larger than the later. We note that the KLOE Collaboration also recently measured mass of the $\eta$ meson to be $547.873 \pm 0.007 \pm 0.031$ MeV which was presented at the 2007 Lepton-Photon conference [23].
Table VIII For each $\eta$ decay channel, absolute branching fraction measurements for $J/\psi \rightarrow e^+e^-$ and $J/\psi \rightarrow \mu^+\mu^-$ combined, with statistical and systematic uncertainties (middle column), as determined in this work. The last column shows the PDG fit result \[1\]. All but $\gamma\gamma$ are first measurements.

<table>
<thead>
<tr>
<th>Channel</th>
<th>this work (%)</th>
<th>PDG [1] (%)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\gamma\gamma$</td>
<td>38.45 ± 0.40 ± 0.36</td>
<td>39.38 ± 0.26</td>
</tr>
<tr>
<td>$3\pi^0$</td>
<td>34.03 ± 0.56 ± 0.49</td>
<td>32.51 ± 0.28</td>
</tr>
<tr>
<td>$\pi^+\pi^-\pi^0$</td>
<td>22.60 ± 0.35 ± 0.29</td>
<td>22.7 ± 0.4</td>
</tr>
<tr>
<td>$\pi^+\pi^-\gamma$</td>
<td>3.96 ± 0.14 ± 0.14</td>
<td>4.69 ± 0.11</td>
</tr>
<tr>
<td>$e^+e^-\gamma$</td>
<td>0.94 ± 0.07 ± 0.05</td>
<td>0.60 ± 0.08</td>
</tr>
</tbody>
</table>

5. Summary

I have presented confirmation of BaBar’s observation of $Y(4260)$ in di-pion transition to $J/\psi$ along with a new observation through neutral di-pion transition. Our precision measurement on $M(D^0)$ calls for more precise measurement on $M(X(3872))$. With 3M $\psi(2S)$ sample, we have results on two-, three-, and four-body decays of $\chi_cJ$ states in which many substructures were seen in three- and four-body modes. Dalitz plot analyses were done for the case of 3-body decays. More detailed analyses can be done with the full 27M $\psi(2S)$ sample. Using the 27M sample, we performed precision measurements on $B(\eta \rightarrow X)$ and $M(\eta)$.

References

Figure 11: Comparison of the results obtained in this analysis with the most precise measurements from other experiments [1,21] (top), and the PDG 2006 global fits [1] (bottom).
Experimental charmonium decay results from BES

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Based on 14 million $\psi(2S)$ and 58 million $J/\psi$ events collected by the BESII detector, the leptonic decay of $\psi(2S)$ into $\tau^+\tau^-$, $\psi(2S)$ multi-body decays, $\chi_{cJ}$ decays, and $J/\psi$ hadronic decays are studied, and the branching fractions of these decays are reported. These results may shed light on the understanding of QCD.

I. INTRODUCTION

The Beijing spectrometer (BES) is a general purpose solenoidal detector at the Beijing Electron Positron Collider (BEPC) storage ring, which operates at center-of-mass energies from 2 to 5 GeV. The BES detector (BESIII) is described in Ref. [1].

In this paper, we focus on studies of the $\psi(2S)$ leptonic decay, $\psi(2S)$ radiative decays, $\psi(2S)$ hadronic decays, $\chi_{cJ}$ decays, and $J/\psi$ decays based on 14 million $\psi(2S)$ and 58 million $J/\psi$ events collected by the BESII detector.

II. $\psi(2S)$ DECAYS

A. $\psi(2S) \rightarrow \tau\tau$

The $\psi(2S)$ provides an opportunity to compare the three lepton generations by studying the leptonic decays $\psi(2S) \rightarrow e^+e^-$, $\mu^+\mu^-$, and $\tau^+\tau^-$. The leptonic decay widths are approximately described by the relation $B_{ee} \approx B_{\mu\mu} \approx B_{e\tau}/0.3885$, which is in good agreement with the BESI $B_{e\tau}$ measurement [2]. Based on 14 million $\psi(2S)$ events, the branching fraction for $\psi(2S) \rightarrow \tau^+\tau^-$ is remeasured [3]. The $\tau^+\tau^-$ pairs are reconstructed from $\tau^+\tau^- \rightarrow \mu^+\nu_\mu e^-\bar{\nu}_e$ and $\tau^+\tau^- \rightarrow e^+\bar{\nu}_e\mu^-\nu_\mu$. At the $\psi(2S)$ resonance, 1015 signal events are observed, while 516 events are estimated to be from $e^+e^- \rightarrow \tau^+\tau^-$, determined using a data sample taken at $\sqrt{s} = 3.65$ GeV. The branching fraction is calculated to be $(0.310 \pm 0.021 \pm 0.038)\%$, where the first error is statistical and the second is systematic. Compared with the BESI result, the number of events is much bigger and the QED contribution and the efficiency and background estimations are improved.

B. Radiative decays

Besides the conventional meson and baryon states, QCD also predicts a rich spectrum of glueballs, meson hybrids, and multi-quark states in the 1.0 to 2.5 GeV/$c^2$ mass region. Therefore, searches for evidence of these exotic states play an important role to test QCD. Such studies of these states have been performed in $J/\psi$ radiative decays for a long time [4, 5], while studies in $\psi(2S)$ radiative decays have been limited due to low statistics in previous experiments [5, 6]. The radiative decays of $\psi(2S)$ to hadrons are expected to contribute about 1% to the total $\psi(2S)$ decay width [7]. However, the measured channels only sum up to about 0.05% [6].

Recently BESII measured the decays of $\psi(2S)$ into $\gamma p\bar{p}$, $\gamma 2(\pi^+\pi^-)$, $\gamma K^0S\bar{K}^0 + c.c.$, $\gamma K^+K^-\pi^+\pi^-$, $\gamma K^+K^-\pi^+\pi^-$, $\gamma K^0K^+\pi^+\pi^-$, $\gamma K^+K^-\pi^+\pi^-$, $\gamma 3(\pi^+\pi^-)$, and $\gamma 2(\pi^+\pi^-)K^+K^-$. The invariant mass of the hadrons ($m_{h\gamma}$) less than 2.9 GeV/$c^2$ for each decay mode [8]. The differential branching fractions are shown in Fig. 1. In the mass distribution of $m_{p\bar{p}}$ for $\psi(2S) \rightarrow \gamma p\bar{p}$, although there is some excess of events near $p\bar{p}$ threshold as shown in Fig. 1(a), no significant narrow structure due to the X(1859), observed in $J/\psi \rightarrow \gamma p\bar{p}$, is seen. The upper limit on the branching fraction is measured to be $B[\psi(2S) \rightarrow \gamma X(1859) \rightarrow \gamma p\bar{p}] < 5.4 \times 10^{-6}$ at the 90% C.L. There is a wide peak in the $m_{2(\pi^+\pi^-)}$ distribution located at $1.4 \sim 2.2$ GeV/$c^2$, but its structure can not be resolved due to the low statistics. No obvious structure is observed in the other final states. The branching fractions below $m_{h\gamma} < 2.9$ GeV/$c^2$ are given in Table I, and they sum up to 0.26% of the total $\psi(2S)$ decay width. This indicates that a larger data sample is needed to search for more decay modes and to resolve the substructure of $\psi(2S)$ radiative decays.

C. Hadronic decays

In perturbative QCD (pQCD), hadronic decays of both $\psi(2S)$ and $J/\psi$ proceed dominantly via the annihilation of $c\bar{c}$ quarks into three gluons or one photon, followed by a hadronization process. This yields the so-called pQCD “12% rule”, i.e.

$$Q_h = \frac{B_{\psi'\rightarrow h}}{B_{J/\psi\rightarrow h}} = \frac{B_{\psi'\rightarrow ee}}{B_{J/\psi\rightarrow ee}} \approx 12\%.$$  

A large violation of this rule was firstly observed in decays to $\rho\pi$ and $K^*K + c.c.$ by Mark II [9], known as the $\rho\pi$ puzzle. Since then there have been many more measurements of $\psi(2S)$ decays by BES and recently
by the CLEO collaboration for the study of the 12% rule. Table II summarizes recent measurements on ψ(2S) decays by BES. For the ψ(2S) decays listed, the $Q_h$ ratios are in agreement with the 12% rule within $1 \sim 2\sigma$, except for the obviously suppressed channel $\psi(2S) \rightarrow \omega f_2(1270)$.

The branching fractions of $\psi(2S)$ decays into octet baryons have also been measured; the results are listed in Table II. They are in agreement with the results published by the CLEO collaboration within $2\sigma$ for $p\bar{p}$ and within $1\sigma$ for the $\Lambda\Lambda$, $\Sigma^0\Sigma^0$, and $\Xi^-\Xi^+$ channels. For $\psi(2S) \rightarrow N\bar{N}\pi$, the ratios of the measured branching fractions of the $p\bar{p}\pi^0$ isospin partners are given by $B(\psi(2S) \rightarrow p\bar{p}\pi^0) : B(\psi(2S) \rightarrow p\bar{n}\pi^-) : B(\psi(2S) \rightarrow \bar{p}n\pi^+) = 1 : 1.86 \pm 0.27 : 1.91 \pm 0.27$, which is consistent with the isospin symmetry prediction $1 : 2 : 2$.

### III. $\chi_{cJ}$ DECAYS

#### A. $\chi_{cJ} \rightarrow \phi\phi$

Decays of $\chi_{cJ} \rightarrow K^+K^-K^+K^-$ are measured using 14 million $\psi(2S)$ decays [10]. The branching fractions including intermediate states are given in Table III. The decay $\chi_{cJ} \rightarrow \phi K^+K^-$ is observed for the first time, and the precision of the measurements $\chi_{cJ} \rightarrow \phi\phi$ and $K^+K^-K^+K^-$ are improved compared with PDG values.

The branching fractions of $\chi_{cJ} \rightarrow \phi\phi$ together with previous BES measurements on $\chi_{cJ} \rightarrow \omega\omega$ [12] and $\chi_{cJ} \rightarrow K^*(892)\bar{K}^*(892)$ [13] are used to predict the decay branching fractions of $\chi_{cJ}$ to other vector meson pairs, like $pp$ and $\omega\phi$ [14], where a large double OZI suppressed amplitude is expected.
TABLE III: Summary of $\chi_{cJ}$ hadronic decays. Upper limits are given at the 90% C.L. For $\chi_{cJ} \to K^0_S K^- + c.c.$ and $\eta \pi \pi$, branching fractions of $\text{Br}(\psi' \to \gamma \chi_{c0}) = (8.6 \pm 0.7)\%$, $\text{Br}(\psi' \to \gamma \chi_{c1}) = (8.4 \pm 0.8)\%$, and $\text{Br}(\psi' \to \gamma \chi_{c2}) = (6.4 \pm 0.6)\%$ are used. For other decays, branching fractions of $\text{Br}(\psi(2S) \to \gamma \chi_{cJ})$ from CLEOc [11] are used.

<table>
<thead>
<tr>
<th>Decay mode X</th>
<th>$\text{Br}(\chi_{c0} \to X) \times 10^{-3}$</th>
<th>$\text{Br}(\chi_{c1} \to X) \times 10^{-3}$</th>
<th>$\text{Br}(\chi_{c2} \to X) \times 10^{-3}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>2$(K^+ K^-)$</td>
<td>3.48 $\pm$ 0.23 $\pm$ 0.47</td>
<td>0.70 $\pm$ 0.13 $\pm$ 0.10</td>
<td>2.17 $\pm$ 0.20 $\pm$ 0.31</td>
</tr>
<tr>
<td>$\phi K^+ K^-$</td>
<td>1.03 $\pm$ 0.22 $\pm$ 0.15</td>
<td>0.46 $\pm$ 0.16 $\pm$ 0.06</td>
<td>1.67 $\pm$ 0.26 $\pm$ 0.24</td>
</tr>
<tr>
<td>$\eta \pi \pi$</td>
<td>0.94 $\pm$ 0.21 $\pm$ 0.13</td>
<td>---</td>
<td>1.70 $\pm$ 0.30 $\pm$ 0.25</td>
</tr>
<tr>
<td>$K^0_S K^+ \pi^- + c.c.$</td>
<td>$&lt;0.35$</td>
<td>$4.0 \pm 0.3 \pm 0.5$</td>
<td>$0.8 \pm 0.3 \pm 0.1$</td>
</tr>
<tr>
<td>$\eta \pi^+ \pi^-$</td>
<td>$&lt;1.1$</td>
<td>$5.9 \pm 0.7 \pm 0.8$</td>
<td>$&lt;1.7$</td>
</tr>
</tbody>
</table>

B. $\chi_{cJ} \to K^0_S K^+, \eta \pi \pi$

Decays of $\chi_{c0}$ and $\chi_{c2}$ into three pseudo-scalars are highly suppressed by the spin-parity selection rule. BES measured the branching fractions of $\chi_{c1}$ decays into $K^0_S K^+ \pi^- + c.c.$ and $\eta \pi^+ \pi^-$, including the intermediate states involved [15]. The branching fractions or upper limits at the 90% C.L. are summarized in Table III.

The $K^0_S K^+ \pi^- + c.c.$ events are mainly produced via $K^+(892)$ intermediate states, and the $\eta \pi^+ \pi^-$ events via $f_2(1270)\eta$ or $a_0(980)\pi$. The branching fractions [15] for these resonances are

$$\text{Br}(\chi_{c1} \to K^+(892)K^0 + c.c.) = (1.1 \pm 0.4 \pm 0.2) \times 10^{-3},$$

$$\text{Br}(\chi_{c1} \to K^+(892)K^- + c.c.) = (1.6 \pm 0.7 \pm 0.3) \times 10^{-3},$$

$$\text{Br}(\chi_{c1} \to f_2(1270)\eta) = (3.0 \pm 0.7 \pm 0.5) \times 10^{-3},$$

$$\text{Br}(\chi_{c1} \to a_0(980)K^+ \pi^- + c.c. \to \eta \pi^+ \pi^-) = (2.0 \pm 0.5 \pm 0.5) \times 10^{-3}.$$
Charmonium Spectroscopy Below Open Flavor Threshold

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Latest experimental results in the charmonium spectroscopy below \( \bar{D}D \) breakup threshold are reviewed.

1. Introduction

Charm quark has large mass (~1.5 GeV) compared to the masses of \( u, d, s \) quarks. Velocity of the charm quarks in hadrons is not too relativistic \((v/c)^2 \sim 0.2\). Strong coupling constant \( \alpha_s(m_c) \) is small (~0.3). Therefore charmonium spectroscopy is a good testing ground for the theories of strong interactions: quantum chromodynamics (QCD) in both perturbative and nonperturbative regimes, QCD inspired pure phenomenological potential models, nonrelativistic QCD (NRQCD) and lattice QCD.

There are 8 bound states of charmonium below the \( \bar{D}D \) breakup threshold (Fig. 1). These are spin triplets \( J/\psi (1^1S_1), \psi (2S)(2^3S_1), \chi_{c0,1.2}(1^3P_0,1.2) \) and spin singlets \( \eta_c(1^1S_0), \eta_c(2S)(2^3S_0), h_c(1^1P_1) \). Only \( J/\psi \) and \( \psi (2S) \) can be produced directly in \( e^+e^- \) annihilation. A lot is known about these triplet states. Spin singlet states population via radiative transitions from the vector states is either very weak (M1 transitions for \( \eta_c(1S), \eta_c(2S) \)), or \( C \)-forbidden \( (h_c(1^1P_1)) \). Only \( J/\psi \) and \( \psi (2S) \) can be produced directly in \( e^+e^- \) annihilation. A lot is known about these triplet states. Spin singlet states population via radiative transitions from the vector states is either very weak (M1 transitions for \( \eta_c(1S), \eta_c(2S) \)), or \( C \)-forbidden \( (h_c(1^1P_1)) \).

Accordingly, little is known about these singlet states.

Table I Status of charmonium states.

<table>
<thead>
<tr>
<th>Spin</th>
<th>Mass (MeV)</th>
<th>Width (MeV)</th>
<th>Number of Decays PDG 2002</th>
<th>2004</th>
<th>2007*</th>
</tr>
</thead>
<tbody>
<tr>
<td>Triplets</td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>( J/\psi )</td>
<td>3096.92±0.01</td>
<td>93.4±2.1 (keV)</td>
<td>134</td>
<td>135</td>
<td>162</td>
</tr>
<tr>
<td>( \psi (2S) )</td>
<td>3686.09±0.03</td>
<td>327±11 (keV)</td>
<td>51</td>
<td>62</td>
<td>115</td>
</tr>
<tr>
<td>( \chi_{c0} )</td>
<td>3414.75±0.35</td>
<td>10.4±0.7</td>
<td>17</td>
<td>17</td>
<td>51</td>
</tr>
<tr>
<td>( \chi_{c1} )</td>
<td>3510.66±0.07</td>
<td>0.89±0.05</td>
<td>12</td>
<td>13</td>
<td>35</td>
</tr>
<tr>
<td>( \chi_{c2} )</td>
<td>3556.20±0.09</td>
<td>2.05±0.12</td>
<td>18</td>
<td>19</td>
<td>37</td>
</tr>
<tr>
<td>Singlets</td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>( \eta_c(1S) )</td>
<td>2979.8±1.2</td>
<td>26.5±3.5</td>
<td>20</td>
<td>21</td>
<td>31</td>
</tr>
<tr>
<td>( \eta_c(2S) )</td>
<td>3637±4</td>
<td>14±7</td>
<td>3</td>
<td>4</td>
<td>4</td>
</tr>
<tr>
<td>( h_c )</td>
<td>3525.93±0.27</td>
<td>&lt;1</td>
<td>3</td>
<td>3</td>
<td>4</td>
</tr>
</tbody>
</table>

It is obvious from Table I that the parameters of \( spin-triplet \) states are measured with precision, and the number of measured decay channels is large (notice the marked improvements after 2004). This is not valid for \( spin-singlet \) states. A lot remains to be done for precision measurements of their parameters and decay channels.

Some new and recent experimental developments on charmonium spectroscopy below open flavor threshold will be reviewed.

2. Observation of \( \eta_c(2S) \)

It is important to identify the spin-singlet states in order to determine the hyperfine, or \( spin-spin \) interaction, which is responsible for singlet-triplet splitting of \( q\bar{q} \) states. Identification of \( \eta_c(2S) \) is important to know the possible variation of spin-spin interaction from Coulombic \( (J/\psi, \eta_c(1S)) \) to Confinement \( (\psi (2S), \eta_c(2S)) \) regions of the \( q\bar{q} \) interaction. Most potential model calculations predicted \( M(\eta_c(2S))=3594-3626 \) (MeV).

Prior to 2002 there were several unsuccessful attempts to identify \( \eta_c(2S) \) in \( p\bar{p} \), \( \gamma\gamma \)-fusion, inclusive photon analysis.

Finally, \( \eta_c(2S) \) was first observed in \( B \) decays by Belle [1]. It was followed by its observation in \( \gamma\gamma \)-fusion by CLEO [2] (see Fig. 2 left), and BaBar [3].

Figure 1: Spectra of the states of charmonium below \( \bar{D}D \) breakup threshold.

The status of charmonium states below \( \bar{D}D \) breakup threshold is summarized in Table I. The masses and widths from PDG 2007, as well as a number of measured decay channels from PDG 2002, 2004 and 2007 are presented separately for \( spin-triplet \) and \( spin-singlet \) states.
3. Observation of $h_c(^1P_1)$

The observation and the measurement of the parameters of $h_c$ is important to determine the hyperfine splitting of P-states $\Delta M_{hf}(1P) \equiv M(<^3P_1>) - M(^1P_1)$, which is expected to be zero from the lowest order pQCD calculations.

Recently CLEO Collaboration has unambiguously identified $h_c$ in their data for 3 million $\psi(2S)$ events. Data have been analyzed for the reaction $\psi(2S) \to \pi^0h_c$, $h_c \to \gamma h_c$ in both inclusive and exclusive studies. In inclusive analysis, photon energy or $h_c$ mass (recoil against $\pi^0\gamma$) was constrained. In exclusive analysis seven known $h_c$ decay channels with a total branching fraction of $\sim 10\%$ were measured. Results of inclusive and exclusive analysis are consistent. The final plot of the $\pi^0$ recoil spectrum from exclusive analysis is shown in Fig. 2 (right).

The overall results are:

- $M(h_c) = (3524.4 \pm 0.6 \pm 0.4) \text{ MeV}$;
- $\Delta M_{hf}(1P) = (+1.0 \pm 0.6 \pm 0.4) \text{ MeV}$, using $M(<^3P_1>) = (3525.4 \pm 0.1) \text{ MeV}$;
- $B(\psi(2S) \to \pi^0h_c) \times B(h_c \to \gamma h_c) = (4.0 \pm 0.8 \pm 0.7) \times 10^{-4}$;
- Significance level of the $h_c$ signal is $> 6\sigma$.

The conclusions are that a) the lowest order pQCD expectation $\Delta M_{hf}(1P) = 0$ is not strongly violated, and b) the magnitude and the sign of $\Delta M_{hf}(1P)$ is not well determined.

The Fermilab E835 $p\bar{p}$ annihilation experiment has also claimed $h_c$ observation at $\sim 3\sigma$ level in the reaction $p\bar{p} \to h_c \to \gamma h_c$ and reported $\Delta M_{hf}(1P) = 0.4\pm0.2\pm0.2 \text{ MeV}$.

CLEOc now has new data with 24 million $\psi(2S)$ events, and a $h_c$ peaks with $\sim 250$ and $\sim 1000$ counts are expected in exclusive and inclusive analysis respectively, which will reduce the error on mass measurement more than a factor of two.

4. Measurements of the $\psi(2S)$ Widths

Using $p\bar{p}$ annihilation to form charmonium $c\bar{c}$ states, Fermilab experiment E835 achieved unprecedented precision in measuring masses and widths of charmonium resonances. This happens due to taking advantage of stochastically cooled antiproton beams, with FWHM energy spreads of 0.4–0.5 MeV in the center-of-mass frame.

Recently new precision measurement of the $\psi(2S)$ total width was performed from excitation curves obtained in $p\bar{p}$ annihilations from 1.64 $pb^{-1}$ scan data in the $\psi(2S)$ region, collected by E835 in 2000 [8]. The channels analyzed were $p\bar{p} \to e^+e^-$ and $p\bar{p} \to J/\psi X \to e^+e^- + X$. New technique of “complementary scans”, based on precise beam revolution-frequency and orbit-length measurements was used. Resonance parameters were extracted from a maximum-likelihood fit to the excitation curves. The total width of the $\psi(2S)$ and the combination of partial widths were measured:

- $\Gamma_{tot}(\psi(2S)) = (290 \pm 25 \pm 4) \text{ keV}$,
- $\Gamma_{e^+e^-}/\Gamma_{tot} = 579 \pm 38 \pm 36 \text{ MeV}$.

These represent the most precise measurements to date (Fig. 3).

BES has also measured recently the $\psi(2S)$ total width using $e^+e^-$ annihilation scan data in the $\psi(2S)$ and $\psi(3770)$ regions, collected by BES II in 2003 [8] (Fig. 3). They have analyzed the channel $e^+e^- \to $
Recent measurements of the \( \psi(2S) \) widths. Resonance parameters were extracted from simultaneous fit of cross section curves covering energy ranges of both \( \psi(2S) \) and \( \psi(3770) \) resonances: 
\[
\Gamma_{\text{tot}}(\psi(2S)) = (331 \pm 58 \pm 2) \text{ keV}, \\
\Gamma_{ee}(\psi(2S)) = (2.330 \pm 0.036 \pm 0.110) \text{ keV}.
\]

5. Measurements of the \( J/\psi \) Widths

Using \( 281 \, pb^{-1} \) CLEOc \( \psi(3770) \) data and looking for radiative return events to \( J/\psi \), CLEO has measured the widths of the \( J/\psi \) \[10\]. They selected \( \mu^+\mu^- (\gamma) \) events, each with a dimuon mass in the region of the \( J/\psi \), and counted the excess over non-resonant QED production. Resulting cross section is proportional to \( Br_{\mu\mu} \times \Gamma_{ee}(J/\psi) \). Assuming lepton universality and dividing by CLEO’s own measurement of \( Br_{ll}(J/\psi) \) \[11\], they obtained \( \Gamma_{ee}(J/\psi) \). Dividing once more by \( Br_{ll} \) they obtained \( \Gamma_{\text{tot}}(J/\psi) \): 
\[
Br_{\mu\mu} \times \Gamma_{ee}(J/\psi) = (0.3384 \pm 0.0058 \pm 0.0071) \text{ keV}; \\
\Gamma_{ee}(J/\psi) = (5.68 \pm 0.11 \pm 0.13) \text{ keV}; \\
\Gamma_{\text{tot}}(J/\psi) = (95.5 \pm 2.4 \pm 2.4) \text{ keV}.
\]
These represent the most precise measurements to date.

6. Measurements of the \( \eta_c(1S) \) and \( \chi_{c0,c2}(1P) \) Parameters in Two-Photon Fusion Reaction

The masses and widths of the spin-triplet \( \chi_{cJ}(1P) \) states are measured with high precision at Fermilab \( p\bar{p} \) experiments E760/E835. But the mass and width of the spin-singlet \( \eta_c(1S) \) state are known with only \(~1 \text{ MeV} \) and \(~3 \text{ MeV} \) precision respectively.

Using \( 395 \, fb^{-1} \) data sample accumulated with the Belle detector, measurements of the \( \eta_c(1S) \) and \( \chi_{c0,c2}(1P) \), produced in two-photon collisions and decaying to four-meson final states (4\( \pi \), 2\( K2\pi \), 4\( K \)) were performed \[12\] (Fig. 4).

6.1. Mass and width of the \( \eta_c(1S) \) and \( \chi_{c0,c2}(1P) \)

The measured values of the mass and width of \( \eta_c(1S) \) and \( \chi_{c0,c2}(1P) \), and the number of signal events used in analysis are presented in Table \[III\]. The values of the mass and width for \( \chi_{c0}(1P) \) and \( \chi_{c2}(1P) \) are consistent within errors with the previous high precision measurements. The precision of the measured mass and width of \( \eta_c(1S) \) is comparable to other available precision measurements.

Table III Mass and width measurements of the \( \eta_c(1S) \), \( \chi_{c0}(1P) \) and \( \chi_{c2}(1P) \) from \[12\].

<table>
<thead>
<tr>
<th>Resonance</th>
<th>Mass (MeV)</th>
<th>Width (MeV)</th>
<th>N(events)</th>
</tr>
</thead>
<tbody>
<tr>
<td>( \eta_c(1S) )</td>
<td>2986.1\pm1.0\pm2.5</td>
<td>28.1\pm3.2\pm2.2</td>
<td>7616\pm553</td>
</tr>
<tr>
<td>( \chi_{c0}(1P) )</td>
<td>3414.2\pm0.5\pm2.3</td>
<td>10.6\pm1.9\pm2.6</td>
<td>5459\pm319</td>
</tr>
<tr>
<td>( \chi_{c2}(1P) )</td>
<td>3555.3\pm0.6\pm2.2</td>
<td>-</td>
<td>2503\pm158</td>
</tr>
</tbody>
</table>

6.2. Two-photon widths of \( \eta_c(1S) \) and \( \chi_{c0,c2}(1P) \)

The two photon decay of the positive C-parity charmonium states in the lowest order is a pure QED process. The measurements of the two photon partial widths of these states can shed light on higher order relativistic and QCD radiative corrections.
The values of the two photon partial widths of \( \eta_c(1S) \) and \( \chi_{c0,2}(1P) \), evaluated from the measurements of the \( \Gamma_{\gamma\gamma} \times Br \) by Belle [12], are presented in Table IV. They are compared to the values obtained from the PDG 2007. The Belle value of \( \Gamma_{\gamma\gamma}(\eta_c) \) is \( \sim 2.7 \) times smaller (\( \sim 4\sigma \) difference) then the PDG 2007 value. The values of \( \Gamma_{\gamma\gamma}(\chi_{c0}) \) and \( \Gamma_{\gamma\gamma}(\chi_{c2}) \) are consistent within errors with those from the PDG 2007.

The ratio \( R \equiv \Gamma_{\gamma\gamma}(\chi_{c2})/\Gamma_{\gamma\gamma}(\chi_{c0}) \) is an interesting quantity, because it allows us to evaluate the reliability of the first order radiative corrections, which are often very large, by calculating \( \alpha_s \) from them

\[
R \equiv \frac{\Gamma_{\gamma\gamma}(\chi_{c2})}{\Gamma_{\gamma\gamma}(\chi_{c0})} = \frac{(4|\Psi(0)|^2\alpha_s^2/m_s^4)/(1-1.7\alpha_s)}{(15|\Psi(0)|^2\alpha_s^2/m_s^4)/(1+0.06\alpha_s)} = 0.267(1 - 1.76\alpha_s)
\]

The Belle value \( R=0.221\pm0.041 \) leads to \( \alpha_s=0.098\pm0.085 \), which is obviously underestimate of \( \alpha_s(m_c) \), which is known to be \( \sim 0.3 \). This makes questionable the reliability of nearly 50% first order correction factor for \( \Gamma_{\gamma\gamma}(\chi_{c2}) \).

### Table IV

Two photon partial widths \( \Gamma_{\gamma\gamma} \) of \( \eta_c(1S) \), \( \chi_{c0}(1P) \) and \( \chi_{c2}(1P) \) [12]. \( \Gamma_{\gamma\gamma} \) values are evaluated from measured \( \Gamma_{\gamma\gamma} \times Br \) using branching fractions from PDG 2007. Results of \( 4\pi, 2K2\pi \) and \( 4K \) channels are combined.

<table>
<thead>
<tr>
<th>Resonance</th>
<th>( \Gamma_{\gamma\gamma} ) (keV), Belle</th>
<th>( \Gamma_{\gamma\gamma} ) (keV), PDG 07</th>
</tr>
</thead>
<tbody>
<tr>
<td>( \eta_c(1S) )</td>
<td>2.46\pm0.60</td>
<td>6.7\pm0.9</td>
</tr>
<tr>
<td>( \chi_{c0}(1P) )</td>
<td>1.98\pm0.24</td>
<td>2.90\pm0.43</td>
</tr>
<tr>
<td>( \chi_{c2}(1P) )</td>
<td>0.438\pm0.062</td>
<td>0.539\pm0.050</td>
</tr>
<tr>
<td>( R \equiv \Gamma_{\gamma\gamma}(\chi_{c2})/\Gamma_{\gamma\gamma}(\chi_{c0}) )</td>
<td>0.221\pm0.041</td>
<td>0.186\pm0.032</td>
</tr>
</tbody>
</table>

### 7. Summary

All Charmonium states below open flavor threshold have now been firmly identified.

The spectroscopy of spin-triplet states is now well in hand, but a lot still needs to be done for spin-singlet states. Masses, widths, particularly of \( \eta_c(2S) \) and \( h_c(1^1P_1) \) need to be better determined. Many more decay channels need to be investigated for each.

A large number of investigations, based on the world’s largest sample of \( \psi(2S) \) acquired by CLEOc, are currently in progress, and results are expected soon. These include:
- Precision results for mass, width and branching fractions of \( h_c(1^1P_1) \);
- Results for many decay channels of \( \eta_c(1S) \);
- Results for attempt to identify \( \eta_c(2S) \) in radiative decay of \( \psi(2S) \);
- Results of studies for \( p\bar{p} \) threshold enhancement in radiative decays of \( J/\psi, \psi(2S) \);
- Results of search for tensor glueball, \( \xi(2230) \);
- Hadronic and radiative decays of \( \psi(2S) \) and \( J/\psi \);
- Two-body and multi-body decays of \( \chi_{cJ}(1P) \) states, and others.

### Acknowledgments

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### References

Recent Results in Bottomonium
Ties to Charmonium

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Recent results in studies of bottomonium especially relevant to charmonium are reviewed. This report covers dipion transition matrix elements, production in \( \Upsilon \) transitions, \( \Upsilon \) decays to invisible particles, a search for a non-standard-model pseudoscalar Higgs in \( \Upsilon \) radiative decays, and \( \Upsilon \) radiative decays to \( f_2(1270) \), \( \eta \), and \( \eta' \).

1. Introduction

In describing bottomonium results at a charm conference, we implicitly invoke heavy quark symmetry. The QCD is supposed to be the same, except that the bottom quark mass is about three times the charm quark mass, the magnitude of the electric charge is half as big, and the running strong coupling constant \( \alpha_s \) is about 30\% smaller. These differences will affect what portion of the effective potential is explored, how well the non-relativistic approximation works, decay multiplicities, and the number of bound quarkonium states, but the changes should, in principle, be calculable. This makes bottomonium a different laboratory to study the same physics as seen in charmonium. In this report, I will emphasize studies where bottomonium either extends, checks, or suggests new studies that can be done in charmonium.

Figure 1 compares the bound spectra of charmonium and bottomonium to illustrate some of these ideas. You can see that the bottomonium spectrum is richer, with more bound states and a wider variety of decay scenarios. It is also true that some of the fundamental states, including the ground state, have not yet been observed.

2. Relevant Experiments

The first \( \Upsilon \) states were discovered in hadron collisions, and there is still interesting work being done in bottomonium at the colliders. In particular, D0 has recently measured polarization in hadroproduction of bottomonium. However, Jonas Rademacker discussed this in detail in his report [1], so I will not cover it here.

Direct production of bottomonia in \( e^+e^- \) collisions has been a fruitful method for studying their properties. CLEO has 6 million \( \Upsilon(3S) \), 9 million \( \Upsilon(2S) \), and 21 million \( \Upsilon(1S) \) events. Belle collected 11 million \( \Upsilon(3S) \). Of course, both Belle and BaBar have hundreds of millions of \( \Upsilon(4S) \) events.

With so much luminosity, Belle and BaBar also produce tens of millions of \( \Upsilon(1S) \), \( \Upsilon(2S) \), and \( \Upsilon(3S) \) events with initial state radiation, although these events are somewhat more difficult to use effectively.

3. Dipion Transition Matrix Element

For two decades, there has been a puzzle in the description of the dipion mass distribution in bottomonium decays, as illustrated in Fig. 2. While the dipion mass distribution for the charmonium transition \( \psi(2S) \rightarrow \pi^+\pi^- J/\psi \), and the two bottomonium transitions \( \Upsilon(2S) \rightarrow \pi^+\pi^- \Upsilon(1S) \) and \( \Upsilon(3S) \rightarrow \pi^+\pi^- \Upsilon(2S) \) are well described by a single term in the matrix element [2], the transition \( \Upsilon(3S) \rightarrow \pi^+\pi^- \Upsilon(1S) \) shows a more complicated, two hump structure. More recently, data from Belle [3] and Babar [4] shown in Fig. 3 add inputs to the puzzle, showing that the decay \( \Upsilon(4S) \rightarrow \pi^+\pi^- \Upsilon(1S) \) shows the simple behavior, while \( \Upsilon(4S) \rightarrow \pi^+\pi^- \Upsilon(2S) \) is more complex.

CLEO has recently attempted to approach the problem by analyzing its \( \Upsilon \) transition data in a two-dimensional Dalitz-like fitting procedure [5]. Brown and Cahn [6] describe the relevant matrix element using PCAC and current algebra in the general form

\[
\mathcal{M} = A(e' \cdot \epsilon)(q^2 - 2m_{\pi}^2) + B(e' \cdot \epsilon)E_1E_2 + C[(e' \cdot q_1)(\epsilon \cdot q_2) + (\epsilon \cdot q_1)(e' \cdot q_2)],
\]

where \( A, B, \) and \( C \) are complex parameters, \( \epsilon \) and \( e' \) are the \( \Upsilon \) and \( \Upsilon' \) polarization vectors, \( q^2 = M_{\pi\pi}^2 \), \( E_1 \) and \( E_2 \) are the pion energies, and \( q_1 \) and \( q_2 \) their four-momenta.

Instead of fitting in a single dimension \( M_{\pi\pi}^2 \), which is equivalent to allowing only the \( A \) term in the matrix element, CLEO fits in \( M_{\pi\pi}^2 - M_{\pi\pi}'^2 \) space. They assume that the complex-valued \( A, B, \) and \( C \) terms do not vary significantly over the phase space of any of the decays. In the fits they find that the \( C \) term is not needed, which is not unexpected because it involves flipping the spin of the heavy \( b \) quark. CLEO’s limit is not very stringent, \( |C/A| < 1.09\% \) at 90\% confidence level, because the shapes of the \( C \) and \( B \) terms are very similar. If they set \( C = 0 \) they fit values for the ratio \( B/A \) given in Table 1 [6].

Dubynskiy and Voloshin [1] argue that the CLEO parametrization is too naive because \( B \) cannot possibly be constant over the Dalitz plot. So CLEO’s fits...
Figure 1: The spectra of bound charmonium and bottomonium. The vertical scales are the same, offset to align the ground states. States which have not yet been observed are rendered as dashed lines.

Figure 2: Dipion mass distribution in $\psi' \rightarrow \pi^+\pi^- J/\psi$ (upper left), $\Upsilon(2S) \rightarrow \pi^+\pi^- \Upsilon(1S)$ (upper right), $\Upsilon(3S) \rightarrow \pi^+\pi^- \Upsilon(2S)$ (lower left), and $\Upsilon(3S) \rightarrow \pi^+\pi^- \Upsilon(1S)$ (lower right). All distributions are well described by the Yan formulation, which keeps only the $A$ term of Eq. 1, except the last which has a distinctly different shape.

Figure 3: Dipion mass distributions observed in $\Upsilon(4S) \rightarrow \pi^+\pi^- \Upsilon(1S)$ by Belle [3] (left), and $\Upsilon(4S) \rightarrow \pi^+\pi^- \Upsilon(2S)$ by BaBar [4] (right). The former transition follows the Yan formulation, but the latter is more complex. The shaded histogram in the Belle plot is an estimate of the background from $\Delta M$ sidebands.

Table I. Fitted values for B/A in $\Upsilon$ dipion transitions

<table>
<thead>
<tr>
<th>Initial $\Upsilon$</th>
<th>Final $\Upsilon$</th>
<th>Re(B/A)</th>
<th>Im(B/A)</th>
</tr>
</thead>
<tbody>
<tr>
<td>3S</td>
<td>1S</td>
<td>$-2.52 \pm 0.04$</td>
<td>$1.19 \pm 0.06$</td>
</tr>
<tr>
<td>2S</td>
<td>1S</td>
<td>$-0.75 \pm 0.15$</td>
<td>$0.00 \pm 0.11$</td>
</tr>
<tr>
<td>3S</td>
<td>2S</td>
<td>$-0.40 \pm 0.32$</td>
<td>$0.00 \pm 1.10$</td>
</tr>
</tbody>
</table>

4. Pseudoscalar Transitions

In the charmonium system $\psi(2S) \rightarrow \eta J/\psi$ is surprisingly large considering the small amount of available phase space, with about a 3% branching fraction [5]. Yan [2] predicted in 1980 that the rate in bottomonium should scale like $\Gamma \propto (p^*)^3/m_\eta^4$. Kuang [6] has more recently refined this procedure and predicts $B(\Upsilon(2S) \rightarrow \eta \Upsilon(1S)) = (8.1 \pm 0.8) \times 10^{-4}$ and...
CLEO has sought $\Upsilon(2S) \rightarrow \eta \Upsilon(1S)$ in the final state where the $\Upsilon$ decays $\Upsilon(1S) \rightarrow \mu\mu$ or $ee$ and the $\eta$ decays $\eta \rightarrow \gamma\gamma$ or $\pi^+\pi^-\pi^0$.

CLEO plots the yield as a function of the kinetic energy of the $\eta$ candidate. In this preliminary analysis CLEO sees a 5 standard deviation peak in the kinetic energy of the $\gamma\gamma$ from $\eta$ decay as shown in Fig. 4. This leads to a (preliminary) branching fraction of $B(\Upsilon(2S) \rightarrow \eta \Upsilon(1S)) = (2.5 \pm 0.7 \pm 0.5) \times 10^{-4}$, somewhat smaller but in the same general range as Kuang’s prediction. CLEO also sees three events, with nearly zero expected background, in the $\eta$ decay channel $\eta \rightarrow \pi^+\pi^-\pi^0$, which corresponds to a consistent branching fraction.

Using the same technique, CLEO also seeks $\Upsilon(2S) \rightarrow \pi^0\Upsilon(1S)$, but sees no significant excess over background, setting the (preliminary) 90% confidence level upper limit $B(\Upsilon(2S) \rightarrow \pi^0\Upsilon(1S)) < 2.1 \times 10^{-4}$. This is consistent with the expectation, obtained using Yan’s simple scaling prediction, that the $\pi^0$ rate should be 0.16 of the $\eta$ rate.

5. $\Upsilon(1S)$ Decays to Invisible Particles

The decays of charmonium or bottomonium states to undetectable particles are a window on physics beyond the Standard Model. McElrath [10] has predicted that the neutralino $\chi$, a dark matter candidate, could be produced in $\Upsilon(1S) \rightarrow \chi\chi$ with a branching fraction of 0.41%. Possible production of new gauge bosons or a light gravitino was described by Fayet [11]. While it is true that $\Upsilon(1S) \rightarrow \nu\nu$ via $Z^0$ production is a potential background, it is calculated to be tiny enough as not to present a problem.

But how can you “see” these invisible decays? The trick is to produce a higher $\Upsilon$ state which decays via a two-pion transition to the $\Upsilon(1S)$. The experimentalist then uses the two pions to both trigger the detector and to signal the presence of the $\Upsilon(1S)$ with the missing mass recoiling against the dipion. The remainder of the detector must be completely empty.

Figures 5 and 6 show the results of two searches. Belle [12] uses $\Upsilon(3S)$ events so the transition pions have enough energy to trigger the detector reliably. The top of Fig. 5 shows the dipion mass spectrum when the $\Upsilon(1S) \rightarrow \mu\mu$ decay is observed, to demonstrate the expected shape of a possible signal. The bottom part of the figure shows the dipion mass spectrum when the rest of the detector shows no tracks and less than 3 GeV of neutral energy.

The CLEO data sample consists of 9 million $\Upsilon(2S)$ decays, almost as large as Belle’s 11 million $\Upsilon(3S)$ decays, and with the advantage of a more favorable dipion branching rate to $\Upsilon(1S)$. Unfortunately, their two track trigger had to be prescaled by a factor 20 to prevent saturating the data acquisition system. CLEO’s results are shown in Fig. 6. The top half shows the inclusive dipion mass spectrum, and the bottom half shows what remains when the rest of the detector is required to show no tracks and no photons of energy more than 250 MeV.

A small peak is visible at the $\Upsilon(1S)$ mass in the bottom parts of both figures. In both cases, the observed peaking is consistent with what is expected from Monte Carlo simulations where the decay products of the $\Upsilon$ traveled down the beam line, thus escaping the detector. Both experiments thus set upper limits to the production of invisible particles in $\Upsilon(1S)$ decays $B(\Upsilon(1S) \rightarrow \text{“invisible”}) < 0.25\%$ (Belle), and $B(\Upsilon(1S) \rightarrow \text{“invisible”}) < 0.39\%$ (CLEO).

Each of these limits is an order of magnitude better than the previous best limit. Together, they set a limit on the branching fraction about half the McElrath prediction for neutralino production, and better
the previous gravitino mass limit by a factor of four, to $m_{3/2} > 1.2 \times 10^{-7}$ eV.

Similar searches to these can be performed in the charmonium system, where a much larger number of $\psi'$ events is available. Of course, the mass range that can be explored is more limited, and the predicted branching fractions tend to be smaller, but such searches might be fruitful for charmonium experiments to pursue.

6. Radiative Decays of $\Upsilon(1S)$

6.1. Higgs Search

In an effort to explain why the Higgs hasn’t yet been seen, Dermisek, Gunion, and McElrath \cite{14} propose adding a non-Standard-Model-like pseudoscalar Higgs $a_0$ to the Minimal Supersymmetric Standard Model (MSSM) to make it the “Nearly MSSM” (NMSSM). This $a_0$ must have mass less than twice the $b$ quark mass, so that it can’t decay to a pair of $b$ quark jets. This proposal explains the failure of the LEP experiments to see the Higgs at masses up to 100 GeV, since the daughters of the Higgs decay can’t make the $b$ jets those experiments sought. Yet the hypothesis is natural, in the sense that it avoids fine tuning of parameters to explain observations.

The $a_0$ should decay predominantly into $\tau\tau$ if it has enough mass, and should be observable in $\Upsilon(1S) \rightarrow \gamma a_0$.

CLEO has sought these new Higgses by looking for monochromatic photons in events likely to contain taus. They tag $\Upsilon(1S)$ from $\Upsilon(2S) \rightarrow \pi^+\pi^- \Upsilon(1S)$ to help eliminate the copious QED backgrounds from $e^+e^- \rightarrow \gamma\tau\tau$. They flag the presence of $\tau$ pairs by seeking two 1-prong $\tau$ decays, one of which must be to a lepton, and by demanding missing energy in the event. The spectrum of photons they observe in such events is shown in Fig. 7 and leads to upper limits shown in Fig. 8. These upper limits improve on older measurements by an order of magnitude or more, and rule out much of the parameter space for NMSSM models.

6.2. $\Upsilon(1S) \rightarrow \gamma f_2(1270)$

In the charmonium system, radiative decays are common and many have been observed. The decay $J/\psi \rightarrow \gamma f_2(1270)$ is one of the most common. In bottomonium, few exclusive radiative decays are measured, but now CLEO has observed $\Upsilon(1S) \rightarrow \gamma f_2(1270)$...
Figure 7: The spectrum of photons in $\tau$-enriched $\Upsilon(1S)$ decays observed at CLEO.

Figure 8: Upper limits on the branching fraction of $\Upsilon(1S) \rightarrow \gamma a_0$ vs. photon energy (bottom scale) and $a_0$ mass (top scale).

$\gamma f_2(1270)$ as can be seen in Fig. 9.

We test the simple heavy quark symmetry relation

$$B(J/\psi \rightarrow \gamma f_2)/B(\Upsilon(1S) \rightarrow \gamma f_2) = \frac{(q_c/q_b)^2(m_b/m_c)^2(\Gamma_{bb}/\Gamma_{cc})}{\approx 20}$$

using the CLEO observations [15].

$B(\Upsilon(1S) \rightarrow \gamma f_2) = (10.2 \pm 0.8 \pm 0.7) \times 10^{-5}$ ($\pi^+\pi^-$)

$B(\Upsilon(1S) \rightarrow \gamma f_2) = (10.5 \pm 1.6 \pm 1.9) \times 10^{-5}$ ($\pi^0\pi^0$)

$B(\Upsilon(1S) \rightarrow \gamma f_2) = (10.23 \pm 0.97) \times 10^{-5}$ (combined).

The observed ratio

$B(J/\psi \rightarrow \gamma f_2)/B(\Upsilon(1S) \rightarrow \gamma f_2) = 14.0 \pm 1.7$

in satisfactory agreement with the expectations from scaling.

6.3. $\Upsilon(1S) \rightarrow \gamma \eta'$ and $\gamma \eta$

Does this success of scaling in radiative decay to $f_2$ carry over to other radiative decays? Another prominent decay in the charmonium system is $B(J/\psi \rightarrow \gamma \eta' = (4.7 \pm 0.3) \times 10^{-3}$ [15]. Using the observed charm system decay rate ratio $B(J/\psi \rightarrow \gamma \eta')/B(J/\psi \rightarrow \gamma f_2) = (3.4 \pm 0.4)$, and the relative rates of $\Upsilon$ and $J/\psi$ to $f_2$, we can predict the radiative decay rates for $\Upsilon(1S)$ to $\eta$ and $\eta'$. The expectation is that these decays should be easily visible.

We already know that $\eta'$ is unconventional. In radiative $J/\psi$ decay its branching fraction is five times as large as that for $\eta$. There have been speculations...

Figure 10: CLEO seeks $\Upsilon(1S) \to \gamma \eta'$ with the $\eta'$ decaying to $\pi^+\pi^-\eta$ and the daughter $\eta$ decaying in any of three modes. The blue arrows indicate where an expected $\eta'$ signal should be visible. Two candidates are seen in the mode where $\eta \to \pi^+\pi^-\pi^0$, but none are visible in the two all-neutral $\eta$ decay modes, leading to the upper limit quoted in the text.

7. Conclusion

Bottomonium remains an active field of research at Fermilab, CLEO, Belle and Babar. I have presented new results in dipion transitions among $\Upsilon$ states, $\eta$ and $\pi^0$ transitions in the $\Upsilon$ system, searches for invisible particles and a new type of Higgs, and radiative transitions to $\varphi(1270)$, $\eta$, and $\eta'$. However, bottomonium studies are continuing, and more new results can be expected next year.

Acknowledgments

Richard Galik was in the midst of preparing this talk when external circumstances forced him to turn it over to me. He had already chosen most of the topics, prepared a few of the slides, and done much of the research that went into the preparation of this note. I want to thank him for that work and the many helpful consultations we had.

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References

Light Higgses and Dark Matter at Bottom and Charm Factories

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University of California-Davis

Neither Dark Matter nor scalar particles in the Higgs sector are ruled out at energies accessible to bottom and charm factories. In Dark Matter searches, the error on the mass of Dark Matter is $\sim 4$ GeV in the best LHC studies. For light Dark Matter this could represent a 100% (or more) error. In Higgs searches, the presence of a light singlet Higgs can make the LHC Higgs search difficult, if not impossible. If Dark Matter or a Higgs scalar is light, it will require a low-energy machine to precisely determine the couplings. We review the models, modes of discovery and rate expectations for these new particle searches at bottom and charm factories. We also discuss the options for new runs at bottom and charm factories relevant for these searches.

1. Introduction

The two major new particles expected at colliders are Dark Matter and the Higgs boson. While some models are now ruled out at energies accessible to bottom and charm factories, it is by no means proven that these cannot be light. In fact there exist many attractive models containing light Higgses, for instance supersymmetric models which solve the $\mu$ problem via an extended Higgs sector. Furthermore the problem of light Dark Matter and light Higgses are related, as light Dark Matter particle $\chi$, in its simplest incarnation, requires a new light particle $U$ with $m_U \simeq 2m_\chi$ to serve as an $s$-channel annihilation mediator. A promising possibility for $U$ is that it is a pseudo-scalar higgs, which can be naturally light due to new symmetries which can protect its small mass. In order to ensure discovery, we should look everywhere that is practical for solutions to these problems, and $b$- and $c$-factories can perform an important set of new-particle searches.

Apart from the Dark Matter question, in the MSSM it was rigorously shown that an extremely light neutralino is experimentally unconstrained if one drops the assumption of gaugino unification, and the requirement that the neutralino relic density be equal to the Dark Matter relic density. For instance, the Dark Matter problem could be solved in another manner, such as with the QCD axion, rendering the neutralino an insignificant contributor to the relic density of the universe.

2. Dark Matter

In Dark Matter searches, the error on the mass of Dark Matter is $\sim 4$ GeV at the LHC in the best studies using optimistic models with large cross sections. Ultimately this is due to the resolution of the hadronic calorimeter, since to determine the mass scale, a missing energy event at the LHC can either use the missing transverse momentum, which is a hadronic observable or if purely leptonic observables exist, one can use the changes in slopes and shapes as a function of overall mass scale, which are only weakly correlated.

The only truly fundamental limit on the mass of Dark Matter comes from the Cosmic Microwave Background, which tells us the fraction of the universe that was non-relativistic at the time that photons decoupled, a measurement which includes Dark Matter. Dark Matter must have been non-relativistic at a temperature of about 0.3 eV, therefore the smallest possible mass consistent with the Standard Cosmological Model is about 0.3 eV.

This means that bottom and charm factories are capable of exploring 10 orders of magnitude in the Dark Matter mass. The LHC can expand to the range 5 GeV – 1 TeV, but has no precision below approximately 4 GeV.

We feel that the most compelling motivation for Dark Matter searches at bottom and charm factories is the demonstrable wisdom of a model independent approach. Indeed, the reason $M < 45$ GeV was ignored for so long is due to heavy reliance on models. In particular the Minimal Supersymmetric Standard Model cannot support Dark Matter this light because it would require another charged or colored particle to be lighter than other limits. The secondary particle is necessary to get the annihilation cross section large enough. Nearly all models which cannot support light Dark Matter cannot do so because of limits on particles other than the Dark Matter candidate itself. Trivial extensions of these models can generically support light Dark Matter by adding a mediator which is mostly singlet under the Standard Model. Several models demonstrate this explicitly.

The most minimal model possible for Dark Matter is to add only the dark matter candidate $\chi$ itself.

1. These studies all use a large value $\sim 100$ GeV for the Dark Matter mass. As this mass is brought closer to zero, the resolution on it at the LHC worsens.
However, these models generate very heavy Dark Matter candidates, outside the reach of $b$- and $c$-factories.

The second most minimal models adds the mediator $U$ as well, which is flavor neutral and couples both to the Standard Model and Dark Matter. These are the models testable at $b$- and $c$-factories.

$U$ can be a new gauge boson as proposed in Refs. [7, 9], or a scalar as proposed in [1]. If $U$ is a vector, it is necessarily anomalous, so building a consistent model requires even more matter than the $U$ and $\chi$. We are not aware of any such model in the literature. This is not because it is impossible, but rather because the resulting models are ugly, requiring several symmetry breaking scales and associated Higgses, as well as extra matter to cancel anomalies. If $U$ is a scalar, it can only couple to Standard Model fermions by mixing with the Higgs bosons due to gauge invariance, making its couplings proportional to mass. In the author’s opinion, a scalar or pseudo-scalar is a more natural candidate for $U$. Though as we will see in the next section, a vector $U$ may be easier to discover.

Treating the relic density as a constraint, acceptable models are achieved for (at least) two values of the Dark Matter mass as a function of the mediator’s mass, $\Omega h^2 = \frac{M_U}{2} \pm \epsilon$, as can be seen in Fig. 1. This is because the process controlling the annihilation of Dark Matter in the early universe is an $s$-channel annihilation diagram, rather than $t$-channel diagrams. In order for a $t$-channel diagram to dominate the annihilation, the new particle in the $t$-channel must be charged or colored. This occurs in the MSSM (e.g. “stau co-annihilation”) and generates one of the most promising regions of parameter space.

There is experimental evidence that Dark Matter may be light from the INTEGRAL satellite [10], which has detected an anomalously large population of positrons in the galactic center, as suggested in Ref. [9]. If this is from Dark Matter annihilation, it requires $M_\chi < \sim 3$ MeV [11].

Another source of evidence is from the DAMA annual modulation signal. As shown in Ref. [12], this is consistent with light Dark Matter due to the lower threshold of Sodium, as compared to heavier elements such as gallium (CDMS) and xenon (XENON).

There are two major modes of discovery for light Dark Matter: invisible meson decay [13, 14] and radiative decay [1, 15]. These are described respectively in the following subsections.

## 2.1. Invisible Quarkonium Decay

In invisible meson decay, one can make a naive calculation of the branching ratio for a meson. One can

\[ \frac{\text{BR}}{\text{meson}} = \frac{1}{2} \left( \frac{M_U}{2} \right)^2 \]

For the masses we consider, $M_\chi < 5$ GeV, annihilation to neutral Higgses or Z-bosons is kinematically disallowed.

\[ \frac{\text{BR}}{\text{meson}} = 10^{-12} \left( \frac{M_U}{5 \text{ GeV}} \right)^2 \]

Calculation here is expanded and corrected relative to Ref. [13], however the uncertainties and approximations made introduce much larger errors than the difference between the two calculations. This is an order-of-magnitude estimate only and little significance can be attached to achieving or exceeding
Where $\Omega_X$ by the average energy of the gas factor, assuming that the per-particle energy is given to separate
get an order of magnitude estimate for the annihilation cross section using
\[ \Omega_X h^2 \simeq \frac{0.1 \text{pb} \cdot c}{\langle \sigma v \rangle}. \] (1)
Where $\Omega_X = \rho_X / \rho_c$ is the relic density for species $X$ relative to the critical density $\rho_c$, $h$ is the Hubble constant, and $\langle \sigma v \rangle$ is the thermally averaged annihilation cross section of the DM into Standard Model particles. Using the central value of the WMAP result for $\Omega_X h^2 = 0.113$, we can invert this equation and solve for the required annihilation cross section for light relics
\[ \langle \sigma v \rangle = 0.88 \text{ pb} \cdot c. \] (2)
The velocity $v$ appearing here is the Möller velocity, which we approximate by the relative velocity in the center-of-mass frame, $v_{\text{rel}} = |v_1 - v_2|$, using $\langle v_{\text{rel}}^2 \rangle = 6 / x_{\text{FO}}$. The approximate temperature at freeze-out is $T = m_X / x_{\text{FO}}$ where $m_X$ is the mass of the DM and $x_{\text{FO}}$ is an expansion parameter evaluated at the freeze-out temperature that is $x_{\text{FO}} \sim 20 - 25$ depending on the model. By approximating that $\langle \sigma v \rangle = \sigma \sqrt{\langle v_{\text{rel}}^2 \rangle}$ we can remove the kinematic velocity factor, assuming that the per-particle energy is given by the average energy of the gas $\frac{3}{2} kT$.

We can expand $\langle \sigma v \rangle$ in the velocity at freeze-out to separate $s$-wave and $p$-wave components, $\langle \sigma v \rangle = a + b v^2$. Since the Dark Matter annihilates through the U-boson and not the meson we’re interested in, the freeze-out may in general occur at a different energy than the invisibly-decaying meson mass. Therefore we also remove the extra $v^2$ term and solve for $b$ in the $p$-wave case. These manipulations remove the kinematic factors of the initial state, giving us a cross section that essentially assumes the Dark Matter is massless with respect to our invisibly-decaying meson; that the meson mass is much larger than the center-of-mass energy at freeze-out.

For these assumptions with $x_{\text{FO}} = 25$ at freeze-out we have:
\[ \sigma(\chi\chi \to SM) = a / v_{\text{rel}} \simeq 1.8 \text{ pb}, \quad (s \text{-wave}) \] (3)
\[ \sigma(\chi\chi \to SM) = b / v_{\text{rel}} \simeq 7.5 \text{ pb}, \quad (p \text{-wave}) \]

The invisible branching ratio of a hadron can then be estimated by assuming that the time-reversed reaction is the same, $\sigma(f \bar{f} \to \chi\chi) \simeq \sigma(\chi\chi \to f \bar{f})$. Since the meson decays by the meson mixing with the $U$ boson, $p$-wave suppression factors are not reintroduced for the reverse reaction. We assume that the DM mediator is not flavor changing and that annihilation occurs in the $s$ channel. Therefore, the best-motivated hadrons to have an invisible width are same-flavor quark-antiquark bound states (quarkonia) with narrow widths.

The invisible width of a hadron composed dominantly of $q\bar{q}$ is given approximately by:
\[ \Gamma(H \to \chi\chi) = f_H^2 M_H \sigma(q\bar{q} \to \chi\chi) \] (4)

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure2.png}
\caption{The branching ratio for $\Upsilon(1s) \to \gamma\chi\chi$ via 3-body decay (i.e. either $m_{A_1} < 2m_{\chi_1}$ or $m_{A_1} > m_{\Upsilon}$) is plotted vs. the LSP mass (left) and relic density $\Omega h^2$ (right). All points shown are consistent with all LEP constraints. Points marked by an $x$ are excluded by one of: $\Upsilon \to \gamma\chi\chi$ (3-body decay) (that which is plotted); $\Upsilon \to \gamma A_1$ (2-body decay) with $A_1 \to \chi_1\chi_1$ (2-body decay); or $\Upsilon \to \gamma A_1$ (2-body decay) where the $A_1$ decays visibly.}
\end{figure}
where $f_H$ is the hadronic form factor (wave function at the origin) for the state $H$, and $M_H$ is the hadron’s mass. Here we ignore final state kinematic and spin factors.

We can predict an approximate expectation for the branching ratios for narrow states. Some of the most promising are shown in Table I. Branching ratios for scalars and pseudo-scalars tend to be smaller since those states are wider.

We emphasize again that this is only an order-of-magnitude calculation. A more precise calculation requires inclusion of kinematic and spin factors, as well as consideration of which fermions the mediator $U$ couples to. Furthermore, the freeze-out of light Dark Matter occurs in the middle of the QCD phase transition, and is much more sensitive to uncertainties due to QCD than heavier Dark Matter. This kind of dark matter is also annihilating through a narrow pole, which must be treated carefully. Narrow poles arise due to the $U$ boson itself, as well as numerous QCD resonances.

Several of these measurements have now been performed including $\Upsilon(1S) \rightarrow \chi\chi$, $\eta \rightarrow \chi\chi$ and $\eta' \rightarrow \chi\chi$; and now $J/\Psi \rightarrow \chi\chi$.

### 2.2. Radiative Decay

Radiative decay refers to meson decays into something visible as well as something invisible. This can be flavor changing, such as $b \rightarrow s\chi\chi$, in which case this is a next-to-leading-order effect requiring a loop of $W^\pm$ bosons to induce flavor changing. The authors of Ref. [14] found that this radiative decay can be as much as 50 times larger than the similar process radiating neutrinos. Other modes include $\Upsilon \rightarrow \gamma\chi\chi$ and $J/\Psi \rightarrow \gamma\chi\chi$.

In Fig. 3 we show the branching ratio of $J/\Psi$ and $\Upsilon$.

Figure 3: Branching ratio of the $J/\Psi$ (top) and $\Upsilon$ (bottom) into a photon and lightest pseudo-scalar Higgs $a_1$ in the NMSSM. The $a_1$ may then decay into Dark Matter (neutralinos) or visible Standard Model particles. The quantity $\cos \theta_A$ parameterizes how singlet-like the $a_1$ is. $\cos \theta_A = 1$ indicates that the $a_1$ is identical to the MSSM $A$. The bottom panels show $\tan \beta = 10$ and $M_{1,2,3} = 100, 200, 300$ GeV at scale $M_Z$. The plots are for $\tan \beta = 10$ and $M_{1,2,3} = 100, 200, 300$ GeV at scale $M_Z$. The bottom left plot comes from simply scanning in $\chi_1$, $\chi_2$, holding $\mu_{\chi \chi} = 150$ GeV fixed. The bottom right plot shows results for the $F < 15$ scenarios among the orange-cross, i.e. $m_{\chi_1} < 2m_{\chi_2}(pole)$, points of Fig. 1 of Ref. [23].

Table I Estimated branching ratios for the narrowest mesons. The two columns correspond to the assumption that the Dark Matter annihilation in the early universe occurs in either the $s$-wave or $p$-wave. Neutrino branching ratios are from Ref. [17]. All mesons have a branching ratio (even if tiny) to neutrinos.

<table>
<thead>
<tr>
<th>mode</th>
<th>$s$-wave</th>
<th>$p$-wave</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\mathrm{BR}(\Upsilon(1S) \rightarrow \chi\chi)$</td>
<td>$4.2 \times 10^{-4}$</td>
<td>$1.8 \times 10^{-3}$</td>
</tr>
<tr>
<td>$\mathrm{BR}(\Upsilon(1S) \rightarrow \nu\bar{\nu})$</td>
<td>$9.0 \times 10^{-6}$</td>
<td>$1.0 \times 10^{-3}$</td>
</tr>
<tr>
<td>$\mathrm{BR}(J/\Psi \rightarrow \chi\chi)$</td>
<td>$2.5 \times 10^{-5}$</td>
<td>$1.0 \times 10^{-4}$</td>
</tr>
<tr>
<td>$\mathrm{BR}(J/\Psi \rightarrow \nu\bar{\nu})$</td>
<td>$2.7 \times 10^{-8}$</td>
<td>$1.0 \times 10^{-5}$</td>
</tr>
<tr>
<td>$\mathrm{BR}(\eta \rightarrow \chi\chi)$</td>
<td>$3.4 \times 10^{-5}$</td>
<td>$1.4 \times 10^{-4}$</td>
</tr>
<tr>
<td>$\mathrm{BR}(\eta' \rightarrow \chi\chi)$</td>
<td>$3.7 \times 10^{-7}$</td>
<td>$1.5 \times 10^{-6}$</td>
</tr>
<tr>
<td>$\mathrm{BR}(\rho_0 \rightarrow \chi\chi)$</td>
<td>$1.3 \times 10^{-7}$</td>
<td>$5.3 \times 10^{-7}$</td>
</tr>
<tr>
<td>$\mathrm{BR}(\chi_{i,0}(1P) \rightarrow \chi\chi)$</td>
<td>$2.7 \times 10^{-8}$</td>
<td>$1.2 \times 10^{-7}$</td>
</tr>
<tr>
<td>$\mathrm{BR}(\phi \rightarrow \chi\chi)$</td>
<td>$1.9 \times 10^{-8}$</td>
<td>$7.8 \times 10^{-8}$</td>
</tr>
<tr>
<td>$\mathrm{BR}(\omega \rightarrow \chi\chi)$</td>
<td>$7.2 \times 10^{-8}$</td>
<td>$3.0 \times 10^{-8}$</td>
</tr>
</tbody>
</table>
$\Upsilon$ into $\gamma a_1$ in the NMSSM. The relevance of these plots for Dark Matter are that the $a_1$ may decay into the neutralino if the bino mass $M_1$ is decreased to make this mode kinematically allowed (without affecting this branching ratio). This can be done such that it is compatible with all collider constraints including $Z \to \text{invisible}$ as described in Ref. [1].

These branching ratios have little to do with the model assumptions of the NMSSM and can be parameterized only with $\theta_A$ and $\beta$:

$$BR(\Upsilon \to \gamma a_1) \propto \cos \theta_A \tan \beta$$  \hspace{1cm} (5)

$$BR(J/\Psi \to \gamma a_1) \propto \cos \theta_A \cot \beta$$ \hspace{1cm} (6)

so that experimentally, the only thing one needs worry about is that $M_{a_1}$ is small enough that the mode is kinematically allowed, and any limit can be interpreted in the $\cos \theta_A \tan \beta$ vs. $M_{a_1}$ plane. This covers a vast array of model space including any model with a light Higgs having some singlet admixture. Such Higgses appear in Refs. [1, 2, 3].

3. Higgses

Likewise, the existence of light Higgses can completely destroy the observability of Standard Model Higgs signals at the LHC via decays such as $h_2 \to h_1 h_1$ if it is dominant, where $h_2$ is a SM-like Higgs and $h_1$ is a lighter, mostly-singlet Higgs. Even if substantial backgrounds at the LHC can be overcome, the LHC will be unable to get a precise measurement of the lighter $h_1$ mass. If $h_1$ decays to $\tau^+ \tau^-$ the missing energy makes the mass measurement imprecise. If the $h_1$ decays to charm, strange, or gluons this renders the dominant Higgs decay entirely hadronic, and likely unobservable at the LHC due to hadronic backgrounds.

By contrast, bottom and charm factories can obtain precise measurements of the mass via the energy of a recoiling photon in the process $Y \to \gamma h_1$. The $h_1$ may have branching fractions to both Standard Model matter and Dark Matter.

The mode $Y \to \gamma H(A)$ was first suggested by Wilczek [24]. This mode is subject to significant radiative and threshold corrections, a comprehensive list of which can be found in Ref. [25]. It was vigorously pursued until about 1995, when it became clear that the LEP accelerator, searching for SM or MSSM Higgses in the Higgsstrahlung modes $e^+ e^- \to Zh$ and $e^+ e^- \to Ah$ was superior. The best measurements on the Upsilon were made by CLEO [26] however these remain about a factor of 10 away from being sensitive to a Standard Model Higgs. Existing data can reach the sensitivity required, as CLEO and Belle have approximately 20 times more data collected than that used in these limits. It will be necessary to reach and exceed the Standard Model limits to have sensitivity to Higgses with some singlet admixture.

The LEP measurements told us that no new particles with masses below $M_Z$ have a significant coupling to the $Z$. However, they tell us little about particles which have small coupling to the $Z$, and cannot rule out the existence of light particles. Particles with small $Z$ coupling are still allowed and can have interesting couplings to Higgses and fermions.

It is perhaps surprising that a light Higgs could still exist at low energies, and be compatible with all existing direct and indirect limits. However numerous studies have borne this out in a variety of models. All relevant experimental limits have been checked and light Higgses remain consistent with them. Some examples are: In the context of the Two Higgs Doublet Model, $(g - 2)_\mu$ (the anomalous magnetic moment of the muon) was examined [27, 28], as well as $BR(b \to s \gamma)$, $R_b$, $A_b$, $BR(\Upsilon \to A\gamma)$, $BR(\eta \to A\gamma)$ [28]. In the context of the NMSSM, $BR(\Upsilon \to \gamma + X)$ [1] was examined and found to be compatible.

We should note also that there is experimental evidence that this decay exists, from considerations of the excess seen at LEP near $M_b = 100$ GeV, fine tuning in the NMSSM [29], as well as some anomalous events at the HyperCP experiment which seem to indicate a $\sim 250$ MeV pseudo-scalar decaying to muons [30] that can be verified using radiative decays.

To allow a Higgs to be light, one must reduce its coupling to the $Z$ boson. In the MSSM this is proportional to $\sin(\beta - \alpha)$ for the CP-even state, and zero (at tree level) for the CP-odd state. Thus, by tuning the Higgs mixing angle $\alpha$ to be close to the ratio of the vacuum expectation values $\tan \beta$, this can be achieved. In the MSSM, however, the relationships among masses, $\alpha$, and $\beta$ is too constrained to allow only one of the Higgses to be lighter than $M_Z$ while simultaneously satisfying the Higgsstrahlung constraints. Basically, one of the CP-even Higgses has a mass related to the CP-odd Higgs, and the other is related to $M_Z$. So one cannot bring the $h$ light while simultaneously keeping the $A$ heavy. The $A$ becomes light as well, and generates a large cross section for $e^+ e^- \to h A$.

This difficulty comes from the fact that there is not enough freedom in the Higgs mass matrices, and as such is a theoretical constraint caused by one’s assumptions, and not experimental proof that there is no light Higgs. In models with more Higgs particles or more freedom in the Higgs self-couplings, the $ZZh$ coupling is more complex, and can be made small. The expansion of the Higgs sector in this manner is well motivated from the need to break any extra gauge symmetries such as a $U(1)''$ or $SU(2)'$, or to solve the MSSM’s $\mu$ problem [31]. Such particles may also generically be associated with SUSY breaking.

In a more general Two Higgs Doublet Model (2HDM), small coupling to the $Z$ can be achieved with light Higgses because $\beta$ and $\alpha$ are essentially free pa-
rameters. There remains some interesting parameter space in the 2HDM accessible at b- and c-factories, but it is small.\cite{16}

Finally there now exists “Gauge-Phobic Higgs” models in which electroweak symmetry breaking occurs by a combination of an elementary Higgs and breaking by by boundary conditions in an extra dimension.\cite{32} In such models, the Higgses become de-coupled from the Z. Their mass again becomes a free parameter and can be light, depending on how much of the symmetry breaking occurs due to the Higgs and how much due to the extra dimension.

4. The Future

As the b-Factories come to the end of their lives, much attention has been given to possible runs off the $\Upsilon(4S)$ resonance. This is an extremely promising idea. A small amount (e.g. weeks to months) of run time at a different energy may provide powerful physics results. Spending that same time on the $\Upsilon(3S)$ will provide only a negligible improvement over the already precise flavor physics results returned by these machines.

The promising options for future runs are on the $\Upsilon(3S)$, $\Upsilon(1S)$, $\Psi(2S)$, and $J/\Psi$. We have argued for the $\Upsilon(3S)$ due to the existence of the radiative decay $\Upsilon(3S) \rightarrow \pi \pi \Upsilon(1S)$,\cite{13} which can be used as a powerful constraint to remove backgrounds for invisible searches\cite{13} and Higgs searches\cite{22}. In addition to the CLEO data collected in the 1990’s, Belle has already collected 2.9 fb$^{-1}$\cite{20} on the $\Upsilon(3S)$.

No new studies have yet been published on the radiative decays $\Upsilon \rightarrow \gamma + X$ or $J/\Psi \rightarrow \gamma + X$. To improve on the capabilities of the CLEO datasets when searching for rare decays, it is necessary to further reject backgrounds. The single photon signal itself (ignoring the rest of the event) has sizeable backgrounds from direct production of 3-body final states $f \bar{f} \gamma$ which is a background to $\Upsilon \rightarrow \gamma a_1 \rightarrow \gamma f \bar{f}$. This argument was presented for $\Upsilon \rightarrow \gamma \tau^+ \tau^-$ in Ref.\cite{23}, but holds just as well for other fermions. Due to the high luminosity, BaBar and Belle have higher photon backgrounds in general than CLEO did.

Therefore, for all searches described here, we believe that new runs on the $\Upsilon(3S)$ and $\Psi(2S)$ will be the most significant. A Super-B factory can be even more powerful.\cite{33}

References

Extraction of $\alpha_s$ and $m_Q$ from Onia

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We briefly review how precise determinations of the strong coupling constant and of the heavy quark masses may be obtained from heavy quarkonium. Such determinations are competitive with theoretical predictions and the corresponding experimental measurements, we should be able to extract the values of the coupling constant $\alpha_s$ and of the heavy quark masses $m$. Therefore, from the theoretical predictions of physical observables and the corresponding experimental measurements, we should be able to extract the values of the coupling constant $\alpha_s$ and of the quark masses $m$. However, everything is complicated by the fact that QCD is a strongly coupled theory in the low energy region. At the scale $\Lambda_{\text{QCD}}$, nonperturbative effects become dominant and $\alpha_s$ becomes large. The nonperturbative QCD dynamics originates the confinement of quarks that in turn is the reason for which the quark mass loses its intuitive definition. Quarks are confined inside hadrons and thus we cannot directly measure their masses. The mass of the quark is a parameter defined in some renormalisation scheme at some renormalisation scale. Systems made by two heavy quarks, quarkonia in the following, are characterized by a quark mass scale $m_Q$ which is large, bigger than $\Lambda_{\text{QCD}}$. Then $\alpha_s(m_Q)$ is small and perturbative expansions may be performed at this scale. This introduces a great simplification and hints at a factorization between high and low energy contributions for quarkonia. For these systems, however, things are even more interesting [1]. They are nonrelativistic systems characterized by another small parameter, the heavy-quark velocity $v$, and by a hierarchy of energy scales: $m_{Q}$ (hard), the relative momentum $p \sim m_Q v$ (soft), and the binding energy $E \sim m_Q v^2$ (ultrasoft). For energy scales close to $\Lambda_{\text{QCD}}$, perturbation theory breaks down and one has to rely on nonperturbative methods. Regardless of this, the nonrelativistic hierarchy $m_Q \gg m_Q v \gg m_Q v^2$ will persist also below the $\Lambda_{\text{QCD}}$ threshold. While the hard scale is always larger than $\Lambda_{\text{QCD}}$, different situations may arise for the other two scales depending on the considered quarkonium system. The soft scale, proportional to the inverse radius $r$, may be a perturbative ($\gg \Lambda_{\text{QCD}}$) or a nonperturbative scale ($\sim \Lambda_{\text{QCD}}$) depending on the physical system. Finally, only for $t\bar{t}$ threshold states the ultrasoft scale may still be perturbative. Heavy quark-antiquark states probe confinement and nonperturbative physics [2] at different scales and are thus an ideal and to some extent unique laboratory where our understanding of nonperturbative QCD, its interplay with perturbative QCD and the behaviour of the perturbative bound state series may be tested and understood in a controlled framework. In particular in some regimes nonperturbative effects will appear in the form of local or nonlocal gluon condensates and will be suppressed in the computation of physical observables. In this framework quarkonia become very appropriate systems to be used for the study of the transition region from high to low energy, for information on the QCD vacuum structure and for precision determinations of the QCD parameters. Precisely this last point is the subject of this paper. In the next Sections we will discuss the systematic framework offered by Non Relativistic Effective Field Theories (NR EFT) [3] for the description of quarkonia and how one can take advantage of the accurate EFT calculations to make precise determinations of the QCD parameters. For some reviews of NR EFTs see [1, 2, 4, 5].

II. EFFECTIVE FIELD THEORIES

It is possible to take advantage from the existence of a hierarchy of scales in quarkonia to introduce NR EFTs, which are simpler but equivalent to QCD. A hierarchy of EFTs may be constructed by systematically integrating out modes associated to high energy scales not relevant for the quarkonium system. Such integration is made in a matching procedure that enforces the complete equivalence between QCD and the EFT at a given order of the expansion in $v$ ($v^2 \sim 0.1$ for $b\bar{b}, v^2 \sim 0.3$ for $c\bar{c}, v \sim 0.1$ for $t\bar{t}$). The EFT realizes a factorization at the Lagrangian level between the high energy contributions carried by matching coefficients and the low energy contributions carried by the dynamical degrees of freedom. The Poincaré symmetry remains intact in a nonlinear realization at the level of the NR EFT imposing exact relations among
the EFT matching coefficients $[7, 10]$. By integrating out the hard modes one obtains Non-Relativistic QCD $[8, 9, 10]$. NRQCD is making explicit at the Lagrangian level the expansions in $m\nu/m$ and $m\nu^2/m$. It is is similar to HQET, but with a different power counting. It also accounts for contact interactions between quarks and antiquark pairs (e.g. in decay processes) and hence has a wider set of operators. In NRQCD soft and ultrasoft scales are left dynamical and their mixing may complicate calculations, power counting and the consideration of the nonperturbative effects. In the last few years the problem of systematically treating the remaining dynamical scales in an EFT framework has been addressed by several groups $[11, 12, 13]$ and has now reached a good level of understanding. Therefore one can go down one step further and integrate out also the soft scale in a matching procedure to the lowest energy and simplest EFT that can be introduced for quarkonia, where only ultrasoft degrees of freedom remain dynamical. We call such EFT potential NonRelativistic QCD (pNRQCD) $[12, 13]$ (an alternative EFT is in $[14]$). pNRQCD is making explicit at the Lagrangian level the expansion in $m\nu^2/m\nu$. This EFT is close to a Schrödinger-like description of the bound state and hence as simple. The bulk of the interaction is carried by potential-like terms, but non-potential interactions, associated with the propagation of low-energy degrees of freedom ($Q\bar{Q}$ colour singlets, $Q\bar{Q}$ colour octets and low energy gluons), are generally present. They start to contribute at NLO in the multipole expansion of the gluon fields and are typically related to nonperturbative effects $[15, 16]$. In this EFT frame, it is important to establish when $A_{\text{QCD}}$ sets in, i.e. when we have to resort to non-perturbative methods. For low-lying resonances, it is reasonable to assume $m\nu^2 \gtrsim A_{\text{QCD}}$. The system is weakly coupled and we may rely on perturbation theory, for instance, to calculate the potential. In this case, we deal with weak coupling pNRQCD. The theoretical challenge here is performing higher-order calculations and the goal is precision physics. This is the case that we will consider in this paper.

## A. The QCD potential and the Static Energy

The masses may be extracted from a calculation of the energy levels and to obtain the energy levels we need the potential. The $Q\bar{Q}$ potential is a Wilson coefficient of pNRQCD $[17]$ obtained by integrating out all degrees of freedom but the ultrasoft ones. It is given by a series of contributions in an expansion in the inverse of the mass of the quark. If the quarkonium system is small, the soft scale is perturbative and the potentials can be entirely calculated in perturbation theory $[6]$. As matching coefficients the potentials undergo renormalization, develop a scale dependence and satisfy renormalization group equations, which eventually allow to resum potentially large logarithms $[15]$. The static singlet potential (the contribution at zero order in the mass expansion) is known at three loops apart from the constant term $[17, 18, 25]$. The first log related to ultrasoft effects arises at three loops. Such logarithm contribution at N$^3$LO and the single logarithm contribution at N$^4$LO may be extracted respectively from a one-loop and two-loop calculation in the EFT and have been calculated in $[17, 19]$. The static energy given by the sum of a constant, the static potential and the ultrasoft corrections is free from renormalon ambiguities. By comparing it (at the NNLL) with lattice calculations one sees that the QCD perturbative series converges very nicely to and agrees with the lattice result $[24]$ in the short range and that no nonperturbative linear (“stringy”) contribution to the static potential exist $[8, 20]$. This is an example of how precise calculations may be performed in this framework. Once the renormalon contribution has been cancelled, in this case between the static potential and the pole mass $[15, 21, 22]$, we are left with a well behaved perturbative series and we can unambiguously define power corrections. It is possible to make predictions of physical quantities (in this case the $Q\bar{Q}$ static energy) at high order in the perturbative expansion and with a small error (including nonperturbative corrections which are suppressed in the power counting) and to make a connection with the lattice results. It is remarkable that the dependence on the lattice spacing can be predicted in perturbation theory.

## B. The QCD perturbative series of the $Q\bar{Q}$ energies and the nonperturbative contributions

In weak coupling pNRQCD the soft scale is perturbative and the potentials are purely perturbative objects. Nonperturbative effects enter energy levels and decay calculations in the form of local or nonlocal electric and magnetic nonlocal correlators may be related to the gluon dressing functions $[23, 24]$. We still lack a precise and systematic knowledge of such nonperturbative purely glue dependent objects. It would be important to have for them lattice determinations or data extraction (see e.g. $[27]$). The leading electric and magnetic nonlocal correlators may be related to the gluon dressing functions $[13]$ and to some existing electric (quenched) determinations $[23, 28]$. However, since the nonperturbative contributions are suppressed in the power counting it is possible to obtain good determinations of the masses of the lowest quarkonium resonances with purely perturbative calculations in the cases in which the perturbative series is convergent (after that the appropriate subtractions of renormalons have been performed) and large logarithms are resummed. In this framework power corrections are unambiguously defined.
III. \textit{m}_c AND \textit{m}_b EXTRATION

The lowest heavy quarkonium states are suitable systems to obtain a precise determination of the mass of the heavy quarks \textit{b} and \textit{c}. Perturbative determinations of the \textit{Y}(1S) and \textit{J/\psi} masses have been used to extract the \textit{b} and \textit{c} masses. These determinations are competitive with those coming from different systems and different approaches (for the \textit{b} mass see e.g. \[56\]).

Determinations of the quark masses from the perturbative calculation of \textit{Y} and \textit{J/\psi} 1S masses differ for the order of the perturbative calculation considered, for the order of the resummation of the logarithms in \textit{v} and the way in which nonperturbative corrections are taken into account. Higher order terms and the residual scale dependence of the result give the theoretical error on the mass. The main uncertainty in these determinations comes from nonperturbative contributions (local and nonlocal gluon condensates) together with possible effects due to subleading renormalons. We report some example of such determinations in Tab. 1.

<table>
<thead>
<tr>
<th>reference</th>
<th>order</th>
<th>\textit{m}_b/\textit{m}_c (GeV)</th>
</tr>
</thead>
<tbody>
<tr>
<td>[36]</td>
<td>NNNLO*</td>
<td>4.210 ± 0.090 ± 0.025</td>
</tr>
<tr>
<td>[30]</td>
<td>NNLO +charm</td>
<td>4.190 ± 0.020 ± 0.025</td>
</tr>
<tr>
<td>[38]</td>
<td>NNLO</td>
<td>4.24 ± 0.10</td>
</tr>
<tr>
<td>[37]</td>
<td>NNLO*</td>
<td>4.346 ± 0.070</td>
</tr>
<tr>
<td>[39]</td>
<td>NNLO*</td>
<td>4.20 ± 0.04</td>
</tr>
<tr>
<td>[40]</td>
<td>NNLO*</td>
<td>4.241 ± 0.070</td>
</tr>
<tr>
<td>[41]</td>
<td>NNLL*</td>
<td>4.19 ± 0.06</td>
</tr>
<tr>
<td>[30]</td>
<td>NNLO</td>
<td>1.24 ± 0.020</td>
</tr>
<tr>
<td>[38]</td>
<td>NNLO</td>
<td>1.19 ± 0.11</td>
</tr>
</tbody>
</table>

TABLE I: Different recent determinations of \textit{m}_b/\textit{m}_c and \textit{m}_c/\textit{m}_c in the \textit{\overline{MS}} scheme from the bottomonium and the charmonium systems. The displayed results use either a direct calculation of the lowest energy level in perturbation theory or non-relativistic sum rules. The * indicates that the theoretical input is only partially complete at that order. For the detailed discussion about how the error has been computed see the original references, for a review see [3].

Once the quark masses have been obtained, the renormalon subtraction and the same calculational approach have been exploited also to obtain the energy levels of the lowest resonances. In \[29\] a prediction of the \textit{B}_c mass has been obtained. The NNLO calculation with finite charm mass effects \[30\] predicts a mass of 6307(17) MeV that well matches the CDF measurement \[31\] and the lattice determination \[32\]. The same procedure seems to work at NNLO even for higher states (inside the theory errors that grow) \[31\]. Including logs resummation at NLL, it is possible to obtain a prediction for the mass of \textit{\eta}_b = 9421 ± 11(\text{th}) +9/-8(\delta\alpha_s) \text{MeV} (where the second error comes from the uncertainty in \alpha_s) and for the \textit{B}_c hyperfine separation \Delta = 65 ± 24^{+19}_{-16} \text{MeV} \[34\]. A NLO calculation reproduces in part the 1P fine splitting \[54\].

A compilation of values of the \textit{b} and \textit{c} mass has been presented by the Quarkonium Working Group in Chapter 6 of [1] and is reported in Figures 1 and 2. The mass determinations presented in such Figures include (relativistic and nonrelativistic) sum rule results, lattice QCD results, semileptonic \textit{B} decays as well as \textit{Y}(1S) and \textit{J/\psi} 1S determinations. One can see that the determinations from quarkonium are competitive with respect to determinations coming from other systems (heavy-light, \textit{B} decays). The original works to which the results in such Figures refer are explicitly given and discussed in [1]. We refer to [1] also for an extended review of the different mass schemes, the different heavy quark mass extractions approaches and the renormalon subtraction.

From these determinations the QWG reported the following values for the \textit{\overline{MS}} masses:

\[
\begin{align*}
\text{\overline{m}}_b(\text{\overline{m}}_b) &= 4.22 ± 0.05 \text{GeV} \\
\text{\overline{m}}_c(\text{\overline{m}}_c) &= 1.28 ± 0.05 \text{GeV}.
\end{align*}
\]

which are displayed by the darker gray area in Figures 1 and 2. For the details of the calculation of these averages and ranges see [1].

We see that the QWG values for the \textit{b} and \textit{c} mass attribute to them an error of 1% and 4% respectively. This is a smaller error than the one given in the PDG [42].

More recent and more accurate mass determinations (from lattice unquenched calculation \textit{\overline{m}}_b(\text{\overline{m}}_b) = 4.4 ± 0.030 \text{GeV} \[44\]; from semileptonic \textit{B} decays, \textit{\overline{m}}_c(\text{\overline{m}}_c) = 1.224 ± 0.017 ± 0.054 \text{GeV} \[43\]; from low momentum sum rules \textit{\overline{m}}_b(\text{\overline{m}}_b) = 4.164 ± 0.025 \text{GeV} \[45\]; \textit{\overline{m}}_c(\text{\overline{m}}_c) = 1.286 ± 0.013 \text{GeV} \[45\]) and a new preliminary calculation of the mass of the \textit{b} in the potential subtracted scheme with unquenched lattice Fermilab action \[46\] would call for a new critical analysis and discussion of such extractions and errors and an updated mass compilation.

A. \textit{m}_t from \textit{t}\textit{t} bar systems

In \[41, 48\] the total cross section for top quark pair production close to threshold in \textit{e}^+\textit{e}^- annihilation is
FIG. 1: Collection of recent bottom quark mass determinations. The circles represent sum rule results, the triangles Upsilon 1S determinations, the squares lattice QCD results and the upside down triangle a determination from semileptonic B decays. The full diamond gives the QWG global average for $m_b(m_b)$. The darker and lighter shaded areas represent the QWG error estimates corresponding to a 1σ error and a range respectively. This Table is taken from Chapter 6, pag. 360 of [1]. For a detailed discussion and the explicit references to the original works see [1].

investigated at NNLL in the weakly coupled EFT. The summation of the large logarithms in the ratio of the energy scales significantly reduces the scale dependence. Studies like these will make feasible a precise extractions of the strong coupling, the top mass and the top width at a future ILC. The present theoretical uncertainties for top mass extraction at the ILC is about 100 MeV [3, 47].

IV. $\alpha_s$ EXTRACTION FROM QUARKONIA

The summary of values of $\alpha_s(M_Z)$ from various processes as reported by the PDG 2006 [42] is given in Fig. 4. We see that the value of $\alpha_s$ as determined from quarkonium is considerably smaller than the other determinations. The effect is seen also in Fig. 4 where the values of $\alpha_s(\mu)$ are reported at the values of $\mu$ where they are measured. The determination of $\alpha_s$ from Υ decays is one of the few ones at a relatively low energy with a relatively small error. It follows from theory calculations of ratio of hadronic and leptonic Υ decays [57] and use of sum rules for the Υ system [58, 59], the smaller error being obtained in the first case. Here we will report about a determination for $\alpha_s$ from the Υ decays [62] that has recently solved this inconsistency.

Heavy quarkonium leptonic and non-leptonic inclusive decay rates have historically provided ways to extract $\alpha_s$ and served as additional confirmation of the validity of QCD. Ratios of these quantities are very sensitive to $\alpha_s$ if the data are sufficiently precise. In particular, today the inclusive decay widths of $J/\psi$, $\psi(2S)$ and $\Upsilon(1S)$ are known with a few percent error, the ones of $\Upsilon(2S), \Upsilon(3S)$ with a 10% error and most of the other inclusive decays are known with an error of 15-20%. In the last few years the error on charmonium P-wave inclusive decays have been reduced to half [42]. On the theory side NRQCD [9] and pNRQCD [49] have provided powerful factorization for-
S and P wave quarkonium inclusive decays are today known in the NRQCD factorization up to order $v^7$ in the relativistic expansion \cite{9,50,64} and at different orders in the perturbative expansion of the matching coefficients (see e.g. \cite{65} for a review). In pNRQCD the nonperturbative matrix elements of the four quark operators on the quarkonium states can be further decomposed in the product of quarkonium wave functions (or derivatives of quarkonium wave functions) in the origin and glue dependent operators, with a substantial reduction in the number of nonperturbative (and unknown) contributions \cite{49}. A lattice calculation of such nonlocal gluonic correlators is however still missing.

Thanks to the EFTs factorization between high energy contributions, calculable in QCD perturbation theory, and low energy nonperturbative contributions, it is possible to consider appropriate ratios of inclusive decays at some order of the expansion in $\alpha_s$ and in $v$. In particular, the ratio $\Gamma(H \to \gamma gg) / \Gamma(H \to ggg)$ ($H$ being a quarkonium state) appears particularly promising for the extraction of $\alpha_s$ \cite{51,52,63}, since both the wave function at the origin and the relativistic corrections cancel out. However, the first measurements of $J/\psi$ and $\Upsilon$ inclusive radiative decays delivered a photon spectrum not compatible with the early QCD predictions. Inside the EFT approach it was understood that colour octet contributions, ignored in the early calculations, become very important in the upper end-point region of the spectrum \cite{53}. By considering such octet contributions, using pNRQCD to calculate them and Soft Collinear Effective Theory (SCET) to resum end-point singularities, a good description of the photon spectrum has been achieved recently, at least for the $\Upsilon(1S)$ state \cite{54}. These recent theoretical advances combined with new and more precise data from CLEO on $\Upsilon(1S)$ radiative decay \cite{55}, has made the ratio $R_\gamma \equiv \Gamma(\Upsilon(1S) \to \gamma X) / \Gamma(\Upsilon(1S) \to X)$, (X being hadrons) particularly suitable for the $\alpha_s$ extraction at the bottom mass scale. For the perturbative calculation of the matching coefficients appearing in such ratio see \cite{9,56}. Colour octet contributions also affect the ratio $R_\gamma$ and are parametrically of the same order of the relativistic corrections. They have so far either been ignored \cite{53} or estimated to be small \cite{57} in the available extractions of $\alpha_s$ from this ratio. In \cite{62}, recent determinations of the $\Upsilon(1S)$ colour octet matrix elements both on the lattice \cite{60} and in the contin-
FIG. 3: Summary of values of $\alpha_s(M_Z)$ from various processes, taken from the PDG [42]. The value shown indicate the process and the measured value of $\alpha_s$ extrapolated to $\mu = M_Z$. The error shown is the total error including theoretical uncertainties. The PDG average coming from these measurements and quoted in the text is also shown. Notice that the value of $\alpha_s$ extracted from Υ decays is considerably lower than all the other determinations.

V. FUTURE PROSPECTS FOR MASS AND $\alpha_s$ EXTRCTIONS

The mass and $\alpha_s$ determinations from quarkonium that we have presented are already competitive with the results obtained from other physical systems.

In the near future $\alpha_s(m_c)$ may be extracted from the $R_\gamma$ ratio for the $J/\psi$ provided that a new measurement of the inclusive photon spectrum for radiative $J/\psi$ decays will be performed at BESIII [60]. In a similar way, the discovery and the measurement of the $\eta_b$ mass with a few MeV accuracy will provide a determination of $\alpha_s(M_Z)$ with 3 per mille error from the hyperfine separation calculated at NLL [33].

For an improved determination of $\alpha_s$ from the lattice calculations of the quarkonium spectrum, we need a nonperturbative unquenched determination of $\Lambda_{\overline{MS}}$ and results on the spectrum obtained with different formulations of sea quarks, besides staggered quarks. Also the improvement in the lattice extraction of the masses would require an improved accuracy in the conversion from the bare lattice mass to the $\overline{MS}$ mass. In particular the two loop matching in such conversion would be needed for the Fermilab and the NRQCD actions. A nonperturbative matching would also be desirable.

For what concerns the mass extraction from the $\Upsilon(1S)$ and $J/\psi$ masses in perturbation theory at present, as it has been discussed, the major theoretical error comes from our ignorance of the ultra-soft nonperturbative corrections. A lattice calculation of the nonperturbative chromoelectric correlator together with its matching from lattice to $\overline{MS}$ scheme.
is needed.

Further improvements in the mass determinations from nonrelativistic sum rule would require the full NNLL calculation; a complete NNNLO computation would also be useful to have a better control on the theoretical uncertainties. For low momentum sum rules, improved determinations of the $R$ measurements around bottomonium and charmonium region would be crucial.

We conclude noticing that, within the EFT approach and the factorization scheme, precision calculations in quarkonium may be applied to all the physical observables of the lowest resonances, spectra and decays included. To this respect it is particularly interesting the example of the calculation of M1 transitions for the lowest quarkonia resonances. In this case the Poincarè invariance of the EFT imposes exact relations among matching coefficients that set to zero the nonperturbative corrections at order $v^2$. At this order the M1 transitions may be exactly calculated in perturbation theory [67].

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Charmonium from Lattice QCD

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Charmonium is an attractive system for the application of lattice QCD methods. While the sub-threshold spectrum has been considered in some detail in previous works, it is only very recently that excited and higher-spin states and further properties such as radiative transitions and two-photon decays have come to be calculated. I report on this recent progress with reference to work done at Jefferson Lab.

1. Introduction

Between 3 and 3.7 GeV a number of states exist which are believed to be the bound states of a charm quark and an anti-charm quark and whose widths are narrow owing to their being below the threshold to decay to a pair of open charm mesons coupled with the suppression of annihilation channels at this high mass scale. Because the hadronic contributions to their widths are so small, radiative transitions between them constitute considerable branching fractions, and the rates of these transitions have been measured with some accuracy by a number of experiments (Yao et al. [2006]). Additionally the $C = +$ states can decay to a pair of photons - this process when time-reversed can serve as a production mechanism (two-photon fusion) at $e^+e^-$ machines.

Rates for these radiative processes have been computed in various varieties of quark-model, and are typically fairly successful when one sets parameters using the experimental spectrum, however corrections beyond approximations like non-relativistic dynamics are often uncontrolled in these models.

In the current century the charmonium picture has filled out considerably and new mysteries have arisen owing to the high statistics and new production methods made possible by CLEO-c and the B-factories. The remaining expected sub-threshold states, $\eta_c, h_c$, have been observed, as have radiative transitions from the $\psi(3770)$ down to the $\chi_{cJ}$. The above-threshold spectrum is rapidly being mapped (Swanson [2006]), with some states living up to the expectations of potential models (Uehara et al. [2006]) and others coming as something of a surprise (Choi et al. [2003]). The increasingly complete set of exclusive data in $e^+e^-$ looks set to allow determination of the vector spectrum with some confidence.

In a series of recent works (Dudek et al. [2007], [2006], Dudek and Edwards [2006], [2006]), members of the Jefferson Lab lattice group have investigated the possibility of computing excited spectral and radiative quantities using lattice QCD. These initial studies have been carried out on quenched lattices with rather promising results. In the sections that follow I will briefly describe the work done.

2. Excited and higher spin states

The mass spectrum of a field theory considered in Euclidean space-time can be extracted from the time-dependence of a two-point correlation function,

$$C_{ij}(t) = \sum_{\vec{x}} \langle O_i(\vec{x}, t) O_j(\vec{0}, 0) \rangle,$$

where $O_{i,j}$ are operators that have the right quantum numbers to produce a particular state from the vacuum which are constructed from the fundamental fields of theory. For example in QCD we might try to study the pseudoscalar spectrum by considering an operator $\bar{\psi} \gamma^5 \psi$. The correlator receives contributions from all states in the theory with the appropriate quantum numbers,

$$C_{ij}(t) = \sum_{\alpha} \frac{Z_{i}^{\alpha} Z_{j}^{\alpha}}{2m_{\alpha}} \exp(-m_{\alpha} t). \quad (1)$$

(In a quantum mechanical bound state model we might think of radial excitations being labeled by $\alpha$). In practice extracting anything other than the ground state mass from fits to the time-dependence of a single correlator is difficult and often unstable. This is particularly troublesome in a system like charmonium where there are significant approximate degeneracies, e.g. the $\psi(3686)$ and $\psi(3770)$. These degeneracy problems are made worse on a cubic lattice where states are labeled not by a continuum spin, but by an irreducible representation of the cubic group. The continuum spin content of these various irreps is shown in Table I. This indicates that, for example, components of a $3^{--}$ state would appear in the same correlators as $1^{--}$ states. Since from potential models we expect there to be a $3^{--}(3D_3)$ roughly degenerate with the $1^{--}(3D_1) \psi(3770)$, we anticipate that there should be three roughly degenerate excited states above the ground state in a $T_{1^{--}}$ correlator. Extracting this from a fit to a single correlator is not practical.

Given this one might consider more reliable ways to extract the excited state spectrum. A variational method utilizing a large basis of operators satisfies this need. Its major advantage is that it utilizes the
orthogonality of states in a space of operators - while states might be degenerate and hence hard to separate on the basis of mass, they remain orthogonal and hence easier to separate on the basis of their state vectors.

In [Dudek et al. 2007], an operator basis was constructed based upon operators that in the continuum would have the structure of fermion bilinears with a number of symmetric covariant derivatives

\[ O_{\mu \nu \rho \ldots} = \bar{\psi}(x) \Gamma_\mu D_\nu D_\rho \cdots \psi(x). \]

Including up to two derivatives, these operators give access to almost all continuum \( J^{PC} \) with \( J \leq 3 \). Suitable linear combinations of these operators can be constructed that transform as the irreducible representations of Table I. These are related to the operators used in Liao and Manke [2002].

Once a matrix of correlators, \( C_{ij}(t) \), has been computed (for a given irrep), the mass spectrum follows from solution of a generalized eigenvalue problem that can be shown to be the quantum mechanical variational solution. We solve

\[ C(t)v_\alpha = \lambda_\alpha(t)C(t_0)v_\alpha, \quad (2) \]

for the eigenvalues \( \lambda_\alpha(t) \) which are related to state masses, and for the eigenvectors \( v_\alpha \) which are related to the overlap of our operators onto the mass eigenstates, the \( Z_i^\alpha \) in eqn (1).

We computed correlators on quenched anisotropic lattices with \( a_s \sim 0.1 \text{ fm} \) and \( a_s^{-1} \sim 6 \text{ GeV} \). Full details can be found in Dudek et al. [2007].

We show in figure 1 the mass spectrum extracted for negative parity and charge conjugation. In the \( T_1 \) representation we see precisely the level structure we expected, namely a ground state and three closely spaced excited states above. Looking at the states in the other irreps we see that one possible continuum spin assignment of the states in the first excited "band" would be to have two spin-1 states, one spin-2 state and one spin-3 state\(^1\). We gain a good deal of support for this hypothesis from studying the eigenvalues extracted from eqn (2). Consider the lattice irrep projections of the "\( a_1 \times \vec{\nabla} \)" operator:

\[ O_{T_2}^i = |\epsilon^{ijk}| \bar{\psi}(x) \gamma_5 \gamma_j D_k \psi(x) \]
\[ O_{E}^i = Q^{ijk} \bar{\psi}(x) \gamma_5 \gamma_j D_k \psi(x), \]

where \( |\epsilon^{ijk}|, Q^{ijk} \) are Clebsch-Gordan coefficients for the lattice cubic group. In the continuum we know the form that the overlap of these operators onto a spin-2 state takes, so that

\[ \langle 0|O_{T_2}^i|2^-(\vec{p}, r)\rangle = Z|\epsilon^{ijk}| \bar{e}_{jk}(\vec{p}, r) \]
\[ \langle 0|O_{E}^i|2^-(\vec{p}, r)\rangle = ZQ^{ijk} \bar{e}_{jk}(\vec{p}, r), \]

\[ ^1 \text{In the continuum the appearance of a e.g. spin-2 state} \]

In the continuum the appearance of a e.g. spin-2 state
where $Z$ is common to both. $Z$ can be extracted from the eigenvectors and if it is found to be close in value in the $T_2$ and $E$ cases then we conclude that it is likely that we have a spin-2 state.

We apply this eigenvector inspection method wherever possible and where the result is conclusive we assign the continuum spin shown by the color coding in figure 1. For the states above the first excited band this method gave inconclusive answers and for this reason we do not try to assign a continuum spin.

This first calculation was performed only at one (quenched) lattice spacing and consequently our results are not extrapolated to the continuum. Nevertheless we present our results for continuum spin assigned states in figure 2 along with experimental state masses taken from the PDG [Yao et al. 2006] and potential model masses taken from [Barnes et al. 2005].

It is clear that we are in agreement with the gross structure predicted by potential models, and in particular we appear to have successfully extracted something like the $\psi(3686)/\psi(3770)$ system. We believe that this has not been achieved before in a lattice calculation. Extracted state masses appear to be systematically high with respect to potential models and experiment - our suspicion is that this is due to some combination of computation at finite lattice spacing and the quenched approximation - this hypothesis can be tested with further calculation now that this method has been demonstrated.

Other $PC$ combinations were also considered. In figure 3 we show our results for $J^{++}$. It is clear that again we are observing masses systematically higher than the potential model states. That we miss the spin-4 state near 4 GeV may be related to the fact that our operators, which have a maximum of two spatial derivatives do not have any overlap with spin-4 mesons in the continuum limit. This could be remedied by enlarging the operator basis.

With $PC = +-$, as well as spin-singlets with odd $J$, one also has the possibility of exotic quantum numbers, i.e. those not accessible to a $q\bar{q}$ Fock state. In a quenched heavy-quark calculation these can only arise through non-trivial gluonic excitation giving rise to states usually described as "hybrids". Our extracted mass spectrum is shown in figure 4 where exotic states with $0^{-+}, 2^{-+}$ quantum numbers appear above 4.5 GeV.

With $PC = -+$, the odd-$J$ states are exotic. Our extracted mass spectrum listed by lattice irrep is shown in figure 5. This case demonstrates the difficulty in continuum spin assignment; the set of five levels near 4.3 GeV could, on the basis of their mass degeneracy, be interpreted either as a single $0^{-+}$ and a single (non-exotic) $4^{-+}$ or as two $0^{-+}$ states, an exotic $1^{-+}$ and a $2^{-+}$. In previous cases we used the eigenvector inspection method to break these ambiguities, but unfortunately here the method produces inconclusive results. We display the two possible spectra in figure 6 where we note that the potential model does have a $4^{-+}$ state in this mass range.

It is worth pointing out that previous studies of the $1^{-+}$ state in charmonia have not taken into account the spin ambiguity and hence they may have in fact reported the mass of a non-exotic $4^{-+}$ state. It is clear that further study with more operators and higher statistics is needed in order to make a definitive statement.

We believe that we have demonstrated the power of using a variational solution in a large, carefully constructed operator basis to extract excited states in lattice QCD. Of course there remain numerous issues to deal with, including the effect of multiparticle ($DD$) states when one relaxes the quenched approximation, but given that they too are orthogonal states we should be well-equipped with the method outlined.

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2In particular the problem of scale setting when one has an incorrect running of the coupling. In [Dudek et al. 2007] the effect of finite box size was tested and was found not to be the source of the level raising effect.
Figure 4: Extracted mass spectrum for $PC = +-$ listed by lattice irreducible representation and continuum spin assigned states.

Figure 5: Extracted mass spectrum for $PC = -+$ listed by lattice irreducible representation.

Figure 6: Two possible continuum spin interpretations of extracted mass spectrum for $PC = -+$. 

and an extended basis featuring operators with good overlap on to these multiparticle states.
3. Radiative transitions

Sub-$D\bar{D}$-threshold charmonia have very narrow widths such that radiative transitions between them constitute considerable branching fractions, these have been measured by a range of experiments and their relative magnitudes give us clues to internal structure.

These transition widths can be computed using lattice QCD by considering three-point correlators of the type

$$C_{\mu j}(t) \sim \langle O_i(x, t) \psi(\bar{y}, t) O_j(\bar{0}, 0) \rangle,$$

where $O_{i,j}$ are operators having overlap with meson states and $\psi_{\mu}$ is a lattice representation of the vector current. For example we could extract the $J/\psi \rightarrow \eta_c \gamma$ matrix element from the large Euclidean times value of the correlator

$$C_{\mu \nu}(t) = \sum_{\vec{x}} e^{-i \vec{p}_{f} \cdot \vec{x}} e^{i \vec{p}_{i} \cdot \vec{x}}$$

$$\times \left\langle \left[ \bar{\psi} \gamma_5 \psi \right](\vec{x}, t_f) \left[ \bar{\psi} \gamma_\mu \psi \right](\vec{y}, t) \left[ \bar{\psi} \gamma_\nu \psi \right](\bar{0}, 0) \right\rangle.$$

Correlators of this type, using only point-like operators were evaluated (details can be found in Dudek et al. [2006]) and transition form-factors extracted for a set of transitions between $J^{PC}$ ground states.

The $J/\psi \rightarrow \eta_c \gamma$ transition form-factor shown in figure 7 is the most statistically precise signal, but it suffers from a large systematic issue related to quenching. It is well known that the experimental hyperfine splitting in charmonium is not reproduced well by studies utilizing the quenched approximation. As such we have an ambiguity when computing the phase space that is required to scale a matrix-element to a width (or vice-versa) - should we use the experimental value or the value extracted from the spectrum portion of our lattice calculation? In figure 7 we show the experimental width\(^3\) scaled to a matrix-element by both possibilities and the lattice data fitted with an exponential in photon virtuality, $Q^2$, used to extrapolate back to $Q^2 = 0$.

A transition with reasonable statistical precision and a very small phase-space ambiguity is the electric dipole transition $\chi_{c0} \rightarrow J/\psi \gamma$. Our results are shown in figure 8 where the fit uses a form motivated by the quark model. Note the points at slightly timelike $Q^2$ are not included in the fit - the agreement with the extrapolated curve then lends support to the fitting form used.

Results for other transitions can be found in Dudek et al. [2006] as can comparison of the lattice results to potential model expectations. Work is currently underway combining the excited state technology of the first section with the radiative transition technology to make it possible to study transitions involving excited and high-spin states. This would include experimentally measured transitions like $\psi(3686) \rightarrow \chi_{cJ} \gamma$.

\(^3\)We note that there is ongoing work at CLEO to confirm the single measurement from Crystal Ball
4. Two-photon decays

At first sight it is not clear how one would go about evaluating the matrix element for the process $\eta_c \rightarrow \gamma \gamma$ in lattice QCD. In the previous section we outlined how to extract the matrix element for a radiative transition between two QCD eigenstates from a three-point function evaluated at large Euclidean times. This issue here is that the photon is not an eigenstate of QCD - taking a vector interpolating field to large Euclidean time would not yield a photon state, but instead the lightest QCD vector eigenstate (the $J/\psi$ in this case).

However, all is not lost, for while the photon is not a QCD eigenstate, it can be constructed from a linear superposition of QCD eigenstates. The precise field-theoretic mechanism for this is the LSZ reduction. The connection in Euclidean space-time, for a different physical process, is made in Ji and Jung [2001] and for the process in question an outline appears in Dudek and Edwards [2006]. The end result is that the following relationship connects the matrix element of interest to a Euclidean three-point function computable on the lattice: $\langle \eta_c(p) | \gamma(q_1, \lambda_1) \gamma(q_2, \lambda_2) \rangle \sim$

\[
e^2 \epsilon_{\mu}(q_1, \lambda_1) \epsilon_{\nu}(q_2, \lambda_2) \int dt e^{-i\omega(t_i-t)} \times \left\langle \int d^3 \vec{x} e^{-i\vec{p} \cdot \vec{x}} \mathcal{O}(\vec{x}, t_f) \int d^3 \vec{y} e^{i\vec{q} \cdot \vec{y}} \gamma^\nu(\vec{y}, t) \gamma^\mu(\vec{0}, t_i) \right\rangle \tag{3}
\]

The difference with respect to the radiative transitions between hadrons considered above is that an integral over the Euclidean time position of a vector source is now involved.

The details of the lattice computation of this object can be found in Dudek and Edwards [2006], here we mention only that an isotropic lattice was used. In figure 9(a) we display the integrand of equation 3 having computed with an operator $\bar{\psi} \gamma_5 \psi$ fixed at $t_f = 37$, a conserved vector current insertion at $t = 4, 16, 32$ and a vector interpolating field at all possible source positions, $t_i = 0 \rightarrow 37$. It is clear that provided one is not too close to the dirichlet wall or to the sink position, one can capture the entire integral by summing timeslices. In figure 9(b) the results of summing timeslices to compute the integral for all possible insertion positions and a number of $Q^2$ are shown - clear plateaus are visible at intermediate times indicating dominance of the $\eta_c$ over the possible excited states.

Given the confidence that the integral can be captured on a lattice of this temporal length, one can use a much faster method to compute the transition form-factor that places the sum over timeslices into a “sequential source”, reducing the computation time by a factor of $O(L_1)$. Results using this method are shown in figure 10 along with PDG values and results inferred from Uehara et al. [2007].

Figure 9: (a) Integrand in equation 3 at three values of vector current insertion time ($t = 4, 16, 32$) with pseudoscalar sequential source at sink position $t_f = 37$. (b) Pseudoscalar two-photon form-factor as a function of time slice, $t_i$ from equation 3. First six time slices ghosted out due to the Dirichlet wall truncating the integral.

Of course here the errors displayed on the lattice data are statistical only and must be augmented by an uncertainty due to scaling from our fixed lattice spacing to the continuum and one related to the lack of light-quark loops within the quenched approximation. This is the first demonstration of this method, such controlled studies will doubtless follow now that efficacy has been demonstrated.

5. Summary

Several new techniques have emerged that much expand the range of charmonium quantities that can be considered in lattice QCD. Initial studies with quenched lattices are clearly systematics dominated, but this can be expected to be improved in the near
future by use of dynamical lattices, in particular the anisotropic dynamical lattices being generated under USQCD at Jefferson Lab. These same methods applied to the light quark sector will provide invaluable information for future meson spectroscopy projects like GlueX.

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References


Figure 10: (a) $\eta_c \rightarrow \gamma \gamma^*$ amplitude. (b) $\chi_{c0} \rightarrow \gamma \gamma^*$ amplitude. Fits are one-pole forms as described in Dudek and Edwards [2006].
Determination of Charm Hadronic Branching Fractions at CLEO-c

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Recent results from CLEO-c on measurements of absolute hadronic branching fractions of $D^0$, $D^+$, and $D^{*+}$ mesons are presented.

I. INTRODUCTION

Precise measurements of absolute hadronic branching fractions for $D^0$, $D^+$, and $D^{*+}$ meson decays are important as they serve to normalize decay yields as well as many charm decays.

Results from the CLEO-c experiment at the Cornell Electron Positron Storage Ring based on 281 pb$^{-1}$ recorded at the $\psi(3770)$ are presented here for studies of $D^0$ and $D^+$ decays. In addition, CLEO-c has analyzed 298 pb$^{-1}$ of $e^+e^-$ annihilation data near $E_{\text{cm}} = 4170$ MeV for studies of $D_s$ decays. These samples provide very clean environments for studying decays of $D$ and $D_s$ mesons. The $\psi(3770)$, produced in the $e^+e^-$ annihilation, decays to pairs of $D$ mesons, either $D^+D^-$ or $D^0\bar{D}^0$. In particular, the produced $D_s$ mesons can not be accompanied by any additional pions. At $E_{\text{cm}} = 4170$ MeV $D_s$ mesons are primarily produced as $D_s^{+}D_s^{-}$ and $D_s^{*+}D_s^{*-}$ pairs.

First, I will discuss the determination of the absolute hadronic $D^0$, $D^+$, and $D^{*+}$ branching fractions. Then I will present CLEO-c measurements of inclusive $\eta$, $\eta'$, and $\phi$ decays; the doubly Cabibbo suppressed decay $D^+ \to K^+\pi^0$; studies of $D \to K_S\pi$ and $D \to K_L\pi$; $D_s$ decays to two pseudoscalars; and two-body $D^0$ and $D^+$ decays to pairs of kaons.

II. ABSOLUTE $D^0$ AND $D^+$ HADRONIC BRANCHING FRACTIONS

This analysis makes use of a ‘double tag’ technique initially used by Mark III. In this technique the yields of single tags, where one $D$ meson is reconstructed, and double tags, where both $D$ mesons are reconstructed, are determined. The number of reconstructed single tags, separately for $D$ and $\bar{D}$ decays, are given by $N_i = \epsilon_i \bar{B}_i N_{D\bar{D}}$ and $\bar{N}_j = \overline{\epsilon}_j B_j N_{D\bar{D}}$, respectively, where $\epsilon_i$ and $\bar{B}_i$ are the efficiency and branching fraction for mode $i$. Similarly, the number of double tags reconstructed are given by $N_{ij} = \epsilon_{ij} B_i B_j N_{D\bar{D}}$, where $i$ and $j$ label the $D$ and $\bar{D}$ mode used to reconstruct the event and $\epsilon_{ij}$ is the efficiency for reconstructing the final state. Combining the equations above and solving for $N_{D\bar{D}}$ gives the number of produced $D\bar{D}$ events as

$$N_{D\bar{D}} = \frac{N_i \bar{N}_j \epsilon_{ij}}{\epsilon_i \overline{\epsilon}_j}$$

and the branching fractions

$$B_i = \frac{N_i}{N_{D\bar{D}}} \epsilon_i \epsilon_{ij}.$$

In this analysis we determine all the single tag and double tag yields in data, determine the efficiencies from Monte Carlo simulations of the detector response, and extract the branching fractions and $D\bar{D}$ yields from a combined fit to all measured data yields.

This analysis uses three $D^0$ decay modes ($D^0 \to K^-\pi^+$, $D^0 \to K^-\pi^+\pi^0$, and $D^0 \to K^-\pi^+\pi^-\pi^+$) and six $D^+$ decay modes ($D^+ \to K^-\pi^+\pi^+$, $D^+ \to K^-\pi^+\pi^-\pi^+$, $D^+ \to K^-\pi^+\pi^0$, $D^+ \to K^-\pi^+\pi^-\pi^+$, $D^+ \to K^-\pi^+\pi^0\pi^+$, and $D^+ \to K^-\pi^+\pi^-\pi^+$). The single tag yields are shown in Fig. 1. The combined double tag yields are shown in Fig. 2 for charged and neutral $D$ modes separately. The scale of the statistical errors on the branching fractions are set by the number of double tags and precisions of $\approx 0.8\%$ and $\approx 1.0\%$ are obtained for the neutral and charged modes respectively. The branching fractions obtained are summarized in Table [1]. For the branching fractions we quote three uncertainties. The first is the statistical uncertainty, the second is the systematic uncertainties excluding the uncertainty in the modeling of final state radiation (FSR), and the third error is the FSR uncertainty. For the $D^0 \to K^-\pi^+$ mode the effect of the FSR is a 3.0% correction. We have taken the uncertainty of the FSR correction to be about 30% of the correction. This covers the difference between including or excluding the effect of interference in simulating FSR in the decay $D^0 \to K^-\pi^+$.

III. ABSOLUTE BRANCHING FRACTIONS FOR HADRONIC $D_s$ DECAYS

This analysis uses a sample of 298 pb$^{-1}$ of data recorded at a center-of-mass energy of 4170 MeV. At this energy $D_s$ mesons are produced, predominantly, as $D_s^{+}D_s^{-}$ or $D_s^{-}D_s^{*+}$ pairs. We use the same tagging technique as for the hadronic $D$ branching fractions; we reconstruct samples of single tags and double tags and use this to extract the branching fractions.

In this study eight $D_s$ final states are used ($D_s^{+} \to K_S^0K^+$, $D_s^{+} \to K^+K^-\pi^+$, $D_s^{*+} \to K^+K^-\pi^+$, $D_s^{*+} \to K_S^0K^-\pi^+\pi^+$, $D_s^{*+} \to K_S^0K^-\pi^+\pi^+$, $D_s^{*+} \to K_S^0K^-\pi^+\pi^+$, $D_s^{*+} \to K_S^0K^-\pi^+\pi^+$, $D_s^{*+} \to K_S^0K^-\pi^+\pi^+$). The single tag
event yields are shown in Fig. 3. The double tag yields are extracted by a cut-and-count procedure in the plot of the invariant mass of the $D_s^+$ vs. $D_s^-$. This plot is shown in Fig. 4. Backgrounds are subtracted from the sidebands indicated in the plot and a total of 976 ± 33 double tag events are found.

From these yields we determine the preliminary branching fractions listed in Table III. We do not quote branching fractions for $D_s^+ \to \phi \pi^+$ as the $\phi$ signal is not well defined. In particular, the $\phi$ resonance interferes with the $f_0$ resonance. Instead we report preliminary results for partial branching fractions for $D_s^+ \to K^+ K^- \pi^+$ in restricted invariant mass ranges of $m_{KK}$ near the $\phi$ resonance. These partial branching fractions are summarized in Table III.

### IV. INCLUSIVE MEASUREMENTS OF $\eta$, $\eta'$, AND $\phi$ PRODUCTION IN $D$ AND $D_s$ DECAYS

Using samples of tagged $D$ and $D_s$ decays CLEO-c has measured the inclusive production of $\eta$, $\eta'$, and $\phi$ mesons by looking at the recoil against the tag [4]. The results are summarized in Table IV. The knowledge of inclusive measurements before this CLEO-c measurement was poor, besides limits, only $B(D^0 \to \phi X) = (1.7 \pm 0.8)\%$ was measured. As expected the $\eta$, $\eta'$, and $\phi$ rates are much higher in $D_s$ decays.

### V. THE DOUBLY CABIBBO SUPPRESSED DECAY $D^+ \to K^+ \pi^0$

CLEO-c [5] has reconstructed $D^+ \to K^+ \pi^0$ candidates in the 281 pb$^{-1}$ sample of $e^+e^-$ data recorded at the $\psi(3770)$. We find the branching fraction $B(D^+ \to$
TABLE I: Fitted branching fractions and $D\overline{D}$ pair yields. For $N_{D^0\overline{D}^0}$ and $N_{D^+D^-}$, uncertainties are statistical and systematic, respectively. For branching fractions and ratios, the systematic uncertainties are divided into the contribution from FSR (third uncertainty) and all others combined (second uncertainty). The column of fractional systematic errors combines all systematic errors, including FSR. The last column, $\Delta_{\text{FSR}}$, is the relative shift in the fit results when FSR is not included in the Monte Carlo simulations used to determine efficiencies.

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Fitted Value</th>
<th>Fractional Error</th>
<th>$\Delta_{\text{FSR}}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$N_{D^0\overline{D}^0}$</td>
<td>$1.031 \pm 0.008 \pm 0.013 \times 10^6$</td>
<td>0.8</td>
<td>1.3</td>
</tr>
<tr>
<td>$B(D^0 \rightarrow K^-\pi^+)$</td>
<td>$(3.891 \pm 0.035 \pm 0.059 \pm 0.035)%$</td>
<td>0.9</td>
<td>1.8</td>
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<tr>
<td>$B(D^0 \rightarrow K^-\pi^+\pi^0)$</td>
<td>$(14.57 \pm 0.12 \pm 0.38 \pm 0.05)%$</td>
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<td>2.7</td>
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<tr>
<td>$B(D^0 \rightarrow K^-\pi^+\pi^-\pi^0)$</td>
<td>$(8.30 \pm 0.07 \pm 0.19 \pm 0.07)%$</td>
<td>0.9</td>
<td>2.4</td>
</tr>
<tr>
<td>$N_{D^+D^-}$</td>
<td>$(0.819 \pm 0.008 \pm 0.010) \times 10^6$</td>
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<td>1.2</td>
</tr>
<tr>
<td>$B(D^+ \rightarrow K^-\pi^+\pi^+\pi^-)$</td>
<td>$(9.14 \pm 0.10 \pm 0.16 \pm 0.07)%$</td>
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<td>1.9</td>
</tr>
<tr>
<td>$B(D^+ \rightarrow K^-\pi^+\pi^+\pi^0)$</td>
<td>$(5.98 \pm 0.08 \pm 0.16 \pm 0.02)%$</td>
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<td>2.8</td>
</tr>
<tr>
<td>$B(D^+ \rightarrow K_S^0\pi^+)$</td>
<td>$(1.526 \pm 0.022 \pm 0.037 \pm 0.009)%$</td>
<td>1.4</td>
<td>2.5</td>
</tr>
<tr>
<td>$B(D^+ \rightarrow K_S^0\pi^+\pi^0)$</td>
<td>$(6.99 \pm 0.09 \pm 0.25 \pm 0.01)%$</td>
<td>1.3</td>
<td>3.5</td>
</tr>
<tr>
<td>$B(D^+ \rightarrow K_S^0\pi^+\pi^-\pi^0)$</td>
<td>$(3.122 \pm 0.046 \pm 0.094 \pm 0.019)%$</td>
<td>1.5</td>
<td>3.0</td>
</tr>
<tr>
<td>$B(D^+ \rightarrow K^+K^-\pi^+)$</td>
<td>$(0.935 \pm 0.017 \pm 0.024 \pm 0.003)%$</td>
<td>1.8</td>
<td>2.6</td>
</tr>
<tr>
<td>$B(D^0 \rightarrow K^-\pi^+\pi^0)/B(K^-\pi^+)$</td>
<td>$(3.744 \pm 0.022 \pm 0.093 \pm 0.021)$</td>
<td>0.6</td>
<td>2.6</td>
</tr>
<tr>
<td>$B(D^0 \rightarrow K^-\pi^+\pi^-)/B(K^-\pi^-)$</td>
<td>$(2.133 \pm 0.013 \pm 0.037 \pm 0.002)$</td>
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<td>$B(D^+ \rightarrow K^-\pi^-\pi^+\pi^-)/B(K^-\pi^-\pi^+)$</td>
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<td>2.7</td>
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<tr>
<td>$B(D^+ \rightarrow K_S^0\pi^+)/B(K^-\pi^+\pi^-)$</td>
<td>$(0.1668 \pm 0.0018 \pm 0.0038 \pm 0.0003)$</td>
<td>1.1</td>
<td>2.3</td>
</tr>
<tr>
<td>$B(D^+ \rightarrow K_S^0\pi^+\pi^-)/B(K^-\pi^-\pi^+)$</td>
<td>$(0.764 \pm 0.007 \pm 0.027 \pm 0.005)$</td>
<td>0.9</td>
<td>3.5</td>
</tr>
<tr>
<td>$B(D^+ \rightarrow K_S^0\pi^+\pi^-)/B(K^-\pi^-\pi^+)$</td>
<td>$(0.3414 \pm 0.0039 \pm 0.0003 \pm 0.0004)$</td>
<td>1.1</td>
<td>2.7</td>
</tr>
<tr>
<td>$B(D^+ \rightarrow K^+K^-\pi^+\pi^-)/B(K^-\pi^-\pi^+)$</td>
<td>$(0.1022 \pm 0.0015 \pm 0.0022 \pm 0.0004)$</td>
<td>1.5</td>
<td>2.2</td>
</tr>
</tbody>
</table>

$K^+\pi^0 = (2.24 \pm 0.36 \pm 0.15 \pm 0.08) \times 10^{-4}$, which is in good agreement with the recent BABAR measurement $B(D^+ \rightarrow K^+\pi^0) = (2.52 \pm 0.46 \pm 0.24 \pm 0.08) \times 10^{-4}$.

VI. MODES WITH $K_S^0$ OR $K_L^0$ IN THE FINAL STATES

It has commonly been assumed that $\Gamma(D \rightarrow K_S^0X) = \Gamma(D \rightarrow K_L^0X)$. However, as pointed out by Bigi and Yamamoto, this is not generally true as for many $D$ decays there are contributions from Cabibbo favored and Cabibbo suppressed decays that interfere and contribute differently to final states with $K_S^0$ and
As an example consider $D^0 \to K^0_{S,L}\pi^0$. Contributions to these final states involve the Cabibbo favored decay $D^0 \to \bar{K}^0\pi^0$ as well as the Cabibbo suppressed decay $D^0 \to K^0\pi^0$. However, we don’t observe the $K^0$ and the $\bar{K}^0$ but rather the $K^0_S$ and the $K^0_L$. As these two amplitudes interfere constructively to form the $K^0_S$ final state we will see a rate asymmetry. Based on factorization Bigi and Yamamoto predicted

$$R(D^0) \equiv \frac{\Gamma(D^0 \to K^0_{S,L}\pi^0) - \Gamma(D^0 \to K^0_{S,L}\pi^0)}{\Gamma(D^0 \to K^0_{S,L}\pi^0) + \Gamma(D^0 \to K^0_{S,L}\pi^0)} \approx 2\tan^2\theta_C \approx 0.11.$$
TABLE III: Preliminary partial branching fractions for $D^+_s \rightarrow K^+ K^- \pi^+$ in limited $m(K^- K^+)$ ranges around the $\phi(1020)$ mass.

<table>
<thead>
<tr>
<th>$m(K^- K^+)$ range</th>
<th>Partial branching fraction(%)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$</td>
<td>m(K^- K^+) - m_{\phi}</td>
</tr>
<tr>
<td>$</td>
<td>m(K^- K^+) - m_{\phi}</td>
</tr>
<tr>
<td>$</td>
<td>m(K^- K^+) - m_{\phi}</td>
</tr>
<tr>
<td>$</td>
<td>m(K^- K^+) - m_{\phi}</td>
</tr>
</tbody>
</table>

TABLE IV: Inclusive branching fractions of $D^0$, $D^+$ and $D^+_s$ meson decays to $\eta$, $\eta'$, and $\phi$.

<table>
<thead>
<tr>
<th>Decay</th>
<th>$\mathcal{B}$ (%)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$D^0 \rightarrow \eta X$</td>
<td>9.5 ± 0.4 ± 0.8</td>
</tr>
<tr>
<td>$D^- \rightarrow \eta X$</td>
<td>6.3 ± 0.5 ± 0.5</td>
</tr>
<tr>
<td>$D^+_s \rightarrow \eta X$</td>
<td>23.5 ± 3.1 ± 2.0</td>
</tr>
<tr>
<td>$D^0 \rightarrow \eta' X$</td>
<td>2.48 ± 0.17 ± 0.21</td>
</tr>
<tr>
<td>$D^- \rightarrow \eta' X$</td>
<td>1.04 ± 0.16 ± 0.09</td>
</tr>
<tr>
<td>$D^+_s \rightarrow \eta' X$</td>
<td>8.7 ± 1.9 ± 1.1</td>
</tr>
<tr>
<td>$D^0 \rightarrow \phi X$</td>
<td>1.05 ± 0.08 ± 0.07</td>
</tr>
<tr>
<td>$D^- \rightarrow \phi X$</td>
<td>1.03 ± 0.10 ± 0.07</td>
</tr>
<tr>
<td>$D^+_s \rightarrow \phi X$</td>
<td>16.1 ± 1.2 ± 1.1</td>
</tr>
</tbody>
</table>

VII. $D_s$ DECAYS TO TWO PSEUDOSCALARS

CLEO-c has performed a study of $D_s$ decays to a pair of pseudoscalars. These final states consists of either a $K^+$ or a $\pi^+$ and one of $\eta$, $\eta'$, $\pi^0$, or $K_S^0$. In the analysis presented here the following final states are studied: $D^+_s \rightarrow K^+ \eta$, $D^+_s \rightarrow K^+ \eta'$, $D^+_s \rightarrow K^+ \pi^0$, $D^+_s \rightarrow \pi^+ K_S^0$, and $D^+_s \rightarrow \pi^+ \pi^0$. The final state $D^+_s \rightarrow \pi^+ \pi^0$ violates isospin and is expected to be small. The details of the analysis can be found in Ref. [10]. The signals are observed in the $D_s$ invariant mass distribution as peaks at the $D_s$ mass. Significant signals are observed in all modes except $D^+_s \rightarrow \pi^+ \pi^0$. The observed mass distributions are shown in Fig. 3. We measure the ratio of the branching fractions of the Cabibbo suppressed modes with respect to the Cabibbo favored modes. The results are summarized in Table IV. The observed ratios of branching fractions are consistent with the naive expectation of $|V_{cd}/V_{cs}|^2 \approx 0.05$. In addition, we have looked for a $CP$ asymmetry in rate for $D^+_s$ and $D^-_s$ decays. No evidence for any $CP$ asymmetry was found; the results are summarized in Table VII.

Using tagged $D$ mesons CLEO-c has measured this asymmetry and obtained

$$R(D^0) = 0.108 \pm 0.025 \pm 0.024,$$

which is in good agreement with the prediction.

Similarly, CLEO-c has also measured the corresponding asymmetry in charged $D$ mesons and obtained

$$R(D^+) = \frac{\Gamma(D^+ \rightarrow K_S^0 \pi^+) - \Gamma(D^+ \rightarrow K_S^0 \pi^+)}{\Gamma(D^+ \rightarrow K_S^0 \pi^+)} + \Gamma(D^+ \rightarrow K_S^0 \pi^+)} = 0.022 \pm 0.016 \pm 0.018. $$

Prediction of the asymmetry in charged $D$ decays is more involved. D.-N. Gao predicts this asymmetry to be in the range $0.035$ to $0.044$, which is consistent with the observed asymmetry.
particular, the decays $D^0 \to K^- K^+$, $D^0 \to K^0_SK^0_S$, and $D^+ \to K^+ K^0_S$ have been analyzed. In addition to TABLE VI: $CP$ asymmetries for Cabibbo suppressed $D_s \to PP$ decays.

<table>
<thead>
<tr>
<th>Mode</th>
<th>$(B_+ - B_-)/(B_+ + B_-)$ (%)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$A(D^+_s \to K^+\eta)$</td>
<td>$-20 \pm 18$</td>
</tr>
<tr>
<td>$A(D^+_s \to K^+\eta')$</td>
<td>$-17 \pm 37$</td>
</tr>
<tr>
<td>$A(D^+_s \to \pi^+K^0_S)$</td>
<td>$27 \pm 11$</td>
</tr>
<tr>
<td>$A(D^+_s \to K^+\pi^0)$</td>
<td>$2 \pm 29$</td>
</tr>
</tbody>
</table>

being Cabibbo suppressed, the $D^0 \to K^0(SK_S)$ mode is strongly suppressed due to destructive interference in the SU(3) limit between the two dominating exchange amplitudes for this decay. Figure 5 shows the observed yields in the three channels studied in this analysis.

The preliminary branching fractions are summarized in Table VII. For $D^0 \to K^+K^-$ and $D^+ \to K^+K^0_S$ there is good agreement with previous measurements. However, for $D^0 \to K^0_SK^0_S$ our new measurement is lower than previous measurements.

**IX. SUMMARY**

I have presented results based on 281 pb$^{-1}$ of $e^+e^-$ annihilation data recorded at the $\psi(3770)$ resonance for studies of $D^0$ and $D^+$ decays. Among the results presented here were the final results for the absolute $D^0 \to K^-\pi^+$ and $D^+ \to K^-\pi^+\pi^+$ branching fractions. CLEO-c has also analyzed 298 pb$^{-1}$ of $e^+e^-$ annihilation data recorded at the center-of-mass energy of 4170 MeV. Here we have studied the absolute hadronic branching fractions of $D_s$ mesons. CLEO-c has recorded more than 800 pb$^{-1}$ of data at the $\psi(3770)$ and are planning to double the data sample recorded at $E_{cm} = 4170$ MeV, so there are still many interesting results to come from the CLEO-c data sample.

**Acknowledgments**

This work was supported by the National Science Foundation grant PHY-0202078 and by the Alfred P. Sloan foundation.
FIG. 6: From left to right the yields in the $D^0 \rightarrow K^+ K^-$, $D^0 \rightarrow K^0_S K^0_S$, and $D^+ \rightarrow K^+ K^0_S$ are shown. We observe 4747 ± 74, 96 ± 13, and 1971 ± 51 events respectively in these modes. For the $D^0 \rightarrow K^0_S K^0_S$ analysis we subtract backgrounds, primarily, from $D^0 \rightarrow K^0_S \pi^+ \pi^-$ and find 70 ± 15 signal events.

TABLE VII: Preliminary branching fractions obtained in the study of two-body Cabibbo suppressed decays of $D$ mesons to pairs of kaons.

<table>
<thead>
<tr>
<th>Decay Mode</th>
<th>Our Measurement $(10^{-3})$</th>
<th>PDG 2007 $(10^{-3})$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$B(D^0 \rightarrow K^- K^+)$</td>
<td>4.01 ± 0.07 ± 0.08 ± 0.07</td>
<td>3.85 ± 0.09</td>
</tr>
<tr>
<td>$B(D^0 \rightarrow K^0_S K^0_S)$</td>
<td>0.149 ± 0.034 ± 0.015 ± 0.03</td>
<td>0.36 ± 0.07</td>
</tr>
<tr>
<td>$B(D^+ \rightarrow K^0_S K^+)$</td>
<td>3.35 ± 0.10 ± 0.10 ± 0.12</td>
<td>2.95 ± 0.19</td>
</tr>
</tbody>
</table>

[11] The result presented here represents the final results and are slightly different from the results presented at the workshop.
D and D_s hadronic branching fractions at B factories

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University of Bari and I.N.F.N., 70126 Bari, Italy

Recent measurements of hadronic branching fractions of D and D_s mesons, performed by the BaBar and Belle experiments at the asymmetric e^+e^- B factories colliders PEP II and KEKB, are reviewed.

1. Introduction

Hadronic branching fractions of D and D_s decays are used as references mode in many measurements of branching fractions of D and B-meson decays as well. A precise measurement of such values improves our knowledge of D and B-meson properties, and of fundamental parameters of the Standard Model, such as the magnitude of the Cabibbo-Kobayashi-Maskawa [1] matrix element.

2. Absolute branching fraction of D^0 → K^-π^+

BaBar collaboration measures the absolute branching fraction B(D^0 → K^-π^+)\(^3\) using D^0 → K^-π^+ decays in a sample of D^0 mesons preselected by their production in D^{*+} decays, obtained with partial reconstruction of the decay D^0 → D^{*-}π^-ν, with D^{*-} → D^0π^- [2]. Such measurement is extremely important because many of the past and current D and B branching fraction measurements are indeed systematically limited by the precision of B(D^0 → K^-π^+).

A sample of partially reconstructed B mesons in the channel D^0 → D^{*-}π^-ν is selected by retaining events containing a charged lepton (ℓ = e, μ) and a low momentum pion (soft pion, π^-) which may arise from the decay D^{*-} → D^0π^- [3]. This sample of events is referred to as the “inclusive sample”.

Using conservation of momentum and energy, the invariant mass squared of the undetected neutrino is calculated as

\[
M_\nu^2 = (E_{\text{beam}} - E_\ell - E_\pi)^2 - (\vec{p}_\ell + \vec{p}_\pi)^2,
\]

where \(E_{\text{beam}}\) is half the total center-of-mass energy, \(E_\ell\) (\(E_\pi\)) \(\vec{p}_\ell\) (\(\vec{p}_\pi\)) are the energy and momentum of the lepton (the D^* meson) and the magnitude of the B meson momentum, \(p_\ell\) is considered negligible compared to \(p_\pi\) and \(p_D^*\). Figure 1 shows the \(M_\nu^2\) distribution and the results of a minimum \(\chi^2\) fit aiming to determine the signal and background contribution.

\(1\) Charge conjugation is implied through the paper.

The number of signal events with \(M_\nu^2 > -2\) GeV^2/c^4 results \(N_{\text{incl}} = (2170.64\pm3.04(stat)\pm18.1(syst)) \times 10^3\).

The \(D^0 \rightarrow K^−\pi^+\) decays in the inclusive sample are selected requiring events in the mass range \(1.82 < M_{K\pi} < 1.91\) GeV/c^2 and \(142.4 < \Delta M < 149.9\) MeV/c^2 where \(\Delta M = M(K^−\pi^+\pi^-) - M(K^-\pi^+)\) and \(\pi_s^+\) is the slow pion from \(D^{*-}\) decay. The exclusive selection yields \(N_{\text{excl}} = 33810\pm290\) signal events, where the error is statistical only.

The branching fraction is computed as

\[
B(D^0 \rightarrow K^-\pi^+) = \frac{N_{\text{excl}}}{N_{\text{incl}}\varepsilon_{(K^-\pi^+)}}\zeta,
\]

where \(\varepsilon_{(K^-\pi^+)}\) is the \(D^0\) reconstruction efficiency as computed in the simulation, and \(\zeta\) is the selection bias introduced by the partial reconstruction.

The main systematic uncertainty on \(N_{\text{incl}}\) and \(N_{\text{excl}}\) are respectively due to the non-peaking combinatorial \(B\overline{B}\) background and the charged-track reconstruction efficiency. The complete set of systematic uncertainties is listed in Tab. 1. The absolute branching fraction of \(D^0 \rightarrow K^-\pi^+\) decay results

\[
B(D^0 \rightarrow K^-\pi^+) = (4.007 \pm 0.037 \pm 0.070)\%,
\]

Figure 1: The \(M_\nu^2\) distribution of the inclusive sample, for right-charge (a) and wrong-charge (b) samples. The data are represented by solid points with error. The MC fit results are overlaid to the data, as explained in the figure.

where the first error is statistical and the second error is systematic. This result is comparable in precision with the present world average, and it is consistent with it within two standard deviations.

Table I The relative systematic errors of $\mathcal{B}(D^0 \rightarrow K^- \pi^+)$.  

<table>
<thead>
<tr>
<th>Source</th>
<th>$\delta(\mathcal{B})/\mathcal{B}$(%)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Selection bias</td>
<td>$\pm 0.35$</td>
</tr>
<tr>
<td>$\mathcal{N}^{incl}$</td>
<td>$\pm 0.89$</td>
</tr>
<tr>
<td>Non-peaking combinatorial background</td>
<td></td>
</tr>
<tr>
<td>Peaking combinatorial background</td>
<td>$\pm 0.34$</td>
</tr>
<tr>
<td>Soft pion decays in flight</td>
<td>$\pm 0.10$</td>
</tr>
<tr>
<td>Fake leptons</td>
<td>$\pm 0.08$</td>
</tr>
<tr>
<td>Cascade decays</td>
<td>$\pm 0.08$</td>
</tr>
<tr>
<td>Monte Carlo events shape</td>
<td>$\pm 0.08$</td>
</tr>
<tr>
<td>Continuum background</td>
<td>$\pm 0.05$</td>
</tr>
<tr>
<td>$D^*$ production</td>
<td>$\pm 0.02$</td>
</tr>
<tr>
<td>Photon radiation</td>
<td>$\pm 0.02$</td>
</tr>
<tr>
<td>$\mathcal{N}^{recl}$</td>
<td>$\pm 1.00$</td>
</tr>
<tr>
<td>Tracking efficiency</td>
<td></td>
</tr>
<tr>
<td>$K^+$ identification</td>
<td>$\pm 0.70$</td>
</tr>
<tr>
<td>$D^*$ invariant mass</td>
<td>$\pm 0.56$</td>
</tr>
<tr>
<td>Combinatorial background shape</td>
<td>$\pm 0.30$</td>
</tr>
<tr>
<td>Combinatorial background normalization</td>
<td>$\pm 0.16$</td>
</tr>
<tr>
<td>Soft pion decay</td>
<td>$\pm 0.12$</td>
</tr>
<tr>
<td>Cabibbo-suppressed decays</td>
<td>$\pm 0.10$</td>
</tr>
<tr>
<td>Photon radiation in $D^0$ decay</td>
<td>$\pm 0.07$</td>
</tr>
<tr>
<td>Total</td>
<td>$\pm 1.74$</td>
</tr>
</tbody>
</table>

3. Absolute branching fraction of $D^+_s \rightarrow K^+ K^- \pi^+$

The poor accuracy of the branching fraction $\mathcal{B}(D^+_s \rightarrow K^+ K^- \pi^+) = (5.2 \pm 0.9)\%$ has been a systematic limitation for some precise measurements. In particular, the recent study of the $CP$ violation in $B^0 \rightarrow D^{(*)\pm} \pi^\mp$ decays is restricted by the knowledge of the ratio of two amplitudes that determine the $CP$-asymmetry [4, 5]. The amplitude $B^0 \rightarrow D^{(*)+} \pi^-$ can be calculated from the branching fraction of $B^0 \rightarrow D_s^{(*)+} \pi^-$ decays assuming factorization. On the other hand, the factorization hypothesis can be tested by measuring the ratio of $B^0 \rightarrow D^{(*)-} \pi^+$ and $B^0 \rightarrow D^{(*)-} D^+_s$ decays. Both $B(B^0 \rightarrow D_s^{(*)+} \pi^-)$ and $B(B^0 \rightarrow D^{(*)-} D^+_s)$ measurements can be improved with better accuracy in $D^+_s$ absolute branching fractions.

Belle collaboration measures $B(D_s^{(*)+} \rightarrow K^+ K^- \pi^+)$ using a partial reconstruction of the process $e^- e^- \rightarrow D^+_s D^-_{s1}$ [6]. In this analysis 4-momentum conservation allows to infer the 4-momentum of the undetected part.

The process $e^- e^- \rightarrow D_s^{(*)+} D^-_{s1}$ is reconstructed using two different tagging procedures. The first one (denoted as the $D_{s1}^-$ tag) includes the full reconstruction of the $D^-_{s1}$ meson via $D_{s1}^- \rightarrow T^* K^-$. The second procedure (denoted as the $D_s^{(*)+}$ tag) includes the full reconstruction of $D_s^{(*)+}$ and observation of the photon from $D_s^{(*)+} \rightarrow D^*_s \gamma$, while the $D_s^{(*)+}$ is not reconstructed. The measured signal yield with the $D_{s1}^-$ tag is proportional to the branching fractions of the reconstructed $T^*$ modes. In the second procedure (denoted as the $D_s^{(*)+}$ tag) a full reconstruction of $D_s^{(*)+}$ is required through $D_s^{(*)+} \rightarrow D^+_s \gamma$ and observation of the kaon from $D_{s1}^- \rightarrow T^- K^-$, but the $T^*$ is not reconstructed. Since the $D_s^{(*)+}$ meson is reconstructed in the channel of interest, $D_s^{(*)+} \rightarrow K^+ K^- \pi^+$, the signal yield measured with the $D_s^{(*)+}$ tag is proportional to this $D_s^{(*)+}$ branching fraction. The (efficiency-corrected) ratio of the two measured signal yields is equal to the ratio of well-known $T^*$ branching fractions and the branching fraction of the $D_s^{(*)+}$:

$$B(D_s^{(*)+} \rightarrow K^+ K^- \pi^+) = \frac{N(D_s^{(*)+})}{N(D_{s1}^-)} \cdot \frac{\epsilon(D^-_{s1})}{\epsilon(D_s^{(*)+})},$$

where $B(T^*)$ is the product of $T^*$ branching fraction and those of sub-decays.

The signal is identified by studying the mass recoiling against the reconstructed particle (or combination of particles) denoted as $X$. This recoil mass is defined as:

$$M_{\text{recoil}}(X) = \sqrt{(E_{CM} - E_X)^2 - P_X^2},$$

where $E_X$ and $P_X$ are the center-of-mass (CM) energy and momentum of $X$, respectively; $E_{CM}$ is the CM beam energy. A peak in the $M_{\text{recoil}}$ distribution at the nominal mass of the recoil particle is expected.

Since the resolution in $M_{\text{recoil}}$ is not enough to separate the relevant final states, the recoil mass difference $\Delta M_{\text{recoil}}$ is used to disentangle the contribution of the different final states:

$$\Delta M_{\text{recoil}}(D^-_{s1} \gamma) \equiv M_{\text{recoil}}(D^-_{s1} \gamma) - M_{\text{recoil}}(D^-_{s1} \gamma),$$

$$\Delta M_{\text{recoil}}(D_s^{(*)+} K^-) \equiv M_{\text{recoil}}(D_s^{(*)+} K^-) - M_{\text{recoil}}(D_s^{(*)+} K^-).$$

As the ratio of $D^-_{s1} \rightarrow T^* K^-$ and $D_{s1}^- \rightarrow D^*_s K^0_s$ branching fractions is unknown, the analysis is performed for these two channels separately. Figure 2 shows the $\Delta M_{\text{recoil}}(D^-_{s1} \gamma)$ and $\Delta M_{\text{recoil}}(D_s^{(*)+} K^-)$ distributions used for $D^-_{s1}$ and $D_{s1}^-$ tag procedures respectively. $\Delta M_{\text{recoil}}(D^-_{s1} \gamma)$ peaks at around $\pm 0.14$ GeV/$c^2 \simeq M(D^-_{s1} \gamma) - M(D^-_{s1})$. $\Delta M_{\text{recoil}}(D_s^{(*)+} K^-)$ peaks at around $\pm 0.525$ GeV/$c^2 \simeq M(D_{s1}^-) - M(D_s^{(*)+})$.

Using the measured signal yields $N(D_s^{(*)+})$ and $N(D_{s1}^-)$ with $D_s^{(*)+}$ and $D_{s1}^-$ tags, respectively, and taking into account the efficiency ratio $\epsilon(D_{s1}^-)/\epsilon(D_s^{(*)+})$, the $D_s^{(*)+}$ absolute branching fraction is computed by Eq. 1 for $D_{s1}^- \rightarrow T^* K^-$ and $D_{s1}^- \rightarrow D^*_s K^0_s$. The average value is $B(D_s^{(*)+} \rightarrow K^+ K^- \pi^+) = (4.0 \pm 0.4(\text{stat}) \pm 0.4(\text{syst}))\%$. 


4. Relative branching fraction of $D^0 \to K^-K^+\pi^0$ and $D^0 \to \pi^+\pi^-\pi^0$

The branching ratios of the singly Cabibbo-suppressed decays of $D^0$ meson are anomalous since the $D^0 \to \pi^-\pi^+$ branching fraction is observed to be suppressed relative to the $D^0 \to K^-K^+$ by a factor of almost three, even though the phase space for the former is larger. The branching ratios of the three-body decays $[3]$ have larger uncertainties but do not appear to exhibit the same suppression. This motivates the current study which measures the branching ratios of $D^0 \to \pi^-\pi^+\pi^0$ and $K^-K^+\pi^0$ with respect to the Cabibbo-favored decay $D^0 \to K^-\pi^+\pi^0$. BaBar collaboration measures both branching ratios $[7]$, Belle collaboration only the decay $D^0 \to \pi^-\pi^+\pi^0$ with respect to the decay $D^0 \to K^-\pi^+\pi^0$ $[8]$. By choosing the normalization mode $D^0 \to K^-\pi^+\pi^0$, many sources of systematic uncertainty including the $\pi^0$ detection efficiency and uncertainty in the tracking efficiency cancel out. To reduce combinatorial backgrounds, $D^0$ candidates are reconstructed in decays $D^{*-} \to D^0 \pi^-$ ($\pi^+$ is a soft, low momentum charged pion) with $D^0 \to K^-\pi^+\pi^0$, $\pi^-\pi^+\pi^0$, and $K^-K^+\pi^0$, by selecting events with at least three charged tracks and a neutral pion.

BaBar obtains the following results for the branching ratios:

$$\frac{\mathcal{B}(D^0 \to \pi^-\pi^+\pi^0)}{\mathcal{B}(D^0 \to K^-\pi^+\pi^0)} = (10.59 \pm 0.06 \pm 0.13) \times 10^{-2},$$

$$\frac{\mathcal{B}(D^0 \to K^-K^+\pi^0)}{\mathcal{B}(D^0 \to K^-\pi^+\pi^0)} = (2.37 \pm 0.03 \pm 0.04) \times 10^{-2},$$

while Belle obtains:

$$\frac{\mathcal{B}(D^0 \to \pi^+\pi^-\pi^0)}{\mathcal{B}(D^0 \to K^-\pi^+\pi^0)} = (9.71 \pm 0.09 \pm 0.30) \times 10^{-2}.$$  

Errors are statistical and systematic, respectively. Figure 3 shows the resulting mass distributions. Reflected $K^-\pi^+\pi^0$ events peak in the lower (upper) sideband of $m_{\pi^-\pi^+\pi^0}$ ($m_{K^-K^+\pi^0}$).

Using the world average value for the $D^0 \to K^-\pi^+\pi^0$ branching fraction $[8]$, the absolute branching ratios result:

BaBar

$$\mathcal{B}(D^0 \to \pi^-\pi^+\pi^0) = (1.493 \pm 0.008 \pm 0.018 \pm 0.053)\%,$$

$$\mathcal{B}(D^0 \to K^-K^+\pi^0) = (0.334\pm0.004\pm0.006\pm0.012)\%,$$

Belle

$$\mathcal{B}(D^0 \to \pi^-\pi^+\pi^0) = (1.369\pm0.013\pm0.042\pm0.049)\%,$$

where the errors are statistical, systematic, and due to the uncertainty of $\mathcal{B}(D^0 \to K^-\pi^+\pi^0)$.

The decay rate for each process can be written as:

$$\Gamma = \int d\Phi |\mathcal{M}|^2,$$

where $\Gamma$ is the decay rate to a particular three-body final state, $\mathcal{M}$ is the decay matrix element, and $\Phi$ is the phase space. Integrating over the Dalitz plot
For the three signal decays $\Phi$ is in the ratio

$$\Phi \propto \text{area of the Dalitz plot}.$$

where

$$\Gamma = \langle |M|^2 \rangle \times \Phi,$$

where $\langle |M|^2 \rangle$ is the average value of $|M|^2$ over the Dalitz plot and the three-body phase space, $\Phi$ is proportional to the area of the Dalitz plot. For the three signal decays $\Phi$ is in the ratio $\pi^-\pi^+\pi^0 : K^-\pi^+\pi^0 : K^-K^+\pi^0 = 5.05 : 3.19 : 1.67$. Combining the statistical and systematic errors, it results:

**Babar**

$$\frac{\langle |M|^2 \rangle (D^0 \rightarrow \pi^-\pi^+\pi^0)}{\langle |M|^2 \rangle (D^0 \rightarrow K^-\pi^+\pi^0)} = (6.68 \pm 0.04 \pm 0.08)\% \, (2)$$

$$\frac{\langle |M|^2 \rangle (D^0 \rightarrow K^-K^+\pi^0)}{\langle |M|^2 \rangle (D^0 \rightarrow K^-\pi^+\pi^0)} = (4.53 \pm 0.06 \pm 0.08)\% \, (3)$$

$$\frac{\langle |M|^2 \rangle (D^0 \rightarrow K^-K^+\pi^0)}{\langle |M|^2 \rangle (D^0 \rightarrow \pi^-\pi^+\pi^0)} = (6.78 \pm 0.14 \pm 0.21)\% \, (4)$$

**Belle**

$$\frac{\langle |M|^2 \rangle (D^0 \rightarrow \pi^-\pi^+\pi^0)}{\langle |M|^2 \rangle (D^0 \rightarrow K^-\pi^+\pi^0)} = (6.13 \pm 0.06 \pm 0.19)\% \, (5)$$

To the extent that the differences in the matrix elements are only due to Cabibbo-suppression at the quark level, the ratios of the matrix elements squared for singly Cabibbo-suppressed decays to that for the Cabibbo-favored decay should be approximately $\sin^2 \theta_C \approx 0.05$ and the ratio of the matrix elements squared for the two singly Cabibbo-suppressed decays should be unity. The deviations from this naive picture are less than 35% for these three-body decays. In contrast, the corresponding ratios may be calculated for the two-body decays $D^0 \rightarrow \pi^-\pi^+$, $D^0 \rightarrow K^-\pi^+$, and $D^0 \rightarrow K^-K^+$. Using the world average values for two-body branching ratios, the ratios of the matrix elements squared for two-body Cabibbo-suppressed decays, corresponding to Eqs. (2) (5), are, respectively, $0.034 \pm 0.001$, $0.111 \pm 0.002$, and $3.53 \pm 0.12$. Thus the naive Cabibbo-suppression model works well for three-body decays but not so well for two-body decays.

5. Amplitude analysis of $D$ and $D_s$ decays

The Dalitz plot analysis is the most complete method of studying the dynamics of three-body charm decays. These decays are expected to proceed through intermediate quasi-two-body modes and experi-
mentally this is the observed pattern. Dalitz plot analyses can also provide new information on the resonances that contribute to observed three-body final states. In this kind of analysis the complex quantum mechanical amplitude \( f \) is a coherent sum of all relevant quasi-two-body \( D^0 \to (r \to AB)C \) isobar model [10] resonances, \( f = \sum_i a_i e^{i\phi_i} \), where \( s = m_{AB}^2 \), and \( A \) is the resonance amplitude. The isobar model is expected to fail when there are large and overlapping resonances. In such case the \( \pi\pi \) S-wave is often parameterized through a K-matrix formalism [11, 12].

5.1. \( D_s^+ \to K^+K^-\pi^+ \) Dalitz plot analysis

**BaBar** collaboration reports the study of the three-body \( D_s^+ \) meson decays to \( K^+K^-\pi^+ \) and in particular the measurement of the branching fractions \( \frac{\mathcal{B}(D_s^+ \to \phi\pi^+)}{\mathcal{B}(D_s^+ \to K^+K^-\pi^+)} \) and \( \frac{\mathcal{B}(D_s^+ \to K^+K^-\pi^+)}{\mathcal{B}(D_s^+ \to K^+K^-\pi^+)} \). The decay \( D_s^+ \to \phi\pi^+ \) is frequently used in particle physics as the \( D_s^+ \) reference decay mode. The improvement in the measurements of these ratios is therefore important because it allows the \( D_s^+ \to K^+K^-\pi^+ \) to be used as reference.

A sample of 101k events with a purity of 95% is selected by a likelihood function using vertex separation and \( p^* \), the momentum of \( D_s^+ \) in CM system. A 66% of this final sample consists of \( D_s^+(2112)^+ \) decaying to \( D_s^+ \gamma \) where the variable

\[
\Delta m = m(K^+K^-\pi^+\gamma) - m(K^+K^-\pi^+)
\]

is required to be within \( \pm 2\sigma \) of the PDG value [3].

The selection efficiency is determined from a sample of Monte-Carlo events in which the \( D_s^+ \) decay is generated according to phase-space.

An unbinned maximum likelihood fit is performed in order to use the distribution of events in the Dalitz plot to determine the relative amplitudes and phases of intermediate resonant and non-resonant states.

The best fit results showing fractions, are summarized in Tab. II. The decay results to be dominated by \( K^*(892) \) and \( \phi \). Their branching ratio are:

\[
\frac{\mathcal{B}(D_s^+ \to \phi\pi^+)}{\mathcal{B}(D_s^+ \to K^+K^-\pi^+)} = 0.379 \pm 0.002 \pm 0.018
\]

and

\[
\frac{\mathcal{B}(D_s^+ \to K^*(892)^0 K^+)}{\mathcal{B}(D_s^+ \to K^+K^-\pi^+)} = 0.487 \pm 0.002 \pm 0.016
\]

where errors are statistic and systematic respectively. These measurements are much more precise than the previous ones, based on a Dalitz plot analysis of only 700 events [13].

A \( f_0(980) \) contribution is large but it is affected by large systematic errors as well due to uncertainty on \( f_0(980) \) and \( f_0(1370) \) parameters. The Dalitz plot projections together with the fit results are shown in Fig. 4.

Table II Fit fractions of a Dalitz plot fit of \( D_s^+ \to K^+K^-\pi^+ \) decay. Errors are statistical and systematic respectively.

<table>
<thead>
<tr>
<th>Decay Mode</th>
<th>Decay fraction(%)</th>
</tr>
</thead>
<tbody>
<tr>
<td>( K^*(892)^0 K^+ )</td>
<td>48.7 \pm 0.2 \pm 1.6</td>
</tr>
<tr>
<td>( \phi(1020)\pi^+ )</td>
<td>37.9 \pm 0.2 \pm 1.8</td>
</tr>
<tr>
<td>( f_0(980)\pi^+ )</td>
<td>35 \pm 1 \pm 14</td>
</tr>
<tr>
<td>( K_0^*(1430)^0 K^+ )</td>
<td>2.0 \pm 0.2 \pm 3.3</td>
</tr>
<tr>
<td>( f_0(1710)\pi^+ )</td>
<td>2.0 \pm 0.1 \pm 1.0</td>
</tr>
<tr>
<td>( f_0(1370)\pi^+ )</td>
<td>6.3 \pm 0.6 \pm 4.8</td>
</tr>
<tr>
<td>( K_0^*(1430)^0 K^+ )</td>
<td>0.17 \pm 0.05 \pm 0.3</td>
</tr>
<tr>
<td>( f_2(1270)\pi^+ )</td>
<td>0.18 \pm 0.03 \pm 0.4</td>
</tr>
</tbody>
</table>

---

Figure 4: The \( D^+ \to K^+K^-\pi^+ \) Dalitz plot projections. The data are represented by the points with error bars; the solid histograms are the projections of the fit described in the text. The inset shows an expanded view of the \( \phi(1020) \) region.
Further tests of the fit quality are performed using unnormalized $Y_0^0$ moment projections onto the $K^+K^-$ and $K^−\pi^+$ axes as functions of the helicity angles $\theta_K$ and $\theta_\pi$. For $K^+K^-$, the angle $\theta_K$ is defined as the angle between the $K^-$ for $D_s^-$ (or $K^+$ for $D_s^+$) in the $K^+K^-$ rest frame and the $K^+K^-$ direction in the $D_s^+$ rest frame. The $K^+K^-$ mass distribution is then modified by weighting by the spherical harmonic $Y_L^m(\cos \theta_K)$ ($L=1-4$). A similar procedure is followed for the $K^-\pi^+$ system. The resulting $\langle Y_0^0 \rangle$ distributions are shown in Fig. 6.

\[ \sqrt{\frac{4\pi}{\sin\theta}} \langle Y_0^0 \rangle = 2 |S||P| \cos \phi_{SP} \] (6)

Here $S$ and $P$ are proportional to the size of the $S$- and $P$-wave contributions and $\phi_{SP}$ is their relative phase. So $\langle Y_0^0 \rangle$ results to be related to the $S$-$P$ interference. Due to the presence of strong reflections on the $K^+K^-$ channel from the $K^−\pi^+$ channel and vice versa, Eq. (6) is meaningful only in the threshold regions. Figure 6 shows a large activity in the low $K^+K^-$ mass distribution, suggesting the presence of a large $S$-wave contribution below the $\phi(1020)$. The $\langle Y_0^0 \rangle$ distribution along the $K^−\pi^+$ projection, on the other hand, has a very small activity in the $K^+(892)^0$ suggesting a small $K\pi$ $S$-wave contribution.

5.2. $D^0 \rightarrow K^+K^−\pi^0$ Dalitz plot analysis

The $K^\pm\pi^0$ systems from the decay $D^0 \rightarrow K^−K^+\pi^0$ can provide information on the $K\pi$ $S$-wave amplitude in the mass range 0.6–1.4 GeV/c$^2$, and hence on the possible existence of the $\kappa(800)$, reported to date only in the neutral state ($\kappa^0 \rightarrow K^−\pi^+$) [16]. If the $\kappa$ has isospin 1/2, it should be observable also in the charged states. Results of the present analysis can also be an input for extracting the CP-violating phase $\gamma$ [17, 18].

$D^0$ from $\bar{D}^0$ are identified by reconstructing the decays $D^{*+} \rightarrow D^0\pi^+$ and $D^{*-} \rightarrow \bar{D}^0\pi^−$. The signal efficiency is estimated for each event as a function of its position in the Dalitz plot using simulated $D^0 \rightarrow K^−K^+\pi^0$ events from $\Upsilon(4S)$ decays, generated uniformly in the available phase space.

For $D^0$ decays to $K^\pm\pi^0$ $S$-wave states, three amplitude models are considered: the LASS amplitude for $K^−\pi^+ \rightarrow K^−\pi^+$ elastic scattering [19], the E-791 results for the $K^−\pi^+$ $S$-wave amplitude from an energy-independent partial-wave analysis in the decay $D^+ \rightarrow K^−\pi^+\pi^0$ [20] and a coherent sum of a uniform nonresonant term, and Breit-Wigner terms for the $\kappa(800)$ and $K^0_\Lambda(1430)$ resonances.

The results of an unbinned maximum likelihood are shown in Fig. 6. While the measured fit fraction (Tab. 11) for $D^0 \rightarrow K^+K^-$ agrees well with a phenomenological prediction [21], based on a large SU(3) symmetry breaking, the corresponding results for $D^0 \rightarrow K^−K^+$ and the color-suppressed $D^0 \rightarrow \phi^0$ decays differ significantly from the predicted values. The $K\pi$ $S$-wave amplitude is consistent with that from the LASS analysis, throughout the available mass range. The $K^−K^+$ $S$-wave amplitude, parameterized as either $f_0(980)$ or $a_0(980)^0$, is required. No
higher mass $f_0$ states are found to contribute significantly.

Table III  The results obtained from the $D^0 \rightarrow K^- K^+\pi^0$ Dalitz plot fit. The errors are statistical and systematic, respectively. The $a_0(980)$ contribution, when it is included in place of the $f_0(980)$, is shown in square brackets. LASS amplitude is used to describe the $K\pi$ S-wave states.

<table>
<thead>
<tr>
<th>State</th>
<th>Decay fraction(%)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$K^*(892)^+$</td>
<td>45.2 ± 0.8 ± 0.6</td>
</tr>
<tr>
<td>$K^*(1410)^+$</td>
<td>3.7 ± 1.1 ± 1.1</td>
</tr>
<tr>
<td>$K^+\pi^0(S)$</td>
<td>16.3 ± 3.4 ± 2.1</td>
</tr>
<tr>
<td>φ(1020)</td>
<td>19.3 ± 0.6 ± 0.4</td>
</tr>
<tr>
<td>$f_0(980)$</td>
<td>6.7 ± 1.4 ± 1.2</td>
</tr>
<tr>
<td>$[a_0(980)^0]$</td>
<td>[6.0 ± 1.8 ± 1.2]</td>
</tr>
<tr>
<td>$f_2^0(1525)$</td>
<td>0.08 ± 0.04 ± 0.05</td>
</tr>
<tr>
<td>$K^*\pi(892)^-$</td>
<td>16.0 ± 0.8 ± 0.6</td>
</tr>
<tr>
<td>$K^*(1410)^-$</td>
<td>4.8 ± 1.8 ± 1.2</td>
</tr>
<tr>
<td>$K^-\pi^0(S)$</td>
<td>2.7 ± 1.4 ± 0.8</td>
</tr>
</tbody>
</table>

Neglecting CP violation, the strong phase difference, $\delta_D$, between the $\bar{D}^0$ and $D^0$ decays to $K^*(892)^+K^-$ state and their amplitude ratio, $r_D$, are given by

$$ r_D e^{i\delta_D} = \frac{a_{D^0 \rightarrow K^*+K^-}}{a_{D^0 \rightarrow K^+K^-}} e^{i(\delta_{K^*+K^-} - \delta_{K^+K^-})}. $$

$\bar{B}A\bar{B}AR$ finds $\delta_D = -35.5^\circ \pm 1.9^\circ (\text{stat}) \pm 2.2^\circ (\text{syst})$ and $r_D = 0.599 \pm 0.013 (\text{stat}) \pm 0.011 (\text{syst})$. These results are consistent with the previous measurements [22], $\delta_D = -28^\circ \pm 8^\circ (\text{stat}) \pm 11^\circ (\text{syst})$ and $r_D = 0.52 \pm 0.05 (\text{stat}) \pm 0.04 (\text{syst})$.

5.3. $D^0 \rightarrow K_0^S\pi^+\pi^-$ Dalitz plot analysis

Recently, evidence for $D^0\bar{D}^0$ mixing has been found in $D^0 \rightarrow K^+K^-/\pi^+\pi^-$ [24] and $D^0 \rightarrow K^+\pi^-$ [25] decays. It is important to measure $D^0-\bar{D}^0$ mixing in other decay modes and to search for CP-violating effects in order to determine whether physics contributions outside the SM are present. $Belle$ collaboration reports a measurement of $D^0-\bar{D}^0$ mixing studying $D^0 \rightarrow K_0^S\pi^+\pi^-$ decay. The relevance of this decay is enhanced by its role in determining the angle $\gamma \equiv \arg[-V_{ud}V_{ub}^*V_{cd}V_{cb}^*]$ of the Unitarity Triangle. In fact, various methods [28] have been proposed to extract $\gamma$ using $B^- \rightarrow D^0K^-$ decays, all exploiting the interference between the color allowed $B^- \rightarrow D^0K^-(\propto V_{cb})$ and the color suppressed $B^- \rightarrow \bar{D}^0K^-(\propto V_{ub})$ transitions, when the $D^0$ and $\bar{D}^0$ are reconstructed in a common final state. The symbol $\bar{D}^0$ indicates either a $D^0$ or a $\bar{D}^0$ meson. Among the $D^0$ decay modes studied so far the $K_0^S\pi^+\pi^-$ channel is the one with the highest sensitivity to $\gamma$ because of the best overall combination of branching ratio magnitude, $D^0-\bar{D}^0$ interference and background level. $B\bar{A}\bar{B}AR$ collaboration reports a measurement of the angle $\gamma$ by studying the Dalitz plot of $D^0 \rightarrow K_0^S\pi^+\pi^-$ [27]. In order to estimate the systematic errors due to model, $B\bar{A}\bar{B}AR$ reports a Dalitz plot analysis where the $\pi\pi$ S-wave is parameterized by a K-matrix model [26].

The results of these analyses are summarized in Tab. IV. The decay is dominated by the $K^*(892)^-$ and $\rho(770)$ contribution. In order to improve the quality of fits, doubly Cabibbo suppressed $K^*$ contributions and two Breit-Wigner amplitudes $\sigma_1$ and $\sigma_2$ (whose masses and widths are float parameters) are included. $\sigma_1$ and $\sigma_2$ take into account the poor knowledge of $S$-wave in the low mass spectrum and $f_0(980)$ parameters. The K-matrix model overcomes this problem describing the $\pi\pi$ S-wave at all.

Figure 7 shows the results of unbinned maximum likelihood fit performed by $Belle$ [22].
Table IV Summary of branching ratios of $D^0 \to K_S^0 \pi^+\pi^-$ Dalitz plot fits performed by Belle(Isobar Model) and BaBar(Isobar and K-matrix Model).

<table>
<thead>
<tr>
<th>State</th>
<th>Belle (Isobar Model)</th>
<th>BaBar (Isobar Model)</th>
<th>BaBar (K-matrix Model)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$K^+(892)^-$</td>
<td>62.27</td>
<td>58.1</td>
<td>58.9</td>
</tr>
<tr>
<td>$K^*_+(1430)^-$</td>
<td>7.24</td>
<td>6.7</td>
<td>9.1</td>
</tr>
<tr>
<td>$K_S^-(1430)^-$</td>
<td>1.33</td>
<td>3.6</td>
<td>3.1</td>
</tr>
<tr>
<td>$K^*(1410)^-$</td>
<td>0.48</td>
<td>0.1</td>
<td>0.2</td>
</tr>
<tr>
<td>$K^+(1680)^-$</td>
<td>0.02</td>
<td>0.6</td>
<td>1.4</td>
</tr>
<tr>
<td>$K^+(892)^+$</td>
<td>0.54</td>
<td>0.5</td>
<td>0.7</td>
</tr>
<tr>
<td>$K^*_+(1430)^+$</td>
<td>0.47</td>
<td>0.0</td>
<td>0.2</td>
</tr>
<tr>
<td>$K^*_+(1430)^+$</td>
<td>0.13</td>
<td>0.1</td>
<td>0.0</td>
</tr>
<tr>
<td>$K^+(1410)^+$</td>
<td>0.13</td>
<td>—</td>
<td>—</td>
</tr>
<tr>
<td>$K^+(1680)^+$</td>
<td>0.04</td>
<td>—</td>
<td>—</td>
</tr>
<tr>
<td>$\rho(770)$</td>
<td>21.11</td>
<td>21.6</td>
<td>22.3</td>
</tr>
<tr>
<td>$\omega(782)$</td>
<td>0.63</td>
<td>0.7</td>
<td>0.6</td>
</tr>
<tr>
<td>$f_2(1270)$</td>
<td>1.8</td>
<td>2.1</td>
<td>2.7</td>
</tr>
<tr>
<td>$\rho(1450)$</td>
<td>0.24</td>
<td>0.1</td>
<td>0.3</td>
</tr>
<tr>
<td>$f_0(980)$</td>
<td>4.52</td>
<td>6.4</td>
<td>—</td>
</tr>
<tr>
<td>$f_0(1370)$</td>
<td>1.62</td>
<td>2.0</td>
<td>S-wave</td>
</tr>
<tr>
<td>$\sigma_1$</td>
<td>9.14</td>
<td>7.6</td>
<td>16.2</td>
</tr>
<tr>
<td>$\sigma_2$</td>
<td>0.88</td>
<td>0.9</td>
<td>—</td>
</tr>
<tr>
<td>NR</td>
<td>6.15</td>
<td>8.5</td>
<td>—</td>
</tr>
</tbody>
</table>

References

[27] B. Aubert et al. [BABAR Collaboration], arXiv:hep-ex/0607104.

Acknowledgments

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Flavor Symmetry and Charm Decays

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A wealth of new data in charmed particle decays allows the testing of flavor symmetry and the extraction of key amplitudes. Information on relative strong phases is obtained.

1. Introduction

The application of flavor symmetries, notably SU(3), to charmed particle decays can shed light on some fundamental questions. Often it is useful to know the strong phases of amplitudes in these decays. For example, the relative strong phase in $D^0 \to K^-\pi^+$ and $\bar{D}^0 \to K^-\pi^+$ is important in interpreting decays of $B$ mesons to $D^0X$ and $\bar{D}^0X$ [1, 2]. Such strong phases are non-negligible even in $B$ decays to pairs of pseudoscalar mesons ($P$) despite some perturbative QCD expectations to the contrary, and can be even more important in $D \to PP$ decays. In the present report we shall illustrate the extraction of strong phases from charmed particle decays using SU(3) flavor symmetry, primarily the U-spin symmetry involving the interchange of $s$ and $d$ quarks.

We begin in Section 2 by discussing the overall diagrammatic approach to flavor symmetry. In Section 3 we treat Cabibbo-favored decays, turning to singly-Cabibbo-suppressed decays in Section 4 and doubly-Cabibbo-suppressed decays in Section 5. We note specific applications to $D^0$ and $\bar{D}^0$ decays to $K^-\pi^+$ in Section 6, mention some other theoretical approaches in Section 7, and conclude in Section 8.

2. Diagrammatic amplitude expansion

We use a flavor-topology language for charmed particle decays first introduced by Chau and Cheng [3, 4]. These topologies, corresponding to linear combinations of SU(3)-invariant amplitudes, are illustrated in Fig. 1. Cabibbo-favored (CF) amplitudes, proportional to the product $V_{us}V^*_{cs}$ of Cabibbo-Kobayashi-Maskawa (CKM) factors, will be denoted by unprimed quantities; singly-Cabibbo-suppressed amplitudes proportional to $V_{us}V^*_{cd}$ or $V_{ud}V^*_{cd}$ will be denoted by primed quantities; and doubly-Cabibbo-suppressed quantities proportional to $V_{us}V^*_{cd}$ will be denoted by amplitudes with a tilde. The relative hierarchy of these amplitudes is $1: \lambda: -\lambda: -\lambda^2$, where $\lambda = \tan \theta_C = 0.232 \pm 0.002$ [5]. Here $\theta_C$ is the Cabibbo angle.

3. Cabibbo-favored decays

A detailed discussion of amplitudes and their relative phases for Cabibbo-favored charm decays was given in Ref. [6]. The main conclusions of that analysis were large relative phases of the $C$ and $E$ amplitudes relative to the dominant $T$ term, and an approximate relation $A \approx -E$. The present updated data confirm these results.

In Table I we show the results of extracting amplitudes $A = M_D[8\pi B/(p^*\tau)]^{1/2}$ from the branching ratios $B$ and lifetimes $\tau$, all from Ref. [2] unless otherwise noted. Here $M_D$ is the mass of the decaying charmed particle, and $p^*$ is the final c.m. 3-momentum.

The extracted amplitudes, with $T$ defined to be real, are, in units of $10^{-6}$ GeV:

$$T = 2.71 + 0 \; i \; ; \quad C = -1.77 - 1.01i \; ; \quad \delta(CT) = -150^\circ \; ; \quad E = -0.71 + 1.49i \; ; \quad \delta(ET) = 115^\circ \; ; \quad A = 0.57 - 1.30i \; ; \quad \delta(AT) = -66^\circ .$$

These values update (and are consistent with) those quoted with less precision in Ref. [6]. New (mainly lower) preliminary branching ratios for many $D_s$ decays reported at this Conference [7] will change some of the results slightly once they are incorporated into averages.

The Cabibbo-favored amplitudes are shown on an Argand diagram in Fig. 2. Here $A$ was extracted from $D_s \to \pi^+\eta$ and $D_s \to \pi^+\eta'$; the amplitude $A$ for $D_s \to \overline{K}^0K^+$ is then predicted to be $2.60 \times 10^{-6}$ GeV vs. $(2.60 \pm 0.25) \times 10^{-6}$ GeV observed. Note the importance of the $E$ and $A \approx -E$ amplitudes.

4. Singly-Cabibbo-suppressed decays

4.1. SCS decays involving pions and kaons

We show in Table II the branching ratios, amplitudes, and representations in terms of reduced amplitudes for singly-Cabibbo-suppressed (SCS) charm decays involving pions and kaons. The ratio of primed (SCS) to unprimed (CF) amplitudes is expected to be $\tan \theta_C \approx 0.232$. 

...
The deviations from flavor SU(3) implicit in Table III are well known. We shall discuss amplitudes in units of $10^{-7}$ GeV. If one rescales the CF amplitudes by the factor of $\tan^2 \theta_C$, one predicts $|A(D^0 \to \pi^+\pi^-)| = |A(D^0 \to K^+K^-)| = 5.78$, to be compared with a smaller observed value for $\pi^+\pi^-$ and a larger observed value (by a factor of $\sqrt{2}$) for $K^+K^-$. One can account for some of this discrepancy via the ratios of decay constants $f_K/f_\pi = 1.22$ and form factors $f_K(D \to K)/f_\pi(D \to \pi) > 1$. Furthermore, one predicts $|A(D^0 \to 0\pi^0)| = 4.45$ (larger than observed) and $|A(D^+ \to \pi^+\pi^0)| = 2.25$ (smaller than observed), which means that the $\pi\pi$ isospin triangle [associated with the fact that there are two independent amplitudes with $I = (0,2)$ for three decays] has a different shape from that predicted by rescaling the CF amplitudes. One predicts $|A(D^+ \to K^+K^-)| = |A(D_s \to \pi^+K^0)| = 5.79$; experimental values are (11%,1%) higher. The decay $D^0 \to K^0\pi^+$ is forbidden by SU(3); the branching ratio of $2\mathcal{B}(D^0 \to K^0\pi^0) = (2.98 \pm 0.68 \pm 0.30 \pm 0.60) \times 10^{-4}$ reported by CLEO [8] is more than a factor of two below that quoted in Table III (based on the average in Ref. [5]) and so does offer some evidence for the expected suppression.

4.2 SCS decays involving $\eta, \eta'$

The amplitudes $C$ and $E$ extracted from Cabibbo-favored charm decays imply values of $C' = \lambda C$ and $E' = \lambda E$ which may be used in constructing amplitudes for singly-Cabibbo-suppressed $D^0$ decays involving $\eta$ and $\eta'$. In Table III we write amplitudes multiplied by factors so that they involve unit coefficient of an amplitude $SE'$ describing a disconnected “singlet” exchange amplitude for $D^0$ decays.

Similarly the decays $D^+ \to (\pi^+\eta, \pi^+\eta')$ and $D^+_s \to (K^+\eta, K^+\eta')$ may be described in terms of a
Table II Branching ratios, amplitudes, and decomposition in terms of reduced amplitudes for singly-Cabibbo-suppressed (SCS) charm decays involving pions and kaons.

| Meson mode | Decay | $B$ $(10^{-3})$ | $p^*$ (MeV) | $|A|$ $(10^{-7}$ GeV) | Rep. |
|------------|-------|----------------|-------------|----------------------|------|
| $D^0$      | $\pi^+\pi^-$ | 1.37±0.03      | 922         | 4.57±0.05            | $-(T' + E')$ |
|            | $\pi^0\pi^0$ | 0.79±0.08      | 923         | 3.46±0.18            | $-(C' - E')/\sqrt{2}$ |
|            | $K^+K^-$      | 3.85±0.09      | 791         | 8.26±0.10            | $(T' + E')$ |
|            | $K^0\bar{K}^0$ | 0.72±0.14      | 789         | 3.58±0.35            | 0 |
| $D^+$      | $\pi^+\pi^0$ | 1.28±0.08      | 925         | 2.77±0.09            | $-(T' + C')/\sqrt{2}$ |
|            | $K^+\bar{K}^0$ | 5.90±0.38      | 793         | 6.43±0.21            | $T' - A'$ |
| $D_s^+$    | $\pi^+K^0$   | 2.46±0.40      | 916         | 5.87±0.48            | $-(T' - A')$ |
|            | $\pi^0K^+$   | 0.75±0.28      | 917         | 3.24±0.60            | $-(C' + A')/\sqrt{2}$ |

Figure 2: Construction of Cabibbo-favored amplitudes from observed processes. Here the sides $C+T$, $C+A$, and $E+T$ correspond to measured processes; the magnitudes of other amplitudes listed in Table I are also needed to specify the reduced amplitudes $T$, $C$, $E$, and $A$.

We show in Fig. 3 the construction proposed in Ref. [8] to obtain the amplitude $SA'$. Two solutions are found. In one, $|SA'|$ is uncomfortably large in comparison with the “connected” amplitudes, while in the other $|SA'|$ is smaller, but nonzero. Corresponding studies of the $D^0$ decays listed in Table III [10], which await further analysis by the CLEO Collaboration, will permit determination of the corresponding amplitude $SE'$ if one or more consistent solutions are found.

5. Doubly-Cabibbo-suppressed decays

In Table V we expand amplitudes for doubly-Cabibbo-suppressed decays in terms of the reduced amplitudes $\tilde{T} \equiv -\tan^2\theta_C T$, $\tilde{C} \equiv -\tan^2\theta_C C$, $\tilde{E} \equiv$
Table III Real and imaginary parts of amplitudes for SCS charm decays involving \( \eta \) and \( \eta' \), in units of \( 10^{-7} \text{ GeV} \) as predicted in Ref. [5].

<table>
<thead>
<tr>
<th>Amplitude</th>
<th>Expression</th>
<th>Re</th>
<th>Im</th>
<th>( \mathcal{A}_{\exp} )</th>
</tr>
</thead>
<tbody>
<tr>
<td>(-\sqrt{6}A(D^0 \rightarrow \pi^0\eta))</td>
<td>(2E' - C + SE')</td>
<td>0.82</td>
<td>9.24</td>
<td></td>
</tr>
<tr>
<td>(-\frac{1}{2}A(D^0 \rightarrow \pi^0\eta'))</td>
<td>(\frac{1}{2}(C' + E') + SE')</td>
<td>-2.87</td>
<td>0.56</td>
<td></td>
</tr>
<tr>
<td>(-\frac{1}{2\sqrt{3}}A(D^0 \rightarrow \eta\eta))</td>
<td>(C' + SE')</td>
<td>-4.10</td>
<td>-2.33</td>
<td></td>
</tr>
<tr>
<td>(-\frac{1}{2\sqrt{2}}A(D^0 \rightarrow \eta\eta'))</td>
<td>(\frac{1}{2}(C' + 6E') + SE')</td>
<td>-1.99</td>
<td>2.63</td>
<td></td>
</tr>
<tr>
<td>(\sqrt{3}A(D^+ \rightarrow \pi^+\eta))</td>
<td>(T' + 2C' + 2A' + SA')</td>
<td>0.71</td>
<td>-10.68</td>
<td>8.29\pm0.38</td>
</tr>
<tr>
<td>(-\frac{1}{2}A(D^+ \rightarrow \pi^+\eta'))</td>
<td>(\frac{1}{2}(T' - C' + 2A') + SA')</td>
<td>3.25</td>
<td>-0.92</td>
<td>4.03\pm0.42</td>
</tr>
<tr>
<td>(-\sqrt{3}A(D^+_s \rightarrow \eta\eta))</td>
<td>(-T' + 2C') + SA'</td>
<td>1.92</td>
<td>4.67</td>
<td>9.40\pm1.05</td>
</tr>
<tr>
<td>(-\frac{1}{2}A(D^+_s \rightarrow \eta\eta'))</td>
<td>((2T' + C' + 3A') + SA')</td>
<td>3.10</td>
<td>-2.84</td>
<td>3.88\pm0.66</td>
</tr>
</tbody>
</table>

Table IV Branching ratios and amplitudes for \( D^+ \) and \( D^+_s \) SCS decays involving \( \eta \) and \( \eta' \).

<table>
<thead>
<tr>
<th>Meson</th>
<th>Decay mode</th>
<th>( B )</th>
<th>( p^* )</th>
<th>( A )</th>
</tr>
</thead>
<tbody>
<tr>
<td>( D^+ )</td>
<td>( \pi^+\eta )</td>
<td>3.50\pm0.32</td>
<td>848</td>
<td>4.79\pm0.22</td>
</tr>
<tr>
<td></td>
<td>( \pi^+\eta' )</td>
<td>5.3\pm1.1</td>
<td>681</td>
<td>6.58\pm0.68</td>
</tr>
<tr>
<td>( D^+_s )</td>
<td>( K^+\eta )</td>
<td>1.92\pm0.43</td>
<td>835</td>
<td>5.43\pm0.61</td>
</tr>
<tr>
<td></td>
<td>( K^+\eta' )</td>
<td>2.02\pm0.69</td>
<td>646</td>
<td>6.33\pm1.08</td>
</tr>
</tbody>
</table>

\(-\tan^2 \theta_C E\), and \( \hat{A} \equiv -\tan^2 \theta_C A \).

With \( \tan \theta_C \approx 0.23 \) one predicts \( |A(D^0 \rightarrow K^+\pi^-)| = 1.32 \times 10^{-7} \text{ GeV} \) and \( |A(D^+ \rightarrow K^+(\pi^0, \eta, \eta')| = (0.93, 0.83, 1.27) \times 10^{-7} \text{ GeV} \), in qualitative agreement with experiment.

5.1. \( D^0 \rightarrow (K^0\pi^0, \bar{K}^0\pi^0) \) interference

The decays \( D^0 \rightarrow K^0\pi^0 \) and \( D^0 \rightarrow \bar{K}^0\pi^0 \) are related to one another by the U-spin interchange \( s \leftrightarrow d \), and SU(3) symmetry breaking is expected to be extremely small in this relation [11]. Graphs contributing to these processes are shown in Fig. [4].

The CLEO Collaboration [12] has reported the asymmetry

\[
R(D^0) \equiv \frac{\Gamma(D^0 \rightarrow K_S\pi^0) - \Gamma(D^0 \rightarrow K_L\pi^0)}{\Gamma(D^0 \rightarrow K_S\pi^0) + \Gamma(D^0 \rightarrow K_L\pi^0)} \tag{5}
\]

to have the value \( R(D^0) = 0.122 \pm 0.024 \pm 0.030 \), consistent with the expected value [11, 13] \( R(D^0) = 2 \tan^2 \theta_C \approx 0.108 \). One expects the same \( R(D^0) \) if \( \pi^0 \) is replaced by \( \eta \) or \( \eta' \) [11]. Moreover, by similar arguments, one expects \( A[D^0 \rightarrow K^0(\rho^0, f_0, \ldots)]/A[D^0 \rightarrow \bar{K}^0(\rho^0, f_0, \ldots)] = -\tan^2 \theta_C \).

5.2. \( D^+ \rightarrow (K^0\pi^+, \bar{K}^0\pi^+) \) interference

In contrast to the case of \( D^0 \rightarrow (K^0\pi^0, \bar{K}^0\pi^0) \), the decays \( D^+ \rightarrow (K^0\pi^+, \bar{K}^0\pi^+) \) are not related to one another by a simple U-spin transformation. Amplitudes contributing to these processes are shown in Fig. [5]. Although both processes receive color-suppressed \( (C \text{ or } \bar{C}) \) contributions, the Cabibbo-favored process receives a color-favored tree \( (T) \) contribution, while the doubly-Cabibbo-suppressed process receives an annihilation \( (\bar{A}) \) contribution. In order to calculate the asymmetry between \( K_S \) and \( K_L \) production in these decays due to interference between CF and DCS amplitudes, one can use the determination of the CF amplitudes discussed previously and the relation between them and DCS amplitudes. Thus, we define

\[
R(D^+) \equiv \frac{\Gamma(D^+ \rightarrow K_S\pi^+) - \Gamma(D^+ \rightarrow K_L\pi^+)}{\Gamma(D^+ \rightarrow K_S\pi^+) + \Gamma(D^+ \rightarrow K_L\pi^+)} \tag{6}
\]

and predict

\[
R(D^+) = -2 \text{Re} \frac{\bar{C} + \bar{A}}{\bar{T} + \bar{C}} = 2 \tan^2 \theta_C \text{Re} \frac{C + A}{T + C} = 0.068 \pm 0.007 \tag{7}
\]

This is consistent with (though slightly larger in central value than) the observed value \( R(D^+) = 0.026 \pm 0.016 \pm 0.018 \) [14]. The relative phase of \( C + A \) and \( T + C \) is about 70°, as can be seen from Fig. [2]. The real part of their ratio hence is small. A similar exercise can be applied to the decays \( D^+_s \rightarrow K^+K^0 \) and \( D^+_s \rightarrow K^+\bar{K}^0 \), which are related by U-spin to the \( D^+ \) decays discussed here. The corresponding ratio

\[
R(D^+_s) \equiv \frac{\Gamma(D^+_s \rightarrow K_SK^+) - \Gamma(D^+_s \rightarrow K_LK^+)}{\Gamma(D^+_s \rightarrow K_SK^+) + \Gamma(D^+_s \rightarrow K_LK^+)} \tag{8}
\]

is predicted to be

\[
R(D^+_s) = -2 \text{Re} \frac{\bar{C} + \bar{T}}{\bar{A} + \bar{C}}
\]
Table V Branching ratios, amplitudes, and representations in terms of reduced amplitudes for doubly-Cabibbo-suppressed decays. Amplitudes denoted by (a) involve interference between the doubly-Cabibbo-suppressed process shown and the corresponding Cabibbo-favored decay to $K^+ + X$.

| Meson | Decay mode | $\mathcal{B}$ ($10^{-4}$) | $p^*$ (MeV) | $|A|$ ($10^{-7}$ GeV) | Rep. |
|-------|------------|--------------------------|-------------|---------------------|------|
| $D^0$ | $K^+\pi^-$ | 1.45±0.04 | 861 | 1.54±0.02 | $T + E$ |
|       | $K^0\pi^0$ (a) | 860 | (a) | (a) | $(\bar{C} - \bar{E})/\sqrt{2}$ |
|       | $K^0\eta$ (a) | 772 | (a) | $\tilde{C}/\sqrt{3}$ | $-\tilde{C} + 3\bar{E})/\sqrt{6}$ |
|       | $K^0\eta'$ (a) | 565 | (a) | $\bar{T}/\sqrt{3}$ | $\tilde{C}/\sqrt{3}$ |
| $D^+$ | $K^0\pi^+$ (a) | 863 | (a) | $C + A$ | |
|       | $K^+\pi^0$ | 2.28±0.39 | 864 | 1.21±0.10 | $\bar{T} - \bar{A})/\sqrt{2}$ |
|       | $K^+\eta$ | 1.01±0.37 | 776 | 0.85±0.16 | $-\tilde{C}/\sqrt{3}$ |
|       | $K^+\eta'$ | $< 1.2$ | 571 | $< 1.08$ | $(\bar{T} + 3\tilde{A})/\sqrt{6}$ |
| $D_s^+$ | $K^0K^+$ (a) | 850 | (a) | $T + C$ | |

Figure 4: Graphs contributing to $D^0 \to (K^0\pi^0, \bar{K}^0\pi^0)$.

$$\begin{align*}
\delta &= 2 \tan^2 \theta_C \Re \frac{C + T}{A + \bar{C}} \\
\delta &= 0.019 \pm 0.002.
\end{align*}$$

6. Strong phases in $(D^0, \bar{D}^0) \to K^-\pi^+$

The relative strong phase in the CF decay $D^0 \to K^-\pi^+$ and the DCS decay $\bar{D}^0 \to K^-\pi^+$ is of interest in studying $B$ decays involving neutral $D$ mesons, where these two processes often can interfere. It was shown in Ref. [1] that one could measure this phase by producing a CP eigenstate $D^0_{CP}$, for example by tagging on a state of opposite CP at the $\psi(3770)$. Define decay amplitudes as

$$\langle K^-\pi^+ | D^0 \rangle \equiv A e^{i\delta_R}, \quad \langle K^-\pi^+ | \bar{D}^0 \rangle \equiv A e^{i\delta_W}. \quad (10)$$

The difference $\delta = \delta_R - \delta_W$ of strong phases would vanish in the SU(3) limit. At $\psi(3770)$ with $K^-\pi^+$ produced opposite a state $S_C$ with CP eigenvalue $\zeta$, one would have

$$\Gamma(K^-\pi^+, S_C) \approx A^2 A_{S_C}^2 (1 + 2\zeta r \cos \delta), \quad (11)$$

so by choosing states with $\zeta = \pm 1$ one can measure $(1 + 2r \cos \delta)/(1 - 2r \cos \delta)$, where $r = |A/A| = 0.057 \simeq \tan^2 \theta_C$.

In an analysis of 281 pb$^{-1}$ of CLEO data [13], the error on $\cos \delta$ is not yet conclusively determined, as a result of uncertainty in fits to $D^0, \bar{D}^0$ mixing. For an eventual integrated luminosity at CLEO of 750 pb$^{-1}$ and a cross section of $\sigma(e^+e^- \to \psi(3770) \to D\bar{D}) = 6$ nb one can estimate by rescaling the calculation in Ref. [1] an eventual error of $\Delta \cos \delta < 0.2$.

7. Other theoretical approaches

One can invoke effects of final state interactions to explain arbitrarily large SU(3) violations (if, for example, a resonance with SU(3)-violating couplings dominates a decay such as $D^0 \to \pi^+\pi^- + D^0 \to K^+K^-$). As one example of this approach [12], both resonant and nonresonant scattering can account for the observed ratio $\Gamma(D^0 \to K^+K^-)/\Gamma(D^0 \to \pi^+\pi^-) = 2.8 \pm 0.1$. This same approach predicted $\mathcal{B}(D^0 \to K^0\bar{K}^0) = 9.8 \times 10^{-4}$, a level of SU(3) violation consistent with the world average of Ref. [8] but far in excess of the recent CLEO value [2]. The paper of Ref. [10] may be consulted for many predictions for PV and PS final states in charm decays, where $V$ denotes a vector meson and $S$ denotes a scalar meson. Results for PV decays also may be found in Refs.
The recent discussion of Ref. [19] entails a prediction \( A \approx -0.4E \) (recall we were finding \( A \approx -E \)), essentially as a consequence of a Fierz identity and QCD corrections. Tree amplitudes are obtained from factorization and semileptonic \( D \to \pi \) and \( D \to K \) form factors. The main source of SU(3) breaking in \( /T \) is assumed to come from \( f_K/f_\pi = 1.22 \). Predictions include asymmetries \( R(D^0,\pi) = (2\tan^2 \theta_C, 0.068 \pm 0.007) \), and – via a sum rule for \( D^0 \to K^\mp\pi^\pm \) and \( D^+ \to K^+\pi^0 \) – expectation of \( |\delta| \approx 7-20^\circ \) (to be compared with 0 in exact SU(3) symmetry).

8. Summary

We have shown that the relative magnitudes and phases of amplitudes contributing to charm decays into two pseudoscalar mesons are describable by flavor symmetry. We have verified that there are large relative phases between the color-favored tree amplitude \( T \) and the color-suppressed amplitude \( C \), as well as between \( T \) and \( E \approx -A \).

The largest symmetry-breaking effects are visible in singly-Cabibbo-suppressed (SCS) decays, particularly in the \( D^0 \to (\pi^+\pi^-)/K^+K^- \) ratio which are at least in part understandable through form factor and decay constant effects. Decays involving \( \eta, \eta' \) are mostly describable with small “disconnected” amplitudes, a possible exception being in SCS \( D^+ \) and \( D_s^+ \) decays.

One sees evidence for the expected interference between Cabibbo-favored (CF) and doubly-Cabibbo-suppressed decays in \( D^{0,+,+} \to K_{S,L}\pi^{0,+,+} \) decays. As a result of CLEO’s present data on \( (D^0,\bar{D}^0) \to K^-\pi^+ \), limits are being placed on the relative strong phase \( \delta \) between these amplitudes, and the full CLEO data sample is expected to result in an error equal to or better than \( \Delta (\cos \delta) = 0.2 \).

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References

[12] Q. He et al. [CLEO Collaboration], hep-ex/0607068 contributed to 33rd International Conference on High Energy Physics (ICHEP 06), Moscow, Russia, 26 July - 2 August 2006.
Study of a model-independent method for the measurement of the angle $\phi_3$

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This report shows the latest results on the study of the method to determine the angle $\phi_3$ of the unitarity triangle using Dalitz plot analysis of $D^0 \to K_S^0 \pi^+ \pi^-$ decay from $B^\pm \to DK^\pm$ process in a model-independent way. We concentrate on the case with a limited charm data sample, which will be available from the CLEO-c collaboration in the nearest future, with the main goal to find the optimal strategy for $\phi_3$ extraction. We find that the analysis using decays of $D_{CP}$ only cannot provide a completely model-independent measurement in the case of limited data sample. The procedure involving binned analysis of $B^\pm \to DK^\pm$ and $\psi(3770) \to (K_S^0 \pi^+ \pi^-)D(K_S^0 \pi^+ \pi^-)_D$ decays is proposed, that allows to obtain the $\phi_3$ precision comparable to unbinned model-dependent fit.

1. Introduction

The measurement of the angle $\phi_3$ ($\gamma$) of the unitarity triangle using Dalitz plot analysis of the $D^0 \to K_S^0 \pi^+ \pi^-$ decay from $B^\pm \to DK^\pm$ process, introduced by Giri et al. [1] and Belle collaboration [2] and successfully implemented by BaBar [3] and Belle [4], presently offers the best constraints on this quantity. However, this technique is sensitive to the choice of the model used to describe the three-body $D^0$ decay. Currently, this uncertainty is estimated to be $\sim 10^3$ and due to large statistical error does not affect the precision of $\phi_3$ measurement. As the amount of $B$ factory data increases, though, this uncertainty will become a major limitation. Fortunately, a model-independent approach exists (see [1]), which uses the data of the $\tau$-charm factory to obtain the missing information about the $D^0$ decay amplitude.

In our previous study of the model-independent Dalitz analysis technique [3] we have implemented a procedure proposed by Giri et al. involving the division of the Dalitz plots into bins, and shown that this procedure allows to measure the phase $\phi_3$ with the statistical precision only 30–40% worse than in the unbinned model-dependent case. We did not attempt to optimize the binning and mainly considered a high-statistics limit with an aim to estimate the sensitivity of the future super-B factory.

The data useful for model-independent measurement are presently available from the CLEO-c experiment [6]. The integrated luminosity at the $\psi(3770)$ resonance decaying to $D\bar{D}$ available for the analysis is 400 pb$^{-1}$. By the end of CLEO-c operation this statistics will grow up to 750 pb$^{-1}$ [7]. This corresponds to $\sim 1000$ events where $D$ meson in a $CP$ eigenstate decays to $K_S^0 \pi^+ \pi^-$, and twice as much events of $\psi(3770) \to D^0 \bar{D}^0$ with both $D$ mesons decaying to $K_S^0 \pi^+ \pi^-$. Both of these processes include the information useful for a model-independent $\phi_3$ measurement. In this paper, we report on studies of the model-independent approach with a limited statistics of both $\psi(3770)$ and $B$ data, using both $D_{CP} \to K_S^0 \pi^+ \pi^-$ and $(K_S^0 \pi^+ \pi^-)_D(K_S^0 \pi^+ \pi^-)_D$ final states.

2. Model-independent approach

The density of $D^0 \to K_S^0 \pi^+ \pi^-$ Dalitz plot is given by the absolute value of the amplitude $f_D$ squared:

$$p_D = |D(m^2_+, m^2_-)|^2 = |f_D(m^2_+, m^2_-)|^2$$  \hspace{1cm} (1)

In the case of no $CP$-violation in $D$ decay the density of the $D^0$ decay $\bar{p}_D$ equals to

$$\bar{p}_D = |\bar{f}_D|^2 = p_D(m^2_+, m^2_-).$$  \hspace{1cm} (2)

Then the density of the $D$ decay Dalitz plot from $B^\pm \to DK^\pm$ process is expressed as

$$p_{B^\pm} = |f_D + r_B e^{i(\delta_B + \phi_3)} \bar{f}_D|^2 = p_D + r^2_B \bar{p}_D + 2\sqrt{p_D \bar{p}_D}(x_\pm c + y_\pm s),$$  \hspace{1cm} (3)

where $x_\pm, y_\pm$ include the value of $\phi_3$ and other related quantities, the strong phase $\delta_B$ of the $B^\pm \to DK^\pm$ decay, and amplitude ratio $r_B$:

$$x_\pm = r_B \cos(\delta_B + \phi_3); \hspace{0.5cm} y_\pm = r_B \sin(\delta_B + \phi_3).$$  \hspace{1cm} (4)

The functions $c$ and $s$ are the cosine and sine of the strong phase difference $\Delta \delta_D$ between the symmetric Dalitz plot points:

$$c = \cos(\delta_D(m^2_+, m^2_-) - \delta_D(m^2_-, m^2_+)) = \cos \Delta \delta_D;$$

$$s = \sin(\delta_D(m^2_+, m^2_-) - \delta_D(m^2_-, m^2_+)) = \sin \Delta \delta_D.$$  \hspace{1cm} (5)

The phase difference $\Delta \delta_D$ can be obtained from the sample of $D$ mesons in a $CP$-eigenstate, decaying to $K_S^0 \pi^+ \pi^-$. The Dalitz plot density of such decay is

$$p_{CP} = |f_D| f_D|^2 = p_D + \bar{p}_D + 2\sqrt{p_D \bar{p}_D} c$$  \hspace{1cm} (6)

(the normalization is arbitrary). Decays of $D$ mesons in $CP$ eigenstate to $K_S^0 \pi^+ \pi^-$ can be obtained in the process, e.g. $e^+e^- \to \psi(3770) \to D\bar{D}$, where the
other (tag-side) $D$ meson is reconstructed in the $CP$ eigenstate, such as $K^+K^-$ or $K^{0}_{S}\omega$.

Another possibility is to use a sample, where both $D$ mesons (we denote them as $D$ and $D'$) from the $\psi(3770)$ meson decay into the $K^{0}_{S}\pi^+\pi^-$ state [4]. Since $\psi(3770)$ is a vector, two $D$ mesons are produced in a $P$-wave, and the wave function of the two mesons is antisymmetric. Then the four-dimensional density of two correlated Dalitz plots is

$$p_{\text{corr}}(m_1^2, m_2^2, m_1', m_2') = |f_D f_D' - f_D f_D'|^2 = p_D p_D' + \bar{p}_D p_D' - 2 \sqrt{p_D \bar{p}_D p_D' \bar{p}_D}(cc' + ss'),$$

(7)

This decay is sensitive to both $c$ and $s$ for the price of having to deal with the four-dimensional phase space.

In a real experiment, one measures scattered data rather than a probability density. To deal with real data, the Dalitz plot can be divided into bins. In what follows, we show that using the appropriate binning, it is possible to reach the statistical sensitivity equivalent to the model-dependent case.

### 3. Binned analysis with $D_{CP}$ data

The binned approach was proposed by Giri et al. [5]. Assume that the Dalitz plot is divided into $2N$ bins symmetrically to the exchange $m_1^2 \leftrightarrow m_2^2$. The bins are denoted by the index $i$ ranging from $-N$ to $N$ (excluding 0); the exchange $m_1^2 \leftrightarrow m_2^2$ corresponds to the exchange $i \leftrightarrow -i$. Then the expected number of events in the bins of the Dalitz plot of $D$ decay from $B^{+} \rightarrow D^{0} K^{\pm}$ is

$$\langle N_i \rangle = h_B[K_i + \bar{K}_i(xc_i + ys_i)],$$

(8)

where $K_i$ is the number of events in the bins in the Dalitz plot of the $D^{0}$ in a flavor eigenstate, $h_B$ is the normalization constant. Coefficients $c_i$ and $s_i$, which include the information about the cosine and sine of the phase difference, are given by

$$c_i = \frac{\int_{D_i} \sqrt{p_D p_D'} \cos(\Delta \delta_D(m_1'^2, m_2'^2))dD}{\int_{D_i} dD \int_{D_i} \sqrt{p_D p_D'} dD},$$

(9)

$$s_i = \frac{\int_{D_i} \sqrt{p_D p_D'} \sin(\Delta \delta_D(m_1'^2, m_2'^2))dD}{\int_{D_i} dD \int_{D_i} \sqrt{p_D p_D'} dD},$$

(10)

$s_i$ is defined similarly with cosine substituted by sine. Here $D_i$ is the bin region, over which the integration is performed. Note that $c_i = c_{-i}$, $s_i = -s_{-i}$ and $c_i^2 + s_i^2 \leq 1$ (the equality $c_i^2 + s_i^2 = 1$ being satisfied if the amplitude is constant across the bin).

The coefficients $K_i$ are obtained precisely from a very large sample of $D^{0}$ decays in the flavor eigenstate, which is accessible at $B$-factories. The expected number of events in the Dalitz plot of $D_{CP}$ decay equals to

$$\langle M_i \rangle = h_{CP}[K_i + K_{-i} + 2\sqrt{K_i K_{-i}}c_i],$$

(11)

and thus can be used to obtain the coefficient $c_i$. As soon as the $c_i$ and $s_i$ coefficients are known, one can obtain $x$ and $y$ values (hence, $\phi_3$ and other related quantities) by a maximum likelihood fit using equation (12).

Note that now the quantities of interest $x$ and $y$ (and consequently $\phi_3$) have two statistical errors: one due to a finite sample of $B^{+} \rightarrow D^{0} K^{\pm}$ data, and due to $D_{CP} \rightarrow K^{0}_{S}\pi^+\pi^-$ statistics. We will refer to these errors as $B$-statistical and $D_{CP}$-statistical, respectively.

Obtaining $s_i$ is a major problem in this analysis. If the binning is fine enough, so that both the phase difference and the amplitude remain constant across the area of each bin, expressions (9) reduce to $c_i = \cos(\Delta \delta_D)$ and $s_i = \sin(\Delta \delta_D)$, so $s_i$ can be obtained as $s_i = \pm \sqrt{1 - c_i^2}$. Using this equality if the amplitude varies will lead to the bias in the $x, y$ fit result. Since $c_i$ is obtained directly, and $s_i$ is overestimated by the absolute value, the bias will mainly affect $y$ determination, resulting in lower absolute values of $y$.

Our studies [5] show that the use of equality $c_i^2 + s_i^2 = 1$ is satisfactory for the number of bins around 200 or more, which cannot be used with presently available $D_{CP}$ data. It is therefore essential to find a relatively coarse binning (the number of bins being 10–20) which a) allows to extract $s_i$ from $c_i$ with low bias, and b) has the sensitivity of the $\phi_3$ phase comparable to the unbinned model-dependent case.

Fortunately, both the a) and b) requirements appear to be equivalent. To determine the $B$-statistical sensitivity of a certain binning, let’s define a quantity $Q$ — a ratio of a statistical sensitivity to that in the unbinned case. Specifically, $Q$ relates the number of standard deviations by which the number of events in bins is changed by varying parameters $x$ and $y$, to the number of standard deviations if the Dalitz plot is divided into infinitely small regions (the unbinned case):

$$Q^2 = \frac{\sum_i \left(\frac{1}{\sqrt{N_i}} \frac{dN_i}{dx}\right)^2 + \left(\frac{1}{\sqrt{N_i}} \frac{dN_i}{dy}\right)^2}{\int_D \left[\left(\frac{1}{\sqrt{|f_B|}} \frac{|d(f_B)|^2}{dx}\right)^2 + \left(\frac{1}{\sqrt{|f_B|}} \frac{|d(f_B)|^2}{dy}\right)^2\right] dD},$$

(11)

$$Q^2|_{x=y=0} = \sum_i (c_i^2 + s_i^2) N_i \sum_i N_i$$

(12)

Therefore, the binning which satisfies $c_i^2 + s_i^2 = 1$ (i.e., the absence of bias if $s_i$ is calculated as $\sqrt{1 - c_i^2}$) also has the same sensitivity as the unbinned approach. The factor $Q$ defined this way is not necessarily the
best measure of the binning quality (the binning with higher $Q$ can be insensitive to either $x$ or $y$, which is impractical from the point of measuring $\phi_3$), but it allows an easy calculation and correctly reproduces the relative quality for a number of binnings we tried in our simulation.

The choice of the optimal binning naturally depends on the $D^0$ model. In our studies we use the two-body amplitude obtained in the latest Belle $\phi_3$ Dalitz analysis [4].

From the consideration above it is clear that a good approximation to the optimal binning is the one obtained from the uniform division of the strong phase difference $\Delta \delta_D$. In the half of the Dalitz plot $m^2_D < m^2_\pi$ (i.e., the bin index $i > 0$) the bin $D_i$ is defined by condition

$$2\pi(i - 1/2)/N < \Delta \delta_D(m^2_\pi, m^2_\rho) < 2\pi(i + 1/2)/N,$$

(13)
in the remaining part ($i < 0$) the bins are defined symmetrically. We will refer to this binning as $\Delta \delta_D$-binning. As an example, such a binning with $N = 8$ is shown in Fig. 1 (left). Although the phase difference variation across the bin is small by definition, the absolute value of the amplitude can vary significantly, so the condition $c_i^2 + s_i^2 = 1$ is not satisfied exactly. The values of $c_i$ and $s_i$ in this binning are shown in Fig. 1 (bottom left) with crosses.

Figure 1 (right) shows the division with $N = 8$ obtained by continuous variation of the $\Delta \delta_D$-binning to maximize the factor $Q$. The sensitivity factor $Q$ increases to 0.89 compared to 0.79 for $\Delta \delta_D$-binning.

![Figure 1](image)

Figure 1: Divisions of the $D^0 \rightarrow K^0_S \pi^+ \pi^-$ Dalitz plot. Uniform binning of $\Delta \delta_D$ strong phase difference with $N = 8$ (left), and the binning obtained by variation of the latter to maximize the sensitivity factor $Q$ (right).

We perform a toy MC simulation to study the statistical sensitivity of the different binning options. We use the amplitude from the Belle analysis [4] to generate decays of flavor $D^0$, $D_{CP}$, and $D$ from $B^\pm \rightarrow DK^\pm$ decays to the $K^0_S \pi^+ \pi^-$ final state according to the probability density given by [1], [8], and [8], respectively. To obtain the $B$-statistical error we use a large number of $D^0$ and $D_{CP}$ decays, while the generated number of $D$ decays from the $B^\pm \rightarrow DK^\pm$ process ranges from 100 to 100000. For each number of $B$ decay events, 100 samples are generated, and the statistical errors of $x$ and $y$ are obtained from the spread of the fitted values. A study of the error due to $D_{CP}$ statistics is performed similarly, with a large number of $B$ decays, and the statistics of $D_{CP}$ decays varied. Both errors are checked to satisfy the square root scaling.

The binning options used are $\Delta \delta_D$-binning with $N = 8$ and $N = 20$, as well as “optimal” binnings with maximized $Q$ obtained from these two with a smooth variation of the bin shape. Note that the “optimal” binning with $N = 20$ offers the $B$-statistical sensitivity only 4% worse than an unbinned technique. For comparison, we use the binnings with the uniform division into rectangular bins (with $N = 8$ and $N = 19$ in the allowed phase space, the ones which are denoted as $3 \times 3$ and $5 \times 5$ in [8]).

The $B$- and $D_{CP}$-statistical precision of different binning options, recalculated to 1000 events of both $B$ and $D_{CP}$ samples, as well as their calculated values of the factor $Q$, are shown in Table I. In the present study we use the errors of parameters $x$ and $y$ rather than $\phi_3$ as a measure of the statistical power since they are nearly independent of the actual values of $\phi_3$, strong phase $\delta$ and amplitude ratio $r_B$. The error of $\phi_3$ can be obtained from these numbers given the value of $r_B$. The factor $Q$ reproduces the ratio of the values $\sqrt{1/\sigma_x^2 + 1/\sigma_y^2}$ for the binned and unbinned approaches with the precision of 1–2%. While the binning with maximized $Q$ offers better $B$-statistical sensitivity, the best $D_{CP}$-statistical precision of the options we have studied is reached for the $\Delta \delta_D$-binning. However, for the expected amount of experimental data of $B$ and $D_{CP}$ decays the $B$-statistical error dominates, therefore, slightly worse precision due to $D_{CP}$ statistics does not affect significantly the total precision.

We have considered the choice of the optimal binning only from the point of statistical power. However, the conditions to satisfy low model dependence are quite different. Since the bins in the binning options we have considered are sufficiently large, the requirement that the phase does not change over the bin area is a strong model assumption. We have performed toy MC simulation to study the model dependence. While the binning was kept the same as in the statistical power study (based on the phase difference from the default $D^0$ amplitude), the amplitude used to generate $D^0$, $D_{CP}$ and $B^\pm \rightarrow DK^\pm$ decays was altered in the same way as in the Belle study of the model-dependence in the unbinned analysis [4]. As a result, the same bias of $\Delta \phi_3 \sim 10^\circ$ is observed as in unbinned analysis. The bias in $x$ and $y$ if demonstrated in Fig. 2. We remind that the cause of this bias is a fixed relation between the $c_i$ and $s_i$. Therefore, proposed binning options, although providing good statistical precision, are not flexible enough to provide also a low model dependence. To minimize
Table I Statistical precision of \((x, y)\) determination using different binnings and with an unbinned approach. The errors correspond to 1000 events in both the \(B\) and \(D_{CP}\) \((K_0^{0}\pi^+\pi^-)^2\) samples.

<table>
<thead>
<tr>
<th>Binning</th>
<th>Q</th>
<th>(\sigma_x)</th>
<th>(\sigma_y)</th>
<th>(\sigma_{x})</th>
<th>(\sigma_{y})</th>
<th>(\sigma_{x,y})</th>
</tr>
</thead>
<tbody>
<tr>
<td>(N = 8) (uniform)</td>
<td>0.57</td>
<td>0.0331</td>
<td>0.0600</td>
<td>0.0053</td>
<td>0.0097</td>
<td>0.0145</td>
</tr>
<tr>
<td>(N = 8) ((\Delta\delta_D))</td>
<td>0.79</td>
<td>0.0273</td>
<td>0.0370</td>
<td>0.0042</td>
<td>0.0072</td>
<td>0.0050</td>
</tr>
<tr>
<td>(N = 8) (optimal)</td>
<td>0.89</td>
<td>0.0232</td>
<td>0.0324</td>
<td>0.0058</td>
<td>0.0114</td>
<td>0.0082</td>
</tr>
<tr>
<td>(N = 19) (uniform)</td>
<td>0.69</td>
<td>0.0274</td>
<td>0.0549</td>
<td>0.0042</td>
<td>0.0112</td>
<td>-</td>
</tr>
<tr>
<td>(N = 20) ((\Delta\delta_D))</td>
<td>0.82</td>
<td>0.0266</td>
<td>0.0350</td>
<td>0.0048</td>
<td>0.0074</td>
<td>-</td>
</tr>
<tr>
<td>(N = 20) (optimal)</td>
<td>0.96</td>
<td>0.0223</td>
<td>0.0290</td>
<td>0.0078</td>
<td>0.0110</td>
<td>-</td>
</tr>
<tr>
<td>Unbinned</td>
<td>-</td>
<td>0.0213</td>
<td>0.0279</td>
<td>-</td>
<td>-</td>
<td>-</td>
</tr>
</tbody>
</table>

4. Binned analysis with correlated \(D^0 \to K_0^{0}\pi^+\pi^-\) data

The use of the \(\psi(3770)\) decays where both neutral \(D\) mesons decay to the \(K_0^{0}\pi^+\pi^-\) state allows to significantly increase the amount of data useful to extract phase information in \(D^0\) decay. It is also possible to detect events of \(\psi(3770) \to (K_0^{0}\pi^+\pi^-)_D (K_0^{0}\pi^+\pi^-)_D\), where \(K_0^L\) is not reconstructed, and its momentum is obtained from kinematic constraints. The number of these events is approximately twice that of \((K_0^{0}\pi^+\pi^-)^2\) due to combinatorics. However, it is impossible to simply combine these samples since the phases of the doubly Cabibbo-suppressed components in \(\overline{D}^0 \to K_0^{0}\pi^+\pi^-\) and \(D^0 \to K_0^{0}\pi^+\pi^-\) amplitudes are opposite [8]. In the analysis of \(B\) data only \(K_0^{0}\pi^+\pi^-\) state can be used, but it is possible to utilize \(K_0^{0}\pi^+\pi^-\) data to better constrain the \(\overline{D}^0 \to K_0^{0}\pi^+\pi^-\) amplitude using model assumptions based on \(SU(3)\) symmetry. In what follows, we will consider the use of \(K_0^{0}\pi^+\pi^-\) data only.

In the case of a binned analysis, the number of events in the region of the \((K_0^{0}\pi^+\pi^-)^2\) phase space is

\[
\langle M \rangle_{ij} = h_{corr}[K_i K_{-j} + K_{-i} K_j] - 2\sqrt{K_i K_{-i} K_j K_{-j}(c_i c_j + s_i s_j)].
\]

Here two indices correspond to two \(D\) mesons from \(\psi(3770)\) decay. It is logical to use the same binning as in the case of \(D_{CP}\) statistics to improve the precision of the determination of \(c_i\) coefficients, and to obtain \(s_i\) from data without model assumptions, contrary to \(D_{CP}\) case. The obvious advantage of this approach is its being unbiased for any finite \((K_0^{0}\pi^+\pi^-)^2\) statistics (not asymptotically as in the case of \(D_{CP}\) data).

Note that in contrast to \(D_{CP}\) analysis, where the sign of \(s_i\) in each bin was undetermined and has to be fixed using model assumptions, \((K_0^{0}\pi^+\pi^-)^2\) analysis has only a four-fold ambiguity: change of the sign of all \(c_i\) or all \(s_i\). In combination with \(D_{CP}\) analysis, where the sign of \(c_i\) is fixed, this ambiguity reduces to only two-fold. One of the two solutions can be
chosen based on a weak model assumption (incorrect $s_i$ sign corresponds to complex-conjugated $D$ decay amplitude, which violates causality requirement when parameterized with the Breit-Wigner amplitudes).

Coefficients $c_i$, $s_i$ can be obtained by minimizing the negative logarithmic likelihood function

$$-2 \log L = -2 \sum_{i,j} \log P(M_{ij}, \langle M \rangle_{ij})$$

where $P(M, \langle M \rangle)$ is the Poisson probability to get $M$ events with the expected number of $\langle M \rangle$ events.

The number of bins in the 4-dimensional phase space is $4N^2$ rather than $2N$ in the $D_{CP}$ case. Since the expected number of events in correlated $K_S^0\pi^+\pi^-$ data is of the same order as for $D_{CP}$, the bins will be much less populated. This, however, does not affect the precision of $c_i$, $s_i$ determination since each of the free parameters is constrained by many bins.

The coefficients $c_i$, $s_i$ obtained this way can then be used to constrain $x$, $y$ with the maximum likelihood fit of the $B$ decay data using Eq. [9] To correctly account for the errors of $c_i$, $s_i$ determination, this likelihood should include distributions of these quantities, in addition to Poisson fluctuations in $B$ data bins. A more convenient way is to use the common likelihood function, covering both $B$ and $(K_S^0\pi^+\pi^-)^2$ data:

$$-2 \log L = -2 \sum_{i,j} \log P(M_{ij}, \langle M \rangle_{ij}) - 2 \sum_i \log P(N_i, \langle N \rangle_i),$$

with $x$, $y$, $h_B$, $h_{corr}$, $c_i$ and $s_i$ as free parameters. This approach is also more optimal in the case of large $B$ data sample, since it imposes additional constraints on $c_i$, $s_i$ values.

The toy MC simulation was performed to study the procedure described above. Using the amplitude from the Belle analysis, we generate a large number of $D^0 \rightarrow K_S^0\pi^+\pi^-$ decays and several sets of $(K_S^0\pi^+\pi^-)^2$ decays (according to the probability density given by [4]) and $B$ decays [3]. We use the same binning options as in $D_{CP}$ study with $N = 8$. The combined negative likelihood [10] is minimized in the fit to each toy MC sample. We constrain $c_i^2 + s_i^2 < 1$ in the fit to avoid entering unphysical region with negative number of events in the bin. The number of $(K_S^0\pi^+\pi^-)^2$ and $B$ decays ranges from $10^3$ to $10^5$. The errors of $x$ and $y$ parameters are calculated from the spread of the fitted values. If the number of $(K_S^0\pi^+\pi^-)^2$ decays is comparable or larger than the number of $B$ decays, the $x$ and $y$ errors can be represented as quadratic sums of two errors, each scaled as a square root of $(K_S^0\pi^+\pi^-)^2$ and $B$ statistics, respectively. However if the number of $B$ decays is large, the errors of $c_i$ and $s_i$ depend also on $B$ decay statistics, so separating the total error into $B$- and $(K_S^0\pi^+\pi^-)^2$-statistical errors becomes impossible.

The best $(K_S^0\pi^+\pi^-)^2$-statistical error is obtained for $\Delta\delta_D$-binning and recalculated to 1000 events yields $\sigma_x = 0.0050$, $\sigma_y = 0.0095$, which is only slightly worse than the error obtained with the same amount of $D_{CP}$ data (see Table I for comparison). We also check that significant change of the model used to define the binning does not lead to the systematic bias (although it does decrease the statistical precision). Figure [3] demonstrates the precision of the determination of $c_i$, $s_i$ coefficients in our toy MC study and the absence of the systematic bias for both $x$ and $y$ when the model is varied.

The numbers of $(K_S^0\pi^+\pi^-)^2$ and $D_{CP}$ decays in $\psi(3770)$ data are comparable, so are the statistical errors due to $\psi(3770)$ data sample for the two approaches. The same binning can be used in both approaches, therefore improving the accuracy of $c_i$ determination. The approach based on $(K_S^0\pi^+\pi^-)^2$ data allows to extract both $c_i$ and $s_i$ without additional model uncertainties, so it can be used to check the validity of the constraint $c_i^2 + s_i^2 = 1$ and therefore to test the sensitivity of the particular binning.
5. Conclusion

We have studied the model-independent approach to \(\phi_3\) measurement using \(B^\pm \to DK^\pm\) decays with neutral \(D\) decaying to \(K_S^0\pi^+\pi^-\). The analysis of \(\psi(3770) \to D\bar{D}\) data allows to extract the information about the strong phase in \(D^0 \to K_S^0\pi^+\pi^-\) decay that is fixed by model assumptions in a model-dependent technique. We specially consider the case with a limited \(\psi(3770) \to D\bar{D}\) data sample which will be available from CLEO-c in the nearest future.

In the binned analysis, we propose a way to obtain the binning that offers an optimal statistical precision (close to the precision of an unbinned approach). Two different strategies of the binned analysis are considered: using \(D_{CP}\) \(K_S^0\pi^+\pi^-\) data sample, and using decays of \(\psi(3770)\) to \((K_S^0\pi^+\pi^-)_D(K_S^0\pi^+\pi^-)_D\). The strategy using \(D_{CP}\) decays alone cannot offer a completely model-independent measurement: it provides only the information about \(c_i\) coefficients, while \(s_i\) for low \(D_{CP}\) statistics has to be fixed using model assumptions. However, as the \(D_{CP}\) data sample increases, model-independence can be reached by reducing the bin size. The strategy using the \(\psi(3770)\) to \((K_S^0\pi^+\pi^-)_D(K_S^0\pi^+\pi^-)_D\) sample, in contrast, allows to obtain both \(c_i\) and \(s_i\) with an accuracy comparable to \(D_{CP}\) approach. Both strategies can use the same binning of the \(D^0 \to K_S^0\pi^+\pi^-\) Dalitz plot and therefore can be used in combination to improve the accuracy due to \(\psi(3770)\) statistics.

The expected sensitivity is obtained based on the \(D^0\) decay model from Belle analysis. For the CLEO-c statistics of 750 pb\(^{-1}\) \((\sim 1000 D_{CP}\) and \((K_S^0\pi^+\pi^-)^2\) events) the expected errors of parameters \(x\) and \(y\) due to \(\psi(3770)\) statistics are of the order of 0.01. For \(r_B = 0.1\) it gives the \(\phi_3\) precision \(\sigma_{\phi_3} = \sigma_{x,y}/(\sqrt{2}r_B) \simeq 5^\circ\), which is far below the expected error due to present-day \(B\) data sample. Further improvement of \(\phi_3\) precision will require larger charm dataset, which can be provided by BES-III experiment [10].

In our study, we did not consider the experimental systematic uncertainties e.g. due to imperfect knowledge of the detection efficiency or background composition. We believe these issues can be addressed in a similar manner as in already completed model-dependent analyses.

6. Acknowledgments

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[8] E. White, Q. He, talk at Charm 07, Ithaca, USA.
Impact on $\gamma/\phi_3$ from CLEO-c Using $CP$-Tagged $D \to K_{S,L}\pi^+\pi^-$ Decays

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Precision determination of the CKM angle $\gamma/\phi_3$ depends upon constraints on charm mixing amplitudes, measurements of doubly-Cabibbo suppressed amplitudes and relative phases, and studies of charm Dalitz plots tagged by flavor or $CP$ eigenstates. In this note we describe the technique used at CLEO-c to constrain the $K_{S,L}\pi^+\pi^-$ uncertainty, and its impact on $\gamma/\phi_3$ measurements at $B$-factories presented at the Charm 2007 Workshop.

1. Introduction

Measurement of the CKM angle $\gamma/\phi_3$ is challenging. Several methods have been proposed using $B^\pm \to DK^\pm$ decays; 1) the Gronau-London-Wyler (GLW) method [1] where the $D$ decays to $CP$ eigenstates 2) the Atwood-Dunietz-Soni (ADS) method [2] where the $D$ decays to flavor eigenstates and 3) the Dalitz plot method [3, 5] where the $D$ decays to a three-body final state. This latter method has been used recently by CLEO to measure the $K^+K^0$ strong phase via the three-body decay $D^0 \to K^+K^-\pi^0$ [6]. Uncertainties due to charm contribute to each of these methods. The CLEO-c physics program includes a variety of charm measurements that will improve the determination of $\gamma/\phi_3$ from the $B$-factory experiments, Babar and Belle. The pertinent components of this program are improved constraints on charm mixing amplitudes - important for GLW, measurement of the relative strong phase $\delta$ between $D^0$ and $\bar{D}^0$ decay to $K^+\pi^-$ - important for ADS, and studies of charm Dalitz plots tagged by hadronic flavor or $CP$ eigenstates. The total number of charm mesons accumulated at CLEO-c will be much smaller than the samples already accumulated by the $B$-factories. However, quantum correlations in the $DD$ system from $\psi(3770)$ provides a unique laboratory in which to study charm.

The decay with the largest branching fraction relevant to the determination of $\gamma/\phi_3$, $D^0 \to K_S\pi^+\pi^-$. Recently Babar [7] and Belle [8] have reported $\gamma = (92 \pm 41 \pm 11 \pm 12)^o$ and $\phi_3 = (53^{+15}_{-18} \pm 3 \pm 9)^o$, respectively, where the third error is the systematic error due to modeling of the Dalitz plot.

Both $D^0$ and $\bar{D}^0$ populate the Dalitz plots $K_S\pi^+\pi^-$, (as well as $\pi^+\pi^-\pi^0$, $K^+\pi^-\pi^0$ and $K_S^0\pi^+\pi^-$) and so can be used in the determination of $\gamma/\phi_3$ which exploit the interference between $b \to c\bar{u}s$ ($B^- \to D^0 K^-$) and $b \to u\bar{c}s$ ($B^- \to \bar{D}^0 K^-$) where the former process is real and the latter is parametrized by $\sim e^{-\gamma}$. Studying $CP$ tagged Dalitz plots allows a model independent determination of the relative $D^0$ and $\bar{D}^0$ phase across the Dalitz plot. We describe this technique in the following sections.

2. Determining $\gamma/\phi_3$ From $B$ Decays

Our analysis follows the work outlined in [3], [5], and [4]. We consider the decay process $B^\pm \to DK^\pm$, followed by the three-body decay $D \to K_S\pi^+\pi^-$. Assuming no $CP$ violation, we define the decay amplitudes for the $D^0$ and $\bar{D}^0$ to be

$$A(D^0 \to K_S\pi^+\pi^-; x, y) = f_D(x, y)$$
$$A(\bar{D}^0 \to K_S\pi^-\pi^-; x, y) = f_D(y, x).$$

Sensitivity to the angle $\phi_3$ comes from the interference of the neutral $D$ mesons from $B^\pm \to DK^\pm$. Since the $D$ meson is in a linear combination of flavor states, the amplitude for a $D^0 \to K_S\pi^+\pi^-$ event originating from a $B$ decay is then

$$A(B^- \to (K_S\pi^+\pi^-)_D K^-) \propto f_D + r_B e^{i\theta} f_D^*. \tag{2}$$

up to an overall normalization. The angle $\theta$ is defined as $\theta_{\pm} \equiv \delta_B \pm \phi_3$. Here $\delta_B$ is the strong phase difference between color-suppressed and favored amplitudes, whilst $r_B$ is the ratio between the color-suppressed to favored amplitudes. Theoretical estimates place $r_B$ between 0.1-0.2 [11]. This has been confirmed by BaBar ($r_B = 0.12 \pm 0.08 \pm 0.03$ (syst) $\pm 0.04$ (model), [7]) and Belle ($r_B = 0.16 \pm 0.05 \pm 0.01$ (syst) $\pm 0.05$ (model), [8]).

The $D^0 \to K_S\pi^+\pi^-$ Dalitz plot is divided into $2N$ bins, symmetric under exchange of $x$ and $y$. The bins are indexed from $-i$ to $i$, excluding zero, as in shown in Fig 1. The coordinate exchange $x \leftrightarrow y$ thus corresponds to the exchange of bins $i \leftrightarrow -i$. For simplicity we ignore the effects of efficiency and background in the Dalitz plot. The number of events in the $i$-th bin of the $K_S\pi^+\pi^-$ Dalitz plot from a $D$ decay is then expressed as

$$K_i = A_D \int_{D_i} |f_D(x, y)|^2 dx \ dy = A_D F_i. \tag{3}$$

The interference between the $D^0$ and $\bar{D}^0$ amplitudes is parametrized by the two quantities

$$c_i = \frac{1}{\sqrt{F_i F_{-i}}} \int_{D_i} \text{Re} \left[ f_D(x, y) f_D^*(y, x) \right] dx \ dy. \tag{4}$$
In terms of this amplitude the number of events in the $i$-th bin of a $CP$-tagged Dalitz plot is

$$M_i^\pm = h_{CP\pm} \left( K_i \pm 2c_i \sqrt{|K_{i-1} - K_{-i}|} \right),$$  \hfill (8)

where $h_{CP\pm}$ is a normalization factor.

The expression given above for $M_i^\pm$ can be used to measure $c_i$ directly, even if only one type of $CP$ tag is reconstructed. Care must be taken to use the corresponding value of $h_{CP\pm}$ as defined above. However, if samples of both $CP$ parities are available we can combine the expressions for $M_i^+$ and $M_i^-$ to get the following equation

$$c_i = \frac{1}{2} \frac{(M_{i+} - M_{i-}) (K_i - K_{-i})}{(M_{i+} + M_{i-}) \sqrt{|K_{i-1} - K_{-i}|}}.$$  \hfill (9)

We thus have an expression for measuring $c_i$ simply by counting events within the bins of flavor-tagged and $CP$-tagged Dalitz plots.

At CLEO-c we produce $D^0\bar{D}^0$ pairs from the decay of a $\psi(3770)$ in a definite eigenstate of $C = -1$. Ignoring both the effects of $CP$ violation, the double tag rate for final states $|1\rangle$ and $|2\rangle$ is given by

$$\Gamma(1,2) = |A(1,2)|^2 + \text{(mixing terms)},$$  \hfill (10)

where

$$A(1,2) \equiv \langle 1|D^0\rangle\langle 2|\bar{D}^0\rangle - \langle 1|\bar{D}^0\rangle\langle 2|D^0\rangle.$$  \hfill (11)

For the time being we ignore the effects of correlations and mixing in the $K\pi$ tagged Dalitz plot. This is not expected to make a significant difference for the $K\pi$ mode, as terms proportional to $r_{K\pi} \approx 0.06$ and $r_{K\pi}^2$ are negligible.

### 3.1. Optimized Binning

Although the quantity $s_i$ can only be measured using a $K_S\pi^+\pi^-$ vs. $K_S\pi^+\pi^-$ double Dalitz analysis [4], it can still be approximated from a single Dalitz plot if the binning is fine enough. If the bins are small enough that the phase difference and the amplitude remains constant across each bin, the strong phase parameters become $c_i = \cos(\delta_D)$, $s_i = \sin(\delta_D)$, so that the equality $s_i = \sqrt{1 - c_i^2}$ is true. It has been shown [3] that this equality holds for 200 or more bins, which is clearly not feasible for the number of $D_{CP}$ tags produced at CLEOc. In order to circumvent this problem, Bondar has proposed an alternate, model-dependent method for binning the $K_S\pi^+\pi^-$ Dalitz plot[4]. The optimal choice depends on the $D^0 \to K_S\pi^+\pi^-$ model. In this analysis we use the isobar model amplitude obtained from the most recent Belle $\phi_3$ Dalitz analysis [8].

From the consideration above it is clear that a good approximation to the optimal binning is the one obtained from the uniform division of the strong phase

![Figure 1: Binning of the $D^0 \to K_S\pi^+\pi^-$ Dalitz plot.](image)

and

$$s_i = \frac{1}{\sqrt{f_{D}(x,y)}} \int \text{Im} \ [f_{D}(x,y) f_{D}^*(y,x)] \ dx \ dy,$$  \hfill (5)

where the integral is performed over a single bin. The number of events in the $i$-th bin of the $K_S\pi^+\pi^-$ Dalitz plot from a $B$ decay is then

$$N_i = K_i + r_B^2 K_{-i} - 2r_B \sqrt{K_i K_{-i}} (c_i \cos \theta_\pi - s_i \sin \theta_\pi),$$  \hfill (6)

again up to an overall normalization. It is important to note that $c_i$ and $s_i$ depend only on the $D$ decay. These are the quantities that we measure using CLEO-c data. Although in principle they could be left as free parameters in a $D \to K_S\pi^+\pi^-$ Dalitz plot analysis from $B^\pm$ decays, their values can be more precisely determined from correlated $D_{CP}$ decays produced at CLEO-c.

Thus, we can constrain $\theta_\pi$, and in turn $\gamma/\phi_3$, if we know $K_i$, $c_i$, and $s_i$. The $K_i$ can be easily determined using flavor-tagged $D^0 \to K_S\pi^+\pi^-$ Dalitz plot. In the next section we show how the $c_i$ can be obtained using binned, $CP$-tagged $D^0 \to K_S\pi^+\pi^-$ Dalitz plots.

### 3. Measuring $c_i$ From $CP$-Tagged $D$ Decays

For $D$ mesons that decay into a $CP$ eigenstate, we write the initial state of the $D$ as a linear combination of flavor eigenstates

$$f_{CP\pm}(x,y) = \frac{1}{\sqrt{2}} [f_D(x,y) \pm f_D(y,x)].$$  \hfill (7)
difference $\delta_D$. We thus take the definition of $i$-th bin to be

$$2\pi(i - 1/2)/N \leq \delta_D(x, y) < 2\pi(i + 1/2)/N. \quad (12)$$

An example of such a binning with $N = 8$ is shown in Fig. 2.

![Figure 2: Divisions of the $D^0 \rightarrow K_S\pi^+\pi^-$ Dalitz plot with uniform binning of $\Delta\delta_D$ strong phase difference with $N = 8$.](image)

Figure 2: Divisions of the $D^0 \rightarrow K_S\pi^+\pi^-$ Dalitz plot with uniform binning of $\Delta\delta_D$ strong phase difference with $N = 8$.

4. Event Selection

4.1. Double-Tagged $D \rightarrow K_S\pi^+\pi^-$ Events

This analysis uses a combination of two-body $CP$ and flavor tags. Since the neutral $D$ mesons are produced at $\psi(3770)$ threshold they are correlated in a $C = -1$ state. If mixing is ignored we can determine whether the parent particle was a $D^0$ or $\bar{D}^0$, up to DCS contributions. Similarly, if $CP$ violation is ignored, then the $D$ mesons must be in eigenstates of opposite $CP$ [10].

To determine the flavor of the $D$ meson, we tag $D^0 \rightarrow K_S\pi^+\pi^-$ events with the two-body $D^0 \rightarrow K^+\pi^-$ mode. We use the two $CP$-even tags $K^+K^-$ and $\pi^+\pi^-$, and the two $CP$-odd tags $K_S\pi^0$ and $K_S\eta$.

We introduce two quantities that are reconstructed on both sides of a double-tagged decay. The beam-constrained mass is defined as $M_{bc} = \sqrt{E_b^2 - p_D^2}$, where $E_b$ is the beam energy and $p_D^2$ is the square of the reconstructed 3-momentum of the $D$ meson.

We require that the beam-constrained mass of the reconstructed candidate is within $3\sigma$ of the nominal $D$ mass, which corresponds to a selection criteria of $1.8603 \leq M_{bc} \leq 1.8687$ GeV. The other quantity is the energy difference between the beam and the reconstructed $D$, defined as $\Delta E \equiv E_{beam} - E_D$. We apply a selection criteria of $|\Delta E| \leq 30$ MeV to all $D^0 \rightarrow K_S\pi^+\pi^-$ candidates.

Additional selection criteria are placed on the disgust particles to ensure basic track quality. For example, we select pion track momenta between $0.05 \leq p \leq 2.0$ GeV. Both signal and tagging modes containing a $K_S$ are selected to be within $3\sigma$ of the $K_S$ mass, which corresponds to $\pm 7.5$ MeV from the central $K_S$ mass value of 497.6 MeV.

We only reconstruct $K_S$ particles that decay through the $\pi^+\pi^-$ channel; we do not attempt to reconstruct $K_S \rightarrow \pi^0\pi^0$. Fake $K_S$ candidates can be misreconstructed from combinatoric $\pi^+\pi^-$ pairs. To suppress these events we apply a selection criteria on the flight significance $f_s \geq 0$ to our $K_S$ candidates. Additionally, we require that the $\pi^0$ mass falls within $3\sigma$ of its nominal value.

4.2. Double-Tagged $D \rightarrow K_L\pi^+\pi^-$ Events

For $D^0 \rightarrow K_L\pi^+\pi^-$ decays we require the same selection criteria on charged pions and $\pi^0$ candidates as those described for $D^0 \rightarrow K_S\pi^+\pi^-$ decays. However, because of the large flight distance of the $K_L$, the $K_L\pi^+\pi^-$ signal is reconstructed using a missing mass technique. We require the signal side to have exactly two charged tracks. We also apply $\pi^0$, $\eta$, and $K_S$ vetoes. Using the measured momentum of the tagged $D$, we compute the missing momentum and energy on the signal side. We require that the missing mass squared satisfies the condition $0.21 \leq m^2 \leq 0.29$ GeV$^2$. The background for $D^0 \rightarrow K_L\pi^+\pi^-$ mode is approximately 5%.

4.3. Double-Tagged $K_L\pi^0$ vs. $K_S\pi^+\pi^-$

We can increase our statistics by reconstructing $D^0 \rightarrow K_S\pi^+\pi^-$ events tagged with the $CP$-even mode $K_L\pi^0$. We require zero tracks and exactly one $\pi^0$ candidate on the tag side. We veto events containing $\eta$ candidates, and impose similar criteria on the $K_L$ missing mass as described above.

The final yields for all tag modes are summarized in Table I.

---

1 The inclusion of charge-conjugate modes is implied throughout our analysis.
Table I: Yields for CP-tagged $K_S\pi^+\pi^-$ and $K_L\pi^+\pi^-$ in 398 pb$^{-1}$ data, by tag mode.

<table>
<thead>
<tr>
<th>Tag Mode</th>
<th>$K_S\pi^+\pi^-$</th>
<th>$K_L\pi^+\pi^-$</th>
</tr>
</thead>
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<tr>
<td>$K^+K^-$</td>
<td>61</td>
<td>194</td>
</tr>
<tr>
<td>$\pi^+\pi^-$</td>
<td>33</td>
<td>90</td>
</tr>
<tr>
<td>$K_S\pi^0$</td>
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<td>263</td>
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<td>$K_S\eta$</td>
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<td>21</td>
</tr>
<tr>
<td>$K_L\pi^0$</td>
<td>190</td>
<td>-</td>
</tr>
</tbody>
</table>

5. Combining $K_S\pi^+\pi^-$ and $K_L\pi^+\pi^-$

The tagged $K_L\pi^+\pi^-$ Dalitz plots are included to increase the statistical accuracy of this analysis. However, if we naively combine the Dalitz plots with $K_S$ and $K_L$ we will find our measurement of $c_i$ to be biased. We must first account for the phenomenological differences between the $K_S\pi^+\pi^-$ and $K_L\pi^+\pi^-$ models.

Since the $K_S$ and $K_L$ mesons are of opposite $CP$, the doubly-Cabibbo suppressed amplitudes in each Dalitz plot will contribute with opposite signs. We can see this by inspecting the $D^0$ decay amplitude for each Dalitz plot

$$A(K_S\pi\pi) = \frac{1}{\sqrt{2}} [A(K^0\pi\pi) + A(K^0\pi\pi)]$$

$$A(K_L\pi\pi) = \frac{1}{\sqrt{2}} [A(K^0\pi\pi) - A(K^0\pi\pi)]$$

The effect of this relative minus sign is to introduce a 180$^\circ$ phase for all DCS $K^*$ resonances in the $K_L\pi^+\pi^-$ model. We can use $U$-spin symmetry to relate the amplitudes for resonances of definite $CP$ eigenvalue. We find that these states acquire a factor of $r_K e^{i\delta_K} \simeq -\tan \theta_C$, where $\theta_C$ is the Cabibbo angle.

In our study we multiply all DCS amplitudes in the $K_L\pi^+\pi^-$ model by -1. From this “base” model we fix $r_K = 0.06$ for each $CP$ eigenstate, then vary the phase $\delta_K$ between 0 and $2\pi$. For each bin we then find the largest resulting deviation in $c_i$, and report this value as the systematic uncertainty in the $K_L\pi^+\pi^-$ model.

To better understand the difference between the $K_S\pi^+\pi^-$ and $K_L\pi^+\pi^-$ models, we compare the numerically calculated values of $c_i$ in each Dalitz plot. We find that the value for $c_i$ is systematically larger in each bin for $K_L\pi^+\pi^-$. In Fig. 3 we can see that the difference is significantly larger than the systematic uncertainty in our $K_L\pi^+\pi^-$ model.

6. Results

We report the difference in $c_i$ between $K_S\pi^+\pi^-$ and $K_L\pi^+\pi^-$ as measured in 398 pb$^{-1}$ of data. In Fig. 4 we compare the $c_i$ differences calculated from our model and measured from data. The error bars in this figure represent both statistical and model uncertainty combined. With a reasonable understanding of the $c_i$ between the $K_S\pi^+\pi^-$ and $K_L\pi^+\pi^-$ Dalitz...
plots, we can estimate the final precision with which we expect to measure $c_i$ once $750 \text{ pb}^{-1}$ of data is available. The values of $c_i$ from our study are once again plotted in Fig. 5, but here the error bars represent the statistical uncertainty obtained from $398 \text{ pb}^{-1}$ of data scaled up to $750 \text{ pb}^{-1}$.

![Figure 5: The central values of $c_i$ are computed from our model with expected sensitivity from $750 \text{ pb}^{-1}$ of data. The error bars are determined by scaling the statistical uncertainty obtained from $398 \text{ pb}^{-1}$ of data, then combining the $K_L\pi^+\pi^-$ model uncertainty.](image)

We expect good sensitivity to the measurement of $c_i$ with the entire CLEO-c data. This measurement can reduce the model uncertainty on $\gamma/\phi_3$ to a precision of about $4^\circ$ [3].

**Acknowledgments**

We would like to thank Mats Selen from the University of Illinois, our colleagues David Asner and Paras Naik from Carleton University, and Ed Thorndike from the University of Rochester for helping us prepare for this conference. Also, we would like to thank the organizers of the Charm 2007 Workshop for providing a stimulating environment and a well-organized program of talks.

**References**

Low Mass S-wave $K\pi$ and $\pi\pi$ Systems

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Knowledge of the details of the S-wave $K\pi$ and $\pi\pi$ systems limits the precision of measurements of heavy quark meson properties. This talk covers recent experimental developments in parametrizing and measuring these waves, and examines possibilities for the future.

I. INTRODUCTION

Hadrons were once regarded as a source of complication in the measurement of the dynamics of weak decays of heavy quark states. The realization that their interactions can be used to help in understanding the phases involved is quite a recent development. There are several examples. Interference between hadrons produced in decays of the $D^0$ and $\bar{D}^0$ mesons to the self-conjugate final state $K^0\pi^+\pi^-\, [31]$ has been used to measure $D^0-\bar{D}^0$ mixing parameters. [1, 2] It has also allowed the measurement of the CKM phase $\gamma$ from the decays $B^+ \rightarrow D^0K^+$ and $B^+ \rightarrow \bar{D}^0K^+$. [3–6] Interference between S- and P-waves in the $K^+\pi^-$ systems produced in $B^0 \rightarrow J/\psi K^+\pi^-$ decays has been used to resolve the ambiguity in the sign of $\cos2\beta$, where $\beta$ is the CKM phase involved in these decays.

Central to all these analyses is the need to assign a partial wave composition at each point in the hadron phase space. Currently, models for the hadron interactions based on resonant composition are mostly used to do this (the “isobar model”). Sometimes a $K$-matrix description of the S-wave component is included. Uncertainties and ambiguities or questionable assumptions, especially with respect to the S-wave behaviour lead to significant systematic uncertainties in the results. As the precision of the measurements improves, it becomes a more pressing goal either to develop a good, model-independent strategy for describing the hadron amplitudes or, at least, a better understanding of how to describe the S-waves in the decays of heavy quark mesons.

It is also of intrinsic interest to study the low mass S-wave $K\pi$ and $\pi\pi$ systems to better understand the scalar states, labelled here as $\kappa(800)$ or $\sigma(500)$, that may exist [7]. Measurements of these amplitudes in this region are sparse. There is also merit, therefore, in pursuing the possibility of using heavy quark meson decays to add experimental data in these regions. This talk focusses on recent progress and future prospects on both fronts.

In the first section, information presently available on both $K\pi$ and $\pi\pi$ S-waves is briefly reviewed. The expected relationship between scattering amplitudes and those measured in decays of heavy quark mesons is then examined. The remainder of the talk discusses ways in which experimental observations, mostly of the available $K\pi$ data, are being studied.

II. S-WAVE SCATTERING DATA

Excellent experimental information comes from model-independent analyses of differential cross sections for reactions in which production of $K\pi$ or $\pi\pi$ systems is dominated by a pion exchange mechanism. The amplitudes so determined show clear evidence for the $K\pi^0(1430)$ resonance in the iso-spin $I = 1/2$ $K\pi$ S-wave and for the $f_0(980)$ in the $I = 0$ $\pi\pi$ S-wave. No exotic states are found in the $I = 3/2$ $K\pi$ or $I = 2$ $\pi\pi$ waves. Data at the very low mass regions, where $\kappa(800)$ or $\sigma(500)$ poles may lurk far from the real axis, are relatively poor.

A. The $K\pi$ System

The SLAC E135 (LASS) experiment [8, 9] provides the best information on the $K^-\pi^+$ system. The data are shown in Fig. 1. Note that there are no data below 825 MeV/c$^2$ from this experiment, though some is available from an earlier era of experiments. The LASS collaboration determined that both $I = 1/2$ and $I = 3/2$ amplitudes $T(s)$ are unitary up to $K\eta'$ threshold ($\sqrt{s} = 1454$ MeV/c$^2$)

$$T(s) = \sin(\delta(s))e^{i\delta(s)}$$

where $s$ is the squared invariant mass and $\delta(s)$ is the phase.

The $I = 3/2$ phase, assumed to be unitary, was measured by Estabrooks, etal [10] and is shown in Fig. 1(c) and (d). The LASS collaboration used this information to separate the $I = 1/2$ and $I = 3/2$ components of the amplitude seen in Fig. 1(a) and (b). The fit to their data, shown as the solid curves, are described by Eq. (1) with

$$I = 1/2 : \quad \delta^{1/2}(s) = \delta^{1/2}_R(s) + \delta^{1/2}_I(s)$$
$$\cot\delta^{1/2}_R(s) = (s_0 - s)/(\sqrt{s}\Gamma_0)$$
$$q\cot\delta^{1/2}_I(s) = 1/a_{1/2} + b_{1/2}q^2$$

$$I = 3/2 : \quad q\cot\delta^{3/2}(s) = 1/a_{3/2} + b_{3/2}q^2$$

where $q$ is the momentum of the $K^-$ in the $K\pi$ center of mass. This parametrization is valid in the range

FIG. 1: $S$-wave $K^-\pi^+$ scattering amplitude extracted by the LASS collaboration from their data on the reaction $K^-p \rightarrow K^-\pi^+n$ at 11 GeV/c. The phase is shown in (a) and the magnitude in (b). The $I$-spin 1/2 component is represented by points on a solid curve and $I$-spin 3/2 is shown with no curve. $I = 3/2$ phases measured in data from Ref. [10] are plotted as a function of $K^+\pi^-$ invariant mass of (c) $K^+p \rightarrow K^+\pi^+n$ and (d) $K^-p \rightarrow K^-\pi^-\Delta^+$. 

825 < $\sqrt{s} < 1454$ MeV/c$^2$, and includes one $I = 1/2$ resonance of mass $\sqrt{s_0} \sim 1345$ MeV/c$^2$ and width $\Gamma_0 \sim 275$ MeV/c$^2$. Non-resonant backgrounds in both waves are described by scattering lengths $a$ and effective ranges $b$.

B. The $\pi\pi$ System

Phases for the $I = 0$ component, shown in Fig. 2(a), have been extracted in several analyses of $\pi^-p \rightarrow \pi^-\pi^+n$ data from Grayer, et al [11] It is clear that data below 600 MeV/c$^2$ in the region of the $\sigma(500)$ from these data are poor.

The $I = 2$ amplitude, assumed to be unitary up to $pp$ threshold ($\sim 1500$ MeV/c$^2$), was derived from data on $\pi^-p \rightarrow \pi^+\pi^-n$ interactions at 12.5 GeV/c [12] and $\pi^-d \rightarrow \pi^-pp$ spectator at 9 GeV/c [13]. It was fit to the form shown in Fig. 2(b) [14]. For $I = 0$, the amplitude is unitary up to $KK$ threshold where its elasticity drops suddenly. Slightly below this, the phase rises rapidly indicating the presence of the $f_0(980)$ resonance.

III. ROLE OF SCATTERING IN D DECAYS

In terms of our two goals, it would be useful if we could draw on these measurements to help in reducing ambiguities in models used in analyzing decays of heavy quark mesons. At the same time, we need to learn how to interpret such decays in learning more about the scattering amplitudes at small $s$ values.

Consider the decay $D \rightarrow (AB)C$ in which a two hadron system $f = AB$ and another system $C$ are produced. A simple assumption about the decay amplitude $F(s)$ for such a process is that it can be factorized into short and long-range effects.

$$F_j(s) = T_{jk}(s)Q_k(s)$$

Here $Q_k(s)$ describes the short-range decay of $D$ to $C$ and an intermediate hadron system $k$, and $T_{jk}(s)$ describes the subsequent re-scattering within the system $k$ to produce the final state $f$. We take $Q_k(s)$ to encode not only the relatively real ratio of decay modes for the weak decay of the $D$ that would imply an $s$-dependence of its magnitude, but also any other short-range effects that may occur in the $Ck$ system which might also impart an $s$-dependence of its phase.

If $s$ is below the threshold where $AB$ scattering becomes inelastic ($K\eta'$ for $L = \text{even waves}$ when $AB = K\pi$ and $K^-K^-$ when $AB = \pi\pi$ systems, for example), then $T_{jk}(s)$ is simply equal to the $AB$ elastic scattering amplitude. If the phase of $Q_k(s)$ is independent of $s$, then the phase of $F(s)$ will have the same $s$-dependence as $T(s)$, i.e., that observed in elastic scattering. This is a statement of the (more rigorously derived) Watson theorem [15]. If partial wave expansions of $F(s)$ and $T(s)$ are made, then this condition must hold for each partial wave. It must hold, in particular, for the $S$-waves.

The Watson theorem could provide a very useful constraint, therefore, in the analysis of heavy quark states, so its range of applicability is of great importance. The considerations above lead us to expect the following:

- If $C$ is a di-lepton ($\ell\nu$, for example) then $Ck$ scattering would be unlikely and the phase of $Q_k$ would have little, if any, $s$-dependence.

- The same might also be true if $C$ were a massive hadron (such as a $J/\psi$) with a small interaction radius.
FIG. 2: (a) $I=0$ phase of the $S$-wave $\pi\pi$ scattering amplitude extracted from data for the reactions $\pi^- p \rightarrow \pi^- \pi^+ n$ from Ref. [11]. The phase in degrees is plotted as a function of $\pi\pi$ invariant mass $M$.

(b) The $I=2$ phase shifts plotted as a function of $\pi\pi$ invariant mass $M$. The $I=2$ amplitude is assumed to be unitary up to 1500 MeV/c and is obtained from data for the reactions $\pi^- p \rightarrow \pi^+ \pi^+ n$ (12.5 GeV/c) [12] and $\pi^+ d \rightarrow \pi^- \pi^- n p$ spectator (9 GeV/c) [13].

- In cases where $C$ is a hadron with mass comparable to $A$ or $B$, significant scattering in the $Ck$ system would be likely. This could lead to a dependence of both the magnitude and the phase of $Q_k(s)$ upon $s$.

- The likelihood of $Ck$ scattering probably increases at shorter range.

We should expect, therefore, to find any deviations from the Watson theorem to be most pronounced when the mass of $C$ is comparable with that of $A$ or $B$, and at small $s$ in the $AB$ system. We also note that, the Watson theorem applies to every partial wave, and it makes little sense to impose it only on the $S$-wave alone.

IV. $K\pi$ PRODUCTION FROM HEAVY QUARK MESON DECAYS

A. Decays with Leptons or Massive Hadrons

In a study of semi-leptonic decays of $D$ mesons

$$D^0 \rightarrow K^- \pi^+ \ell\nu. \quad (4)$$

the FOCUS collaboration [16] observed asymmetry in the distribution of $\cos \theta$, where $\theta$ is the angle between $K^-$ and the $\ell\nu$ system in the $K^-\pi^+$ rest frame. The asymmetry results from interference between $S$- and $P$-waves, and is proportional to the cosine of the relative phase $\gamma$ between them. As seen in Fig. 3 this follows the behaviour observed in the LASS data quite closely as the $K^-\pi^+$ invariant mass moves through the $K^*(890)$ region. This conforms to the first of our expectations.

FIG. 3: Asymmetry in cosine of $K^-\pi^+$ helicity angle in $D^0 \rightarrow K^- \pi^+ \ell\nu$ data from the FOCUS collaboration [16] as a function of invariant mass for the $K^-\pi^+$ system. The behaviour observed in the LASS data is indicated by the solid curve.

The BABAR collaboration observed [17] similar behaviour in the asymmetry of the helicity distribution for $K^+\pi^-$ produced in

$$B^0 \rightarrow J/\psi K^+\pi^- \quad (5)$$

decays. Fig. 4 shows the two solutions for the $S$-$P$ phase difference $\gamma$ in the $K^*(890)$ invariant mass region. The relative phase is dominated by the rapid variation in the $P$-wave due to the $K^*(890)$ and only one solution is physically meaningful. It is seen in the figure that this matches the LASS phase variation (though shifted by $+\pi$ radians) very well.

This conforms to the second of our expectations, but it is not obvious why there is an overall phase shift of $+\pi$ radians.
interfering resonances like the Table II. at this invariant mass is ∼ − LASS data [8]. This shows that the phase of Q φ K

FIG. 4: Two solutions (open and closed circles with error bars) for the relative phase γ = δ_S − δ_P between S- and P-wave amplitudes of K⁺π⁻ systems produced in B⁰ → J/ψK⁺π⁻ decays from the BABAR collaboration [17]. Values for γ are plotted as a function of K⁺π⁻ invariant mass. The solid circles are for the only physical solution and follow the LASS data points, plotted as triangles, closely.

B. Decays to Light Hadrons

The most detailed experimental information comes from studies of D⁺ → K⁻π⁺π⁺ decays. These Cabibbo favoured decays are known to contain a large S-wave component.

A study of ∼ 15,000 such decays by the E791 collaboration [18] provides an illustration. The E791 Dalitz plot is shown in Fig. 5 where significant S-P interference is evident from the asymmetry of the K⁺(890) bands. This is plotted in Fig. 5 vs. the Breit-Wigner phase φ_{BW} = \tan^{-1} M_0 \Gamma(M_0/M_0^2 - s), where M_0 is the mass and \Gamma(\sqrt{s}) the invariant mass-dependent width of the K⁺(890). The asymmetry is zero when φ_{BW} = 56°, almost 80° below the phase found in the LASS data [8]. This shows that the phase of Q_K(γ) at this invariant mass is ∼ −80° and is entered in Table II.

1. Isobar Model Fits

In the earliest analyses of these decays, a model with interfering resonances like the K⁺(890), the L = 0 K⁺(1430) and a constant non-resonant “NR” 3-body amplitude could account for the voids and asymmetries observed in the Dalitz plot. With their larger sample, the E791 isobar model analysis showed that additional structure in the S-wave was required to achieve an acceptable fit, and the addition of a k(800) Breit-Wigner isobar, which interfered destructively with the NR term, worked well.

The “isobar model” description of the L = 0, 1 and 2 wave amplitudes F_L in the K⁺π⁺ systems for this fit can be summarized as:

\[ F_0(s) = c_{00} + \alpha_{10} BW_{K^*_0(1430)}(s) + \alpha_{20} BW_{\kappa^0(800)}(s) \]
\[ F_1(s) = \alpha_{11} BW_{K^+_{5/2}(800)}(s) + \alpha_{21} BW_{K^*_{1/2}(1688)}(s) \]
\[ F_2(s) = \alpha_{21} BW_{K^*_{3/2}(140)} \]

where the BW(s) are relativistic Breit-Wigner functions with s-dependent widths, and the \alpha_L are complex coefficients determined in the fit. The overall phase was defined by setting \alpha_{11} = 1.0, and \alpha_{00} was the NR term.

Two further isobar model analyses of this decay mode were recently made, one by FOCUS [19] and the other by CLEOc [20]. Each used samples ∼ 3.5 times larger. The conclusions, and estimates of the resonant fractions [32], of both were in good general agreement with E791.

Parameters for \kappa and K⁺(1430) S-wave Breit-Wigner isobars are compared for the three experiments in Table IV B1. The K⁺(1430) parameters in this model disagree significantly with those obtained by the LASS experiment or with the World average [21]. There is general agreement that this description of the S-wave, described by Eq. (6), with two broad, Breit Wigner resonances, one of which is also near threshold, is theoretically problematic and could
account for this discrepancy. It would be virtually certain that this amplitude would have an s-dependent phase that would differ from the Watson theorem expectation.

TABLE I: Breit Wigner Parameters for the $K^−\pi^+$ S-wave Isobar States. All quantities are in MeV/c².

<table>
<thead>
<tr>
<th></th>
<th>E791</th>
<th>Focus</th>
<th>CLEO c</th>
<th>PDG</th>
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<td>$\Gamma_0$</td>
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<tr>
<td>$K^*_0$</td>
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<td>177 $\pm$ 8</td>
<td>169 $\pm$ 5</td>
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</tr>
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</table>

2. Model-Independent Measurement

A test of the Watson theorem requires a measurement of the phase of $F(s)$ at several values of $s$ in a model-independent way. The first attempt to do this for the $S$-wave for $K^−\pi^+$ produced in this decay mode was made by the E791 collaboration [22].

They replaced the analytical function describing the $S$-wave in Eq. (6) by a set of 38 complex values at discrete values for $s$, using a spline interpolation for other values. The $P$- and $D$-waves were parameterized, as before, by the form in Eqs. (7) and (8) and the coefficients $\alpha_{iL}$ were allowed to float. A fit was then made to determine the best values (magnitude and phase) for each of the 38 $S$-wave points. The result, shown in Fig. 6, is compared with the LASS model for the $I=1/2$ $K^−\pi^+$ system in Eq. (2). Agreement is good in the $K^*_0(1430)$ region (above $\sim$1100 MeV/c²) after a shift in phase of $75^\circ$ and an arbitrary scale factor are applied to the LASS amplitude. A very significant discrepancy is, however, seen for lower values of invariant mass.

This might conform to our 4th expectation, but there are two problems. First, though $I=1/2$ $K^−\pi^+$ production probably dominates, $I=3/2$ production in these decays cannot be excluded. Second, the isobar model form in Eqs. (7) and (8) for the $P$- and $D$-waves, upon which this result depends, is questionable. The $P$-wave contains more than one Breit-Wigner, and both waves are assumed to be dominated by resonant behaviour. More importantly, neither wave is likely to follow the Watson theorem, so a test of the $S$-wave alone cannot be conclusive.

The CLEO collaboration [20] has attempted to overcome the latter difficulty. Using their high purity sample of $\sim$ 60,000 events, they proceeded in the same way as E791, interpolating the $S$-wave between discrete values of $s$, while parametrizing the $P$- and $D$-waves as above. Their results were very similar. They then fixed the $S$-wave and, using a similar procedure, fit the $P$-wave parametrized in the same way. Then they repeated this for the $D$-wave. In each step, only one wave was allowed to float with the others fixed.

This procedure can converge only by simultaneously floating all waves at once, and this was not done. Also, the phases have to be defined in some part of the phase space (as little as possible) in order for this to work.

It seems that a truly model-independent measurement of the $S$-wave decay amplitudes in these hadron decays is desirable, but has yet to be made, and will probably need to await the much larger data samples of the $B$ factories, BES or the Panda experiment.

3. I-spin Analysis by the FOCUS Collaboration

Using a sample of about 50,000 events with 96% purity, the FOCUS collaboration used the $K$-matrix method to separate $I=1/2$ and $I=3/2$ production for this decay [19].

The $P$- and $D$-waves were described as in Eqs. (7) and (8), however, the $S$-wave amplitude $F_0(s)$ in Eq. (6) was replaced by the sum of two terms, one for each $I$-spin. Each was described by

$$F_f(s) = (I - i\rho K(s))^{-1}_{jk}P_k(s)$$

$$T_{jk}(s) = (I - i\rho K(s))^{-1}_{ki}K_{ij}(s)$$

$$Q_f(s) = K^{-1}_f(s)P_k(s).$$

Here, the amplitudes $F$, $T$, $P$ and $Q$ were introduced for each $I$-spin and their indices $K$ and $f$ labelled intermediate and final states ($I = K\pi$, $2=\eta K$). The production vector $P_k(s)$ describing the couplings of the $D$ decay to the two intermediate states was de-
scribed by parameters obtained from the fit, and $\rho_{bf}$ was the phase space matrix for the two channels.

Real values for all elements of the $K$ matrices (2 × 2 for $I = 1/2$ and $1 \times 1$ for $I = 3/2$) were determined by a fit to the LASS measurements of $T$.

The fit required a very large contribution from the $I = 3/2$ $S$-wave (40.50%) interfering destructively with $I = 1/2$ (207.25%) giving a total $S$-wave (83.23%).

The phases of the resulting decay amplitudes, $F(s)$, are shown in Fig. 7, with the $K$-matrix fit to the LASS elastic scattering data for $I = 1/2$ superimposed. The phase of the total amplitude $F(s)$, similar to that found by E791 and by CLEOc, differs significantly from the LASS $I = 1/2$ data. The $I = 1/2$ component is shown alone in Fig. 7(b) and shows better agreement.

The parametrization of the production vector was chosen specifically to allow an s-dependence in its phase, thereby allowing a deviation from the Watson theorem. Some deviations are, indeed, evident in Fig. 7(b), as the invariant mass approaches the $K^*_0(1430)$ region. It is possible that this is due to the $K^*_0(1430)$ pole and the onset of effects from the $K\eta'$ channel. However, another interpretation may be possible. The phase shown in this figure is shifted by an arbitrary amount to achieve good agreement at the lower invariant masses. Were a different shift in phase applied, agreement would, in fact, be good at the high invariant mass end and poor at low mass, consistent with our expectation number 4.

This analysis attempted to make a valid comparison between scattering and decay. It appears premature, however, to conclude that the Watson theorem works for these hadronic decays especially because the same is not required for the $P$- nor $D$-waves in this analysis.

V. $\pi\pi$ PRODUCTION FROM HEAVY QUARK MESON DECAYS

Decays of $D^+ \rightarrow \pi^-\pi^+\pi^+$ result in $S$-wave enriched $\pi^+\pi^-$ systems. Unfortunately, these decays are Cabibbo suppressed. Another decay rich in $S$-wave content is $D^0 \rightarrow K^0\pi^+\pi^-$. This is partly Cabibbo favoured and partly doubly suppressed, so is somewhat complex with many resonances contributing. Consequently, model-independent analyses of these decays have yet to be attempted.

A. $D^+ \rightarrow \pi^-\pi^+\pi^+$ Decays

The largest sample yet studied comes from the CLEO experiment [23]. It consists of only about 4,000 events with a purity of $\sim 95\%$. It has been used to test a variety of models.

An isobar model fit to the Dalitz plot confirms, as does a similar analysis by FOCUS [24], the E791 collaboration conclusion [25] that structure in the low mass $\pi^+\pi^-$ system is well described by a scalar $\sigma(500)$ Breit Wigner destructively interfering with a constant $NR$ amplitude. The fit obtained was of marginal quality.

CLEO tried several variations on the isobar model. They included an $I = 2$ $\pi^+\pi^+$ $S$-wave contribution, slightly improving the fit quality. They also tried variations of $\pi\pi$ $S$-wave isobar model. In one, as suggested in Ref. [26], they replaced the $\sigma$ Breit-Wigner by a simple pole

$$1/(m_0^2 - s - im\Gamma) \rightarrow 1/(s_0 - s)$$

where $s_0 = (0.47 - 0.22i)$. Other $S$-wave models used were an amplitude based on the linear sigma model [27] and an amplitude derived by N.Achamov described in the CLEO paper Ref. [23]. All models provided fits to the data slightly more acceptable than the isobar model, but no clear distinction was obvious.

B. $D^0 \rightarrow K^0\pi^+\pi^-$ Decays

Isobar model fits to large samples have been made by BABAR [4] and Belle [5] collaborations. They each require two $\sigma$ states in the $\pi^+\pi^-$ $S$-wave, one similar to that found in isobar fits to the previous channel by E791, FOCUS and CLEOc, and the other at a mass of $\sim 1.0$ GeV/c$^2$. Both collaborations have also used $K$-matrix parametrizations of the $S$-wave $\pi^+\pi^-$ amplitude, obtaining slightly better fits. These fits involve no assumptions about $\sigma$ states, and they do enforce the Watson theorem in this system. However, the other waves are those defined by the isobar model, with no such constraint. None of the fits so far published have acceptable quality.

Further progress in this system may come from the larger data samples form the $B$ factories, or from the next generation of charm factories.

C. Other Channels

A number of $B$ and $D$ decays in which $K\pi$ systems are produced have now been published. A comparison between the general characteristics of the $S$-wave amplitudes observed and those of the LASS scattering data are summarized in Table II [9]. The two examples above are included and represent the best examples of the validity of the Watson theorem. Hopefully, an underlying pattern for these characteristics will eventually emerge, provided that such decays are studied with this goal in mind. Any pattern is certainly not yet evident.
FIG. 7: Decay amplitudes for $K^-\pi^+$ systems from $D^+ \to K^-\pi^+\pi^+$ decays from Ref. [19] for (a) the total $S$-wave and (b) the $I = 1/2$ contribution. The black bands correspond to one standard deviation limits about the central fit in each of the amplitudes. Scattering amplitudes from a $K$-matrix fit to the LASS data [8] are shown as red continuous lines.

TABLE II: Qualitative comparison between $K\pi$ $S$-wave decay amplitude $F_0(s)$ and the LASS Model in a variety of cases studied. Three characteristics are compared. The $S$-$P$ relative phase at the $K^*(890)$ mass. The observed value minus that observed in LASS data, $\Delta \phi_{SP}$, is tabulated to the nearest 15°. The general shape of $|F_0(s)|$ for $s$ values both below and above $s_0 = 1.0$ GeV/c$^2$ are also tabulated.

| Decay Process | $\phi_{SP}$ approx | $|F_0|\ s < s_0$ | $|F_0|\ s > s_0$ |
|---------------|---------------------|-----------------|-----------------|
| $B^0 \to J/\psi K^+\pi^-$ [17] | +180 | Poorly defined | to LASS |
| $B^+ \to K^+\pi^-\pi^+$ [28] | 0 | Unknown | Similar to LASS |
| $B^+ \to K^+\pi^-\rho$ [29] | +180 | Unknown | Unknown |
| $D^0 \to K^-K^+\pi^0$ [30] | −90 | Similar to LASS | Similar |
| $D^+ \to K^-\pi^+\pi^+$ [22] | −75 | Very different | Similar to LASS |
| $D_s^+ \to K^-K^+\pi^+$ | −90 | Similar to LASS | Similar |
| $D^+ \to K^-\pi^+\ell\nu$ [16] | 0 | Similar to LASS | Similar |

VI. CONCLUSIONS

The most reliable data on $S$-wave amplitudes are still those from LASS or CERN-Munich data on elastic scattering. Data at the lowest energies, however, are somewhat sparse. It would surely help in the understanding of the pole structure relevant to the firm establishment of existence or otherwise of light scalar $\kappa(800)$ or $\sigma(500)$ states to improve on this situation. It is difficult, at present, to see how hadronic decays of $D$ or $B$ mesons will help since $I$-spin considerations, and uncertainties in the $D$ form-factor, make such measurements difficult. It appears, therefore, that such information is most likely to come from high statistics studies of $D$ semi-leptonic decays, or decays of $B$ mesons to $J/\psi K^+\pi^-$ (or $J/\psi\pi^+\pi^-$) from BABAR or Belle or the new facilities in BES or Darmstadt.

Progress in understanding ways to parametrize the $S$-wave decay amplitudes in Dalitz plot analyses is slow, but the body of information is growing. The prize could be less systematic uncertainty resulting from model dependence of such analyses in measurements of important phenomena such as $D^0\bar{D}^0$ mixing or of the CKM angle $\gamma$.

A positive trend is that more techniques beyond the isobar model, with its known problems and strong model-dependence, are being developed.

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[30] Charge conjugate states are implied unless specified otherwise throughout this talk.
[31] The fraction for each resonance or NR sub process is defined in ref. [18].
K-matrix and Dalitz plot analysis from FOCUS

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1. Introduction

Over the last years we have seen a resurrection of Dalitz plot analyses in modern Heavy Flavor experiments. This analysis tool, first applied in the charm sector, has more recently become a standard technique, used for sophisticated studies and searches for new physics in the beauty sector. Paradigmatic examples are $B \rightarrow \rho \pi$ and $B \rightarrow D^{(*)}K^{(*)}$ for the extraction of the $\alpha$ and $\gamma$ angles of the Unitarity Triangle. Indeed, the road to go from the detected final states to the intermediate resonances can be rather insidious and complications arise in both decays. More precisely, the extraction of $\alpha$ in $B \rightarrow \rho \pi$ means, operationally, selecting and filtering the desired intermediate states among all the possible $(\pi \pi)$ combinations, e.g. $\sigma \pi$, $f_0(980)\pi$ etc. The extraction of $\gamma$ in $B \rightarrow D^{(*)}K^{(*)}$ requires, in turn, modeling the $D$ amplitudes. This poses the problem of how to deal with strong-dynamics effects, in particular those regarding the scalar mesons. The $\pi \pi$ and $K \pi$ S-wave are characterized by broad, overlapping states: unitarity is not explicitly guaranteed by a simple sum of Breit–Wigner functions. In addition, independently of the nature of the $\sigma$, it is not a simple Breit–Wigner. The $f_0(980)$ is a Flatté-like function, and its linehape parametrization needs a precise determination of $KK$ and $\pi \pi$ couplings. Recent analyses of CP violation in the $B \rightarrow D K$ channel from the beauty factories have used the Cabibbo-favored mode $K_s\pi^+\pi^-$, which is common to both $D^0$ and $\bar{D}^0$. A set of 16 two-body resonances had to be introduced to describe the $(K\pi)\pi$ and $K_s(\pi\pi)$ states in the $D^0$ amplitude: two ad hoc resonances were required to reproduce the excess of events in the $\pi \pi$ spectrum, one at the low-mass threshold, the other at 1.1 GeV. Masses and widths of the two states, named $\sigma_1$ and $\sigma_2$, were fitted to the data themselves and found to be $M_{\sigma_1} = 484 \pm 9$ MeV, $\Gamma_{\sigma_1} = 383 \pm 14$ and $M_{\sigma_2} = 1014 \pm 7$ MeV, $\Gamma_{\sigma_2} = 88 \pm 13$ in BaBar [1] and $M_{\sigma_1} = 519 \pm 6$ MeV, $\Gamma_{\sigma_1} = 454 \pm 12$ $M_{\sigma_2} = 1050 \pm 8$ MeV, $\Gamma_{\sigma_2} = 101 \pm 7$ in Belle [2]. These scalars were invoked with no reference to those found in other processes, in particular scattering data, and with no assumption as to the correctness of the physics the model embodies. This procedure of “effectively” fitting data invites a word of caution on estimating the systematics of these measurements. A question then naturally arises: in the era of precise measurements, do we know sufficiently well how to deal with strong-dynamics effects in the analyses?

We have faced parametrization problems in the FOCUS experiment and learnt that many difficulties are already known and studied in different fields of physics, such as nuclear and intermediate-energy physics, where broad, multi-channel, overlapping resonances are treated in the K-matrix formalism. The effort we have had to make mainly consisted in building a bridge of knowledge and language to reach the high-energy community; our pioneering work in the charm sector might inspire future accurate studies in the beauty sector.

2. The K-matrix and P-vector formalism

A formalism for studying overlapping and many-channel resonances was proposed long ago and is based on the K-matrix [3, 4] parametrization. The K-matrix formalism provides a direct way of imposing the two-body unitarity constraint, which is not explicitly guaranteed in the simple sum of Breit–Wigners, here referred to as the isobar model. Minor unitarity violations are expected for narrow, isolated resonances but more severe ones exist for broad, overlapping states. This is the real advantage of the K-matrix approach: it heavily simplifies the formalization of any scattering problem since the unitarity of the S matrix is automatically encoded.

Originating in the context of two-body scattering, the formalism can be generalized to cover the case of production of resonances in more complex reactions [5], with the assumption that the two-body system in the final state is an isolated one and that the two particles do not simultaneously interact with the rest of the final state in the production process [4]. The validity of the assumed quasi two-body nature of the process of the K-matrix approach can only be verified by a direct comparison of the model predictions with data. In particular, the failure to reproduce the Dalitz
plot distribution could be an indication of the presence of relevant, neglected three-body effects.

3. The FOCUS results

3.1. The three pion analysis

The FOCUS collaboration has implemented the K-matrix approach in the $D_s$ and $D^+ \rightarrow \pi^+\pi^-\pi^+$ analyses. It was the first application of this formalism in the charm sector. Results and details can be found in [6]. Here I only reproduce plots of the final results. In Fig. 1 and Fig. 2 the Dalitz-plot projections are shown for $D_s$ and $D^+$ into three pions.

![Figure 1: FOCUS $D_s$ Dalitz-plot projections with fit results superimposed. The background shape under the signal is also shown.](image1)

![Figure 2: FOCUS $D^+$ Dalitz-plot projections with fit results superimposed. The background shape under the signal is also shown.](image2)

In Fig. 3 the FOCUS adaptive binning schemes for $D_s$ and $D^+$ are plotted. In this model [5], the production process, i.e., the D decay, can be viewed as consisting of an initial preparation of states, described by the P-vector, which then propagates according to $(I - iK\rho)^{-1}$ into the final one. The K-matrix here is the scattering matrix and is used as fixed input in our analysis. Its form was inferred by the global fit to a rich set of data performed in [7]. It is interesting to note that this formalism, beside restoring the proper dynamical features of the resonances, allows for the inclusion in D decays of the knowledge coming from scattering experiments, i.e., an enormous amount of results and science. No re-tuning of the K-matrix parameters was needed. The confidence levels of the final fits are 3.0% and 7.7% for the $D_s$ and $D^+$ respectively. The results were extremely encouraging since the same K-matrix description gave a coherent picture of both two-body scattering measurements in light-quark experiments as well as charm-meson decay. This result was not obvious beforehand. Furthermore, the same model was able to reproduce features of the $D^+ \rightarrow \pi^+\pi^-\pi^+$ Dalitz plot that would otherwise require an ad hoc $\sigma$ resonance. The better treatment of the S-wave contribution provided by the K-matrix model was able reproduce the low-mass $\pi^+\pi^-$ structure of the $D^+$ Dalitz plot. This suggests that any $\sigma$-like object in the D decay should be consistent with the same $\sigma$-like object measured in $\pi^+\pi^-$ scattering.

![Figure 3: $D_s$ and $D^+$ adaptive binning Dalitz-plots for the three pion FOCUS K-matrix fit.](image3)

Further considerations and conclusions from the FOCUS three-pion analysis were limited by the sample statistics, i.e. 1475 ± 50 and 1527 ± 51 events for $D_s$ and $D^+$ respectively. We considered mandatory to test the formalism at higher statistics. This was accomplished by the $D^+ \rightarrow K^-\pi^+\pi^+$ analysis.

3.2. The $D^+ \rightarrow K^-\pi^+\pi^+$

The recent FOCUS study of the $D^+ \rightarrow K^-\pi^+\pi^+$ channel uses 53653 Dalitz-plot events with a signal fraction of ~ 97%, and represents the highest statis-
tistics, most complete Dalitz plot analysis for this channel. Invariant mass and Dalitz plot are shown in Fig.4.

Figure 4: The $D^+ \rightarrow K^-\pi^+\pi^+$ Dalitz plot (left) and mass distribution (right): signal and sideband regions are indicated in red and blue respectively. The sidebands are at ±(6-8) $\sigma$ from the peak.

Details of the analysis are in [8].

An additional complication in the $K\pi$ system comes from the presence in the $S$-wave of the two isospin states, $I = 1/2$ and $I = 3/2$. Although only the $I = 1/2$ is dominated by resonances, both isospin components are involved in the decay of the $D^+$ meson into $K^-\pi^+\pi^+$. A model for the decay amplitudes of the two isospin states can be constructed from the $2 \times 2$ $K$-matrix describing the $I = 1/2$ $S$-wave scattering in $(K\pi)_1$ and $(K\eta')_2$ (with the subscripts 1 and 2, respectively, labelling these two channels), and the single-channel $K$-matrix describing the $I = 3/2$ $K^-\pi^+\pi^+$ scattering.

The $K$-matrix form we use as input describes the $S$-wave $K^-\pi^+\pi^+$ scattering from the LASS experiment [9] for energy above 825 MeV and $K^-\pi^-\rightarrow K^-\pi^+\pi^-$ scattering from Estabrooks et al. [10]. The $K$-matrix form follows the extrapolation to the $K\pi$ threshold for both $I = 1/2$ and $I = 3/2$ $S$-wave components by the dispersive analysis by Büttiker et al. [11], consistent with Chiral Perturbation Theory [12]. The complete form is given below in Eqs. (4-5) with the parameters listed in Table I [13].

The total $D$-decay amplitude can be written as

$$\mathcal{M} = (F_{1/2})_1(s) + F_{3/2}(s) + \sum_j a_j e^{i\delta_j} B(abc|\rho),$$  

where $s = M^2(K\pi)$, $(F_{1/2})_1$ and $F_{3/2}$ represent the $I = 1/2$ and $I = 3/2$ decay amplitudes in the $K\pi$ channel, $j$ runs over vector and spin-2 tensor resonances, and $B(abc|\rho)$ are Breit–Wigner forms. The $J > 0$ resonances should, in principle, be treated in the same $K$-matrix formalism. However, the contribution from the vector wave comes mainly from the $K^*(892)$ state, which is well separated from the higher mass $K^*(1410)$ and $K^*(1680)$, and the contribution from the spin-2 wave comes from $K_2^*(1430)$ alone. Their contributions are limited to small percentages, and, as a first approximation, they can be reasonably described by a simple sum of Breit–Wigners. More precise results would require a better treatment of the overlapping $K^*(1410)$ and $K^*(1680)$ resonances as well. In accord with SU(3) expectations, the coupling of the $K\pi$ system to $K\eta$ is supposed to be suppressed. Indeed we find little evidence that it is required. Thus $F_{1/2}$ is actually a vector consisting of two components: the first accounting for the description of the $K\pi$ channel, the second of the $K\eta'$ channel: in fitting $D^+ \rightarrow K^-\pi^+\pi^+$ we need, of course, the $(F_{1/2})_1$ element. Its form is

$$(F_{1/2})_1 = (I - iK_{1/2}\rho)^{-1}_1(P_{1/2})_j,$$  

where $I$ is the identity matrix, $K_{1/2}$ is the $K$-matrix for the $I = 1/2$ $S$-wave scattering in $K\pi$ and $K\eta'$, $\rho$ is the corresponding phase-space matrix for the two channels [4] and $(P_{1/2})_j$ is the production vector in the channel $j$.

The form for $F_{3/2}$ is

$$F_{3/2} = (I - iK_{3/2}\rho)^{-1}P_{3/2},$$  

where $K_{3/2}$ is the single-channel scalar function describing the $I = 3/2$ $K^-\pi^+\rightarrow K^-\pi^+$ scattering, and $P_{3/2}$ is the production function into $K\pi$.

Fitting of the real and imaginary parts of the $K^-\pi^+\rightarrow K^-\pi^+$ LASS amplitude, shown in Fig. 5, and using the predictions of Chiral Perturbation Theory to continue this to threshold, gives the $K$-matrix parameters in Table I.

Figure 5: Real and imaginary $K^-\pi^+\rightarrow K^-\pi^+$ amplitudes from the LASS experiment and their $K$-matrix fit results.

The $I = 1/2$ $K$-matrix is a single-pole, two-channel matrix whose elements are given in Eq. (4).

1Higher spin resonances have been tried in the fit with both formalisms but found to be statistically insignificant.
\[ K_{11} = \left( \frac{s - s_0}{s_{\text{norm}}} \right) \left( \frac{g_1 \cdot g_1}{s_1 - s} \right) + C_{110} + C_{111} \hat{s} + C_{112} \hat{s}^2 \]  
\[ K_{22} = \left( \frac{s - s_0}{s_{\text{norm}}} \right) \left( \frac{g_2 \cdot g_2}{s_1 - s} \right) + C_{220} + C_{221} \hat{s} + C_{222} \hat{s}^2 \]  
\[ K_{12} = \left( \frac{s - s_0}{s_{\text{norm}}} \right) \left( \frac{g_1 \cdot g_2}{s_1 - s} \right) + C_{120} + C_{121} \hat{s} + C_{122} \hat{s}^2 \]  

(4)

where the factor of \( s_{\text{norm}} = m_K^2 + m_\pi^2 \) is conveniently introduced to make the individual terms in the above expression dimensionless. \( g_1 \) and \( g_2 \) are the real couplings of the \( s_1 \) pole to the first and the second channel respectively. \( s_{0\perp} = 0.23 \text{ GeV}^2 \) is the position of the Adler zero in the \( I = 1/2 \) ChPT elastic scattering amplitude \(^2\). \( C_{11i}, C_{22i} \) and \( C_{12i} \) for \( i = 0, 1, 2 \) are the three coefficients of a second order polynomial for the diagonal and off-diagonal elements of the symmetric \( K \)-matrix. Polynomials are expanded around \( \hat{s} = s/s_{\text{norm}} - 1 \). This form generates an \( S \)-matrix pole, which is conventionally quoted in the complex energy plane as \( E = M - i\eta/2 = 1.408 - i0.110 \text{ GeV} \). Any more distant pole than \( K_0^*(1430) \) is not reliably determined as this simple \( K \)-matrix expression does not have the required analyticity properties. Nevertheless, it is an accurate description for real values of the energy, where scattering takes place. Numerical values of the terms in Eq. (4) are reported in Table I.

The \( I = 3/2 \) \( K \)-matrix is given in Eq. (5). Its form is derived from a simultaneous fit to LASS data \(^9\) and to \( K^-\pi^- \to K^-\pi^- \) scattering data \(^{10}\). It is a non-resonant, single-channel scalar function.

\[ K_{3/2} = \left( \frac{s - s_0}{s_{\text{norm}}} \right) \left( D_{110} + D_{111} \hat{s} + D_{112} \hat{s}^2 \right). \]  

(5)

In Eq. (5) \( s_{0\perp} = 0.27 \text{ GeV}^2 \) is the Adler zero position in the \( I = 3/2 \) ChPT elastic scattering and the values of the polynomial coefficients are \( D_{110} = -0.22147 \), \( D_{111} = 0.026637 \), and \( D_{112} = -0.00092057 \) \(^{13}\).

When moving from scattering processes to \( D \)-decays, the production \( P \)-vector has to be introduced. While the \( K \)-matrix is real, \( P \)-vectors are in general complex reflecting the fact that the initial coupling \( D^+ \to (K^-\pi^+)_{\text{spectator}} \) need not be real. The \( P \)-vector has to have the same poles as the \( K \)-matrix, so that these cancel in the physical decay amplitude. Their functional forms are:

\[
(P_{1/2})_1 = \frac{\beta g_1 e^{i\theta}}{s_1 - s} + (c_{10} + c_{11} \hat{s} + c_{12} \hat{s}^2) e^{i\gamma_1}
\]

(6)

\[
(P_{1/2})_2 = \frac{\beta g_2 e^{i\theta}}{s_1 - s} + (c_{20} + c_{21} \hat{s} + c_{22} \hat{s}^2) e^{i\gamma_2}
\]

(7)

\[
P_{3/2} = (c_{30} + c_{31} \hat{s} + c_{32} \hat{s}^2) e^{i\gamma_3}.
\]

(8)

\( \beta e^{i\theta} \) is the complex coupling to the pole in the ‘initial’ production process, \( g_1 \) and \( g_2 \) are the couplings as given by Table I. The \( K\pi \) mass squared \( s_c = 2 \text{ GeV}^2 \) corresponds to the center of the Dalitz plot. It is convenient to choose this as the value of \( s \) about which the polynomials of Eqs. (6-8) are expanded, by defining \( \hat{s} = s - s_c \). The polynomial terms in each channel are chosen to have a common phase \( \gamma_i \) to limit the number of free parameters in the fit and avoid uncontrolled interference among the physical background terms. Thus, the coefficients of the second order polynomial, \( c_{ij} \), are real. Coefficients and phases of the \( P \)-vectors, except \( g_1 \) and \( g_2 \), are the only free parameters of the fit determining the scalar components.

Free parameters for vectors and tensors are amplitudes and phases \((a_i \) and \( \delta_i \)). \( K\pi \) scattering determines the parameters of the \( K \)-matrix elements and these are fixed inputs to this \( D \) decay analysis. Table II reports our \( K \)-matrix fit results. It shows quadratic terms in \( (P_{1/2})_1 \) are significant in fitting data, while in both \( (P_{1/2})_2 \) and \( P_{3/2} \) constants are sufficient.

The \( J > 0 \) states required by the fit are listed in Table III.

The \( S \)-wave component accounts for the dominant portion of the decay \( (83.23 \pm 1.50)\% \). A significant fraction, 13.61 \( \pm 0.98\% \), comes, as expected, from \( K^*(892) \); smaller contributions come from two vectors \( K^*(1410) \) and \( K^*(1680) \) and from the tensor \( K_2^*(1430) \). It is conventional to quote fit fractions for each component and this is what we do. Care should be taken in interpreting some of these since strong interference can occur. This is particularly apparent between contributions in the same-spin partial wave. While the total \( S \)-wave fraction is a sensitive measure of its contribution to the Dalitz plot, the separate fit fractions for \( I = 1/2 \) and \( I = 3/2 \) must be treated with care. The broad \( I = 1/2 \) \( S \)-wave component inevitably interferes strongly with the slowly varying \( I = 3/2 \) \( S \)-wave, as seen for instance in \(^{14}\). Fit results on the projections are re shown in Fig. 6. The corresponding adaptive binning scheme is at the top of Fig. 7.

The fit \( \chi^2/d.o.f \) is 1.27 corresponding to a confidence level of 1.2\%. If the \( I = 3/2 \) component is removed from the fit, the \( \chi^2/d.o.f \) worsens to 1.54, corresponding to a confidence level of \( 10^{-5} \).

\(^2\)Chiral symmetry breaking demands an Adler zero in the elastic \( S \)-wave amplitudes in the unphysical region. ChPT at next-to-leading order fixes these positions \( s_{0\perp} \) \(^{11, 12}\).
Table I Values of parameters for the $I = 1/2$ $K$-matrix.

<table>
<thead>
<tr>
<th>pole (GeV$^2$)</th>
<th>coupling (GeV)</th>
<th>$C_{11i}$</th>
<th>$C_{12i}$</th>
<th>$C_{22i}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$s_1 = 1.7919$</td>
<td>$g_1 = 0.31072$</td>
<td>$g_2 = -0.02323$</td>
<td>$C_{110} = 0.79299$</td>
<td>$C_{120} = 0.15040$</td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td>$C_{111} = -0.15099$</td>
<td>$C_{121} = -0.038266$</td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td>$C_{112} = 0.00811$</td>
<td>$C_{122} = 0.0022596$</td>
</tr>
</tbody>
</table>

Table II $S$-wave parameters from the $K$-matrix fit to the FOCUS $D^+ \to K^- \pi^+ \pi^+$ data. The first error is statistic, the second error is systematic from the experiment, and the third is systematic induced by model input parameters for higher resonances. Coefficients are for the unnormalized $S$-wave.

<table>
<thead>
<tr>
<th>coefficient</th>
<th>phase (deg)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\beta = 3.380 \pm 0.152 \pm 0.002 \pm 0.068$</td>
<td>$\theta = 286 \pm 4 \pm 0 \pm 3.0$</td>
</tr>
<tr>
<td>$c_{10} = 1.655 \pm 0.156 \pm 0.010 \pm 0.101$</td>
<td>$\gamma_1 = 304 \pm 6 \pm 0 \pm 5.8$</td>
</tr>
<tr>
<td>$c_{11} = 0.780 \pm 0.096 \pm 0.003 \pm 0.090$</td>
<td>$\gamma_2 = 126 \pm 3 \pm 0 \pm 1.2$</td>
</tr>
<tr>
<td>$c_{12} = -0.954 \pm 0.058 \pm 0.0015 \pm 0.025$</td>
<td>$\gamma_3 = 211 \pm 10 \pm 0 \pm 7.8$</td>
</tr>
<tr>
<td>$c_{20} = 17.182 \pm 1.036 \pm 0.023 \pm 0.362$</td>
<td></td>
</tr>
<tr>
<td>$c_{30} = 0.734 \pm 0.080 \pm 0.005 \pm 0.030$</td>
<td></td>
</tr>
<tr>
<td><strong>Total $S$-wave fit fraction</strong></td>
<td><strong>83.23 \pm 1.50 \pm 0.04 \pm 0.07 %</strong></td>
</tr>
<tr>
<td><strong>Isospin 1/2 fraction</strong></td>
<td><strong>207.25 \pm 25.45 \pm 1.81 \pm 12.23 %</strong></td>
</tr>
<tr>
<td><strong>Isospin 3/2 fraction</strong></td>
<td><strong>40.50 \pm 9.63 \pm 0.55 \pm 3.15 %</strong></td>
</tr>
</tbody>
</table>

Table III Fit fractions, phases, and coefficients for the $J > 0$ components from the $K$-matrix fit to the FOCUS $D^+ \to K^- \pi^+ \pi^+$ data. The first error is statistic, the second error is systematic from the experiment, and the third error is systematic induced by model input parameters for higher resonances.

<table>
<thead>
<tr>
<th>component</th>
<th>fit fraction (%)</th>
<th>phase $\delta_j$ (deg)</th>
<th>coefficient</th>
</tr>
</thead>
<tbody>
<tr>
<td>$K^+(892)\pi^+$</td>
<td>$13.61 \pm 0.98$</td>
<td>0 (fixed)</td>
<td>1 (fixed)</td>
</tr>
<tr>
<td>$\pm 0.01 \pm 0.30$</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>$K^+(1680)\pi^+$</td>
<td>$1.90 \pm 0.63$</td>
<td>$1 \pm 7$</td>
<td>$0.373 \pm 0.067$</td>
</tr>
<tr>
<td>$\pm 0.009 \pm 0.43$</td>
<td>$\pm 0.1 \pm 6$</td>
<td>$\pm 0.009 \pm 0.047$</td>
<td></td>
</tr>
<tr>
<td>$K_2^+(1430)\pi^+$</td>
<td>$0.39 \pm 0.09$</td>
<td>$296 \pm 7$</td>
<td>$0.169 \pm 0.017$</td>
</tr>
<tr>
<td>$\pm 0.004 \pm 0.05$</td>
<td>$\pm 0.3 \pm 1$</td>
<td>$\pm 0.010 \pm 0.012$</td>
<td></td>
</tr>
<tr>
<td>$K^+(1410)\pi^+$</td>
<td>$0.48 \pm 0.21$</td>
<td>$293 \pm 17$</td>
<td>$0.188 \pm 0.041$</td>
</tr>
<tr>
<td>$\pm 0.012 \pm 0.17$</td>
<td>$\pm 0.4 \pm 7$</td>
<td>$\pm 0.002 \pm 0.030$</td>
<td></td>
</tr>
</tbody>
</table>

These results can be compared with those obtained in the effective isobar model, which can serve as the standard for fit quality. Projections are shown in Fig. 8 and the adaptive binning scheme at the bottom of Fig. 7.

Two ad hoc scalar resonances are required, of mass 856 $\pm 17$ and 1461 $\pm 4$ and width 464 $\pm 28$ and 177 $\pm 8$ MeV/$c^2$ respectively. A detailed discussion of the results and the systematics can be found in [8]. The results of the $K$-matrix fit showed that a consistent representation with scattering is possible, the global fit quality being indeed good. However, it deteriorates at higher $K\pi$ mass. This is not surprising since our $K$-matrix treatment only includes two channels $K\pi$ and $K\eta'$. While we have reliable information on the former channel, we have relatively poor constraints on the latter. This means that as we consider $K\pi$ masses far above $K\eta'$ threshold, these inadequacies in the description of the $K\eta'$ channel become increasingly important. This is expected to become worse as yet further inelastic channels open up. Consequently, improvements could be made by using a number of $D$-decay chains with $K\pi$ final state interactions and inputting all these in one combined analysis in which
several inelastic channels are included in the \textit{K-matrix} formalism. In the present single $D^+ \to K^\pm \pi^\mp \pi^\pm$ channel, adding further inelastic modes would be just adding free unconstrained parameters for which there is little justification. It is interesting to note that the adaptive binning scheme shows that both the \textit{K-matrix} and the isobar fit are not able to reproduce data well in the region at 2 GeV$^2$, in the vicinity of the $K\eta'$ threshold. It is also the energy domain where higher spin states live. Vector and tensor fit parameters in the two models are in very good agreement: we do not exclude the possibility that a better treatment of these amplitudes could improve the $\chi^2$. Some isolated spots of high $\chi^2$ could be caused by an imperfect modeling of the efficiency as they are in the same regions in both fits.

A feature of the \textit{K-matrix} amplitude analysis is that it allows an indirect phase measurement of the separate isospin components: it is this phase variation with isospin $I = 1/2$ which should be compared with the same $I = 1/2$ LASS phase, extrapolated from 825 GeV down to threshold according to Chiral Perturbation Theory. This is done in the right plot of Fig. 9. In this model [5] the $P$-vector allows for a phase variation accounting for the interaction with the third particle in the process of resonance formation. It so happens that the Dalitz fit gives a nearly constant production phase. The two phases in Fig. 9b) have the same behaviour up to $\sim 1.1$ GeV. However, approaching $K\eta'$ threshold, effects of inelasticity and differing final state interactions start to appear.

The difference between the phases in Fig. 9a) is due to the $I = 3/2$ component.

These results are consistent with $K\pi$ scattering data, and consequently with Watson’s theorem predictions for two-body $K\pi$ interactions in the low $K\pi$ mass region, up to $\sim 1.1$ GeV, where elastic processes dominate. This means that possible three-body interaction effects, not accounted for in the \textit{K-matrix} parametrization, play a marginal role.

Our results for the total $S$-wave are in general agreement with those from the E791 analysis, in which the $S$-wave modulus and phase were determined in each $K\pi$ slice [15], [16]. What does this analysis contribute to the discussion of the existence and parameters of the $\kappa$? We know from analysis [17] of the LASS data (which in $K^-\pi^+$ scattering only start at 825 MeV) there is no pole, the $\kappa$(900), in its energy range. However, below 800 MeV, deep in the complex plane, there is very likely such a state. Its precise location requires a more sophisticated analytic continuation onto the unphysical sheet than the \textit{K-matrix} representation provided here. This is because of the need to approach close to the crossed channel cut, which is not correctly represented for a robust analytic continuation. However, our \textit{K-matrix} representation fits along the real energy axis inputs on scattering data and Chiral Perturbation Theory in close agreement with those used in the analysis by Descotes-Genon and Moussallam [18] that locates the $\kappa$ with a mass of $(658 \pm 13)$ MeV and a width of $(557 \pm 24)$ MeV by careful continuation. These pole parameters are quite different from those.
implied by the simple isobar fits. We have thus shown that whatever \( \kappa \) is revealed by our \( D^+ \to K^-\pi^+\pi^+ \) results, it is the same as that found in scattering data. Consequently, our analysis supports the conclusions of [18] and [19].

4. Conclusions

Dalitz-plot analysis represents a unique, powerful and promising tool for studying the Heavy Flavor decay dynamics. There is a recent, vigorous effort to perform amplitude analysis: a more robust formalism has been implemented, many channels have been investigated. The beauty community can benefit from charm experience and expertise. The high statistic \( D^+ \to K^-\pi^+\pi^+ \) from FOCUS showed us that \( D^- \) decay can also teach us about \( K\pi \) interaction much closer to threshold than the older scattering results. This serves as a valuable check from experiment [20] of the inputs to the analyses of [18] and [19] based largely on theoretical considerations. Dalitz-plot analysis will definitely keep us company over the next few years. There will be a lot of work for both experimentalists and theorists alike: synergy will be invaluable. Some complications have already emerged, especially in the charm field, others, unexpected, will only become clearer when we delve deeper into the beauty sector. \( B_s \) will be a completely new chapter. The analysis is challenging but there are no shortcuts toward ambitious and high-precision studies and, ultimately, to New Physics searches.

References

Adaptive binning and $\chi^2$ contributions for dp to ppp (data)

$\chi^2_{tot} = 49.375$
Free Par = 17
DOF = 36
CL = 6.79E-02
$\chi^2$/dof = 1.372
Cut = 20
Bins = 53

$\chi^2$ contributions for dp to ppp (data)
Leptonic and semileptonic $D$ and $D_s$ decays at B-factories

L. Widhalm (Belle collaboration)
Institute of High-Energy Physics, Austrian Academy of Sciences, Vienna 1050, Austria

Recent measurements of branching fractions, form factors and decay constants of leptonic and semileptonic decays of $D_{(s)}$-mesons acquired at experiments running at the $T(4S)$ resonance energy are reviewed.

I. INTRODUCTION

One of the important goals of particle physics is the precise measurement and understanding of the Cabibbo-Kobayashi-Maskawa (CKM) Matrix. To interpret results from B-factory experiments such as BaBar [1] and Belle [2], theoretical calculations of form factors and decay constants (usually based on lattice gauge theory, see e.g. [3]) are needed. It is necessary to have accurate measurements in the charm sector to check (and allow further tuning of) theoretical methods and predictions.

Due to their relative abundance and simplified theoretical treatment, (semi)leptonic decays of $D$ or $D_s$ mesons are a favored means of determining the weak interaction couplings of quarks within the standard model.

This review concentrates on experimental results for such decays achieved at experiments running at the $T(4S)$ resonance threshold, namely the BaBar and Belle experiments. It uses adapted excerpts from the cited Belle and BaBar publications.

II. THE EXPERIMENTS

The BaBar detector [1] reconstructs charged particles by matching hits in the 5-layer double-sided silicon vertex tracker (SVT) with track elements in the 40-layer drift chamber (DCH), which is filled with a gas mixture of helium and isobutane. Slow particles which do not leave enough hits in the DCH due to the bending in the 1.5-T magnetic field, are reconstructed in the SVT. Charged hadron identification is performed combining the measurements of the energy deposition in the SVT and in the DCH with the information from the Cherenkov detector (DIRC). Photons are detected and measured in the CsI(Tl) electromagnetic calorimeter (EMC). Electrons are identified by the ratio of the track momentum to the associated energy deposited in the EMC, the transverse profile of the shower, the energy loss in the DCH, and the Cherenkov angle in the DIRC. Muons are identified in the instrumented flux return, composed of resistive plate chambers interleaved with layers of steel and brass.

The Belle detector [2] is a large-solid-angle magnetic spectrometer that consists of a silicon vertex detector (SVD), a 50-layer central drift chamber (CDC), an array of aerogel threshold Cherenkov counters (ACC), a barrel-like arrangement of time-of-flight scintillation counters (TOF), and an electromagnetic calorimeter comprised of CsI(Tl) crystals (ECL) located inside a superconducting solenoid coil that provides a 1.5 T magnetic field. An iron flux-return located outside of the coil is instrumented to detect $K_L^0$ mesons and to identify muons (KLM). The detector is described in detail elsewhere [2]. Two inner detector configurations were used. A 2.0 cm beampipe and a 3-layer silicon vertex detector were used for the first sample of 156 fb$^{-1}$, while a 1.5 cm beampipe, a 4-layer silicon detector and a small-cell inner drift chamber were used to record the remaining 392 fb$^{-1}$ [3].

III. SEMILEPTONIC DECAYS TO SCALAR MESONS

Form factors from $D$ meson semileptonic decay have been calculated using lattice QCD techniques [2, 6, 7]. In the theoretical description, the differential decay width is dominated by the form factor $f_+(q^2)$ [8].

$$\frac{d\Gamma^{K(\pi)}_{D}}{dq^2} = \frac{G_F^2 |V_{cs}(d)|^2}{24\pi^3} |f_+^{K(\pi)}(q^2)|^2 p_K^{3}(q)$$ (1)

where $p_K^{3}(q)$ is the magnitude of the meson 3-momentum in the $D_{(s)}^{0}$ rest frame.

In the modified pole model [9], the form factor $f_+$ is described as

$$f_+(q^2) = \frac{f_+(0)}{(1 - q^2/m^2_{pole})/(1 - \alpha_p q^2/m^2_{pole})}$$ (2)

with the pole masses predicted as $m(D^+_s) = 2.11$ GeV/$c^2$ (for $D^0_{sig} \rightarrow K^+\ell^+\nu$) and $m(D^*_s) = 2.01$ GeV/$c^2$ (for $D^0_{sig} \rightarrow \pi^+\ell^+\nu$). Setting $\alpha_p = 0$ leads to the simple pole model [8].

A model independent description of the form factor has been studied in [10]. The most general expressions of the form factor $f_+(q^2)$ are analytic functions satisfying the dispersion relation:

$$f_+(q^2) = \frac{Res(f_+(q^2))}{m^2_{D^+} - q^2} + \frac{1}{\pi} \int_{t_+}^{\infty} dt \frac{\Im f_+(t)}{t - q^2 - i\epsilon}.$$ (3)
The only singularities in the complex $t \equiv q^2$ plane originate from the interaction of the charm and the strange quarks in vector states. They are a pole, situated at the $D_s^*$ mass squared and a cut, along the positive real axis, starting at threshold $(t_{\text{cut}} = (m_D + m_K)^2)$ for $D^0K^-$ production.

This cut $t$-plane can be mapped onto the open unit disk with center at $t = 0$ using the variable:

$$
z(t, t_0) = \frac{\sqrt{t - t_{\text{cut}} - t_0}}{\sqrt{t - t_{\text{cut}} + t_0}}.
$$

In this variable, the physical region for the semileptonic decay corresponds to the small (real) range between $\pm z_{\text{max}} = \pm 0.051$. The $z$ expansion of $f_+$:

$$
f_+(t) \propto \sum_{k=0}^{\infty} a_k (t, t_0) z^k(t, t_0),
$$

is thus expected to converge quickly.

### A. $D \to (K/\pi)(e/\mu)\nu(e/\mu)$ at Belle

Belle has measured the absolute branching fractions and form factors of $D^0 \to K^-\pi^+\nu$ and $D^0 \to \pi^-\ell^+\nu$ ($l = e, \mu$) [11], using a novel reconstruction method with better $q^2$ resolution than in previous experiments. The analysis is based on data corresponding to a total integrated luminosity of 282 fb$^{-1}$.

To achieve good resolution in the neutrino momentum and $q^2$, the $D^0$ are tagged by fully reconstructing the remainder of the event. The studied events are of the type $e^+e^- \to D^{(*)}_{\text{tag}}D^{(*)}_{\text{sig}}X$ ($D^{(*)}_{\text{sig}} \to D^0_{\text{sig}}\pi^-$), where $X$ may include additional $\pi^\pm$, $\pi^0$, or $K^\pm$ mesons (inclusion of charge-conjugate states is implied throughout this paper). The $D^{(*)}_{\text{tag}}$ is reconstructed in the modes $D^{*+} \to D^0_\pi^+\ell^+\nu$, $D^{*0} \to D^0_\pi^0\ell^\pm\gamma$, with $D^{*0} \to K^-(n\pi)^{++/\pm}$ ($n = 1, 2, 3$). Each $D_{\text{tag}}$ and $D^{(*)}_{\text{tag}}$ candidate is subjected to a mass-constrained vertex fit to improve the momentum resolution. The 4-momentum of $D^{(*)}_{\text{tag}}$ is found by energy-momentum conservation, assuming a $D^{(*)}_{\text{tag}}D^{(*)}_{\text{sig}}X$ event. Its resolution is improved by subjecting it to a fit of the $X$ tracks and the $D^{(*)}_{\text{tag}}$ momentum, constrained to originate at the run-by-run average collision point, while the invariant mass is constrained to the nominal mass of a $D^{*+}$. Candidates for $\pi^\pm$ are selected from among the remaining tracks, and for each the candidate $D^0_{\text{sig}}$ 4-momentum is calculated from that of the $D^{(*)}_{\text{tag}}$ and $\pi^\pm$. The momentum is then adjusted by a kinematic fit constraining the candidate mass to that of the $D^0$. For this fit, the decay vertex of the $D^0_{\text{sig}}$ has been estimated by extrapolating from the collision point in the direction of the $D^0_{\text{sig}}$ momentum assuming the average decay length.

![FIG. 1: Belle: Form factors for (a) $D^0 \to K^-\ell^+\nu$, in $q^2$ bins of 0.067 GeV$^2$/c$^2$ and (b) $D^0 \to \pi^-\ell^+\nu$, in $q^2$ bins of 0.3 GeV$^2$/c$^2$. Overlaid are the predictions of the simple pole model using the physical pole mass (dashed), and a quenched (yellow) and unquenched (purple) LQCD calculation. Each LQCD curve is obtained by fitting a parabola to values calculated at specific $q^2$ points. The shaded band reflects the theoretical uncertainty and is shown within the range of $q^2$ for which calculations are reported.

Background lying under the $D^0_{\text{sig}}$ mass peak (i.e. fake-$D^0_{\text{sig}}$) is estimated using a wrong sign (WS) sample where the tag- and signal-side $D$ candidates have the same flavor ($\bar{D}_{\text{tag}}$ instead of $D_{\text{tag}}$). A MC study (including $Y(4S) \to B\bar{B}$ and continuum $q\bar{q}$, where $q = c, s, u, d$ events) [12] [13] has found that this sample can properly model the shape of background except for a small contribution from real $D^0_{\text{sig}}$ decays ($\approx 2\%$) from interchange between particles used for the tag due to particle misidentification. Background from fake $D^0_{\text{sig}}$ is subtracted normalizing this shape in a sideband region $1.84 - 1.85$ GeV/c$^2$, yielding $56461 \pm 309_{\text{stat}} \pm 830_{\text{sys}}$ signal $D^0_{\text{sig}}$ tags.

Within this sample of $D^0_{\text{sig}}$ tags, the semileptonic decay $D^0_{\text{sig}} \to K^+(\pi^+)\ell^-\nu$ is reconstructed with $\ell^-\nu$ and $\ell^+\nu$ candidates from among the remaining tracks. The neutrino candidate 4-momentum is reconstructed by energy-momentum conservation, and its invariant mass squared, $m_{\nu\ell}^2$, is required to satisfy $|m_{\nu\ell}^2| < 0.05$ GeV$^2$/c$^4$.

Multiple candidates still remain in one third of $D^0_{\text{sig}}$ tags, and in about one quarter of the semileptonic sample. In these cases all candidates are saved and given equal weights such that each event has a total weight of 1.
TABLE I: Belle: Yields in data, estimated backgrounds, extracted signal yields and branching fractions, where for the latter two, the first uncertainty is statistical and second is systematic; small differences in the numbers are due to rounding.

<table>
<thead>
<tr>
<th>channel</th>
<th>full $\bar{D}_0$</th>
<th>$K^+e^-\nu_e$</th>
<th>$K^+\mu^-\nu_\mu$</th>
<th>$\pi^+e^-\nu_e$</th>
<th>$\pi^+\mu^-\nu_\mu$</th>
</tr>
</thead>
<tbody>
<tr>
<td>yield</td>
<td>95250</td>
<td>1349</td>
<td>1333</td>
<td>152</td>
<td>141</td>
</tr>
<tr>
<td>fake $\bar{D}_0$</td>
<td>38789</td>
<td>12.6</td>
<td>12.2</td>
<td>12.3</td>
<td>12.5</td>
</tr>
<tr>
<td>semileptonic</td>
<td>n/a</td>
<td>6.7</td>
<td>10.0</td>
<td>11.7</td>
<td>12.6</td>
</tr>
<tr>
<td>hadronic</td>
<td>n/a</td>
<td>11.9</td>
<td>62.1</td>
<td>1.8</td>
<td>9.7</td>
</tr>
<tr>
<td>signal</td>
<td>56461</td>
<td>1318</td>
<td>1249</td>
<td>126</td>
<td>106</td>
</tr>
<tr>
<td>stat. error</td>
<td>309</td>
<td>37</td>
<td>37</td>
<td>12</td>
<td>12</td>
</tr>
<tr>
<td>syst. error</td>
<td>830</td>
<td>7</td>
<td>25</td>
<td>3</td>
<td>6</td>
</tr>
<tr>
<td>Braching Fraction</td>
<td>(10^{-4})</td>
<td>345 ± 10 ± 19</td>
<td>345 ± 10 ± 21</td>
<td>27.9 ± 2.7 ± 1.6</td>
<td>23.1 ± 2.6 ± 1.9</td>
</tr>
</tbody>
</table>
(\epsilon and \mu channels, average) | 345 ± 7 ± 20 | & 25.5 ± 1.9 ± 1.6 |

The contribution from fake $\bar{D}_0$ in the sample of semileptonic decay candidates is estimated using the $\bar{D}_0$ invariant mass WS shape of the $\bar{D}_0$ tag sample, normalized in the previously defined sideband region. Backgrounds from semileptonic decays with either an incorrectly identified meson or where additional mesons are lost in reconstruction are highly suppressed by the good neutrino mass resolution. For $\bar{D}_0 \to \pi^+\ell^-\nu$ the most significant background is $\bar{D}_0 \to K^+\ell^-\nu$ amounting to 6% - 8% of the total yield. It was estimated using the reconstructed $\bar{D}_0 \to K^+\ell^-\nu$ decays in data, reweighted with the (independently measured) probability of kaons to fake pions. Smaller backgrounds from $\bar{D}_0 \to K^+\ell^-\nu$ and $\bar{D}_0 \to \rho^+\ell^-\nu$ decays amounting to 0.8% - 0.9% were measured by normalizing MC yield to data in the upper sideband region $m_\tau^2 > 0.3$ GeV$^2$/c$^4$, which is dominated by these channels. For $\bar{D}_0 \to K^+\ell^-\nu$, decays of $\bar{D}_0 \to K^+\ell^-\nu$ contribute at the level of 5% - 0.8%, measured using a sideband evaluation as described above, while background from $\bar{D}_0 \to \pi^+\ell^-\nu$ and $\bar{D}_0 \to \rho^+\ell^-\nu$ was found to be negligible (< 0.07% of the total yield). Background from $\bar{D}_0$ decays to hadrons, where a hadron is mis-identified as a lepton, is measured with an opposite sign (OS) sample, where the lepton charge is opposite to that of the $\bar{D}_0$ slow pion. Note that the signal is extracted from the same sign (SS) sample. In contrast to the SS sample, the OS sample has no signal or semileptonic backgrounds; fake $\bar{D}_0$ are subtracted in the same manner described previously. Assigning well identified pion and kaon tracks a lepton mass, pure background $m_\tau^2$ distributions are constructed in both SS and OS, which are labelled $f_m^\SS$ and $f_m^\OS$, $m = K, \pi$. A fit of the weights $a_K$ and $a_\pi$ of the components $f_K^\OS$ and $f_\pi^\OS$ in the $m_\tau^2$ distribution of the OS data sample is performed, and the hadronic background in the SS data sample is calculated as $(a_K f_K^\SS + a_\pi f_\pi^\SS)$, utilizing the fact that the hadron misidentification rate does not depend on the charge correlation defining SS and OS. The method has been validated using MC samples. As the muon fake rate is about an order of magnitude larger than that for electrons, this background is much more significant for muon modes. The signal yields and estimated backgrounds are summarized in the upper part of Table I.

Efficiencies depend strongly on $n_X$, defined as the number of $\pi^\pm(0)$ and $K^\pm$ mesons assigned to $X$ in $e^+e^- \to D_{\ell\nu}^{(*)}D_{\ell\nu}^*X$, and are determined with MC; differences in the $n_X$ distribution between MC and data give rise to a further (+1.9 ± 3.9)% correction. Applying these corrections, the absolute branching fractions (normalized to the total number of $\bar{D}_0$) summarized in the lower part of Table I are obtained.

The resolution in $q^2$ of semileptonic decays is found to be $\sigma_{q^2} = 0.0145 ± 0.0007 \text{stat GeV}^2/c^4$ in MC signal events. This is much smaller than statistically reasonable bin widths, which have been chosen as 0.067 (0.3) GeV$^2$/c$^4$ for kaon (pion) modes, and hence no unfolding is necessary. Bias in the measurement of $q^2$ that may arise due to events where the lepton and meson are interchanged, a double mis-assignment, was checked with candidate $\bar{D}_0$ slow pion to $K^+\ell^-\nu$ decays and found to be negligible. The differential decay width is bin-by-bin background subtracted and efficiency corrected, using the same methods described previously.

The measured $q^2$ distribution is fitted with 2 free parameters to the predicted differential decay width $d\Gamma/dq^2$ of the pole models with $f_p(0)$ being one of the parameters, and either $m_{\text{pole}}$ (setting $\alpha_p = 0$) or $\alpha_p$ (assuming the theoretical pole) the other. Binning effects are accounted for by averaging the model functions within individual $q^2$ bins. The fit to the simple pole model yields $m_{\text{pole}}(K^+\ell^-\nu) = 1.82 ± 0.04_{\text{stat}} ± 0.03_{\text{syst}} \text{GeV}/c^2$ ($x^2$/ndf = 34/28) and $m_{\text{pole}}(\pi^+\ell^-\nu) = 1.97 ± 0.08_{\text{stat}} ± 0.04_{\text{syst}} \text{GeV}/c^2$ ($x^2$/ndf = 6.2/10). While the pole mass for the $\pi\nu$ decay agrees within errors with the predicted value, $m(D^*)$, the more accurate fit of $m_{\text{pole}}(K\ell\nu)$ is several standard deviations below $m(D^*)$. In the mod-
ified pole model, $\alpha_p$ describes this deviation of the real poles from the $m(D^*_0)$ masses. Fixing these masses to their known experimental values, a fit of $\alpha_p$ yields $\alpha_p(D^0 \to K^- e^+ \nu) = 0.52 \pm 0.08_{\text{stat}} \pm 0.06_{\text{syst}}$ ($\chi^2/\text{ndf} = 31/28$) and $\alpha_p(D^0 \to \pi^- e^+ \nu) = 0.10 \pm 0.21_{\text{stat}} \pm 0.10_{\text{syst}}$ ($\chi^2/\text{ndf} = 6.4/10$).

The fitted values for $f_\pm^{K,\pi}(0)$ vary little for the different fits, for the modified pole model the results are $f_+^{K}(0) = 0.695 \pm 0.007_{\text{stat}} \pm 0.022_{\text{syst}}$ and $f_+^{\pi}(0) = 0.624 \pm 0.020_{\text{stat}} \pm 0.030_{\text{syst}}$.

The measured form factors $f_\pm^{K,\pi}(q^2)$ are shown in Figure 1 with predictions of the simple pole model, unquenched [6] and quenched [7] LQCD. To obtain a continuous curve for $f_\pm$ from the LQCD values reported at discrete $q^2$ points, the values were fitted by a parabola, which is found to fit well within the stated theoretical errors and is not associated with any specific model. To quantify the degree of agreement, a $\chi^2/\text{ndf}$ is calculated between this measurement and the interpolated LQCD curve within the $q^2$ range for which LQCD predictions are made. For the kaon modes, $\chi^2/\text{ndf}$ is 28/18 (34/23), for the pion modes 9.8/5 (3.4/5); correlations induced by the fit of the calculated $q^2$ points to a parabola have been considered.

B. $D \to K e \nu_e$ at BABAR

This subsection is an adapted excerpt of BABAR’s publication [14].

The corresponding BABAR analysis [14] is using a total integrated luminosity of 75 fb$^{-1}$ collected during the years 2000-2002. It measures the $q^2$ variation and the absolute value of the hadronic form factor at $q^2 = 0$ for the decay $D^0 \to K^- e^+ \nu_e(\gamma)$. Normalizing to $D^0 \to K^- \pi^+$, it also gives a value for its branching fraction. In contrast to the Belle analysis, a semi-inclusive reconstruction technique is used to select semileptonic decays with less resolution, but much higher efficiency. As a result of this approach, events with a photon radiated during the $D^0$ decay are included in the signal.

$D^0 \to K^- e^+ \nu_e(\gamma)$ decays are reconstructed in $e^+ e^- \to c\bar{c}$ events where the $D^0$ originates from the $D^{*+} \to D^0 \pi^+$. Charged and neutral particles are boosted to the center of mass system (c.m.) and the event thrust axis is determined. The direction of this axis is required to be in the interval $|\cos(\theta_{\text{thrust}})| < 0.6$ to minimize the loss of particles in regions close to the beam axis. A plane perpendicular to the thrust axis is used to define two hemispheres, equivalent to the two jets produced by quark fragmentation. In each hemisphere, pairs of oppositely charged leptons and kaons are searched for. For the charged lepton candidates only electrons or positrons with c.m. momentum greater than 0.5 GeV/c are considered.

Since the $\nu_e$ momentum is unmeasured, a kinematic fit is performed, constraining the invariant mass of the candidate $e^+ K^- \nu_e$ system to the $D^0$ mass. In this fit, the $D^0$ momentum and the neutrino energy are estimated from the other particles measured in the event. The $D^0$ direction is taken as the direction opposite to the sum of the momenta of all reconstructed particles in the event, except for the kaon and the positron associated with the signal candidate. The energy of the jet is determined from the total c.m. energy and from the measured masses of the two jets. The neutrino energy is estimated as the difference between the total energy of the jet and the sum of the energies of all reconstructed particles in the hemisphere. A correction, which depends on the value of the missing energy measured in the opposite jet, is applied to account for the presence of missing energy due to particles escaping detection, even in the absence of a neutrino from the $D^0$ decay.

Background events arise from $\Upsilon(4S)$ decays and hadronic events from the continuum. To reduce the contribution from $B\bar{B}$ events, selection criteria exploiting the topological differences to events with $c\bar{c}$ fragmentation are used. Background events from the continuum arise mainly from charm particles. Because charm hadrons take a large fraction of the charm quark energy, charm decay products have higher average energies and different angular distributions (relative to the thrust axis or to the $D$ direction) compared with other particles in the hemisphere emitted from the hadronization of the $c$ and $\bar{c}$ quarks. Selection criteria based on these considerations are applied to suppress this kind of background.

The remaining background from $c\bar{c}$ events can be divided into peaking (60%) and non-peaking (40%) candidates. Peaking events are those background events whose distribution is peaked around the signal region. These are mainly events with a real $D^{*+}$ in which the slow $\pi^+$ is included in the candidate track combination. Backgrounds from $e^+ e^-$ annihilations into light $u\bar{u}$, $d\bar{d}$, $s\bar{s}$ quarks and $BB$ events are non-peaking. To improve the accuracy of the reconstructed $D^0$ momentum, the nominal $D^{*+}$ mass is added as a constraint in the previous fit and only events with a $\chi^2$ probability higher than 1% are kept, resulting in 85260 selected $D^0$ candidates containing an estimated number of 11280 background events. The non-peaking component comprises 54% of the background. Detailed studies have been performed to understand corrections (and the connected systematics) of various components of the peaking background.

To obtain the true $q^2$ distribution, the measured one has to be corrected for selection efficiency and detector resolution effects. This is done using an unfolding algorithm based on MC simulation of these effects: Using Singular Value Decomposition (SVD) [15] of the resolution matrix, the unfolded $q^2$ distribution.
for signal events, corrected for resolution and acceptance effects, is obtained. This approach provides the full covariance matrix for the bin contents of the unfolded distribution. To verify that the $q^2$ variation of the selection efficiency is well described by the simulation, a control sample of $D^0 \rightarrow K^-\pi^+\pi^0$ is reconstructed as if they were $K^-e^+\nu_e$ events, and indicates no significant bias. With a second control sample of $D^0 \rightarrow K^-\pi^+$, the accuracy of the $D^0$ direction and missing energy reconstruction for the $D^0 \rightarrow K^-e^+\nu_e$ analysis is checked. This information is used in the mass-constrained fits and thus influences the $q^2$ reconstruction. Once the simulation is tuned to reproduce the results obtained on data for these parameters, the $q^2$ resolution distributions agree very well. One half of the measured variation on the fitted parameters from these corrections has been taken as a systematic uncertainty.

Effects from a momentum-dependent difference between data and simulated events on the charged lepton and on the kaon identification have been found to be $<2\%$ and included in the corrections and systematics. Care has also been taken to correctly understand radiative decays where $q^2 = (p_D - p_K)^2 = (p_e + p_\mu + p_\pi)^2$. Corresponding corrections have been applied and the corresponding uncertainties enter in the systematic uncertainty evaluation. Toy simulations have been used to verify that the statistical precision obtained for each binned unfolded value is correct and if biases generated by removing information are under control.

The fit to a model is done by comparing the number of events measured in a given bin of $q^2$ with the expectation from the exact analytic integration of the expression $|\tilde{p}_K(q^2)|^3 |f_+(q^2)|^2$ over the bin range, with the overall normalization left free. The summary of the fits to the normalized $q^2$ distributions is presented in Table II. As long as the form factor parameters are left free in the fit, the fitted distributions agree well with the data and it is not possible to reject any of the parameterizations. As also observed by Belle and other experiments, the simple pole model ansatz, with $m_{pole} = m_{D^0} = 2.112$ GeV/c$^2$ does not reproduce the measurements.

Table II: $\text{BaBar}$: Fitted values of the parameters corresponding to different parameterizations of $f_+(q^2)$. The last column gives the $\chi^2/NDF$ of the fit when using the value expected for the parameter.

<table>
<thead>
<tr>
<th>Theoretical Ansatz</th>
<th>Unit</th>
<th>Parameters</th>
<th>$\chi^2/NDF$ Expectations $\chi^2/NDF$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$z$ expansion</td>
<td></td>
<td>$a_1 = -2.5 \pm 0.2 \pm 0.2$</td>
<td>5.9/7</td>
</tr>
<tr>
<td></td>
<td></td>
<td>$a_2 = 0.6 \pm 6. \pm 5$</td>
<td></td>
</tr>
<tr>
<td>Modified pole</td>
<td></td>
<td>$\alpha_{pole} = 0.377 \pm 0.023 \pm 0.029$</td>
<td>6.0/8</td>
</tr>
<tr>
<td>Simple pole</td>
<td>GeV/c$^2$</td>
<td>$m_{pole} = 1.884 \pm 0.012 \pm 0.015$</td>
<td>7.4/8</td>
</tr>
</tbody>
</table>

$z(q^2)$ (defined above), where $P \times \Phi$ is the theoretical normalization, which constrains this product to unity at $z = z_{max}$ (equivalent to $q^2 = 0$). The data are compatible with a linear dependence, which is fully consistent with the modified pole ansatz for $f_+(q^2)$.

![FIG. 2: $\text{BaBar}$: Measured values for $P \times \Phi \times f_+$ are plotted versus $-z$ and requiring that $P \times \Phi \times f_+ = 1$ for $z = z_{max}$. The straight lines represent the result for the modified pole ansatz, the fit in the center and the statistical and total uncertainty.](image)

The $D^0 \rightarrow K^-e^+\nu_e$ branching fraction is measured relative to the reference decay channel, $D^0 \rightarrow K^-\pi^+$:

$$R_D = \frac{BR(D^0 \rightarrow K^-e^+\nu_e)_{\text{data}}}{BR(D^0 \rightarrow K^-\pi^+)}$$

Using slightly adapted selection criteria, after background subtraction there remain 76283 ± 323 events of $D^0 \rightarrow K^-e^+\nu_e$ in data. To select $D^0 \rightarrow K^-\pi^+$ candidates, care has been taken to do this in the most similar way possible. After background subtraction and the necessary corrections, there are
134537±374 candidates selected in the interval $\delta(m) \in [0.142, 0.149]$ GeV/c$^2$.

A summary of the systematic uncertainties on $R_D$ is given in [14]. The measured relative decay rate is:

$$R_D = 0.9269 \pm 0.0072 \pm 0.0119.$$  \hspace{1cm} (7)

Using the world average for the branching fraction $BR(D^0 \rightarrow K^-\pi^+) = (3.80 \pm 0.07)$%, [16], gives $BR(D^0 \rightarrow K^-e^+\nu_e(\gamma)) = (3.522 \pm 0.027 \pm 0.045 \pm 0.065)$%, where the last quoted uncertainty corresponds to the accuracy on $BR(D^0 \rightarrow K^-\pi^+)$. The value of the hadronic form factor at $q^2 = 0$ can be obtained as

$$f_+(0) = \frac{1}{V_{cs}} \sqrt{\frac{24\pi^3}{G_F} \frac{BR}{\tau_{D^0} I}},$$  \hspace{1cm} (8)

where $BR$ is the measured $D^0 \rightarrow K^-e^+\nu_e$ branching fraction, $\tau_{D^0} = (410.1 \pm 1.5) \times 10^{-15}$ s [16] is the $D^0$ lifetime and $I = \int \frac{d^2q}{q^4} |F_{D^0}^e|^2 dq^2$. To account for the variation of the form factor within one bin, and in particular to extrapolate the result at $q^2 = 0$, the pole mass and the modified pole ansätze have been used; the corresponding values obtained for $f_+(0)$ differ by 0.002. Taking the average between these two values and including their difference in the systematic uncertainty, this gives

$$f_+(0) = 0.727 \pm 0.007 \pm 0.005 \pm 0.007,$$  \hspace{1cm} (9)

where the last quoted uncertainty corresponds to the accuracy on $BR(D^0 \rightarrow K^-\pi^+)$, $\tau_{D^0}$ and $|V_{cs}|$. For the $z$ expansion, this corresponds to $a_0 = (2.98 \pm 0.01 \pm 0.03 \pm 0.03) \times 10^{-2}$.

IV. SEMILEPTONIC DECAYS TO VECTOR MESONS

The differential semileptonic decay rate of a scalar meson to a vector meson, specifically, $D_s^+ \rightarrow \phi e^+\nu_e$, depends on the four variables $q^2$, $\theta_e$, $\theta_V$ and $\chi$ [17], depicted in Fig. 3.

Neglecting the electron mass, the differential decay rate as function of these four variables depends in a given way [18, 19] on three form factors

$$A_{1,2}(q^2) = \frac{A_{1,2}(0)}{1 - q^2/m_A^2},$$  \hspace{1cm} (10)

$$V(q^2) = \frac{V(0)}{1 - q^2/m_V^2},$$  \hspace{1cm} (11)

with the pole masses $m_A = 2.5$ GeV/c$^2$ and $m_V = 2.1$ GeV/c$^2$. Measurements have usually been expressed in terms of the ratios of the form factors at $q^2 = 0$, namely:

$$r_V = V(0)/A_1(0) \quad \text{and} \quad r_2 = A_2(0)/A_1(0).$$  \hspace{1cm} (12)

Based on a prediction by [20], $r_V$ is a constant depending only on particle masses,

$$r_V = \frac{(m_{D^*_s} + m_\phi)^2}{m_{D_s^*}^2 + m_\phi^2} = 1.8.$$  \hspace{1cm} (13)

A. $D_s \rightarrow \phi e\nu_e$ at BABAR

This subsection is an adapted excerpt of BABar’s publication [18].

BABar has presented a study of the hadronic form factors for the vector meson decay $D_s^+ \rightarrow \phi e^+\nu_e$ with $\phi \rightarrow K^+K^-$ [18] (results still preliminary). This analysis is based on a fraction of the total available BABAR data sample, corresponding to integrated luminosities of 78.5 fb$^{-1}$ recorded on the $\Upsilon(4S)$ resonance. It focuses on semileptonic decays of $D_s^+$ mesons which are produced via $e^+e^- \rightarrow \phi$ annihilation. $D_s$ mesons produced in $B\bar{B}$ events are not included and treated as background.

Similar to BABar’s $D^0 \rightarrow K e\nu_e$ analysis presented above, a plane perpendicular to the thrust axis is used to define two hemispheres, equivalent to the two jets produced by quark fragmentation. In each hemisphere, decay products of the $D_s^+$, a charged lepton and two oppositely charged kaons are searched for. Charged leptons are required to have a c.m. momentum larger than 0.5 GeV/c. The unmeasured neutrino momentum is determined in a way similar to the $D^0 \rightarrow K e\nu_e$ analysis presented above.

Figure 4 shows the $K^+K^-$ invariant mass distribution for the selected decays compared to MC simulation and the composition of the background. $\phi$ candidates are defined as $K^+K^-$ pairs with an invariant mass in the interval from 1.01 and 1.03 GeV/c$^2$. Various selection criteria are applied to suppress the background [18]. About 71% of the total background include a true $\phi$ decay combined with an electron from another source, namely $B$ meson decays (41%), charm...
FIG. 4: \textit{BaBar}: $K^+K^-$ invariant mass distribution from data and simulated events. MC events have been normalized to the data luminosity according to the different cross sections. The excess of signal events in the $\phi$ region can be attributed to a different production rate and decay branching fraction of $D_s^+$ mesons in data and in simulated events. Dedicated studies have been done to evaluate the amount of peaking background in real events.

particle decays (25%), photon conversions or Dalitz decays (24%), and the rest are fake electrons. These $\phi$ mesons are expected to originate from the primary vertex, or from a secondary charm decay vertex.

A maximum likelihood fit is performed to the four-dimensional decay distribution in the reconstructed variables $q^2$, $\cos(\theta_V)$, $\cos(\theta_e)$, and $\chi$, using the likelihood function

$$\mathcal{L} = - \sum_{i=1}^{625} \ln \mathcal{P}(n^\text{data}_i | n^\text{MC}_i).$$

In this expression, for each bin $i$, $\mathcal{P}(n^\text{data}_i | n^\text{MC}_i)$ is the Poisson probability to observe $n^\text{data}_i$ events, when $n^\text{MC}_i$ are expected. Considering the typical resolutions and the available statistics, the four variables are divided into 4 bins each, corresponding to a four-dimensional array with a total number of bins of 625.

To determine the expected number of signal events, a dedicated sample of signal events is generated in MC with a uniform decay phase space distribution, and each event is weighted using the differential decay rate divided by $p_{\phi}$. Two of the four variables, $\cos(\theta_V)$ and $\chi$, are integrated (averaged), taking advantage of the fact that the estimated background rate is flat in these variables. The background components are normalized to correspond to the expected rates for the integrated luminosity of the data sample. The absolute normalization for signal events ($N_\text{S}$) is left free to vary in the fit. In each bin $(i)$, the expected number of events is evaluated to be:

$$n^\text{MC}_i = N_\text{S} \sum_{j=1}^{N^\text{signal}_i} w_j(\lambda_k) \frac{n^\text{bckg}_i}{W_{\text{tot}}(\lambda_k)} + n^\text{bckg}_i.$$  

Here $n^\text{signal}_i$ refers to the number of simulated signal events, with reconstructed values of the four variables corresponding to bin $i$. The weight $w_j$ is evaluated for each event, using the generated values of the kinematic variables, thus accounting for resolution effects. $W_{\text{tot}}(\lambda_k) = \sum_{j=1}^{N^\text{signal}_i} w_j(\lambda_k)$ is the sum of the weights for all simulated signal events which have been generated according to a uniform phase space distribution. $N_\text{S}$ and $\lambda_k$ are the parameters to be fitted. Specifically, the free parameters $\lambda_k$ are $r_2$, $r_V$, and...
parameters which define $q^2$ dependence of the form factors. To avoid having to introduce finite ranges for the fit to the pole masses, $m_i$, we define $m_i = 1 + \lambda_i^2$. This expression ensures that $m_i$ is always larger than $q_{max}^2 \approx 0.9$ GeV$^2$.

The fit to the four-dimensional data distribution is performed using simulated signal events generated according to a uniform phase space distribution. Signal MC events are weighted to correct for differences in the quark fragmentation process between data and simulated events.

Differences between data and MC have been measured using $D_s^+ \rightarrow \phi \pi^+$ decays, according corrections were applied. The influence of combinatorial background has been studied and considered in the systematics. Background from $\phi$ mesons produced in $D$ or $B$ decays results in a further correction in MC (to calibrate the $\phi$ production rate) and corresponding systematics. The effect of uncertainties due to finite MC statistics and background estimation on the fit has been studied with toy simulations, and was also included in the final results. Remaining detector effects have been determined with control data samples. Using $D^+ \rightarrow D^0 \pi^+$ and $D^0 \rightarrow K^- \pi^+ \pi^0$ events it has been verified that differences between data and simulated events in the resolution of the variables $q^2$ and $\cos(\theta_e)$ are small compared with other sources of systematic uncertainties. They have been neglected at present.

Using fixed values for the pole masses ($m_A = 2.5$ GeV/c$^2$ and $m_V = 2.1$ GeV/c$^2$), the final fit results including all corrections are:

$$N_S = 12886 \pm 129$$

$$r_V = 1.636 \pm 0.067 \pm 0.038$$

$$r_2 = 0.705 \pm 0.056 \pm 0.029.$$

Keeping $m_V$ fixed, for which there is no sensitivity, and adding $m_A$ as additional free parameter, the fit results in

$$N_S = 12887 \pm 129$$

$$r_V = 1.633 \pm 0.081 \pm 0.068$$

$$r_2 = 0.711 \pm 0.111 \pm 0.096$$

$$m_A = 2.53^{+0.54}_{-0.35} \pm 0.54 \text{ GeV}/c^2.$$

The measurements of the parameters $r_V$ and $r_2$ for the semileptonic decay $D_s^+ \rightarrow \phi e^+ \nu_e$ have an accuracy similar to the one obtained for $D \rightarrow K^* e^+ \nu_e$ decays [21], see Fig. 5.

V. LEPTONIC DECAYS

The purely leptonic decay $D_s^+ \rightarrow \ell^+ \nu_{\ell}$ (the charge conjugate mode is implied throughout this paper) is theoretically a rather clean decay; in the Standard Model (SM), the decay is mediated by a single virtual $W^+$-boson. The decay rate is given by

$$\Gamma(D_s^+ \rightarrow \ell^+ \nu_{\ell}) = \frac{G_F^2}{8\pi} f_{D_s^+}^2 m_{\ell}^2 M_{D_s^+}^2 \left(1 - \frac{m_{\ell}^2}{M_{D_s^+}^2}\right)^2 |V_{cs}|^2,$$

(16)

where $G_F$ is the Fermi coupling constant, $m_{\ell}$ and $M_{D_s^+}$ are the masses of the lepton and of the $D_s$ meson, respectively. $V_{cs}$ is the corresponding CKM-matrix element, while all effects of strong interaction are accounted for in the decay constant $f_{D_s^+}$. Due to helicity conservation, the decay rate is highly suppressed for electrons. Since the detection of $\tau$’s involve additional neutrinos, the muon mode is experimentally the cleanest and most accessible mode.

A. $D_s \rightarrow \mu \nu_{\mu}$ at BABAR

This subsection is an adapted excerpt of BABAR’s publication [22].

BABAR performed a measurement of the ratio $\Gamma(D_s \rightarrow \mu \nu_{\mu})/\Gamma(D_s \rightarrow \phi \pi)$ and the decay constant $f_{D_s}$, based on a total integrated luminosity of 230.2 fb$^{-1}$.

In order to measure $D_s^+ \rightarrow \mu^+ \nu_{\mu}$, the decay chain $D_s^{*-} \rightarrow \gamma D_s^+, D_s^+ \rightarrow \mu^+ \nu_{\mu}$ is reconstructed from $D_s^+$ mesons produced in the hard fragmentation of continuum $e^-\bar{e}$ events. The subsequent decay results in a photon, a high-momentum $D_s^+$ and daughter muon and neutrino, lying mostly in the same hemisphere of the event. Signal candidates are required to lie in the recoil of a fully reconstructed $D^0$, $D^+$, $D_s^+$, or $D_s^{*-}$ meson (the “tag”) reconstructed in a variety of modes [23] wherein the tag flavor is uniquely determined. To eliminate signal from $B$ decays, the minimum tag momentum is chosen to be close to the kinematic limit for charm mesons arising from $B$ decays.

For each event a single tag candidate is chosen and then used in the subsequent analysis. To pick this tag among multiple candidates within an event (there are 1.2 candidates on average in events with at least one candidate) modes of higher purity are preferred. In events where two tag candidates are reconstructed in the same mode, the quality of the vertex fit of the $D$ meson is used as a secondary criterion. After subtracting combinatorial background there are $5 \times 10^3$ charm tagged events with a muon amongst the recoiling particles.

The signature of the decay $D_s^{*-} \rightarrow \gamma D_s^+$ is a narrow peak in the distribution of the mass difference $\Delta M = M(\mu \nu e) - M(\mu \nu $) at 143.5 MeV/c$^2$. The $D_s^+$ signal is reconstructed from a muon and a photon candidate in the recoil of the tag. Muons are identified as non-showering tracks penetrating the IFR. The muon must have a momentum of at least 1.2 GeV/c in the
center-of-mass (CM) frame and have a charge consistent with the tag flavor. Clusters of energy in the EMC not associated with charged tracks and exceeding an EMC energy of 0.115 GeV are identified as photon candidates.

The CM missing energy ($E^*_{\text{miss}}$) and momentum ($\vec{p}^*_{\text{miss}}$) are calculated from the four-momenta of the incoming $e^+e^-$, the tag four-momentum, and the four-momenta of all remaining tracks and photons in the event. The energy of the charged particles that do not belong to the tag is calculated from the track momentum under a pion mass hypothesis. Assigning a mass according to the most likely particle hypothesis has negligible effect on the missing energy resolution.

The neutrino CM four-momentum ($p^\nu_\mu = (|p^\nu_\mu|, \vec{p}^\nu_\mu)$) is estimated from the muon CM four-momentum ($\vec{p}^\nu_{\mu}$) and $\vec{p}^*_{\text{miss}}$, using a technique adopted from Ref. [23]. The difference $|p^\nu_{\mu} - \vec{p}^\nu_{\mu}|$ is minimized, while the invariant mass of the neutrino-muon pair is required to be the known mass of the $D^+_s[24]$. The muon CM four-momentum ($p^\nu_{\mu}$) is combined with $\vec{p}^\nu_{\mu}$ to form the $D^+_s$ candidate. The $D^+_s$ candidate is then combined with a photon candidate to form the $D^*_s$ candidate. The selection requirements on $E^*_{\text{miss}}$, $p_{\text{corr}}$, and other variables have been optimized to maximize the significance $s/\sqrt{s+b}$, where $s$ and $b$ are the signal and background yields expected in the data set.

One class of background are events $e^+e^- \to f\bar{f}$, where $f = u, d, s, b, or \tau$, which do not contain a real charm tag. The contribution of these events is estimated from data using the tag sidebands. In addition there are events $e^+e^- \to \pi\pi$ where the tag is incorrectly reconstructed. Although these events potentially contain the signal decay, they are also subtracted using the tag sidebands. These two sources amount to \( \approx 42\% \) of the background. A second class of background events (\( \approx 26\% \)) are correctly tagged $\pi\pi$ events with the recoil muon coming from a semileptonic charm decay or from $\tau^+ \to \mu^+\nu_\mu\bar{\nu}_\tau$. This includes events $D^+_s \to \gamma D^+_s \to \gamma\tau^+\nu_\tau$, $\tau^+ \to \mu^+\nu_\mu\bar{\nu}_\tau$. To estimate the size and shape of this background contribution, the analysis is repeated, substituting a well-identified electron for the muon. Except for a small phase-space correction, the widths of weak charm decays into muons and electrons are assumed to be equal. QED effects such as bremsstrahlung ($e^+ \to \gamma e^+$) energy losses and photon conversion ($\gamma \to e^+e^-$), where the muon equivalents have a much lower rate, are explicitly removed. In particular, bremsstrahlung photons found in the vicinity of an electron track are combined with the track. The small number of events with an electron from a converted photon that survive the selection are suppressed by a photon conversion veto, using the vertex and the known radial distribution of the material in the detector. The muon selection efficiency as a function of momentum and direction is measured using $e^+e^- \to \mu^+\mu^-\gamma$ events, while radiative Bhabha events are used to quantify the electron efficiency. The ratio of muon to electron efficiencies is applied as a weight to each electron event. The remaining backgrounds are estimated from simulation.

Events that pass the signal selection are grouped into four sets, depending on whether the tag lies in the signal region or the sideband regions, and on whether the lepton is a muon or an electron (Fig. 6). For each lepton type the sideband $\Delta M$ distribution is subtracted. The electron distribution, scaled by the relative phase-space factor (0.97) appropriate to semileptonic charm meson decays and leptonic $\tau$ decays is then subtracted from the muon distribution. The resulting $\Delta M$ distribution is fitted with a function ($N_{\text{Sig}}f_{\text{Sig}} + N_{\text{Bkgd}}f_{\text{Bkgd}}(\Delta M)$, where $f_{\text{Sig}}$ and $f_{\text{Bkgd}}$ describe the simulated signal and background $\Delta M$ distributions. The function $f_{\text{Sig}}$ is a double Gaussian distribution. The function $f_{\text{Bkgd}}$ consists of a double and a single Gaussian distribution describing the two peaking background components, and a function [22] describing the flat background component. The relative sizes of the background components, along with all parameters except $N_{\text{Sig}}$ and $N_{\text{Bkgd}}$ are fixed to the values estimated from simulation. The $\chi^2$ fit yields $N_{\text{Sig}} = 489 \pm 55(stat)$ signal events and has a fit probability of 8.9% (Fig. 7).

The branching fraction of $D^+_s \to \mu^+\nu_\mu$ cannot be determined directly, since the production rate of $D^+_s(\mu^+\nu_\mu)$ mesons in $\pi\pi$ fragmentation is unknown. Instead the partial width ratio $\Gamma(D^+_s \to \mu^+\nu_\mu)/\Gamma(D^+_s \to \phi\pi^+)$ is measured by reconstructing $D^+_s \to \gamma D^+_s \to \gamma\phi\pi^+$ decays. The $D^+_s \to \mu^+\nu_\mu$ branching fraction is evaluated using the measured branching fraction for $D^+_s \to \phi\pi^+$.

Candidate $\phi$ mesons are reconstructed from two kaons of opposite charge. The $\phi$ candidates are combined with charged pions to form $D^+_s$ meson candidates. Both times a geometrically constrained fit is employed, and a minimum requirement on the
fit quality is made. The φ and the $D_s^+$ candidate masses must lie within 2σ of their nominal values, obtained from fits to simulated events and data. Photon candidates are then combined with the $D_s^+$ to form $D_s^+$ candidates. The same requirements on the CM photon energy and $D_s^+$ momentum as in the $D_s^+$ → $\mu^+\nu_\mu$ signal selection are made. Data events that pass the selection are grouped into two sets: the tag signal and sideband regions. After the tag sideband has been subtracted from the tag signal $\Delta M$ distribution, the remaining distribution is fitted with the product of $N_{\phi}\phi$ and $N_{Bkgd}Bkgd$ $(\Delta M)$, where $f_{\phi}$ is a double Gaussian, describing the simulated $D_s^+$ → $\gamma D_s^+$ → $\gamma \phi \pi^+$ signal, and $f_{\phi Bkgd}$ consists of a broad Gaussian centered at 70 MeV/c$^2$ and a function \cite{24} describing the simulated background $\Delta M$ distributions. The Gaussian describes the background $D_s^+$ → $\pi^0 D_s^*$ → $\pi^0 \phi \pi^+$ where the photon candidate originates from the $\pi^0$. The relative sizes of the background components, along with all parameters except $N_{\phi}$, $N_{\phi Bkgd}$, and the mean of the peak are fixed to the values estimated from simulation. The $\chi^2$ fit yields $N_{\phi} = 2093 \pm 99$ events and has a probability of 25.0% (Fig. 7). From simulation 48 ± 23 events $D_s^+$ → $\gamma D_s^*$ → $\gamma f_0(980)(K^+ K^-)\pi^+$ are expected to contribute to the signal, where the error is mostly from the uncertainty in the $D_s^*$ → $f_0(980)(K^+ K^-)\pi^+$ branching ratio.

Precise knowledge of the efficiency of reconstructing the tag is not important, since it mostly cancels in the calculation of the partial width ratio. However, the presence of two charged kaons in $D_s^+$ → $\phi \pi^+$ events leads to an increased number of random tag candidates, compared to $D_s^+$ → $\mu^+\nu_\mu$ events, which decreases the chances that the correct tag is picked. The size of the correction for this effect to the efficiency ratio ($\epsilon_{\phi}/\epsilon_{Sig}$) is determined to be -1.4% in simulated events.

To measure the effect of a difference between the $D_s^+$ momentum spectrum in simulated and data events, $D_s^+$ → $\gamma D_s^*$ → $\gamma \phi \pi^+$ events are selected in data with the $D_s^+$ momentum correction removed. The sample is purified by requiring the CM momentum of the charged pion to be at least 0.8 GeV/c. The efficiency-corrected $D_s^+$ momentum distribution in data is compared to that of $D_s^*$ in simulated $D_s^+$ → $\gamma D_s^*$ → $\gamma \phi \pi^+$ events. A harder momentum spectrum is observed in data. The detection efficiencies for signal and $D_s^*$ → $\gamma D_s^*$ → $\gamma \phi \pi^+$ events are re-evaluated after weighting simulated events to match the $D_s^*$ momentum distribution measured in data. The correction to the efficiency ratio is +1.5%.

With both corrections applied, the partial width ratio is determined to be $\Gamma_{\mu\nu}/\Gamma_{\phi\pi} = (N/\epsilon)_{Sig}/(N/\epsilon)_{Bkgd} B(\phi \rightarrow K^+ K^-) = 0.143 \pm 0.018 (stat)$, with $B(\phi \rightarrow K^+ K^-) = 49.1%$ \cite{24}.

A detailed discussion of the systematics can be found in \cite{24}. Using the $\bar{B}\bar{A}\bar{B}$ average for the branching ratio $B(D_s^* \rightarrow \phi \pi^+) = (4.71 \pm 0.46)%$ \cite{24}, the branching fraction $B(D_s^* \rightarrow \mu^+\nu_\mu) = (6.74 \pm 0.83 \pm 0.26 \pm 0.66) \times 10^{-3}$ and the decay constant $f_D = (283 \pm 17 \pm 7 \pm 14)$ MeV are obtained. The first and second errors are statistical and systematic, respectively; the third is the uncertainty from $B(D_s^* \rightarrow \phi \pi^+)$.}

B. $D_s \rightarrow \mu\nu$, at Belle

The Belle analysis uses data corresponding to 548 fb$^{-1}$ to study the decay $D_s^+ \rightarrow \mu^+\nu_\mu$, using the full-reconstruction recoil method first established in the study of semileptonic $D$ mesons described above.

This analysis uses fully reconstructed events of the type $e^+ e^- \rightarrow D_s^* D^{\pm,0} K^{\pm,0} X$, where $X$ can be any number of additional pions from fragmentation, and up to one photon. \cite{31} The tag side consists of a $D$- and a $K$ meson (in any charge combination) while the signal side is a $D_s^*$ meson decaying to $D_s\gamma$. Recon-
structing the tag side, and allowing any possible set of particles in \( X \), the signal side is reconstructed in the recoil, using the known beam momentum.

Tracks are detected with the CDC and the SVD. They are required to have at least one associated hit in the SVD and an impact parameter with respect to the interaction point in the radial direction of less than 2 cm and in the beam direction of less than 4 cm. Tracks are also required to have momenta in the laboratory frame greater than 92 MeV/c. A likelihood ratio for a given track to be a kaon or pion, which is required to be larger than 50%, is obtained by utilising specific ionisation energy loss measurements made with the CDC, light yield measurements from the ACC, and time-of-flight information from the TOF. Lepton candidates are required to have momentum in the lab frame larger than 500 MeV/c. For electron identification we use position, cluster energy, shower shape in the ECL, combined with track momentum and \( dE/dx \) measurements in the CDC and hits in the ACC. For muon identification, we extrapolate the CDC track to the KLM and compare the overlap the CDC track to the KLM and compare the primary \( K \) meson decay into (\( K_L \) or \( K_S \)). The number is reduced by requiring the presence of a photon that is consistent with \( p_T \) and \( m_T \) of the nominal \( K^0 \) mass. Neutral kaon candidates are reconstructed using photon pairs with invariant mass within \( \pm 10 \) MeV/c^2 of the nominal \( \pi^0 \) mass. Neutral kaon candidates are reconstructed using charged pion pairs within \( \pm 30 \) MeV/c^2 of the nominal \( K^0 \) mass.

Tag-side \( D \) mesons (both charged and neutral) are reconstructed in \( D \rightarrow K\pi \) with \( n = 1, 2, 3 \) (total branching fraction \( \approx 25\% \)). Mass windows have been optimised for each channel separately, and a combined mass- and vertex fit (requiring a confidence level greater than 0.1%) is applied to the \( D \) meson to improve the momentum resolution. \( D_s^* \)-candidates are not directly reconstructed, but searched for in the recoil of \( D K \), using the known beam momentum, by applying a mass window cut of \( \pm 150 \) MeV/c^2 around the nominal \( D_s^* \) mass. Since at this point in the reconstruction \( X \) can be any set of remaining pions and photons, there are usually a large number of combinatorial possibilities. This number is reduced by requiring the presence of a photon that is consistent with the decay \( D_s^* \rightarrow D_s \gamma \) where the \( D_s \) has its nominal mass within a mass window of \( \pm 150 \) MeV/c^2. Further selection criteria are applied on the momentum of the primary \( K \) meson in the \( e^+e^- \) rest frame, \( p_X \), which should be smaller than 2 GeV/c; the momentum of the \( D \) meson in the \( e^+e^- \) rest frame, \( p_D \), should be larger than 2 GeV/c; the momentum of the \( D_s \) meson in the \( e^+e^- \) rest frame, \( p_{D_s} \), is required to be larger than 3 GeV/c and the energy of the photon in \( D_s^* \rightarrow D_s \gamma \), \( E_\gamma \), in the lab frame, is required to be larger than 150 MeV/c^2, irrespective of \( \theta_\gamma \). To further improve the recoil momentum resolution, inverse \( k \) mass-constrained vertex fits are then performed for the \( D_s^* \) and \( D_s \), requiring a confidence level greater than 1%. After all these selections are applied, the average number of combinatorial reconstruction possibilities is \( \approx 2 \) per event. The sample is further divided into a right- (RS) and wrong-sign (WS) part, depending on the relative charges of the primary \( K \) meson, the \( D \) meson and that of the \( K \) meson the \( D \) meson decays into (\( K_D \)), compared to the charge of the \( D_s^* \) meson, which is fixed by the total charge of the \( X \) assuming overall charge conservation for the event.

Within this sample of \( D_s \)-tags, decays of the type \( D_s \rightarrow \mu^+\nu \) are selected by requiring another charged track that is identified as a muon and has the same charge as the \( D_s \) candidate. No additional charged particles are allowed in the event, and additional energy corresponding to neutral particles is required to be smaller than \( 1/n \) GeV where \( n \) is the number of additional neutral particles. After these selections, in almost all cases only one combinatorial reconstruction possibility remains. Figure 9 shows the mass spectra for \( D_s \)-tags and neutrino candidates.

\( n_X \) is defined as the number of primary particles in the event, where primary means that the particle is not a daughter of any particle reconstructed in the event. The minimal value for \( n_X \) is three corresponding to a \( e^+e^- \rightarrow D_s^*DK \) event without any further particles from fragmentation. The upper limit for \( n_X \) is determined by the reconstruction efficiency; Monte-Carlo (MC) shows that the number of reconstructed signal events is negligible for \( n_X > 10 \). As the efficiency very sensitively depends on \( n_X \), it is crucial to use MC that correctly reflects the \( n_X \) distribution observed in data. Unfortunately, the details of fragmentation processes are not very well understood, and standard MC [12] shows notable differences compared to data. Furthermore, the true (generated) \( n_X \) value differs from the reconstructed \( n_X^R \), as particles can be lost or wrongly assigned. Thus the measured (reconstructed) \( n_X^R \) distribution has to be deconvoluted so that the analysis can be done in bins of \( n_X^R \) to avoid bias in the results.

To extract the number of \( D_s \)-tags as function of \( n_X^T \) in data from the background, 2-dimensional histograms in \( n_X^R \) (ranging from 3 to 8) and the invariant recoil mass \( m_{D_s} \) are used. The signal shapes for different values of \( n_X^T \) (ranging from 3 to 9 [33]) of the signal are modelled with generic MC (GMC) [13], which has been filtered on the generator-level for events of the type \( e^+e^- \rightarrow D_s^*DKX \). The weights of these components, \( w_{D_s}^i, i = 1..6 \), are free parameters in the fit to data. As a model for the background in RS, the WS data sample is used. Each slice of \( n_X^T \) is fitted separately, adding another 6 free parameters. Since the WS-sample contains some signal (\( \approx 10\% \) of the RS signal), these signal components (in slices of \( n_X^T \)) are also included in the fit as independent parameters (yielding a negative weight to compensate for the WS signal present in the data shapes). The fit is performed simultaneously with all these free parameters. As a crosscheck, the fit has also been performed us-
ing generic MC RS-sample backgrounds, which gives a negligible change in the results. A further cross-
check involved dividing the MC sample randomly into two halves, using the shapes of the first half to fit the
signal in the second. The result as function of \( n_X^T \), normalized to the amount of signal in the first
half, fits to a flat line as 0.990 \pm 0.046, which agrees well with the expectation of 1. The total number of recon-
structed \( D_s \)-tags in data is calculated as

\[
N_{D_s}^{\text{rec}} = \sum_{i=1}^{6} w_i^{D_s} N_{D_s}^{\text{GMC},i},
\]

(17)

where \( N_{D_s}^{\text{GMC},i} \) represents the total number of reconstructed filtered GMC events that were generated in the
i-th bin of \( n_X^T \) (regardless of the reconstructed \( n_X^R \)) and \( w_i^{D_s} \) is the fitted weight of this component.

To fit the number of \( D_s \rightarrow \mu \nu \)-events as function of \( n_X^T \), 2-dimensional histograms in \( n_X^R \) and the missing
mass squared \( m_{\text{miss}}^2 \) are used. The shape of the signal is modelled with signal MC. As MC studies show, the
background under the \( \mu \nu \)-signal peak consists primarily of non-\( D_s \) decays, semileptonic \( D_s \) decays (where the additional hadrons have low momenta and remain undetected) and leptonic \( \tau \) decays (where the \( \tau \) decays to a muon and two neutrinos). Hadronic \( D_s \) decays (with one hadron misidentified as muon) are a rather small background component. Except for hadronic decays, which are negligible, all backgrounds are common to the \( e\nu \)-mode, which is suppressed by a factor of \( O(10^{-5}) \). Thus, the \( e\nu \)-sample provides a good model of the \( \mu \nu \) background that has to be corrected only for
kinematical and efficiency differences. Including this

FIG. 9: Belle: Invariant mass spectrum of \( D_s \)-tags (left plot), and missing mass spectrum squared of \( D_s \rightarrow \mu \nu \) candidates (right plot) in data with the selection criteria described in the text (points with error bars show statistical errors only). The red shaded areas show the fitted background, the cyan shaded bands show the fit with systematic uncertainties. The vertical lines indicate the signal regions.

corrected shape in the fit, the total number of fitted \( \mu \nu \)-events in data is given by

\[
N_{\mu\nu}^{\text{rec}} = \sum_{i=1}^{6} w_i^{\mu\nu} N_{\mu\nu}^{\text{SMC},i},
\]

(18)

where \( N_{\mu\nu}^{\text{SMC},i} \) represents the total number of reconstructed signal MC events that were generated in the
i-th bin of \( n_X^T \) (regardless of the reconstructed \( n_X^R \)) and \( w_i^{\mu\nu} \) is the fitted weight of this component.

The numerical result for \( N_{\mu\nu}^{\text{rec}} \) is 32100 \pm 870(stat) \pm 1210(syst), that for \( N_{\mu\nu}^{\text{rec}} \) is 169 \pm 16(stat) \pm 8(syst). The statistical uncertainties are due to statistics in the data signal, the systematic uncertainties due to statistics of the data background samples and those of the MC samples used. These errors include the non-negligible correlations between the \( n_X^T \) bins.

As the branching fraction of \( D_s \rightarrow \mu \nu \) used for the generation of generic MC is known, the branching fraction in data can be determined using the following formula:

\[
B(D_s \rightarrow \mu \nu) = \frac{N_{\mu\nu}^{\text{rec}}}{N_{\mu\nu}^{\text{GMCexp}}} \cdot B_{\text{generated}}(D_s \rightarrow \mu \nu),
\]

(19)

where \( B_{\text{generated}}(D_s \rightarrow \mu \nu) = 0.0051 \) and \( N_{\mu\nu}^{\text{GMCexp}} \) is the number of reconstructed \( \mu \nu \)-events in the generic
MC, weighted according to the fit to data, i.e.

\[
N_{\mu\nu}^{\text{GMCexp}} = \sum_{i=1}^{6} w_i^{\mu\nu} N_{\mu\nu}^{\text{GMC},i}.
\]

(20)

where \( N_{\mu\nu}^{\text{GMC},i} \) represents the total number of reconstructed \( \mu \nu \)-events filtered from GMC that were gen-
erated in the i-th bin of \( n^T_X \) (regardless of the reconstructed \( n^T_X \)).

Figure 10 shows the branching fraction determined in bins of \( n^T_X \) (using correlated fit results). The result is stable within errors in \( n^T_X \); note that the errors shown for the \( n_X \) bins are the total errors, including correlation. The final result is:

\[
\mathcal{B}(D_s \to \mu \nu) = (6.44 \pm 0.76(\text{stat}) \pm 0.52(\text{syst})) \cdot 10^{-3} \tag{21}
\]

The statistical error reflects the statistics of the signal sample. The systematic error is dominated by statistical uncertainties due to background samples from data and MC samples (0.29) and the statistical uncertainty on \( N_{\mu \mu}^{\text{MC}} \cdot n^T_X \) (0.41). Since the branching fraction is determined by calculating a ratio of the signal yield to the number of \( D_s \)-tags, systematics in the reconstruction of the tag side cancel; the only remaining systematics are due to the tracking and identification of the muon, which have been estimated as 2%, contributing 0.13 to the total systematic error. As a crosscheck, also the branching fraction for \( n^T_X \leq 6 \) has been determined as \((6.54 \pm 0.76(\text{stat}) \pm 0.54(\text{syst})) \cdot 10^{-3}\), which agrees nicely with the result given above including all available \( n^T_X \) bins.

The decay constant \( f_{D_s} \), using Eqn. (10) and recent values from PDG \[16\] yields

\[
f_{D_s} = 275 \pm 16(\text{stat}) \pm 12(\text{syst})\text{MeV} \tag{22}
\]

VI. SUMMARY: OVERVIEW AND COMPARISON

Studies in the charm sector are notoriously difficult for experiments running at much higher than threshold energy; still both Belle and \(\text{B}A\text{B}AR\) present an interesting variety of results on (semi)leptonic charm decays. While Belle concentrated on fully reconstructed (and consequently tagged) events, \(\text{B}A\text{B}AR\) preferred methods with partially reconstructed events, and uses tag information only for its \(D_s \to \mu \nu \mu\) analysis.

The advantage of Belle’s approach is a very effective background suppression and an excellent neutrino momentum resolution which can compete with results achieved at experiments operating at threshold energy like Cleo-c \[27\]. It also allows absolute measurements, by normalization to the number of \(D_s\) tags.

However, \(\text{B}A\text{B}AR\)’s approach is significantly more efficient in terms of event statistics. In the case of the \(D^0 \to K\ell\nu\) analysis, despite higher backgrounds, this results in a much better accuracy of measurements. For the \(D_s \to \mu \nu \mu\) analysis, the advantage of higher statistics is more or less equalled by the disadvantage of larger backgrounds, which eventually gives somewhat larger errors than Belle’s.

In any case, it is a valuable crosscheck to have rather different experimental approaches at different experiments. Table III summarizes all results by Belle and \(\text{B}A\text{B}AR\).

Within the semileptonic channels discussed in this review, only the mode \(D^0 \to K\ell\nu\) has been studied by both \(\text{B}A\text{B}AR\) and Belle, and can be compared. The measured branching fractions agree well within errors, larger differences are seen in the form factor measurements. However, these differences do not exceed 1.3\(\sigma\), and could be due to the more complicated systematics of the fits involved. The results are much more precise than those of previous experiments, and also well compatible with other recent results \[28, 29\]. Obviously, the untagged, partial reconstruction of \(\text{B}A\text{B}AR\) has clearly smaller errors, even though it suffers a large uncertainty due to the normalizing channel used, which is the dominating part of its systematic error. Belle does an absolute measurement, but also is clearly limited by systematics. Thus in neither case a significant further improvement of the measurements can be expected with more data accumulated, without also further developing the experimental methods.

The other mode where results can be compared is \(D_s \to \mu \nu \mu\). The results agree very well with each other. Here Belle profils from its full reconstruction method, and has somewhat, but not dramatically smaller errors. Both results are well compatible with other recent results, and still compatible with theoretical predictions, which tend to be lower by \(\approx 2.5\sigma_{\exp}\), but bear some uncertainties as well \[30\]. In both experiments, statistical and systematic error are of the same order. Considering the fact that part of the
TABLE III: Overview of results obtained by BaBar and Belle experiments, compared to theoretical expectations (where available); errors are statistical (first) and systematic (second).

<table>
<thead>
<tr>
<th>decay mode</th>
<th>parameter</th>
<th>prediction</th>
<th>BaBar result</th>
<th>Belle result</th>
<th>difference</th>
</tr>
</thead>
<tbody>
<tr>
<td>$D \rightarrow K e\nu_e$</td>
<td>BF (%)</td>
<td>$3.522 \pm 0.027 \pm 0.079$</td>
<td>$3.45 \pm 0.10 \pm 0.19$</td>
<td>n/a</td>
<td>0.3\sigma</td>
</tr>
<tr>
<td></td>
<td>z-expansion, $a_0$</td>
<td>$2.98 \pm 0.01 \pm 0.04$</td>
<td>n/a</td>
<td>n/a</td>
<td>0.3\sigma</td>
</tr>
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<td></td>
<td>z-expansion, $a_1$</td>
<td>$-2.5 \pm 0.2 \pm 0.2$</td>
<td>n/a</td>
<td>n/a</td>
<td>0.3\sigma</td>
</tr>
<tr>
<td></td>
<td>z-expansion, $a_2$</td>
<td>$0.6 \pm 6.0 \pm 5.0$</td>
<td>n/a</td>
<td>n/a</td>
<td>0.3\sigma</td>
</tr>
<tr>
<td></td>
<td>$m_{\text{pole}}$ (GeV/c²)</td>
<td>$2.112 (= m_{D^*}^3)$</td>
<td>$1.884 \pm 0.012 \pm 0.015$</td>
<td>$1.82 \pm 0.04 \pm 0.03$</td>
<td>1.2\sigma</td>
</tr>
<tr>
<td></td>
<td></td>
<td>$\alpha$</td>
<td>$0.50 \pm 0.04$</td>
<td>$0.377 \pm 0.023 \pm 0.029$</td>
<td>$0.52 \pm 0.08 \pm 0.06$</td>
</tr>
<tr>
<td></td>
<td></td>
<td>$f_+(0)$</td>
<td>$0.73(3)(7)$</td>
<td>$0.727 \pm 0.007 \pm 0.009$</td>
<td>$0.695 \pm 0.007 \pm 0.022$</td>
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<td>$D \rightarrow K \mu \nu_\mu$</td>
<td>BF (%)</td>
<td>n/a</td>
<td>$3.45 \pm 0.10 \pm 0.21$</td>
<td>n/a</td>
<td>0.3\sigma</td>
</tr>
<tr>
<td></td>
<td>$m_{\text{pole}}$ (GeV/c²)</td>
<td>$2.112 (= m_{D^*}^3)$</td>
<td>n/a</td>
<td>included in $K e\nu_e$</td>
<td>n/a</td>
</tr>
<tr>
<td></td>
<td></td>
<td>$\alpha$</td>
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<td>n/a</td>
<td>results shown</td>
</tr>
<tr>
<td></td>
<td></td>
<td>$f_+(0)$</td>
<td>$0.73(3)(7)$</td>
<td>n/a</td>
<td>above</td>
</tr>
<tr>
<td>$D \rightarrow \pi e\nu_e$</td>
<td>BF (%)</td>
<td>n/a</td>
<td>$0.279 \pm 0.027 \pm 0.016$</td>
<td>n/a</td>
<td>0.3\sigma</td>
</tr>
<tr>
<td></td>
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<td>$\alpha$</td>
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<td>n/a</td>
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<td></td>
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<td>$0.64(3)(6)$</td>
<td>n/a</td>
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<tr>
<td>$D \rightarrow \pi \mu \nu_\mu$</td>
<td>BF (%)</td>
<td>n/a</td>
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<td>n/a</td>
<td>0.3\sigma</td>
</tr>
<tr>
<td></td>
<td>$m_{\text{pole}}$ (GeV/c²)</td>
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<td>n/a</td>
<td>included in $\pi e\nu_e$</td>
<td>n/a</td>
</tr>
<tr>
<td></td>
<td></td>
<td>$\alpha$</td>
<td>$0.44 \pm 0.04$</td>
<td>n/a</td>
<td>results shown</td>
</tr>
<tr>
<td></td>
<td></td>
<td>$f_+(0)$</td>
<td>$0.64(3)(6)$</td>
<td>n/a</td>
<td>above</td>
</tr>
</tbody>
</table>

systematic error is due to the size of control samples which will get larger with more statistics, there is some room for further improvements once the full data sets of the experiments are available.

Acknowledgments

We wish to thank Arantza Oyanguren, Justine Serrano and Paul Jackson from BABAR for providing material to compile this talk. Text excerpts from BABAR publications covering the described analyses have been used to compile this review. As far as my own collaboration Belle is concerned, the thanks go to the KEKB group for the excellent operation of the accelerator, the KEK cryogenics group for the efficient operation of the solenoid, and the KEK computer group and the National Institute of Informatics for valuable computing and Super-SINET network support. We acknowledge support from the Ministry of Education, Culture, Sports, Science, and Technology of Japan and the Japan Society for the Promotion of Science; the Australian Research Council and the Australian Department of Education, Science and Training; the National Science Foundation of China under contract No. 10575109 and 10775142; the Department of Science and Technology of India; the BK21 program of the Ministry of Education of Korea, the CHEP SRC program and Basic Research program (grant No. R01-2005-000-1089-0) of the Korea Science and Engineering Foundation, and the Pure Basic Research Group program of the Korea Research Foundation; the Polish State Committee for Scientific Research; the Ministry of Education and Science of the Russian Federation and the Russian Federal Agency for Atomic Energy; the Slovenian Research Agency; the Swiss National Science Foundation; the National Science Council and the Ministry of Education of Taiwan; and the U.S. Department of Energy.
[31] It has been found that events with additional kaons or more than one photon have a poor signal/background ratio and have been therefore excluded.
[32] The fit is called inverse since it uses information from the mother and sister particles, rather than information about daughter particles as is usually the case.
[33] The upper limit of 9 is chosen because the bin $n_X^R = 9$ has some overlap with $n_X^B \leq 8$. 
We present recent results from the CLEO Collaboration on leptonic decay rates of $D$ and $D_s$ near $c\bar{c}$ production threshold. From these decay rates, we extract the decay constants, $f_{D^+} = (222.6 \pm 16.7^{+2.8}_{-2.4})$ MeV, $f_{D_s^+} = (274 \pm 10 \pm 5)$ MeV, and the ratio $f_{D_s^+}/f_{D^+} = 1.23 \pm 0.11 \pm 0.03$.

\[ \Gamma(D_{(s)}^+ \rightarrow l^+\nu) = \frac{G_F^2}{8\pi} f_{D_{(s)}^+}^2 m_t^2 M_{D_{(s)}^+} \left(1 - \frac{m_l}{M_{D_{(s)}^+}}\right)^2|V_{cd(s)}|^2, \]

where $G_F$ is the Fermi constant, $M_{D_{(s)}^+}$ is the $D^+$ ($D_{s}^+$) mass, $m_t$ is the final state lepton’s mass, and $V_{cd(s)}$ are the relevant CKM matrix elements. The quantity $f_{D}$ is the decay constant and represents the hadronic matrix element discussed above. A critical input to $B$ mixing and CP violation measurements in the $B$ sector is the $B$ decay constant, $f_B$. Due to the difficulty in measuring $f_B$, we take the value from theory, usually lattice QCD. To have confidence in the theoretical number, a stringent theoretical test is provided by a precision measurement of the $D$ decay constant, $f_D$. Such a measurement provides a critical test of any theory or model that makes predictions for decay constants.

The CLEO experiment, operating near $c\bar{c}$ threshold, is well positioned to measure these decay rates, and hence $f_{D^+}$ and $f_{D_{s}^+}$. Charge conjugate finals states are implied throughout unless otherwise noted.

2. Measurement of $f_{D}^+$

To measure $f_{D}^+$ [9], we use 281 pb$^{-1}$ of data collected at the $\psi(3770)$ resonance. The proximity to the production threshold implies that the $\psi(3770)$ decays to $D\bar{D}$ with no additional particles. We exploit this clean final state, along with the hermiticity of the detector to reconstruct the neutrino from the missing momentum in the event. Specifically, we fully reconstruct a $D^-$ meson (the tag) in six hadronic final states, comprising $N_{tag} = 158,354 \pm 496$ tags.

To search for $D^+ \rightarrow \mu^+\nu$, we require a single extra charged particle with an energy deposition in the central calorimetry (CC), $E_{CC} < 300$ MeV, and veto events with any additional photon candidates with energy larger than 250 MeV. From this subsample of events, we compute the square of the missing mass ($MM^2$) recoiling against the $D^-\mu^+$ system. For $D^+ \rightarrow \mu^+\nu$, a peak at zero is obtained with a resolution of $\sigma(MM^2) \sim 0.025$ GeV$^2$. The $MM^2$ distribution is shown in Fig. 1 for data. The clear excess near zero is the $D^+ \rightarrow \mu^+\nu$ signal. Some $D^+ \rightarrow K_{S,L}\pi^+$ events pass the selection requirements and appear as a prominent, but well-separated peak near $MM^2 \sim 0.25$ GeV$^2$.

Figure 1: Missing-mass squared distribution for $D^+ \rightarrow \mu^+\nu$ candidates. The peak near zero corresponds to signal events, and is expanded in the inset. The larger peak at $MM^2 \sim 0.25$ GeV$^2$ corresponds to $D^+ \rightarrow K_{S,L}\pi^+$ events which pass the selection requirements.
The branching fraction is computed using:

\[
\mathcal{B} = \frac{N_{\text{cand}} - N_{\text{back}}}{N_{\text{tag}}\epsilon\text{CC}},
\]

where \(N_{\text{cand}} = 50\) is the number of signal candidates in the region \(|MM^2| < 0.050\ \text{GeV}^2\), \(N_{\text{back}} = 2.81 \pm 0.30 \pm 0.27\) is the expected number of background events, \(N_{\text{tag}} = 158,354 \pm 496\) is the number of fully-reconstructed \(D^-\) tags, \(\epsilon_{\mu} = 69.4\%\) is the efficiency for reconstructing and identifying the muon, and \(\epsilon_{\text{CC}} = 96.1\%\) is the fraction of events that do not have any additional photon candidates with energy larger than 250 MeV. An additional correction of (1.5 \pm 0.4 \pm 0.5)\% is applied to account for the higher efficiency for reconstructing a \(D^-\) tag in \(D^+ \to \mu^+\nu\) events than in generic hadronic events.

The resulting branching fraction is

\[
\mathcal{B}(D^+ \to \mu^+\nu) = (4.40 \pm 0.66^{+0.99}_{-0.12}) \times 10^{-4}.
\]

Using Eq. (3) we determine the decay constant to be:

\[
f_{D^+} = (222.6 \pm 16.7^{+2.8}_{-3.4})\ \text{MeV}.
\]

3. Measurement of \(f_{D^+_s}\)

The measurements of \(f_{D^+_s}\) at CLEO require higher energy running in order to produce the \(D_s\bar{D}_s\) pair. A scan of the energy region from 3970 to 4260 MeV was performed, and it was determined that the optimal energy for \(D_s\) physics was 4170 MeV [13], where \(D_s\bar{D}_s^*\) is dominant, e.g., \(\sigma(D_s\bar{D}_s^*) = (916 \pm 50)\ \text{pb}\) and \(\sigma(D_s^*\bar{D}_s) = (35 \pm 19)\ \text{pb}\). A slight complication with using \(D_s\bar{D}_s^*\) is the additional (~150 MeV) photon(s) from the \(D_s^*\) decay. Two independent analyses have been carried out. The first analysis is similar to the \(D^+ \to \mu^+\nu\) measurement described previously, where, in addition to measuring \(\mathcal{B}(D_s^+ \to \mu^+\nu)\), we also measure \(\mathcal{B}(D_s^+ \to \tau^+\nu)\), where, \(\tau^+ \to \pi^+\nu\bar{\nu}\). In the second analysis, we measure \(\mathcal{B}(D_s^+ \to \tau^+\nu)\), \(\tau^+ \to \nu\bar{n}\nu\bar{u}\).

3.1. Measurement of \(\mathcal{B}(D_s^+ \to (\mu^+\tau^+)\nu)\) using Missing Mass

We use 314 \(\text{pb}^{-1}\) of data collected at \(E_{\text{cm}} = 4170\ \text{MeV}\) for this analysis. We search for final states consistent with either \(D_s^+ \to \mu^+\nu\) or \(D_s^+ \to \tau^+\nu\). The branching fraction is obtained from:

\[
\mathcal{B} = \frac{N_{\text{cand}} - N_{\text{back}}}{N_{\text{tag}}\epsilon}
\]

where \(N_{\text{tag}}\) is the number of reconstructed \(D_s\bar{D}_s^*\) events and \(\epsilon\) is the efficiency for reconstruction and identification of the \(\mu^+\) for \(D_s^+ \to \mu^+\nu\), or the \(\pi^+\) for \(D_s^+ \to \tau^+\nu\), \(\tau^+ \to \pi^+\nu\bar{\nu}\). We therefore absorb the full reconstruction of the \(D_s^*\) into the denominator, and do not rely on Monte Carlo simulation for the efficiency of the ~100 MeV photon.

To determine \(N_{\text{tag}}\), we first fully reconstruct a hadronic \(D_s^-\) tag in eight tag modes, from which we obtain \(31,302 \pm 472\) \(D_s^+\) tags. To identify \(D_s\bar{D}_s^*\) events, we combine a \(D_s^*\) tag with any additional photon candidate in the event and form the missing-mass squared \((MM^2)\) recoiling against the \(\gamma\) system, \(MM^2 = (E_{\text{cm}} - E_{D_s} - E_{\gamma})^2 - (\vec{p}_{D_s} - \vec{p}_{\gamma})^2\).

This quantity peaks at \(M_{D_s^*}^2\), regardless of whether the photon came from the \(D_s^*\) (the tag) or from the \(D_s^*\). The distribution of \(MM^2\) is shown in Fig 2 for all eight tag modes combined. A fit to this distribution yields \(18645 \pm 426 \pm 1081\) \(D_s\bar{D}_s\) events within ±2.5 standard deviations of \(M(D_s)\).

![Figure 2: Square of the missing mass recoiling against a \(\gamma\) candidate.](image-url)
GeV$^2$, and falls smoothly to zero at $MM^2 = -0.05$ and extends to $MM^2 \sim 0.8$ GeV$^2$. We thus define signal samples as follows. (i)-$\mu$: For $D_\pm^+ \to \mu^+ \nu$, we require an energy deposition, $E^\text{trk}_{\text{CC}} < 300$ MeV, and $|MM^2| < 0.05$ GeV$^2$. For $D^+_s \to \tau^+ \nu$, $\tau^+ \to \pi^+ \nu \bar{\nu}$, we define two subsamples - (i)-$\tau$: $E^\text{trk}_{\text{CC}} < 300$ MeV and $0.05 < MM^2 < 0.20$ GeV$^2$, and (ii)-$\tau$: $E^\text{trk}_{\text{CC}} > 300$ MeV and $-0.05 < MM^2 < 0.20$ GeV$^2$. The upper cutoff in $MM^2$ is to avoid background from $D^+_s \to K^0 \pi^+$. We also consider a third sample, (iii)-$\tau$, for $D^+_s \to e^+ \nu$ by requiring the track’s energy deposition to be consistent with its momentum and $|MM^2| < 0.050$ GeV$^2$.

The $MM^2$ distributions are shown in Fig. 3 where cases (i)-$\mu$ and (i)-$\tau$ are combined. In the $D^+_s \to \mu^+ \nu$ signal region of $|MM^2| < 0.05$ GeV$^2$, we find 92 events with an expected background of 3.5 $\pm$ 1.4 events. This sample is mostly of $D^+_s \to \mu^+ \nu$, with some cross-feed from $D^+_s \to \tau^+ \nu$, $\tau^+ \to \pi^+ \nu \bar{\nu}$. We evaluate the branching fraction using the number of $\mu^+ \nu$ events, $N_{\mu \nu}=92$, in the signal region:

$$N_{\mu \nu} = N_{\text{det}} - N_{\text{back}}$$

$$= N_{\text{tag}} \cdot \epsilon \cdot \mathcal{B}(D^+_s \to \mu^+ \nu)$$

$$+ \epsilon'' \mathcal{B}(D^+_s \to \tau^+ \nu, \tau^+ \to \pi^+ \nu \bar{\nu}),$$

where $\epsilon = 80.1\%$ is the efficiency of reconstructing the charged particle in a $D^+_s \to \mu^+ \nu$ event, and includes the veto on events with additional photons with $E > 250$ MeV. The quantity, $\epsilon' = 91.4\%$, is the product of the muon identification efficiency (99\%) and the $MM^2 < 0.05$ GeV$^2$ requirement (92.3\%). The cross-feed efficiency, $\epsilon''=7.9\%$, which is the product of the efficiency of the pion depositing less than 300 MeV in the CC (60\%) and the $MM^2 < 0.05$ GeV$^2$ requirement (13.2\%). One can re-express $\mathcal{B}(D^+_s \to \tau^+ \nu, \tau^+ \to \pi^+ \nu \bar{\nu})$ as:

$$\mathcal{B}(D^+_s \to \tau^+ \nu, \tau^+ \to \pi^+ \nu \bar{\nu}) = R \cdot \mathcal{B}(\tau^+ \to \pi^+ \nu \bar{\nu})$$

$$\times \mathcal{B}(D^+_s \to \mu^+ \nu) = 1.059 \cdot \mathcal{B}(D^+_s \to \mu^+ \nu)$$

where we use the Standard Model ratio for $R$:

$$R = \frac{\Gamma(D^+_s \to \tau^+ \nu)}{\Gamma(D^+_s \to \mu^+ \nu)} = \left(\frac{m_{\tau^+}}{m_{\mu^+}}\right)^2 \left(1 - \frac{m_{\tau^+}^2}{m_{\mu^+}^2}\right)^2 = 9.72.$$  

We thus find:

$$\mathcal{B}(D^+_s \to \mu^+ \nu) = (0.594 \pm 0.066 \pm 0.031)\%,$$

where the 5.2\% systematic error is dominated by the 5\% uncertainty on $N_{\text{tag}}$.

We also compute $\mathcal{B}(D^+_s \to \tau^+ \nu, \tau^+ \to \pi^+ \nu \bar{\nu})$ using cases (i)-$\tau$ and (ii)-$\tau$. For these two cases, we find yields of 31 and 25 events, and expected backgrounds of $3.5^{+1.7}_{-1.1}$ and $5.1^{+1.6}_{-1.3}$ events, respectively. The fraction of $D^+_s \to \tau^+ \nu, \tau^+ \to \pi^+ \nu \bar{\nu}$ events in the respective $MM^2$ regions are 32\% and 45\%. We thus find:

$$\mathcal{B}(D^+_s \to \tau^+ \nu, \tau^+ \to \pi^+ \nu \bar{\nu}) = (8.0 \pm 1.3 \pm 0.4)\%.$$

With the measured branching fractions, $\mathcal{B}(D^+_s \to \mu^+ \nu)$ and $\mathcal{B}(D^+_s \to \tau^+ \nu, \tau^+ \to \pi^+ \nu \bar{\nu})$, we measure the ratio of partial widths, $R=13.4 \pm 2.6 \pm 0.2$ (defined in Eq. 5), which is consistent with the Standard Model value of 9.72.

We may improve on the precision of $\mathcal{B}(D^+_s \to \mu^+ \nu)$ by combining the $D^+_s \to \mu^+ \nu$ and $D^+_s \to \tau^+ \nu, \tau^+ \to \pi^+ \nu \bar{\nu}$
π⁺ν̅ν candidates. We can still use Eq. 7 except ε’ and ε’’ increase from 91.4% and 7.9% to 96.2% and 45.2%, respectively. We thus find an effective branching fraction:

\[ \mathcal{B}^{\text{eff}}(D_s^+ \to \mu^+ν) = (0.638 ± 0.059 ± 0.033)\% . \quad (11) \]

Again, the dominant systematic uncertainty (5%) is on the number of \( D_s^+ \) tags.

The \( MM^2 \) distribution for all selected \( D_s^+ \to \mu^+ν \) and \( D_s^+ \to τ^+ν, \ τ^+ \to π^+ν̅ν \) candidates is shown in Fig. 4. Overlayed is a curve that represents the expected shape, normalized to the event yield in the \( MM^2 \) region below 0.2 GeV². We find good agreement between the shape in data and expectations.

![Figure 4: Square of the missing mass recoiling against \( γD_s^0, μ^+ν \) (or \( π^+ν \)) candidates. The curve is the expected shape from simulation, normalized to the number of events with \( MM^2 < 0.2 \) GeV.](image)

We also search for the decay \( D_s^+ \to e^+ν \). The helicity suppression in this decay is much larger, and the expected rate is \( ∼50,000 \) times smaller than in \( D_s^+ \to μ^+ν \). We find no \( D_s^+ \to e^+ν \) candidates and set the upper limit, \( \mathcal{B}(D_s^+ \to e^+ν) < 1.3 \times 10^{-4} \) at the 90% confidence level.

Using the more precise value for \( \mathcal{B}(D_s^+ \to μ^+ν) \) from Eq. 11, we compute the decay constant, \( f_{D_s^+} \):

\[ f_{D_s^+} = 274 ± 13 ± 7 \text{ MeV} \quad (12) \]

Combining this with our previous result for \( f_{D_s^+} = (222.6 ± 16.7^{+3.4}_{-3.0}) \text{ MeV} \) we determine the ratio:

\[ \frac{f_{D_s^+}}{f_{D_s^+}} = 1.23 ± 0.11 ± 0.04 . \quad (13) \]

### 4. Measurement of \( D_s^+ \to τ^+ν, \ τ^+ \to e^+ν̅ν \)

In the second measurement of \( \mathcal{B}(D_s^+ \to τ^+ν) \), we use 298 pb⁻¹ of data collected at \( E_{cm} = 4170 \) MeV. We utilize the decay \( τ^+ \to e^+ν̅ν \), where we benefit from the large value of \( \mathcal{B}(τ^+ \to e^+ν̅ν) \sim 18\% \), and the excellent electron identification capabilities of the CLEO-c detector. We fully reconstruct the three hadronic decay channels: \( D_s^+ \to φπ^0K^- \) and \( K_S^0K^- \). Charged hadrons are identified using standard selection criteria [12], and the intermediate resonances, \( φ \to K^+K^- \), \( K^0 \to K^-π^+ \), and \( K_S^0 \to π^-π^- \), are required to have an invariant mass within \( ±10 \) MeV, \( ±75 \) MeV and \( ±12 \) MeV of their known values [12]. Signal candidates are required to a reconstructed invariant mass, \( M(D_s) \) within \( ±20 \) MeV of the known \( D_s \) mass \( (m_{D_s}) \). We also define sideband regions, \( 35 < |M(D_s) - m_{D_s}| < 55 \) MeV, to study the combinatorial background. The invariant mass distributions of the three \( D_s^- \) tag channels are shown in Fig. 5.

![Figure 5: Invariant mass distributions of \( D_s^- \) candidates from data. The points are data, the solid line is a fit, and the dashed line is the background.](image)

To ensure we have \( D_s\bar{D}_s^* \), we compute the mass recoiling against the reconstructed \( D_s \), and require it to be within \( ±55 \) MeV of the \( D_s^* \) mass [12]. We then select the subset of events with a single additional charged track with \( p > 200 \) MeV that has opposite charge to the \( D_s \) tag and is consistent with being a positron. The discriminating variable we use to identify \( D_s^+ \to τ^+ν, \ τ^+ \to e^+ν̅ν \) is \( E_{\text{extra}} \), the total energy remaining in the calorimeter after all showers associated with the tag and the positron are removed. In signal events, the only additional particles beyond the \( D_s \) tag and the positron are the two neutrinos and either a photon from \( D_s^* \to γD_s \), or a π⁰ from \( D_s^* \to π⁰D_s \). Kinematically, these photons populate the energy regions from 114-170 MeV (for \( γD_s^+ \)) and 39-117 MeV (from \( π⁰D_s^+ \)).

The distribution of \( E_{\text{extra}} \) in data is shown in Fig. 6. The large excess at low values of \( E_{\text{extra}} \)
is the $D^+_s \rightarrow \tau^+\nu$, $\tau^+ \rightarrow e^+\nu\bar{\nu}$ signal. The broad background which peaks near 1 GeV is predominantly semi-leptonic decays, such as $D^+_s \rightarrow \phi e^+\nu$, $\eta e^+\nu$, $\eta' e^+\nu$ and $K^* e^+\nu$ and $K^{*0} e^+\nu$. The Cabibbo-suppressed decay, $K^0_L e^+\nu$, produces a small peaking component in the signal region. The shape of this background is taken from Monte Carlo simulation, and is normalized to our measured rate for $D^+_s \rightarrow K^0_L e^+\nu$ of $B(D^+_s \rightarrow K^0_L e^+\nu) = (0.27 \pm 0.10)%$. We choose the signal region as $E_{\text{extra}} < 400$ MeV, which is chosen based on optimizing the signal significance. The expected non-peaking background in the signal region is estimated by scaling the number of data events with $E_{\text{extra}} > 600$ MeV by the MC ratio of events in the sideband ($E_{\text{extra}} > 600$ MeV) to signal region ($E_{\text{extra}}^{\text{MC}} < 400$ MeV). The yields of $D^-_s$ tags and $D^+_s \rightarrow \tau^+\nu$, $\tau^+ \rightarrow e^+\nu\bar{\nu}$ signal events are shown in Table I. The scale factor, $s$ shown in Table I, is a correction to account for slight differences in the expected number of background events in the signal and sideband regions. Using the efficiency to reconstruct the final state, $D^+_s \rightarrow \tau^+\nu$, $\tau^+ \rightarrow e^+\nu\bar{\nu}$ of $(71.4 \pm 0.4)%$ and the $B(\tau^+ \rightarrow e^+\nu\bar{\nu}) = (17.84 \pm 0.05)%$, we find:

$$B(D^+_s \rightarrow \tau\nu) = (6.24 \pm 0.71 \pm 0.36)\%.$$  \hspace{1cm} (14)

The 5.8% systematic uncertainty is dominated by the 4.3% contribution from the simulation of $K^0_L$ showering in the calorimeter.

When this result is combined with the result in Eq. 12 we obtain:

$$f_{D^+_s} = 274 \pm 10 \pm 5 \text{ MeV} \hspace{1cm} (15)$$

5. Summary

We have presented measurements of the branching fractions $B(D^+ \rightarrow \mu^+\nu$, $D^+_s \rightarrow \mu^+\nu$ and $D^+_s \rightarrow \tau^+\nu$, $\tau^+ \rightarrow \tau^+\nu\bar{\nu}$ with the CLEO-c detector. The results are the most precise measurements of these leptonic decay rates to date. Using Eq IV we extract the decay constants:

$$f_{D^+} = (222.6 \pm 16.7^{+2.8}_{-3.4}) \text{ MeV} \hspace{1cm} (16)$$

$$f_{D^+_s} = (274 \pm 10 \pm 5) \text{ MeV} \hspace{1cm} (17)$$

$$f_{D^+_s}/f_{D^+} = 1.23 \pm 0.11 \pm 0.03 \hspace{1cm} (18)$$

Our measurement of $f_{D^+_s}$ is consistent with and significantly more precise than the recent measurement by BaBar [14]. The only other measurement of $f_{D^+_s}$ was reported by BES based on 1 signal candidate. Recent lattice QCD predictions [15, 16] of both $f_{D^+_s}$ and $f_{D^+_s}$ are typically $\sim 10\%$ lower than our measurements, whereas the ratio of $f_{D^+_s}/f_{D^+}$ is in good agreement with our measurement.

We gratefully acknowledge the effort of the CESR staff in providing us with excellent luminosity and running conditions. We also thank the National Science Foundation for support of this research.

References


Table I Summary of $D_s^-$ tagged events (yield, background from sidebands, sidebands scale factor ($s$), and sideband-subtracted yield), and $D_s^+ \rightarrow \tau^+ \nu$, $\tau^+ \rightarrow e^+ \nu \bar{\nu}$ events (yield, background from $D_s^-$ sidebands, background from $D_s^+$ semileptonic decays, and sideband-subtracted yield).

<table>
<thead>
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<th>Mode</th>
<th>$D_s^-$ Tags</th>
<th>$D_s^+ \rightarrow \tau^+ \nu$, $\tau^+ \rightarrow e^+ \nu \bar{\nu}$</th>
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<td></td>
<td>Yield</td>
<td>Back</td>
</tr>
<tr>
<td>$D_s^- \rightarrow \phi \pi^-$</td>
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</tr>
<tr>
<td>$D_s^- \rightarrow K^- K^{*0}$</td>
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<td>3618</td>
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<tr>
<td>Total</td>
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</table>

Weak Charm Decays with Lattice QCD

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In this paper I review the status of lattice QCD calculations of $D$ and $D_s$ meson decay constants and of $D$ meson semileptonic decay form factors. I restrict my discussion to results obtained from simulations with $n_f = 2 + 1$ sea quarks.

1. Introduction and Motivation

Lattice QCD is the only systematically improvable calculational tool we have for quantitatively understanding nonperturbative QCD effects. Accurate theoretical calculations of nonperturbative QCD effects are essential for the experimental flavor physics program. One set of goals of the experimental program are accurate determinations of the CKM matrix elements. This is illustrated for the weak decay process $D \rightarrow K \nu$. The experimentally measured (differential) decay rate can be written as

$$\frac{d\Gamma}{dq^2} = (\text{known})|V_{cs}|^2 f_+^2(q^2)$$

where $f_+(q^2)$ is one of the hadronic form factors which parameterize the hadronic matrix element for this process, $\langle K | V_{\mu}^\dagger | D \rangle$. Hence, to determine $|V_{cs}|$ from experimental measurements, we need a theoretical calculation of the form factor with matching precision. Another set of goals is to constrain beyond the standard model theories and to search for new physics signals. This effort complements the experiments at the high energy frontier. Accurate theoretical calculations are again essential.

Since lattice QCD calculations are complicated and time consuming, a third important goal specifically of the charm physics experiments is to test lattice QCD methods. For example, we can use Eq. 1 to determine the form factors from experimental measurements after taking $|V_{cs}|$ from other sources. These tests are important to establish lattice methods for the $B$ meson system, where the CKM matrix elements are less well known, and where input from lattice QCD is essential. The leptonic and semileptonic $D$ meson decays discussed in this talk are ideal for this. They are not expected to be sensitive to new physics, and the corresponding hadronic matrix elements are straightforward to calculate. Once established, lattice QCD together with the experimental measurements can then be used to improve the determinations of the CKM angles $V_{cd}$ and $V_{cs}$. All of these goals require accurate measurements and calculations.

1.1. Introduction to Lattice QCD

In lattice field theory, the space-time continuum is replaced by a discrete lattice. (For reviews of lattice QCD see Ref. [1] ) This implies that derivatives are replaced by discrete differences, which in turn introduces discretisation errors into physical quantities. These errors generally vanish with a positive power of the lattice spacing ($a$).

Nonperturbative calculations in lattice QCD can be performed using Monte Carlo methods. Lattice artifacts can be removed by reducing the lattice spacing used in numerical calculations. However, the computational cost increases as $1/a^2$ (keeping the other parameters fixed). Alternatively, one can reduce discretisation errors by adding higher-dimensional operators to the action. This is called improvement. With improved actions the computational effort needed to perform reliable lattice QCD calculations can potentially be significantly reduced. This idea is behind much of the important progress made in lattice QCD calculations in recent years and has been an increasing part of research in lattice field theory.

The main obstacle for obtaining quantitative results (at the few percent level) from numerical simulations of lattice QCD has always been the computational effort associated with the proper inclusion of sea quark effects. Several years ago, substantial progress was made on this problem, in large part due to the development of an improved staggered fermion action [2]. For the first time, computationally efficient lattice simulations with realistic sea quark effects have become possible.

1.2. Light Quark Methods

The simplest lattice quark action replaces the covariant derivative in the continuum action by a discrete difference operator. This so-called naive action suffers from the doubling problem. For every continuum quark flavor, it has 15 additional unphysical flavors, called tastes. The staggered quark action combines four of these tastes into one Dirac field, by staggering the quark fields on a hypercube. This leaves four unphysical flavors (tastes). This action suffers from large $O(a^2)$ lattice spacing artifacts due to taste
changing interactions. The Asqtad action is an improved staggered action where all tree-level discretisation errors are removed [2]. Its leading lattice spacing errors are therefore of \( O(\alpha_s a^2) \) and greatly reduced compared to the original staggered action. The Asqtad action is the computationally most efficient light quark action available. However, in order to use it for the sea quarks in numerical simulations, the unphysical flavors must be removed. The sea quarks are present in the fermion determinant of the path integral. To simulate two degenerate light (up and down) sea quarks, the MILC collaboration simply takes the square root of the light quark fermion determinant. For the strange sea quark, they take the fourth root of the determinant. This procedure is still somewhat controversial, but there is a growing body of evidence that its effects are controllable and disappear in the continuum limit [3]. The Asqtad action with the square root trick has been extensively tested in numerical simulations, most prominently in Ref. [4].

The HISQ (highly improved staggered quark) action is another version of an improved staggered action [3]. Like the Asqtad action, it removes all tree-level \( O(a^2) \) errors. The \( O(\alpha_s a^2) \) errors in the Asqtad action are rather large, due to taste changing interactions which appear at one-loop order. The HISQ action reduces the \( O(\alpha_s a^2) \) taste-changing effects by roughly a factor of three over the Asqtad action. The HISQ action has not yet been used to generate \( n_f = 2 + 1 \) sea quark ensembles. Its computational cost is expected to be about a factor of two larger than the Asqtad action.

Other light quark methods include the Wilson action and its improvements [6], Domain Wall Fermions [7] and Overlap fermions [8], with increasing computational cost. The Wilson action solves the doubling problem by adding a dimension five operator which breaks chiral symmetry. Domain Wall Fermions solve the doubling problem by adding a fifth dimension, while keeping chiral symmetry almost exact. Overlap fermions have exact chiral symmetry, but a complicated operator structure.

1.3. Heavy Quark Methods

On the lattice, heavy quarks with \( am_Q \) large, are best treated within an effective field theory framework (NRQCD or HQET). One can start with an effective field theory, and discretise it as in Ref. [11], for example. Alternatively, one can start with a relativistic lattice action and analyze its mass dependent discretisation errors using effective field theories. The charm quark is too light for a straightforward implementation of the former approach, so we will focus on the latter.

The Fermilab approach [12] starts with the improved relativistic Wilson action [6] and the observation that the Wilson action has the same heavy quark limit as QCD. With a simple prescription, the Wilson action can be used for heavy quarks without errors that grow with the heavy quark mass, \((am_Q)^n\). This approach can be used for both charm and and beauty quarks. With the improved Wilson action, the leading discretisation errors are \( O(\alpha_s \Lambda/m_Q) \) and \( O(\Lambda/m_Q)^2 \).

The HISQ action is so highly improved that it can be used for charm quarks with an additional tuning of a parameter in the action, provided that the lattice spacing is small enough [3]. The leading mass dependent discretisation errors are formally of order \( O(\alpha_s(m_Q)^2) \) and \( O(m_Q)^4 \).

1.4. Systematic Errors

The most important sources of systematic error in lattice QCD calculations are sea quark effects; using unphysically large masses for the up and down quarks; discretisation effects; finite volume effects; and renormalisation effects.

In order to be phenomenologically relevant, a lattice QCD calculation must use gauge configurations that include the effects of three light sea quarks. Since the masses of the up and down quarks are generally large, are \( \rho \), and \( \pi \) mesons. Furthermore, it can be extended to include the leading light quark discretisation errors. Indeed, this has been done for the taste changing interactions of the Asqtad action and is called staggered ChPT (SChPT) [10].
The ChPT extrapolations are a significant but controllable source of systematic error. In order to keep this error at the few percent level or less, one needs to include simulations with a range of light sea masses, keeping \( m_l < m_s / 2 \). The lattice QCD calculations described here include light sea quarks with masses in the range \( m_l = 1/10 m_s - 1/2 m_s \). Hence, the lightest light sea quark masses are only a factor of 2–3 away from their physical value.

As described in the previous sections, the lattice actions give rise to discretisation errors. They can usually be estimated \textit{a priori} using power counting arguments. However, even with improved actions, it is important to study and possibly remove these errors by repeating the calculation at several lattice spacings.

### 1.5. Simulation Parameters

The simulation parameters of the \( n_f = 2 + 1 \) sea quark ensembles generated by the MILC collaboration using the Asqtad action (with the square root trick) are listed below and are shown in Figure 1. Each ensemble contains between 450–800 configurations. The ensembles contain one sea quark with a mass near the strange quark mass, \( m_s \), and two degenerate light sea quarks with masses, \( m_l \).

- \( a = 0.15 \text{ fm}; \ m_l = 0.1 m_s, 0.2 m_s, 0.4 m_s, 0.6 m_s \).
- \( a = 0.12 \text{ fm}; \ m_l = 0.125 m_s, 0.25 m_s, 0.5 m_s, 0.75 m_s \).
- \( a = 0.09 \text{ fm}; \ m_l = 0.1 m_s, 0.2 m_s, 0.4 m_s \).

![Figure 1: Simulation parameters for the MILC ensembles with \( n_f = 2 + 1 \). showing \( m_l/m_s \) vs. lattice spacing \( a \) (red squares). The physical point is at \( m_l/m_s = 1/25 \) (pink burst).](image)

### 2. Semileptonic \( D \) Meson Form Factors

The semileptonic decays \( D \to K(\pi)\ell \nu \) are mediated by weak vector currents. The hadronic matrix elements for semileptonic decays are parameterized in terms of form factors. In our case there are two form factors, conventionally \( f_1(q^2) \) and \( f_0(q^2) \). The form factors are functions of the virtual \( W \) boson momentum transfer, \( q^2 \), or, equivalently, the recoil momentum of the daughter meson. This introduces additional lattice spacing errors:

\[
(K|V_\mu|D)^\text{lat} = (K|V_\mu|D)^\text{cont} + O(\alpha p_K)^n
\]

Hence, discretisation errors are smallest, when \( p_K \) is small and \( q^2 \approx q_{\text{max}}^2 = (m_D - m_K)^2 \).

The finite lattice volume provides an infrared cut-off, and therefore a minimum value for finite momentum, \( p_{\text{min}} = \frac{2 \pi}{a} \). Lattice three-momenta can be written in terms of \( p_{\text{min}} \) as \( \vec{p} = p_{\text{min}}(n_x, n_y, n_z) \), where \( n_x, n_y, n_z \) are integers. For example, for \( a = 0.1 \text{ fm} \), \( L = 20, p_{\text{min}} = 620 \text{ MeV} \).

To date, the only lattice results for semileptonic \( D \) meson form factors with \( n_f = 2 + 1 \) are from the Fermilab Lattice and MILC collaborations [13]. They use the MILC \( a = 0.12 \text{ fm} \) lattices with light sea quark masses in the range \( m_l = 1/8 m_s - 3/4 m_s \), the Asqtad action for the light valence quarks and the Fermilab action for the charm quark. Staggered chiral perturbation theory is used to extrapolate to the physical light quark masses and to remove the leading discretisation errors due to taste violations.

Figure 2 shows a comparison of the lattice QCD result for the normalization \( f_0^D(0) \) for \( D \to K\ell\nu \) with experimental determinations (where \( V_{cs} \) is taken from other sources). The results are in very good agreement; however, the lattice QCD result has much larger errors than the experimental determinations. The comparison between lattice theory and experiment for \( f_0^D(0) \) is similar [14].

The shape of the form factor can also be determined in lattice QCD. However, in Ref. [13] the form factors were calculated at a few values of recoil momentum. Then the BK [10] parameterisation was used to determine the \( q^2 \) dependence of the form factors. Since the errors increase with recoil momentum, the shape of the form factors is fixed mainly from the form factors near \( q_{\text{max}}^2 \) and from using the BK parameterisation [13]. The lattice QCD result appeared before the new measurements by the FOCUS [17] and Belle [18] collaborations were announced, so it is one of a very few lattice QCD predictions. Figure 3 [13] shows a comparison of the lattice prediction for the \( q^2 \) dependence with experimental data from the Belle collaboration [18]. The agreement is excellent. However, a quantitative comparison between the BK shape parameter determined from experiment and lattice theory is difficult to interpret, as eloquently argued by Richard.
Figure 2: $f_+^K(0)$ in comparison from Ref. [14].

Figure 3: The shape of $f_K^+ (q^2)$ in comparison. The lattice prediction for the shape is from Ref. [12].

Hill [19]. A model independent parameterisation of the shape based on the $z$-expansion would avoid this difficulty [20]. The $z$-expansion is being used by the Fermilab Lattice and MILC (FNAL/MILC) collaboration to parameterize the $q^2$ dependence of the form factors for $B \to Dl\nu$ [23]. This works quite well, as shown in Figure 4. Any new lattice analysis of semileptonic $D$ meson decay form factors will (should) adopt the $z$-expansion to determine the shape.

A number of additional improvements are possible in future calculations. Twisted boundary conditions can be used to adjust the lattice momenta to arbitrarily small values [22], which would improve the shape determination. Double ratio methods similar to what has been developed for lattice studies of $B \to Dl\nu$ [23] can be adapted to semileptonic $D$ meson decays [24]. This may lead to reduced statistical errors as well as improvements of some of the systematic errors.

3. Leptonic Decay Constants $f_D$ and $f_{Ds}$

Charm leptonic decays provide another important test of lattice QCD. The lattice methods for calculating decay constants in the charm and beauty meson systems are the same. Indeed, with the Fermilab approach one uses the same heavy quark action in both systems and the heavy quark discretisation errors are expected to be larger for $D$ mesons than for $B$ mesons.

There are now results from two groups (FNAL/MILC and HPQCD). Both use the MILC ensembles at $a = 0.09$ fm, 0.12 fm, 0.15 fm. The first FNAL/MILC results came out in 2005 [25], just days before CLEO-c announced its first precise determination of $f_D^+$ [26]; the two results were in good agreement.

The HPQCD collaboration announced their results for decay constants with much reduced errors this summer [27] and FNAL/MILC presented updated results at the Lattice 2007 conference [28], also with reduced errors. The new FNAL/MILC analysis was done “blind”, where an overall unknown offset was added to the lattice data. The final results were unblinded shortly before they were presented at the Lattice 2007 conference, making this the first (intentionally) blind lattice analysis. Table III compares the main
features of the two calculations. More details about
the HPQCD and FNAL/MILC calculations, including
discussion of the error analysis and plots of chiral and
lattice spacing extrapolations can be found in Refs. [29]
and [28] respectively. The FNAL/MILC analysis includes
more lattice ensembles, more valence quark masses per ensemble, and uses staggered chiral
perturbation theory (Staggered ChPT) to remove the
leading light quark discretisation errors. The HPQCD
collaboration considers only the case \( m_q = m_l \), where
\( m_q \) denotes the light valence quark mass and \( m_l \)
denotes the light sea quark mass. They use continuum
ChPT with generic \( O(a^2) \) terms added in the chiral
fits.

The main difference between the two calculations
is the valence quark actions. The HPQCD collabora-
tion uses the HISQ action for all (charm, strange and
light) valence quarks, whereas the FNAL/MILC collabora-
tion uses the Fermilab action for the charm quarks and
the Asqtad action for the strange and light
valence quarks. Since the HISQ action is more
improved than the Fermilab action, the HPQCD result
has much smaller heavy quark discretisation errors.
This is the main reason for the difference in the total
errors between the two results.

Table II compares the error budgets for the 2005
FNAL/MILC calculation with the FNAL/MILC Lattice
2007 one. The error reduction is mainly due to
including three MILC ensembles at \( a = 0.09 \) fm (and
8-12 different valence masses). This reduces the heavy
quark and light quark discretisation errors, and better
constrains the staggered ChPT.

Figures 5, 6, and 7 compare the lattice results for
\( f_{D^+}, f_{D_s}, \) and \( f_{D_s}/f_{D^+} \), respectively, to the corre-
sponding experimental averages. The experimental
averages are from Ref. [14]. The new CLEO-c result
\( f_{D_s} = 275 \pm 10 \pm 5 \) presented at this conference by
Steven Blusk [30], is very similar to Ref. [14].

The FNAL/MILC results agree with the experimen-
tal averages at the one sigma level. The HPQCD results
agree very well with the FNAL/MILC results.
There is a hint of disagreement between the HPQCD
result for \( f_{D_s} \) and the experimental average at the
two sigma level. However, the experimental determina-
tions of the decay constants must assume a value for
the CKM angle \( V_{cs} \) from other sources. We are
approaching a level of precision, where tests of lattice
QCD should be performed on CKM free quanti-
ties such as the ratio of semileptonic to leptonic decay
rates suggested in Ref. [31].

4. Conclusions and Outlook

With the generation of the MILC ensembles, the
stakes for lattice QCD calculations have risen. We are
now able to calculate the simplest quantities to a few
percent accuracy. As always, repetition is desirable
to test different lattice methods against each other.
To date, all lattice calculations that include realistic
sea quark effects use the MILC ensembles with rooted
Asqtad sea quarks. As mentioned in section 1.2 the
Asqtad action carries the smallest computational cost
of any light quark action. Nevertheless, recently other
collaborations have started to generate ensembles with
different sea quark actions. An overview is given in
Figure 8. It shows that the other ensembles are be-
ing generated at similar values of lattice spacing and
light quark masses as the MILC ensembles. The MILC
collaboration continues to generate new ensembles at
even smaller lattice spacings. They are also generating
additional configurations for the existing ensembles to
further reduce statistical errors. As in experiment, in
lattice QCD smaller statistical errors give better con-

![Figure 5: Comparison of lattice QCD results for \( f_{D^+} \) with experiment.](image1)

![Figure 6: Comparison of lattice QCD results for \( f_{D_s} \) with experiment.](image2)
Table I Comparison of the main features of the HPQCD and Fermilab Lattice/MILC calculations.

<table>
<thead>
<tr>
<th>FNAL/MILC</th>
<th>HPQCD</th>
</tr>
</thead>
<tbody>
<tr>
<td>Fermilab action for charm quark</td>
<td>HISQ action for charm and light valence quarks</td>
</tr>
<tr>
<td>Asqtad action for strange and light valence</td>
<td></td>
</tr>
<tr>
<td>a (fm)</td>
<td>m_l/m_s sea quark</td>
</tr>
<tr>
<td>0.09</td>
<td>1/10, 1/5, 2/5</td>
</tr>
<tr>
<td>0.12</td>
<td>1/8, 1/4, 1/2, 3/4</td>
</tr>
<tr>
<td>0.15</td>
<td>1/10, 1/5, 2/5, 3/5</td>
</tr>
</tbody>
</table>

8−12 light valence quark masses per ensemble 1 valence quark mass/ensemble, m_{valence} = m_{sea}

Partial nonperturbative renormalisation Nonperturbative renormalisation from PCAC
Staggered ChPT fits to all valence Continuum ChPT +O(a^2) terms fit to all ensembles together
Blind analysis for Lattice 2007

Table II Comparison of the error budget of the 2005 FNAL/MILC results with the Lattice 2007 results. Numbers are given in percent.

<table>
<thead>
<tr>
<th></th>
<th></th>
<th></th>
</tr>
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<tbody>
<tr>
<td>source</td>
<td>f_D^- f_D^+</td>
<td>f_D^- f_D^+</td>
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<tr>
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<td>light quark + Chiral fits</td>
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<tr>
<td>inputs (a, m_c, m_s)</td>
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<td>0.6</td>
</tr>
<tr>
<td>higher order PT</td>
<td>1.3</td>
<td>1.3</td>
</tr>
<tr>
<td>+ other small sources (finite volume, . . .)</td>
<td></td>
<td></td>
</tr>
<tr>
<td>total systematic</td>
<td>8.5</td>
<td>5.4</td>
</tr>
</tbody>
</table>

Figure 7: Comparison of lattice QCD results for f_D^-/f_D^+ with experiment.

Acknowledgments

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References

Figure 8: Simulation parameters for ensembles with $n_f = 2 + 1$ showing $m_l/m_s$ vs. lattice spacing $a$. Filled symbols denote existing ensembles. Unfilled symbols denote ensembles which are currently being generated or planned. Red squares: MILC [9], blue diamonds: RBC/UKQCD [32], purple left triangles: BMW (improved Wilson) [33], pink right triangles: PACS-CS (nonperturbatively improved Wilson) [34], green circles: JLQCD (Overlap) [35]. The physical point is at $m_l/m_s = 1/25$ (pink burst).


[34] N. Ukita et al. [PACS-CS Collaboration], arXiv:0710.3462 [hep-lat].

Precision Lattice Calculation of D and Ds decay constants

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We present a determination of the decay constants of the D and Ds mesons from lattice QCD, each with a total error of about 2%, approximately a factor of three better than previous calculations. We have been able to achieve this through the use of a highly improved discretization of QCD for charm quarks, coupled to gauge configurations generated by the MILC collaboration that include the full effect of sea u, d, and s quarks. We have results for a range of u/d masses down to $m_u/5$ and three values of the lattice spacing, which allow us to perform accurate continuum and chiral extrapolations. We fix the charm quark mass to give the experimental value of the $\eta_c$ mass, and then a stringent test of our approach is the fact that we obtain correct (and accurate) values for the mass of the D and Ds mesons. We compare $f_D$ and $f_{Ds}$, with $f_K$ and $f_\pi$, and using experiment determine corresponding CKM elements with good precision.

I. INTRODUCTION

Precision calculations in lattice QCD play a crucial role in testing our non-perturbative theoretical tools, by comparing the results of the calculation with precisely measured quantities. In addition accurate calculations of non-perturbative QCD quantities are very important in the extraction of information from analysis of experimental data, for example in the determination of the Cabibbo-Kobayashi-Maskawa (CKM) matrix elements.

This is most clearly seen in the case of "gold-plated" processes, for example the leptonic decay of $D_s$, $D$, $\pi$ and $K$ mesons. In this process the corresponding matrix elements are of order $a^2$, with quark content $\bar{a}b$ (or $\bar{a}b$) annihilating weakly into a W boson, with a width given, up to calculated electromagnetic corrections [1,2], by:

$$\Gamma(P \to l\eta(\gamma)) = \frac{G_F^2|V_{ab}|^2}{8\pi} f_P^2 m_P^2 (1 - \frac{m_P^2}{m_{\eta^2}})^2. \quad (1)$$

$V_{ab}$ is the corresponding element of the CKM matrix, and the decay constant $f_P$ parametrizes the amplitude for W annihilation. By combining a measurement of $\Gamma$ with an accurate calculation of $f_P$ (1) can be used to determine $V_{ab}$. If $V_{ab}$ is known from elsewhere we can use (1) to get a value for $f_P$.

The decay constant $f_P$ is conventionally defined to be a property of the pseudoscalar meson, calculable in QCD without QED effects, and is given by:

$$\langle 0|\bar{\psi}\gamma_{\mu}\gamma_5\psi|P(p)\rangle = f_P p_\mu. \quad (2)$$

The calculation of $f_P$ is a hard experimental problem, which at present can only be done fully with lattice QCD. There are very precise experimental measurements for the leptonic decay rates in the case of the $\pi$ and $K$, and new results are appearing for $D$ and $D_s$, which make the calculations a highly non-trivial test of lattice QCD, and ultimately of QCD itself. This tests are important to give us confidence in similar lattice QCD predictions of matrix elements in B systems, for which experimental results are much harder to obtain.

II. IMPROVED STAGGERED QUARKS

We use HISQ staggered quarks in the valence sector, whereas the sea quarks are ASQTAD staggered quarks with the fourth root trick [3,4,5].

Staggered quark actions suffer the doubling problem: there are four “tastes” (non-physical flavours) of fermions in the spectrum, which couple through taste-changing interactions. These are lattice artifacts of order $a^2$, involving at leading order the exchange of a gluon of momentum $q \approx \pi/a$. Although quite large in the original one-link (Kogut-Susskind) staggered action, such interactions are perturbative for typical values of the lattice spacing, and can be corrected systematically with the Symonzik. By judiciously smearing the gauge field we can remove the coupling between quarks and high momentum gluons.

The most widely used improved staggered action is called ASQTAD, and removes all tree-level $a^2$ discretization errors in the action [6,7,8].

The HISQ (highly improved staggered quarks) staggered Dirac operator involves two levels of smearing with an intermediate projection onto $SU(3)$. It is designed so that, as well as eliminating all tree-level $a^2$ discretization errors, it further reduces the one-loop taste-changing errors (see [9] for a more detailed discussion.) This action has been shown to substantially reduce the errors associated with the taste-changing interactions [8,10,11].

When we put massive quarks on the lattice, the discretization errors grow with the quark mass as powers of $am$. Therefore to obtain small errors we would need $am \ll 1$. For heavy quarks this would require very small lattice spacings. On the other hand, to keep our lattice big enough to accommodate the light degrees of freedom, we need $La \gg m^{-1}$. The fact that we have two very different scales in the problem makes difficult a direct solution. What we can do instead is to take advantage of the fact that $m$ is large, by using an effective field theory (NRQCD, HQET). This program has been very successful for b quarks [12,13,14].
The charm quark is in between the light and heavy mass regime. It is quite light for an easy application of NRQCD, but quite large for the usual relativistic quark actions, \( a m_c \lesssim 1 \). However, if we use a very accurate action (HISQ) and fine enough lattices (fine MILC ensembles), it is possible to get results accurate at the few percent level. A non-relativistic analysis shows that for HISQ charm quarks the largest remaining source of error is due to the quark’s energy, and can be further suppressed by powers of \( v/c \), where \( v \) is the typical velocity of the quark in the system of interest, simply by retuning the overall coefficient of a term called Naik term to impose the correct relativistic dispersion relation \( c^2(p) = 1 \) for low lattice momentum \( p \).

One advantage of the use of a relativistic action is the existence of a partially conserved current, which implies the non-renormalization of the lattice result for \( f_P \). We can extract \( f_P \) from the PCAC relation for zero momentum meson \( P \):

\[
f_P m_P^2 = (m_u + m_b) \langle 0 |\bar{a} \gamma_\mu b| P \rangle
\]

### III. RESULTS

<table>
<thead>
<tr>
<th>Lattice/sea</th>
<th>valence ( u ) mass, ( u ) mass, ( a ) mass, ( a ) mass, ( c ) mass, ( 1 + \epsilon )</th>
<th>( r_1/a )</th>
</tr>
</thead>
<tbody>
<tr>
<td>16(^3) \times 48</td>
<td>0.0194, 0.0484</td>
<td>0.0264, 0.066, 0.85, 0.66</td>
</tr>
<tr>
<td></td>
<td>0.0097, 0.0484</td>
<td>0.0132, 0.066, 0.85, 0.66</td>
</tr>
<tr>
<td>20(^3) \times 64</td>
<td>0.02, 0.05</td>
<td>0.0278, 0.0525, 0.648, 0.79</td>
</tr>
<tr>
<td></td>
<td>0.01, 0.05</td>
<td>0.01365, 0.0546, 0.66, 0.79</td>
</tr>
<tr>
<td>24(^3) \times 64</td>
<td>0.005, 0.05</td>
<td>0.0067, 0.0537, 0.65, 0.79</td>
</tr>
<tr>
<td>28(^3) \times 96</td>
<td>0.0124, 0.031</td>
<td>0.01635, 0.03635, 0.427, 0.885</td>
</tr>
<tr>
<td></td>
<td>0.0062, 0.031</td>
<td>0.00705, 0.0366, 0.43, 0.885</td>
</tr>
</tbody>
</table>

**TABLE I: MILC configurations and mass parameters used for this analysis.** The \( 16^3 \times 48 \) lattices are ‘very coarse’, the \( 20^3 \times 64 \) and the \( 24^3 \times 64 \), ‘coarse’ and the \( 28^3 \times 96 \), ‘fine’. The sea asqtad quark masses \((l = u/d)\) are given in the MILC convention where \( u_0 \) is the plaquette tadpole parameter. Note that the sea s masses on fine and coarse lattices are above the subsequently determined physical value 7. The lattice spacing values in units of \( r_1 \) after ‘smoothing’ are given in the rightmost column 14-18. The third column gives the HISQ valence \( u/d \) and s masses along with the coefficient of the Naik term, \( 1 + \epsilon \), used for \( c \) quarks 9.

We use gluon field configurations including \( 2 + 1 \) flavours of sea quarks generated by the MILC collaboration 13-18. The parameters of the ensembles we have used for both the sea and the valence sectors are in table 11. The lattice results are converted to physical units through the heavy quark potential parameter \( r_1 \), as determined by the MILC collaboration (table 11 10). The physical value of \( r_1 \) is determined from the \( \Upsilon \) spectrum calculated in NRQCD with \( b \) quarks on the same MILC ensembles 10, with the result \( r_1 = 0.321(5) \) fm, \( r_1^{-1} = 0.615(10) \) GeV.

We use multiple precessing random wall sources, which gives a 3-4-fold reduction in statistical errors with respect to conventional local sources.

The mass of the charm quark is fixed by adjusting the mass of the “goldstone “ \( \eta_c \) to its experimental value. The light \( (u/d) \) and strange quark masses are fixed using the experimental values for the masses of \( \pi \) and \( K \). Our results use masses for the \( u \) and \( d \) quarks that are substantially larger (by a factor of around three) than the real ones. In order to get physical answers we extrapolate to the correct \( u/d \) mass using chiral perturbation theory. Once the masses have been thus fixed, there is no remaining freedom to change any parameters, and in particular the results we obtain for the masses of heavy-light mesons are a stringent test of our method. In figure 1 we show the spectrum of charmonium. We obtain an hyperfine splitting of 111(5) MeV (experiment 117(1) MeV) We have made no attempt as yet to optimize the calculation of the excited states.

In addition to the chiral extrapolation, we have systematic errors coming from a variety of sources 10, among them from the finite lattice spacing. Because we have three different lattice spacings and very precise data, we can extrapolate to the continuum limit. This extrapolation is linked to the chiral extrapolation through discretization errors in the light quark action. We therefore perform a simultaneous bayesian fit for both chiral and continuum extrapolations, allowing for expected functional forms in both. We tested the validity of the method by fitting hundreds of fake
FIG. 2: Masses of the $D^+$ and $D_s$ meson as a function of the $u/d$ quark mass in units of the $s$ quark mass at three different values of the lattice spacing. The very coarse results are above the coarse and the fine are the lowest. The lines give the simultaneous chiral fits and the dashed line the continuum extrapolation as described in the text. Our final error bars, including the overall scale uncertainty, are given by the shaded bands. These are offset from the dashed lines by an estimate of electromagnetic, $m_u \neq m_d$ and other systematic corrections to the masses. The experimental results are marked at the physical $m_d/m_s$.

Datasets generated using staggered chiral perturbation theory with random couplings. We fit simultaneously to the masses and the decay constants, that is, we fit $m_\pi, m_K, f_\pi$ and $f_K$ simultaneously, and similarly for $m_D, m_Ds, f_D$ and $fDs$. We present some of the results in figures 2 and 3.

We get an excellent agreement with experiment for the masses: $m_{D_s} = 1.963(5)$ GeV (experiment 1.968 GeV), and $m_D = 1.869(6)$ GeV (experiment 1.869 GeV). Our calculation also reproduces correctly the difference in binding energies between a heavy-heavy ($\eta_c$) and a heavy-light ($m_D$ and $m_{D_s}$) state: $(2m_{D_s} - m_{\eta_c})/(2m_D - m_{\eta_c}) = 1.249(14)$ (experiment 1.260(2)). Our charm quark action is the first one to be accurate enough to do this calculation (which also cannot be done, for example, in potential models.)

We also have agreement with experiment for the light-light decay constants [19]. The result for the ratio is very accurate, $f_K/f_\pi = 1.189(7)$, and shows tiny discretization effects (figure 4). Combining this ratio with experimental leptonic branching fractions [17, 20], we get $V_{us} = 0.2262(13)(4)$, where the first error is theoretical and the second experimental. This gives the unitarity relation $1 - V_{ud}^2 - V_{us}^2 - V_{ub}^2 = 0.0006(8)$.

Our results for the heavy-light decay constants are 4-5 times more accurate than previous lattice QCD results and existing experimental measurements: $f_{D_s} = 241(3)$ MeV, $f_D = 208(4)$ MeV, and a ratio of $f_{D_s}/f_D = 1.162(9)$ (see figure 5). For the double ratio $(f_{D_s}/f_D)/(f_K/f_\pi)$, which is estimated to be close to

FIG. 3: Results for the $D, D_s, K$ and $\pi$ decay constants on very coarse, coarse and fine ensembles, as a function of the $u/d$ quark mass in units of the $s$ quark mass. The chiral fits are performed simultaneously with those of the corresponding meson masses, and the resulting continuum extrapolation curve is given by the dashed line. The shaded band gives our final result. At the left are experimental results from CLEO-c [22, 24] (on the left with the $\tau$ decay result above the $\mu$ decay result above the $\mu$ decay result for $D_s$) and BaBar [23] ($D_s$ only) and from the Particle Data Tables [2] for $K$ and $\pi$. For the $K$ we have updated the result quoted by the PDG to be consistent with their quoted value of $V_{us}$. 

FIG. 4: Ratio of decay constants $f_K/f_\pi$ on very coarse, coarse and fine ensembles, as a function of the $u, d$ quark mass in units of the $s$ quark mass.
tor\n\nthe difference between Data T ables [2].

ond experimental. The result for 4.42(4)(41). The first error is theoretical and the sec-

quark. The high statistical accuracy of our data com-
delivering very precise results on systems with a charm
relativistic action on fine enough lattices is capable of
finite value of the mass of the charm and bottom.

ences are small compared to the absolute masses of
heavy quarks, for states with a ¸quark and systems with
light quarks, we can also obtain a very precise value
through the calculation for both the charm and the

1 from low order chiral perturbation theory [21], we
get a value of 0.977(10).

The experimental leptonic branching rates, together
with CKM matrix elements determined from other
processes (assuming $V_{cs} = V_{ud}$) give $f_{D_s} = 264(17)$
MeV for $\mu$ decays and 310(26) MeV for $\tau$ decay from
CLEO-c [22] and 283(23) MeV from BaBar [23], and
for $f_0$ 223(17) MeV from CLEO-c for $\mu$ decay [24].
Using our results for $f_{D_s}$ and $f_{D_s}/f_D$ and the exper-
imental values from CLEO-c [22] for $\mu$ decay (since
the electromagnetic corrections are well-known in that
case) we can directly determine the corresponding
CKM elements: $V_{cs} = 1.07(1)(7)$ and $V_{cs}/V_{cd} = 4.42(4)(41)$. The first error is theoretical and the sec-
ond experimental. The result for $V_{cs}$ improves on the
direct determination of 0.96(9) given in the Particle
Data Tables [2].

Our calculation is precise enough that we can see
the difference between $m_B(m_t) - m_B(m_t)$ in the bot-
tom sector and the similar quantity in the charm sec-
mor $m_{D_s}(m_t) - m_{D_s}(m_t)$ (figure 6). These mass differ-
ces are small compared to the absolute masses of the
states, and should be the same in the infinitely
heavy quark limit. We can see that our calculation
correctly reproduces the small difference due to the
finite value of the mass of the charm and bottom.

IV. CONCLUSIONS AND OUTLOOK

We have shown that the use of a highly improved
relativistic action on fine enough lattices is capable of
delivering very precise results on systems with a charm
quark. The high statistical accuracy of our data com-
bined with calculations at several values of the lattice
spacing and light quark masses allows us to make a
controlled joint chiral and continuum extrapolation.

We can calculate accurately the mass of heavy-light
systems, which provide a stringent test of the calculation.
We can calculate precise values for the decay constants of pseudoscalar heavy-light mesons (as well
as light-light mesons), and especially for the ratio of
such decay constants.

The very precise calculation of the masses of heavy-
heavy pseudoscalar mesons should make possible a di-
rect lattice determination of the mass of the charm
quark. Because we use the same relativistic action
through the calculation for both the charm and the
light quarks, we can also obtain a very precise value
for the ratio $m_c/m_s$, and therefore if $m_c$ is deter-
mined through another method use the ratio to get
$m_s$. We are also working (in collaboration with the
Karlsruhe group) on a new method for the determina-
tion of $m_c$ by combining continuum perturbation
results with lattice data.

Another quantity which we plan to calculate in the
near future is the leptonic decay width $\Gamma_{e^+e^-(\psi)}$, as
well as the semileptonic form factors for $D \rightarrow \pi l\nu,
D \rightarrow K l\nu$.

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Update on Semileptonic Charm Decays

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A brief update is given on recent developments in the theory of exclusive semileptonic charm decays. A check on analyticity arguments from the kaon system is reviewed. Recent results on form factor shape measurements are discussed.

1. Introduction

Semileptonic meson decays provide a valuable arena to study the weak and strong interactions. On the one hand, once the effects of the strong interaction are under control, weak mixing parameters (\(V_{ub}, V_{us}, \ldots\)) can be extracted from the overall normalization. On the other hand, if these mixing parameters are taken from other processes, then the semileptonic rates probe complicated underlying strong dynamics, thus yielding an important test for lattice methods [1] and for our understanding of nonperturbative QCD.

Apart from the overall normalization, the spectral shapes are also interesting [2]. The energy spectra (the “\(q^2\) dial”) are governed by quark-hadron duality, and are constrained by dispersion relations and analyticity. Shape parameters for different quark masses (the “\(m\) dial”) can be interpreted by means of effective field theory descriptions in the appropriate regimes of validity. These parameters provide inputs to other processes governed by the same effective field theory. For example, the form factors measured in \(B \rightarrow \pi \ell \nu\) constrain \(B \rightarrow \pi \pi\).

The charm quark holds a privileged position in this scheme. Abundant experimental data exist for \(D \rightarrow K\) and for \(D \rightarrow \pi\) transitions. The charm mass is heavy enough to be treated using heavy-quark methods on the lattice, but light enough so that the full range of physically allowed momentum transfers are accessible in present simulations [3]. Charm decays thus provide a powerful test of lattice QCD methods that can be applied to other heavy-meson systems.

Charm decays are important for testing another important, but perhaps less well-known aspect of QCD, namely the constraints imposed by analyticity [4, 5, 6, 7, 8]. These constraints imply a convergent expansion in powers of a small parameter that measures the distance between the physically allowed kinematic region, and the region of resonances and production thresholds. While the existence of this small parameter is well known, the usefulness of the expansion has been practically negated by appeals to “unitarity bounds”. It has only recently become clear with the advent of rigorous power-counting arguments that we can “take seriously” the expansion provided by analyticity [7]. Charm decays provide a valuable illustration and crosscheck of these arguments.

Table I

| Process | CKM element | \(|z|_{\text{max}}\) |
|---------|-------------|------------------|
| \(\pi^+ \rightarrow \pi^0\) | \(V_{ub}\) | \(3.5 \times 10^{-5}\) |
| \(B \rightarrow D\) | \(V_{cb}\) | 0.032 |
| \(K \rightarrow \pi\) | \(V_{us}\) | 0.047 |
| \(D \rightarrow K\) | \(V_{cs}\) | 0.051 |
| \(D \rightarrow \pi\) | \(V_{cd}\) | 0.17 |
| \(B \rightarrow \pi\) | \(V_{ub}\) | 0.28 |

An overview of these ideas for general semileptonic meson transitions has been presented previously in [2]; further discussion and references may be found there. The focus in the present report is on some recent illustrations from kaon physics that are relevant to charm physics, and on an update of shape measurements in the charm system.

The remainder of the talk is organized as follows. Section 2 outlines the constraints of analyticity, and describes the explicit test of power-counting afforded by \(K_{\ell 3}\) decays. Section 3 tabulates recent results on the shape parameters in \(D \rightarrow K\) and \(D \rightarrow \pi\) semileptonic transitions. Section 4 concludes with a discussion of the relevance of these shape parameters for testing lattice QCD, and for applying factorization in \(B\) decays.

2. Analyticity and simplicity

The mere existence of a field theory description of the physical hadrons, and their weak current probes, implies powerful constraints on form factors. In particular, singularities in hadronic transition amplitudes are determined by kinematics. Analyticity translates into a convergence expansion in a small variable once the domain of analyticity is mapped onto a standard region.

In practical terms, this mapping is simply a rearrangement of the series expansion of the form factor,

\[
F(t) = F(0) + \ldots \tag{1}
\]

For example, a simple pole model of the form factor
would “resum” into the form
\[
F(t)/F(0) = 1 + t/m^2 + (t/m^2)^2 + \ldots = \frac{1}{1 - q^2/m^2}.
\]
Without knowing what the form factor “resums” into, analyticity implies that the series has to rearrange itself into the form
\[
F(t)/F(0) = 1 + a_1 z(t) + a_2 z^2(t) + \ldots,
\]
where \(a_k\) are \(O(1)\) in a rigorous sense (\(\sum a_i^2\) is also \(O(1)\)), and \(z\) is a variable bounded by the distance of the physical region from singularities.

The smallness of \(|z|\) in (3) implies that terms beyond linear order are highly suppressed. There is thus an essentially unique choice of shape parameter that unites semileptonic transitions from various decay modes—the slope of the form factor, say at \(q^2 = 0\). It turns out that the numerical value of this parameter is in itself an interesting quantity. It provides a quantitative test of lattice versus experimental shape; it is an input to related hadronic processes; and it probes an unsolved question in the QCD dynamics governing form factors [2].

The tools of analyticity are well-known, but their usefulness has not been appreciated, due to a reliance on unitarity as the only means to bound the coefficients \(a_k\). In fact, rigorous power counting arguments show that both \(a_k\) and \(\sum a_k^2\) are order unity, even when large ratios of scales, such as \(m_Q/\Lambda_{\text{QCD}}\), are present (\(m_Q\) a heavy-quark mass, and \(\Lambda\) a hadronic scale). It is important to test and make full use of this expansion when dealing with complicated QCD systems.

Analysis of \(K_{\ell 3}\) decays \((K \to \pi \ell \nu)\) provides a revealing confirmation of these ideas. The analysis also provides an important constraint on the determination of \(|V_{us}|\) [9]. Figure 1 compares the unitarity bound to the actual value of the sum of squares of coefficients, as calculated from ALEPH \(\tau\) decay data [10]. The result is plotted for a range of values of the OPE parameter \(Q\) \((Q\) corresponds to the virtuality of the produce meson pair in the crossed channel; large \(Q\) implies that the OPE is increasingly reliable). It can be seen that the unitarity bound begins to wildly overestimate the possible size of the coefficients for even moderately large \(Q\). The convergence of the series improves as \(Q\) becomes larger, since the sensitivity to the \(K^*\) pole is lessened; reliance on the unitarity bound would however force us to work at small \(Q\).

### 3. Shape parameters

\[
\begin{align*}
\delta &\equiv 1 - \frac{m_H^2 - m_L^2}{F_+(0)} \left( \frac{dF_+}{dt} \bigg|_{t=0} - \frac{dF_0}{dt} \bigg|_{t=0} \right) = \frac{F_+(0) + F_-(0)}{F_+(0)}. 
\end{align*}
\]

Here \(F_+\) is the vector form factor that dominates for massless leptons in pseudoscalar-pseudoscalar transi-
At leading power \[17\], where \(E\) is the energy of the outgoing light hadron. To extract the heavy-quark mass dependence, and organize form-factor contributions in powers of \(\Lambda/\sqrt{m}\), and massive leptons, or in lattice simulations. Results on residual shape uncertainty from allowing \(a_2/\alpha_1\) to vary, and with previous CLEO, FOCUS and BELLE measurements, these values are displayed in Figure 2. A naive average of these results yields \(1 + 1/\beta - \delta = 0.97 \pm 0.05 \, (\chi^2 = 2.5 \text{ for 5-1 d.o.f.})\). From the CLEOc results on \(D \to \pi\),

\[
D \to \pi : \quad 1 + 1/\beta - \delta = 1.27 \pm 0.13 \pm 0.27 \quad (9)
\]

Together with previous CLEO and BELLE measurements, these values are displayed in Figure 3. A naive average of these results yields \(1 + 1/\beta - \delta = 1.19 \pm 0.23 \, (\chi^2 = 0.4 \text{ for 3-1 d.o.f.})\).

4. Conclusions

Semileptonic charm decays provide an important arena in which to test lattice QCD. Unfortunately, it is not possible to definitively test the experimental shape predictions against unquenched lattice simulations at present, since the lattice results are so far reported only in terms of model parameters, that need not agree between theory and experiment, or between experiments with different acceptances and systematics. It is interesting to note that while some model parameters show poor agreement between experiments \[18\], there is no obvious discrepancies among the physical results shown in Figure 2.

The actual value of the shape parameters are also of interest. For example, the parameter \(\delta\) is an important input to factorization analyses of \(B \to \pi\pi\) decays \[19\], usually phrased in terms of an (inverse) moment of the B meson wavefunction, \(\lambda\):

\[
\delta(m_\pi, m_B) = \frac{6\pi C_F}{N_c m_B \lambda B F_V(0)} + \ldots \quad (10)
\]

A plausible conjecture of monotonicity for \(\delta\) implies that \(\delta(m_\pi, m_B) < 0.35 \pm 0.03\). A much larger value of \(\delta(m_\pi, m_B)\) in \(B \to \pi\), e.g. \(\delta \approx 1\) \[20\], would require a dramatic behavior of this physical quantity as a function of quark mass, a behavior that should already be apparent in the charm system. For example, we would expect \(\delta(m_\pi, m_D) > \delta(m_K, m_D)\) (see Figure 8 of \[2\]). It is exciting that current lattice simulations can probe this quantity in both the charm and bottom systems, not by evaluating moments of \(B\) meson wavefunctions, but by computing form factor slopes. Experimental measurements are sensitive to the combination \(1/\beta - \delta\), so that the scalar form factor slope (denoted \(1/\beta\)) is the most urgent requirement from the lattice. Again, definitive conclusions must await a model-independent presentation of simulation results.
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Recent results on 4-body, charm semileptonic decays

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We summarize recent data on 4-body charm semileptonic decay concentrating on $D^+_s \rightarrow K^+K^-e^+\nu$ and $D^+ \rightarrow K^-\pi^+e^+\nu$. We begin with giving some motivation for the study of these decays. We discuss several of the models traditionally used to describe these decays and conclude by presenting a non-parametric analysis of $D^+ \rightarrow K^-\pi^+\mu^+\nu$ and its possible extension into non-parametric studies of $D^+_s \rightarrow K^+K^-e^+\nu$ and $D^+ \rightarrow K^-\pi^+\mu^+\nu$.

I. INTRODUCTION

Figure 1 shows a cartoon of the $D^0 \rightarrow K^-\ell^+\nu$ decay process. All of the hadronic complications for this process is contained in $q^2$ dependent form factors that are computable using non-perturbative methods such as LQCD. Although semi-leptonic process can in principle provide a determination of charm CKM elements, one frequently uses the (unitarity constrained) CKM measurements, lifetime, and branching fraction to measure the scale of charm semileptonic decay constants and compare them to LQCD predictions. The $q^2$ dependence of the semileptonic form factor can also be directly measured and compared to theoretical predictions.

The hope is that charm semileptonic decays can provide high statistics, precise tests of LQCD calculations and thus validate the computational techniques for charm. Once validated, the same LQCD techniques can be used in related calculations for $B$-decay and thus produce CKM parameters with significantly reduced theory systematics.

Although recent, unquenched LQCD calculations are unavailable for $D \rightarrow$ vector $\ell^+\nu$ processes, owing to the instability of the vector parent, I hope that the 4-body will provide additional tests of LQCD for a variety of spin states which will further help calibrate the lattice, and provide confidence in analogous decays for the beauty sector.

I find it remarkable that 4-body semileptonic decays such as $D^+_s \rightarrow K^+K^-e^+\nu$ and $D^+ \rightarrow K^-\pi^+\mu^+\nu$ are so heavily dominated by the vector decays $D^+_s \rightarrow \phi \ e^+\nu_e$ and $D^+ \rightarrow \overline{K}^0 \mu^+\nu$. Figure 2 illustrates this dominance by showing data from FOCUS[1] and recent data from BaBar[2]. The absence of a substantial

![Diagram 1](image1.png)

**FIG. 1:** Diagrams for the semileptonic decay of charmed mesons. The hadronic, QCD complications are contained in $q^2$ dependent form factors.

![Diagram 2](image2.png)

**FIG. 2:** We show the $m(K^+K^-)$ spectra obtained in $D^+_s \rightarrow K^+K^-\ell^+\nu$ by BaBar[2] and $m(K^-\pi^+)$ by FOCUS[1]. The curve on the FOCUS $m(K^-\pi^+)$ spectra is a $K^*^0$ line shape both with ($A = 0.36$) and without ($A = 0$) a small s-wave, non-resonant component which was found through an interference in the decay intensity and is described later. The $m(K^+K^-)$ spectra obtained by BaBar is very strongly dominated by the $\phi$ resonance along with a few known backgrounds.

non-resonant, or higher spin resonance component to these decays means the decay angular distribution can be described in terms of three, $q^2$-dependent helicity basis form factors that describe the coupling of the lepton system to the three helicity states of the vector meson according to Eq. 1:
The three decay angles describing the $D^+ \rightarrow K^- \pi^+ \ell^+ \nu$ decay, referenced in Eq. (1), are illustrated by Fig. 3.

![Fig. 3: Definition of kinematic variables.](image)

\[
|\mathcal{A}|^2 \approx \frac{q^2}{8} \left| \begin{array}{c}
(1 + \cos \theta_i) \sin \theta_V e^{i\chi} H_+(q^2) \\
-(1 - \cos \theta_i) \sin \theta_V e^{-i\chi} H_-(q^2) \\
-2 \sin \theta_i \cos \theta_V H_0(q^2)
\end{array} \right|^2
\]  

(1)

where $K$ is the momentum of the $K^-\pi^+$ system and $m_{K\pi}$ is its mass.

Eq. (3) provides considerable insight into the expected analytic form for semileptonic form factors. It uses a dispersion relation obtained using Cauchy's Theorem under the assumption that a form factor is an analytic, complex function apart from some known singularities. Fig. 4 illustrates the Cauchy’s Theorem contour for the case of the $f_+(q^2)$ form factor describing $D^0 \rightarrow K^- \ell^+ \nu$.

The form factor singularities will consist of a sum of simple poles at the $D$ meson-kaon vector bound states (e.g. $D^{*+}$) plus a cut beginning at the $D-$kaon continuum in the cross process: $\nu \ell^+ \rightarrow D$ kaon. The dispersion relation gives the form factor (F(q^2)) as a sum over the spectroscopic poles plus an integral over the cut.

\[
F(q^2) = \frac{\mathcal{R}}{m_{D^*}^2 - q^2} + \frac{1}{\pi} \int_{m_D + m_K}^{\infty} \frac{\text{Im} \{f_+(s)\}}{s - q^2 - i\varepsilon} \, ds
\]  

(3)

Both the cuts and poles are generally beyond the physical $q^2_{\text{max}}$ and thus can never be actually realized.

Spectroscopic pole dominance (SPD) was an early parameterization for the form factors relevant to both $D \rightarrow \text{vector } \ell^+ \nu$ and $D \rightarrow \text{pseudoscalar } \ell^+ \nu$. SPD ignores the cut integral entirely and approximates F(q^2) using just the first term of Eq. (3). The advantage of SPD approach is that it requires only a single unknown fitting parameter $\mathcal{R}$ to describe each F(q^2) since the positions of the bound states are well known. SPD entirely predicts shape of $D \rightarrow \text{pseudoscalar } \ell^+ \nu$ decay intensity and predicts that the shape for the

**II. ANALYTIC MODELS FOR FORM FACTORS**

We begin by describing the three form factors relevant to $D \rightarrow \text{vector } \ell^+ \nu$ although there is strong evidence $[1,3]$ for a non-resonant, s-wave component to $D^+ \rightarrow K^- \pi^+ \ell^+ \nu$. A new, fifth form factor $H_T(q^2)$ is also required for $D^+ \rightarrow K^0 \mu^+ \nu$ to describe the suppressed coupling of the $K^0$ to a left-handed $\mu^+$.

The $H_+(q^2), H_-(q^2), H_0(q^2)$ form factors are linear combinations of two axial and one vector form factor $[5]$ according to Eq. (2):

\[
H_{\pm}(q^2) = (M_D + m_{K\pi}) A_1(q^2) \mp 2 \frac{M_D K}{M_D + m_{K\pi}} V(q^2),
\]

\[
H_0(q^2) = \frac{1}{2m_{K\pi}\sqrt{q^2}} \left[ (M_D^2 - m_{K\pi}^2 - q^2)(M_D + m_{K\pi}) A_1(q^2) - 4 \frac{M_D^2 K^2}{M_D + m_{K\pi}} A_2(q^2) \right]
\]  

(2)

![Fig. 4: Each form factor is assumed to be an analytic function with pole singularities at the masses of bound states, and cuts that start at the start of the continuum. We illustrate the case of $D^0 \rightarrow K^- \ell^+ \nu$. One can use Cauchy’s theorem with the indicated contour to write an dispersion expression for each form factor in the physical range $0 < q^2 < (m_D - m_K)^2$.](image)
\[ D^+ \rightarrow \overline{K}^0 \ell^+ \nu_\ell \] can be fit by just two parameters which are traditionally taken to be the axial and vector form factor ratios at \( q^2 = 0 \): \( r_v = V(0)/A_1(0) \) and \( r_2 = A_2(0)/A_1(0) \).

BaBar \[2\] has recently published an interesting SU(3) test based on SPD applied to \( D_s^+ \rightarrow \phi \, e^+ \nu_e \). Figure 5 compares the \( r_v \) and \( r_2 \) parameters measured for \( D_s^+ \rightarrow \phi \, e^+ \nu_e \) to those previously measured for \( D^+ \rightarrow \overline{K}^0 \ell^+ \nu_\ell \). By SU(3) symmetry and explicit calculation, the \( r_v \) and \( r_2 \) form factor ratios for \( D^+ \rightarrow \overline{K}^0 \ell^+ \nu_\ell \) and \( D_s^+ \rightarrow \phi \, \ell^+ \nu_\ell \) decays are expected to be very close to each other. This is true for \( r_v \), but previous to the recent measurement by the FOCUS Collaboration \[6\], \( r_2 \) for \( D_s^+ \rightarrow \phi \, \ell^+ \nu_\ell \) was measured to be roughly a factor of two larger than that for \( D^+ \rightarrow \overline{K}^0 \ell^+ \nu_\ell \). BaBar \[2\] has confirmed the expected consistency between the form factor ratios obtained for \( D_s^+ \rightarrow \phi \, \ell^+ \nu_\ell \) and \( D^+ \rightarrow \overline{K}^0 \ell^+ \nu_\ell \) with unparalleled statistics.

Several experiments have tested SPD by measuring an “effective” pole mass \( (m_{\text{pole}})_{\text{eff}} \) in \( D^0 \rightarrow K^- e^+ \nu_e \) decay, where the pole mass is defined using \( f_+(q^2) \propto 1/(m_{\text{pole}}^2 - q^2) \). As Fig. 6 from Reference \[4\] shows, as errors have improved over the years, it becomes clear that effective pole is significantly lower than the spectroscopic pole, underscoring the importance of the cut integral contribution for this decay.

Several parameterizations have been proposed to include the cut integral in Eq. (3) as well as the spectroscopic poles. Becirevic and Kaidalov (1999) \[7\] proposed a new parameterization for the \( D \rightarrow \) pseudoscalar \( \ell^+ \nu \) factor \( f_+(q^2) \) that replaces the cut integral by an effective pole where the heavy quark symmetry and other theoretical ideas are used to relate the residue and effective pole position. These constraints leads to a modified pole form with a single additional parameter \( \alpha \) that describes the degree to which the single spectroscopic pole fails to match \( f_+(q^2) \) for a given process.

\[
f_+(q^2) = \frac{f_+(0)}{(1-q^2/m_{D^*}^2)(1-\alpha q^2/m_{D^*}^2)}
\]

S. Fajfer and J. Kamenik \[8\] have recently extended the effective pole approach to the three helicity form factors relevant to \( D \rightarrow \) vector \( \ell^+ \nu \) decays.

R.J. Hill \[9\] has proposed an alternative way of viewing form factors which is illustrated in Fig. 7. The basic idea is to devise a transformation of a form factor from the complex \( q^2 \) plane to a complex \( z \) plane.
This transformation is devised to (1) remove the spectroscopic poles and (2) put the cuts far away from the physical $z$ region. After the transformation, since the singularities have been removed or diminished, each form factor can be well represented by a low order Taylor series in $z$. The transformation approach is known to work very well in $B$-decays where the physical $q^2$ region gets very close to the singularities for pseudo-scalar $B$ semileptonic decay. It also works well for pseudoscalar charm pseudoscalar semi-leptonic decay.[8]

III. $D^+ \to K^-\pi^+\ell^+\nu$ DECAYS

Although historically $D^+ \to K^0\ell^+\nu$ have been the most accessible semileptonic decays in fixed target experiments owing to their ease of isolating a signal, they are significantly more complicated to analyze than $D \to$ pseudoscalar $\ell^+\nu$. One problem is that a separate helicity form factor is required for each of the three helicity states of vector meson. The $q^2$ dependence of these form factors cannot be simply measured from the $q^2$ dependence of the decay rate as is the case in $D \to$ pseudoscalar $\ell^+\nu$ but rather must be entangled from the $q^2$ dependence of the angular distribution such as that given by Eq. (1).

Another complication is that since $D^+ \to K^-\pi^+\ell^+\nu$ states result in a multihadronic final state, the $D^+ \to K^0\ell^+\nu$ final states can potentially interfere with $D^+ \to K^-\pi^+\ell^+\nu$ processes with the $K^-\pi^+$ in various angular momentum waves with each wave requiring its own form factor. Because the $m_{K\pi}$ distribution in $D^+ \to K^-\pi^+\ell^+\nu$ was an excellent fit to the $K^0$ Breit-Wigner as shown in Fig. [2], it was assumed for many years that any non-resonant component to $D^+ \to K^-\pi^+\ell^+\nu$ must be negligible. In 2002, FOCUS observed a strong, forward-backward asymmetry in $\cos \theta_V$ for events with $m_{K\pi}$ below the $K^0$ pole with essentially no asymmetry above the pole as shown in Figure [3]. The simplest explanation for this asymmetry is the presence of a linear $\cos \theta_V$ term in the decay intensity due to interference between the $D^+ \to K^0\mu^+\nu$ and a non-resonant, $s$-wave amplitude. This interference is the second-to-last term in Eq. (5), which is basically an expanded version of Eq. (1), integrated over acoplanarity $\chi$. We also explicitly include the $K^0$ Breit-Wigner amplitude (BW). Note that all other interference terms (such as a possible $H_+(q^2) \times H_-(q^2)$ contribution) vanish because of the $\int_0^{2\pi} d\chi \exp(i\Delta\chi)$ integration. Only “same” helicity contributions can interfere in the acoplanarity averaged intensity. We will argue shortly that an appropriate $\delta$ can create the asymmetry pattern shown in Fig. [3].

Finally we introduce an additional form factor ($h_0(q^2)$) in Eq. (5) to describe the coupling to the $0.8 < M(K\pi) < 0.9 \text{ GeV}/c^2$ and a non-resonant, $s$-wave amplitude.

\[
\int |A|^2 d\chi = \frac{1}{8} q^2 \left\{ \begin{array}{l}
((1 + \cos \theta) \sin \theta \nu)^2 |H_+(q^2)|^2 |BW|^2 \\
+ ((1 - \cos \theta) \sin \theta \nu)^2 |H_-(q^2)|^2 |BW|^2 \\
+ (2 \sin \theta \cos \theta \nu)^2 |H_0(q^2)|^2 |BW|^2 \\
+ 8 (\sin \theta \cos \theta \nu) H_0(q^2) h_0(q^2) \text{Re} \{ Ae^{-i\delta} BW \} \\
+ O(A^2) \end{array} \right\} \tag{5}
\]

Since the helicity intensity contributions are proportional to $q^2 H_\pm^2(q^2)$, according to Eq. (5), the $H_\pm$ intensity contributions vanish in this limit, while $q^2 H_0^2(q^2)$ will approach a constant.

Figure [3] explains why this is true. As $q^2 \to 0$, the $e^+$ and $\nu$ become collinear with the virtual $W^+$. For $H_+(q^2)$ and $H_-(q^2)$, the virtual $W^+$ must be in the $|1, \pm 1\rangle$ state which means that the $e^+$ and $\nu$ must

---

**A. Asymptotic Forms**

Assuming that $A_{1,2}(q^2)$ and $V(q^2)$ approach a constant in the low $q^2$ limit, as expected in spectroscopic pole dominance, Eq. (2) shows $q^2 \to 0$, both $H_+(q^2)$ and $H_-(q^2)$ approach a constant as well. By way of constraint, $H_0(q^2)$ will diverge in the low $q^2$ limit according to Eq. (2), owing to the $1/\sqrt{q^2}$ prefactor.
both appear as either right-handed or left-handed thus violating the charged current helicity rules. Hence $q^2 H_\pm(q^2)$ vanishes at low $q^2$. For $H_0(q^2)$, the $W^+$ is in $|1,0\rangle$ state thus allowing the $e^+$ and $\nu$ to be in their (opposite) natural helicity state. Hence at low $q^2$, $q^2 H_0(q^2) \to \text{constant}$ which allows for $D^+ \to K^0 \mu^+ \nu$ decays as $q^2 \to 0$. Presumably $h_0(q^2) \to 1/\sqrt{q^2}$ as well since it also describes a process with $W^+$ in the $|1,0\rangle$ state

This technique was initially developed by the FOCUS Collaboration\textsuperscript{[13]} and applied to CLEO\textsuperscript{[3]} data. As shown in Eq. (5), after integrating over acoplanarity, the decay intensity is just a sum over four terms that consist of a form factor product times a characteristic angular distribution in $\theta_V$ and $\theta_t$. The acoplanarity integration has significantly simplified the problem by eliminating the five of the possible six interference terms between the four form factor amplitudes with different helicities. We begin by making a binned version of Eq. (4) given by Eq. (6), where for simplicity we only write three of the terms:

$\vec{D}_i = f_+(q_i^2) \vec{m}_+ + f_-(q_i^2) \vec{m}_- + f_0(q_i^2) \vec{m}_0$ (6)

We use 25 joint $\Delta \cos \theta_V \times \Delta \cos \theta_t$ angular bins: 5 evenly spaced bins in $\cos \theta_V$ times 5 bins in $\cos \theta_t$ and 6 bins in $q^2$ ($i = 0 \to 6$). The number of $D^+ \to K^- \pi^+ \ell^+ \nu$ events observed in each of the 25 angular bins is packed into a twenty-five component $\vec{D}_i$ “data” vector.

The $f_\pm(q_i^2)$ and $f_0(q_i^2)$ are proportional to $H_{\pm}(q_i^2)$, $H_0^2(q_i^2)$ averaged over the $q_i^2$ bin along with all phase space and efficiency factors. The $\vec{m}_\pm$ and $\vec{m}_0$ are the angular distributions due to each individual form factor product packed into a 25-vector for each of the six $q^2$ bins. The acceptance and phase space corrected $m$-vectors are obtained directly from a Monte Carlo simulation where a given form factor product is turned on and all others are turned off. We can write Eq. (6) as the “component equation” shown in Eq. (7) by forming the dot product with each of the three $m$-vectors:

$$
\begin{pmatrix}
\vec{m}_+ \cdot \vec{D}_i \\
\vec{m}_- \cdot \vec{D}_i \\
\vec{m}_0 \cdot \vec{D}_i
\end{pmatrix} =
\begin{pmatrix}
\vec{m}_+ \cdot \vec{m}_+ & \vec{m}_+ \cdot \vec{m}_- & \vec{m}_+ \cdot \vec{m}_0 \\
\vec{m}_- \cdot \vec{m}_+ & \vec{m}_- \cdot \vec{m}_- & \vec{m}_- \cdot \vec{m}_0 \\
\vec{m}_0 \cdot \vec{m}_+ & \vec{m}_0 \cdot \vec{m}_- & \vec{m}_0 \cdot \vec{m}_0
\end{pmatrix}
\begin{pmatrix}
f_+(q_i^2) \\
f_-(q_i^2) \\
f_0(q_i^2)
\end{pmatrix}
$$

The solution to Eq. (7) can be written as:

$$
f_+(q_i^2) = i \vec{P}_+ \cdot \vec{D}_i, \quad f_-(q_i^2) = i \vec{P}_- \cdot \vec{D}_i, \quad f_0(q_i^2) = i \vec{P}_0 \cdot \vec{D}_i
$$

where $i \vec{P}_\alpha$ vectors are given by Eq. (9).

\[
\begin{pmatrix}
  i \vec{P}_+ \\
  i \vec{P}_- \\
  i \vec{P}_0
\end{pmatrix}
= \begin{pmatrix}
  \vec{m}_+ \cdot \vec{m}_+ & \vec{m}_+ \cdot \vec{m}_- & \vec{m}_0 \cdot \vec{m}_0 \\
  \vec{m}_- \cdot \vec{m}_+ & \vec{m}_- \cdot \vec{m}_- & \vec{m}_0 \cdot \vec{m}_- \\
  \vec{m}_0 \cdot \vec{m}_+ & \vec{m}_0 \cdot \vec{m}_- & \vec{m}_0 \cdot \vec{m}_0
\end{pmatrix}^{-1}
\begin{pmatrix}
  \vec{m}_+ \\
  \vec{m}_- \\
  \vec{m}_0
\end{pmatrix}
\]

It is useful to think of forming the dot products in Eq. (8) by making a weighted histogram:

\[
\vec{P}_+ \cdot \vec{D} = \left[ \vec{P}_+ \right]_1 n_1 + \left[ \vec{P}_+ \right]_2 n_2 + \cdots + \left[ \vec{P}_+ \right]_{25} n_{25}
\]

Eq. (10) demonstrates the product \(\vec{P}_+ \cdot \vec{D}\) is equivalent to weighting the \(n_1\) events in angular bin 1 by \(\left[ \vec{P}_+ \right]_1\), weighting the \(n_2\) events in angular bin 2 by \(\left[ \vec{P}_+ \right]_2\), etc. Hence each form factor product such as \(f_+(q^2)\) can be obtained by simply weighting the data by \(\left[ \vec{P}_+ \right]_i\), where \(i\) is the angular bin of the given datum. The acceptance and phase space factors can be easily included the projective weights as well in order to directly produce each form factor product. Hence the (arbitrarily normalized) form factor products \(H_2^2(q^2)\), \(H_2^2(q^2)\), and \(H_2^2(q^2)\) can then be obtained by making three weighted histograms using the efficiency rescaled \(i \vec{P}_+, i \vec{P}_-\), and \(i \vec{P}_0\) weights respectively.

The same, basic projective weighting approach has been recently applied by the FOCUS Collaboration\[11\] for a non-parametric analysis of the \(K^{-}\pi^+\) amplitudes in the hadronic decay \(D^+ \to K^{-}\pi^+\pi^+\). To whet the appetite, Fig.\[10\] shows the \(K^{-}\pi^+\) amplitudes obtained in that analysis. The s-wave amplitude shown in Fig.\[10\] (a) and begs comparison with the s-wave amplitude obtained in a K-matrix analysis\[12\] of \(D^+ \to K^{-}\pi^+\pi^+\) described by S. Malvezzi in these proceedings.

V. A NON-PARAMETRIC ANALYSIS OF THE HELICITY FORM FACTORS IN \(D^+ \to K^{-}\pi^+\ell^+\nu\)

Figure\[11\] shows the four weighted histograms from an analysis of 281 pb\(^{-1}\) \(\psi(3770)\) CLEO data\[3\]. Figure\[11\] shows the expected behavior discussed in Section III A. In particular, \(H_\pm q^2\) \(\to\) constant as \(q^2\) \(\to\) 0 while the zero-helicity form factors, \(H_0(q^2)\) and \(h_0(q^2)\), diverge as \(1/\sqrt{q^2}\). It is interesting to note that although the non-resonant, s-wave amplitude is too small to see in the \(K^{-}\pi^+\) mass spectrum (Fig. 2), its form factor is measured with roughly the same precision as \(H_2^2(q^2)\) or \(H_2^2(q^2)\). The curves give the helicity form factors according to Eq. (9), using spectroscopic pole dominance and the \(r_0\), \(r_2\), and s-wave parameters measured by FOCUS\[14\]. Apart from the \(h_0(q^2)\) \(H_0(q^2)\) interference form factor product, the spectroscopic pole dominance model is a fairly good match to the CLEO non-parametric analysis. This suggests that the ad-
helicity basis form factors by plotting the intensity coefficients of each of the form factor products. This is the form factor product multiplied by \( q^2 \). Since \( q^2 H_0^2(q^2) \) dominates, we normalized form factors such that \( q^2 H_0^2(q^2) = 1 \) at \( q^2 = 0 \) but use the same scale factor for the other three form factors. As expected, both \( q^2 H_0^2(q^2) \) and \( q^2 H_2^2(q^2) \) rise from zero with increasing \( q^2 \) and they both appear to approach \( q^2 H_0^2(q^2) \) at \( q^2_{\text{max}} \) — although \( q^2 H_2^2(q^2) \) seems slightly lower than \( q^2 H_0^2(q^2) \) at \( q^2_{\text{max}} \).

![Figure 12](image)

**Figure 12:** Non-parametric form factor products obtained for the data sample (multiplied by \( q^2 \)). The reconstructed form factor products are shown as the points with error bars, where the error bars represent the statistical uncertainties. The solid curves in the histograms represent a form factor model described in Ref. [14]. The reconstructed form factor products are the points with error bars. The three plots on the right are run with the axial and vector pole masses, while the three plots on the right are run with the axial and vector pole masses taken to infinity.

What can we learn about the pole masses? Unfortunately Fig. 13 shows that the present data is insufficient to learn anything useful about the pole masses. On the left of Figure 13 the helicity form factors are compared to a model generated with the FOCUS form factor ratios [14] and the standard pole masses of 2.1 GeV for the vector pole and 2.5 GeV for the two axial poles. On the right side of Fig. 13 the form factors are compared to a model where the pole masses are set to infinity meaning that the axial and vector form factors are constant. Both models fit the data equally well.

The data of Fig. 13 is consistent with the spectroscopic pole dominance albeit with essentially no sensitivity to the pole masses. Fig. 14 shows that it is also consistent with the expected behavior under a Hill transformation, illustrated earlier in Fig. 7.

![Figure 13](image)

**Figure 13:** Non-parametric form factor products obtained for data (multiplied by \( q^2 \)). The solid curves are based on the s-wave model and measurements described in Reference [14]. The reconstructed form factor products are the points with error bars. The three plots on the right are the usual model with the spectroscopic pole masses; while the three plots on the right are run with the axial and vector pole masses taken to infinity.

It is interesting to note that the FOCUS analysis was based on a sample of 11400 \( D^+ \to K^- \pi^+ \mu^+ \nu \) events, while the CLEO analysis was based on a sample of only 2470 \( D^+ \to K^- \pi^+ e^+ \nu \) events. The error bars in Fig. 14 for FOCUS data are much larger than those for the much smaller CLEO data set and only four FOCUS \( q^2 \) bins are reported on. This is because of the much poorer \( q^2 \) resolution in fixed target semileptonic decay compared to the order-of-magnitude better \( q^2 \) resolution obtainable for semileptonic analyses in charm threshold data from e+e− colliders where the neutrino can be reconstructed using energy-momentum balance. This is especially relevant for \( D^+ \to K^- \pi^+ \ell^+ \nu \) since the 1 GeV \( q^2 \) range for \( D^+ \to K^- \pi^+ \ell^+ \nu \) is a factor of two smaller than that in \( D^0 \to K^- \ell^+ \nu \). Error inflation due to deconvolution grows dramatically once the bin-to-bin separation, \( \Delta q^2 \), approaches the r.m.s. resolution, \( \sigma(q^2) \), which was typically 0.18 GeV^2 in the four bins reported on by FOCUS [13].

What can we learn about the phase of the s-wave contribution? Recall in Figure 8 the asymmetry created by the interference between the s-wave and \( D^+ \to K^0 \ell^+ \nu \) only appeared below the \( K^0 \) pole in FOCUS data and thus the s-wave phase was such that
it was orthogonal with the $m_{K\pi} > m(K_0^{*0})$ half of the Breit-Wigner amplitude or $\langle BW_\pi \rangle$. Since the asymmetry is “negative” according to the convention of Eq. (5), in that favors the backward over the forward $\cos \theta$ direction, it must be anti-collinear to $\langle BW_\pi \rangle$ as well. Hence it must have roughly the phase of $40^\circ$ as illustrated by Fig. 15. FOCUS measured the $s$-wave phase to be $\delta = (39 \pm 4 \pm 3)^\circ$.

Finally, is there evidence for higher $K^{-}\pi^+$ angular momentum amplitudes in $D^+ \rightarrow K^{-}\pi^+ \ell^+\nu$? We searched for possible additional interference terms such as a (zero helicity) $d$-wave contribution: $4 \sin^2 \theta_{V} (3 \cos^2 \theta_{V} - 1) H_0(q^2) h_0^{(d)}(q^2) \Re \{A e^{-i\delta} \langle BW \rangle \}$ or an $f$-wave contribution: $4 \sin^2 \theta_{V} \sin \theta_{V} (5 \cos^2 \theta_{V} - 3) H_0(q^2) h_0^{(f)}(q^2) \Re \{A e^{-i\delta} \langle BW \rangle \}$. As shown in Figure 17 there is no evidence for such additional contributions which should diverge as $1/q^2$ at low $q^2$.

**FIG. 15:** Illustration of $s$-wave phase

As Figure 16 shows, the same thing happens in CLEO data. The effective $h_0(q^2) H_0(q^2)$ disappears above the $K_0^{*0}$ pole and is very strong below the pole. The amplitude $A$ of the $s$-wave piece is arbitrary since using interference we can only observe the product $A H_0(q^2) h_0(q^2)$. This means any change in $A$ scale can be compensated by a change of scale in $h_0(q^2)$. The fact that the $h_0(q^2) H_0(q^2)$ data was a tolerable match (at least in the low $q^2$ region) to the FOCUS curve in Figure 14 does imply, however, that the $s$-wave amplitude observed in CLEO is consistent with that of FOCUS. A more formal fit of the $s$-wave parameters in CLEO data is in progress.

**FIG. 16:** The $s$-wave interference term for events below the $K_0^{*0}$ pole (left) and above the pole (right). The interference term depends on the $s$-wave phase relative to the phase average phase of each half of the Breit-Wigner. All of the $H_0(q^2)$ interference observed by FOCUS was also below the $K_0^{*0}$ pole as shown in Fig. 8.

**VI. FUTURE DIRECTIONS**

It will be interesting to pursue the non-parametric $D^+ \rightarrow K^{-}\pi^+ \ell^+\nu$ analysis with more data. One motivation is will be to further study the $h_0(q^2)$ form factor which appears to be somewhat different than $H_0(q^2)$. It would also be interesting to pursue tighter limits on possible $d$-wave and $f$-wave non-resonant contributions to $D^+ \rightarrow K^{-}\pi^+ \ell^+\nu$ and make more stringent tests of SPD. CLEO is slated to increase their luminosity at the $\psi(3770)$ from the 280 pb$^{-1}$ reported here to 750 pb$^{-1}$. In addition Surik Mehrabyan and I, are studying $D^+ \rightarrow K^{-}\pi^+ \mu^+\nu$ as well as $D^+ \rightarrow K^{-}\pi^+ e^+\nu$ in CLEO data. This is a some-
what challenging project since the CLEO muon detector was designed for higher energy B-meson running and the muons from charm semileptonic decay tend to range out before being identified. Hence special care must be exercised to reduce backgrounds. Besides increasing our statistics, the $D^+ \to K^- \pi^+ \mu^+ \nu$ should allow us to make the first measurements of the $H_T(q^2)$ form factor which is suppressed by a factor of $m_c^2/q^2$. Since this is a zero helicity factor, it can interfere with $H_0(q^2)$ and hence two new projectors will be required: one for the $H^2_T(q^2)$ term and one for $H_0(q^2) \times H_T(q^2)$ interference. At present the prognosis for making these measurements looks good.

VII. SUMMARY

Progress in understanding $D \to \text{vector } \ell^+ \nu$ decays was reviewed. These have historically been analyzed under the assumption of spectroscopic pole dominance (SPD). A recent result from BaBar was reviewed that used SPD to show that the form factors for $D^+ \to \phi \ell^+ \nu$ are consistent with those from $D^+ \to K^- \pi^+ \ell^+ \nu$ as expected from SU(3) symmetry. Experiments have obtained consistent results with the SPD assumption, but as of yet there have been no incisive tests of spectroscopic pole dominance. We concluded by describing a first non-parametric look at the $D^+ \to K^- \pi^+ \ell^+ \nu$ form factors. Although the results were very consistent with the traditional pole dominance fits, the data was not precise enough to incisively measure $q^2$ dependence of the axial and vector form factors and thus test SPD. This preliminary analysis did confirm the existence of an $s$-wave effect first observed by FOCUS [1], but was unable to obtain evidence for $d$ and $f$-waves.

Spectroscopy and Strong Decays of Charmed Baryons

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Spectroscopy and strong decays of the charmed baryons are reviewed. Possible spin-parity quantum numbers of several newly observed charmed baryon resonances are discussed. Strong decays of charmed baryons are analyzed in the framework of heavy hadron chiral perturbation theory in which heavy quark symmetry and chiral symmetry are synthesized.

1. Introduction

In the past years many new excited charmed baryon states have been discovered by BaBar, Belle and CLEO. In particular, $B$ factories have provided a very rich source of charmed baryons both from $B$ decays and from the continuum $e^+e^-\to c\bar{c}$. A new era for the charmed baryon spectroscopy is opened by the rich mass spectrum and the relatively narrow widths of the excited states. Experimentally and theoretically, it is important to identify the quantum numbers of these new states and understand their properties. Since the pseudoscalar mesons involved in the strong decays of charmed baryons are soft, the charmed baryon system offers an excellent ground for testing the ideas and predictions of heavy quark symmetry of the heavy quarks and chiral symmetry of the light quarks.

2. Spectroscopy

Charmed baryon spectroscopy provides an ideal place for studying the dynamics of the light quarks in the environment of a heavy quark. The charmed baryon of interest contains a charmed quark and two light quarks, which we will often refer to as a diquark. Each light quark is a triplet of the flavor SU(3). Since the orbital angular momentum between the diquark and the charmed quark. The lowest-lying orbitally excited baryon states are the $p$-wave charmed baryons. Denoting the quantum numbers $L_{\rho}$ and $L_{\lambda}$ as the eigenvalues of $L_{\rho}^2$ and $L_{\lambda}^2$, respectively, the $p$-wave heavy baryon can be either in the $(L_{\rho} = 0, L_{\lambda} = 1)$ $\lambda$-state or the $(L_{\rho} = 1, L_{\lambda} = 0)$ $\rho$-state. It is obvious that the orbital $\lambda$-state ($\rho$-state) is symmetric (antisymmetric) under the interchange of two light quarks $q_1$ and $q_2$. The total angular momentum of the diquark is $J_\ell = S_\rho + L_\rho$ and the total angular momentum of the charmed baryon is $J = S_c + J_\ell$. In the heavy quark limit, the spin of the charmed quark $S_c$ and the total angular momentum of the two light quarks $J_\ell$ are separately conserved.

There are seven lowest-lying $p$-wave $\Lambda_c$ arising from combining the charmed quark spin $S_c$ with light constituents in $J_\ell^{P_s} = 1^-$ state: three $J^P = \frac{1}{2}^-$ states, three $J^P = \frac{3}{2}^-$ states and one $J^P = \frac{5}{2}^-$ state. They form three doublets $\Lambda_c1(1^-), \Lambda_c1(3^-), \Lambda_c2(5^-)$ and one singlet $\Lambda_c0(1^-)$ in the notation $\Lambda_c(J^P)$, where we have used a tilde to denote the multiplets antisymmetric in the orbital wave functions under the exchange of two light quarks. Quark models [1] indicate that the untilded states for $\Lambda$- and $\Sigma$-type charmed baryons with symmetric orbital wave functions lie about 150 MeV below the tilde ones. The two states in each doublet with $J = J_\ell = \frac{3}{2}$ are nearly degenerate; their masses split only by a chromomagnetic interaction.

The next orbitally excited states are the positive parity excitations with $L_{\rho} + L_{\lambda} = 2$. There are two multiplets for the first positive-parity excited $\Lambda_c$ with the symmetric orbital wave function, corresponding to $L_{\lambda} = 2$, $L_{\rho} = 0, L = 2$ and $L_{\rho} = 2, L_{\lambda} = 2, L = 2$, see Table [1](for other charmed baryons, see [2] for details). For the case of $L_{\lambda} = L_{\rho} = 1$, the total orbital angular momentum $L_\ell$ of the diquark is 2, 1 or 0. Since the orbital states are antisymmetric under the interchange of two light quarks, we shall use a tilde to denote the $L_{\lambda} = L_{\rho} = 1$ states. The Fermi-Dirac statistics for baryons yields seven more multiplets for positive-parity excited $\Lambda_c$ states.

The observed mass spectra and decay widths of charmed baryons are summarized in Table [1]. For the experimental status of charmed baryons, see [3]. In the following we discuss some of the new excited charmed baryon states:

2.1. $\Lambda_c$

It is known that $\Lambda_c(2595)^+$ and $\Lambda_c(2625)^+$ form a doublet $\Lambda_c1(\frac{1}{2}^-, \frac{3}{2}^-)$ [4]. The dominant decay mode is $\Sigma_c\pi$ in an $S$ wave for $\Lambda_c1(\frac{1}{2}^-)$ and $\Lambda_c\pi\pi$ in a $P$ wave for $\Lambda_c1(\frac{3}{2}^-)$. (The two-body mode $\Sigma_c\pi$ is a $D$-wave in $\Lambda_c(\frac{3}{2}^-)$ decay.) This explains why the width
of \( \Lambda_c(2625)^+ \) is narrower than that of \( \Lambda_c(2595)^+ \).

\( \Lambda_c(2765)^+ \) is a broad state \(( \Gamma \approx 50 \text{ MeV}) \) first seen in \( \Lambda_c^+ \pi^+ \pi^- \) by CLEO [8]. It appears to resonate through \( \Sigma_c \) and probably also \( \Sigma_c^* \). However, whether it is a \( \Lambda_c^+ \) or a \( \Sigma_c^+ \) or whether the width might be due to overlapping states are not known. According to PDG [6], this state has a nickname, namely, \( \Lambda_c(2765)^+ \). The Skyrme model [2] and the quark model [1] suggest a \( J^P = \frac{3}{2}^+ \), \( \Lambda_c \) state with a mass 2742 and 2775 MeV, respectively. Therefore, \( \Lambda_c(2765)^+ \) could be a first positive-parity excitation of \( \Lambda_c \). However, two recent studies based on the relativistic quark model advocate a different assignment: a radial excitation \( \frac{1}{2}^+ \) by [8] and a negative-parity state with \( J^P = \frac{1}{2}^- \) by [9].

The state \( \Lambda_c(2880)^+ \) first observed by CLEO [3] in \( \Lambda_c^+ \pi^+ \pi^- \) was also seen by BaBar in the \( D^0p \) spectrum [10]. It was originally conjectured that, based on its narrow width, \( \Lambda_c(2880)^+ \) might be a \( \Lambda_c^0(\frac{1}{2}^-) \) state [8]. Recently, Belle has studied the experimental constraint on the \( J^P \) quantum numbers of \( \Lambda_c(2880)^+ \) [11]. The angular analysis of \( \Lambda_c(2880)^+ \rightarrow \Sigma_c^0 \pi^+ \pi^0 \) indicates that \( J = \frac{3}{2} \) is favored over \( J = \frac{1}{2} \) or \( \frac{5}{2} \). In the quark model, the candidates for the spin-\( \frac{3}{2} \) state are \( \Lambda_{c2}(\frac{3}{2}^+) \), \( \Lambda_{c3}(\frac{3}{2}^+) \), \( \Lambda'_{c2}(\frac{3}{2}^+) \), \( \Lambda'_{c3}(\frac{3}{2}^+) \), \( \Lambda''_{c2}(\frac{3}{2}^+) \) and \( \Lambda''_{c3}(\frac{3}{2}^+) \) (see Table 1). And only one of them has odd parity.

Belle has also studied the resonant structure of \( \Lambda_c(2880)^+ \rightarrow \Lambda_c^+ \pi^+ \pi^- \) and found the existence of the \( \Sigma_c^* \pi \) intermediate states [11]. The ratio of \( \Sigma_c^* \pi/\Sigma_c \pi \) is measured to be

\[
R = \frac{\Gamma(\Lambda_c(2880) \rightarrow \Sigma_c^* \pi)}{\Gamma(\Lambda_c(2880) \rightarrow \Sigma_c \pi)} = (24.1 \pm 6.4^{+1.1}_{-1.5})\%.
\]

Table I. The first positive-parity excitations of \( \Lambda \) charmed baryons and their quantum numbers. States with antisymmetric orbital wave functions (i.e. \( L_\rho = L_\lambda = 1 \)) under the interchange of two light quarks are denoted by a tilde. There are two multiplets \( \Lambda_{c2} \) and \( \Lambda_{c3} \) with symmetric orbital wave functions arising from the orbital states \( L_\rho = 0, L_\lambda = 2 \) and \( L_\rho = 2, L_\lambda = 0 \), respectively. We use a hat to distinguish between them.

For \( J^P = \frac{3}{2}^- \), \( \Lambda_c(2880) \) decays to \( \Sigma_c^* \pi \) and \( \Sigma_c \pi \) in a \( D \) wave and we obtain

\[
\frac{\Gamma(\Lambda_c(2880) \rightarrow \Sigma_c^* \pi)}{\Gamma(\Lambda_c(2880) \rightarrow \Sigma_c \pi)} = \frac{7}{2} \frac{\rho_s^2(\Lambda_c(2880) \rightarrow \Sigma_c^* \pi)}{\rho_s^0(\Lambda_c(2880) \rightarrow \Sigma_c \pi)} = 1.45,
\]

where the factor of 7/2 follows from heavy quark symmetry. Hence, the assignment of \( J^P = \frac{3}{2}^- \) for \( \Lambda_c(2880) \) is disfavored. For \( J^P = \frac{5}{2}^+ \), \( \Lambda_c, \hat{\Lambda}_c, \hat{\Lambda}'_{c2} \) and \( \hat{\Lambda}''_{c2} \) with \( J_F = 2 \) decay to \( \Sigma_c \pi \) in a \( F \) wave and \( \Sigma_c^* \pi \) in \( F \) and \( P \) waves. Neglecting the \( P \)-wave contribution for the moment,

\[
\frac{\Gamma(\Lambda_c(2880) \rightarrow \Sigma_c^* \pi)}{\Gamma(\Lambda_c(2880) \rightarrow \Sigma_c \pi)} = \frac{4}{5} \frac{\rho_s^2(\Lambda_c(2880) \rightarrow \Sigma_c^* \pi)}{\rho_s^0(\Lambda_c(2880) \rightarrow \Sigma_c \pi)} = 0.23.
\]

At first glance, it appears that this is in good agreement with experiment. However, the \( \Sigma_c^* \pi \) channel is available via a \( P \)-wave and is enhanced by a factor of \( 1/\rho_s^0 \) relative to the \( F \)-wave one. Unfortunately, we cannot apply heavy quark symmetry to calculate the contribution of the \( \Sigma_c^* \pi \) channel to the ratio \( R \) as the reduced matrix elements are different for \( P \)-wave and \( F \)-wave modes. In this case, one has to reply on a phenomenological model to compute the ratio \( R \). At any event, the \( \Sigma_c^* \pi \) mode produced in \( \Lambda_c(2880) \) is a priori not necessarily suppressed relative to \( \Sigma_c \pi \) channel.

Therefore, if \( \Lambda_c(2880)^+ \) is one of the states \( \Lambda_{c2}, \hat{\Lambda}_c, \hat{\Lambda}'_{c2} \) and \( \hat{\Lambda}''_{c2} \), the prediction \( R = 0.23 \) is not robust as it can be easily upset by the contribution from the \( P \)-wave \( \Sigma_c^* \pi \).

As for \( \hat{\Lambda}'_{c3}(\frac{3}{2}^+) \), it decays to \( \Sigma_c \pi, \Sigma_c \pi \) and \( \Lambda_c \pi \) all in \( F \) waves. Since \( J_L = 3, L_L = 2 \), it turns out that

\[
\frac{\Gamma(\hat{\Lambda}'_{c3}(\frac{3}{2}^+) \rightarrow \Sigma_c \pi)}{\Gamma(\hat{\Lambda}'_{c3}(\frac{3}{2}^+) \rightarrow \Sigma_c \pi)} = \frac{5}{4} \frac{\rho_s^2(\Lambda(2880) \rightarrow \Sigma_c \pi)}{\rho_s^2(\Lambda(2880) \rightarrow \Sigma_c \pi)} = 0.36.
\]

Although this deviates from the experimental measurement [11] by \( 1\sigma \), it is a robust prediction. This has motivated Chun-Khiang Chua and me to conjecture that that the first positive-parity excited charmed baryon \( \Lambda_c(2880)^+ \) could be an admixture of \( \Lambda_{c2}(\frac{3}{2}^+), \Lambda_{c3}(\frac{3}{2}^+) \) and \( \hat{\Lambda}'_{c3}(\frac{3}{2}^+) \) [2].

It is worth mentioning that very recently the Peking group [12] has studied the strong decays of charmed baryons based on the so-called \( ^3P_0 \) recombination model. For the \( \Lambda_c(2880) \), Peking group found that (i) the possibility of \( \Lambda_c(2880) \) being a radial excitation is ruled out as its decay into \( D^0p \) is prohibited in the \( ^3P_0 \) model if \( \Lambda_c(2880) \) is a first radial excitation of \( \Lambda_c \), and (ii) the only possible assignment is \( \Lambda_{c3}(\frac{3}{2}^+) \)
Table II Mass spectra and decay widths (in units of MeV) of charmed baryons taken from [2]. Except for the parity of the lightest \( \Lambda_c^0 \) and the spin-parity of \( \Lambda_c(2880)^+ \), none of the other \( J^P \) quantum numbers given in the table has been measured. One has to rely on the quark model to determine the spin-parity assignments.

<table>
<thead>
<tr>
<th>State</th>
<th>( J^P )</th>
<th>( S_{\pi} )</th>
<th>( L_{\pi} )</th>
<th>( J^P_{\pi} )</th>
<th>Mass</th>
<th>Width</th>
<th>Principal decay modes</th>
</tr>
</thead>
<tbody>
<tr>
<td>( \Lambda_c^+(2595) )</td>
<td>( \frac{3}{2}^+ )</td>
<td>0</td>
<td>0</td>
<td>0</td>
<td>2286.46 ± 0.14</td>
<td>weak</td>
<td>( \Sigma_c^+ ), ( \Lambda_c \pi )</td>
</tr>
<tr>
<td>( \Lambda_c(2625)^+ )</td>
<td>( \frac{3}{2}^+ )</td>
<td>0</td>
<td>1</td>
<td>1</td>
<td>2595.4 ± 0.6</td>
<td>3.6^{+2.0}_{-1.3}</td>
<td>( \Sigma_c^+ ), ( \Lambda_c \pi )</td>
</tr>
<tr>
<td>( \Lambda_c(2765)^+ )</td>
<td>( \frac{5}{2}^+ )</td>
<td>0</td>
<td>1</td>
<td>1</td>
<td>2628.1 ± 0.6</td>
<td>&lt; 1.9</td>
<td>( \Lambda_c \pi ), ( \Sigma_c \pi )</td>
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<tr>
<td>( \Lambda_c(2880)^+ )</td>
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<td>0</td>
<td>0</td>
<td>0</td>
<td>2766.6 ± 2.4</td>
<td>50</td>
<td>( \Sigma_c^+ ), ( \Lambda_c \pi \pi )</td>
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<tr>
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<td>0</td>
<td>0</td>
<td>2881.5 ± 0.3</td>
<td>5.5 ± 0.6</td>
<td>( \Sigma_c^+ ), ( \Lambda_c \pi \pi ), ( D^0 p )</td>
</tr>
<tr>
<td>( \Sigma_c(2455)^{1+} )</td>
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<td>1</td>
<td>0</td>
<td>1</td>
<td>2938.8 ± 1.1</td>
<td>13.0 ± 5.0</td>
<td>( \Sigma_c^+ ), ( \Lambda_c \pi \pi ), ( D^0 p )</td>
</tr>
<tr>
<td>( \Sigma_c(2455)^{0+} )</td>
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<td>0</td>
<td>1</td>
<td>2452.9 ± 0.4</td>
<td>&lt; 4.6</td>
<td>( \Lambda_c \pi )</td>
</tr>
<tr>
<td>( \Sigma_c(2520)^{1+} )</td>
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<td>0</td>
<td>1</td>
<td>2453.76 ± 0.18</td>
<td>2.2 ± 0.4</td>
<td>( \Lambda_c \pi )</td>
</tr>
<tr>
<td>( \Sigma_c(2520)^{0+} )</td>
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<td>0</td>
<td>1</td>
<td>2518.4 ± 0.6</td>
<td>14.9 ± 1.9</td>
<td>( \Lambda_c \pi )</td>
</tr>
<tr>
<td>( \Sigma_c(2800)^{1+} )</td>
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<td>1</td>
<td>2</td>
<td>2801.4^{+4.4}_{-6.6}</td>
<td>75^{+22}_{-17}</td>
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<td>0</td>
<td>2</td>
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<td>62^{+50}_{-30}</td>
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<td>0</td>
<td>2</td>
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<td>61^{+28}_{-18}</td>
<td>( \Lambda_c \pi ), ( \Sigma_c^+ ), ( \pi \pi ), ( \Lambda_c \pi \pi )</td>
</tr>
<tr>
<td>( \Xi_c^0 )</td>
<td>( \frac{1}{2}^+ )</td>
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<td>0</td>
<td>0</td>
<td>2467.9 ± 0.4</td>
<td>weak</td>
<td>( \Xi_c^+ )</td>
</tr>
<tr>
<td>( \Xi_c^+ )</td>
<td>( \frac{1}{2}^+ )</td>
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<td>0</td>
<td>1</td>
<td>2575.7 ± 3.1</td>
<td>weak</td>
<td>( \Xi_c^+ )</td>
</tr>
<tr>
<td>( \Xi_c^0 )</td>
<td>( \frac{1}{2}^+ )</td>
<td>1</td>
<td>0</td>
<td>1</td>
<td>2578.0 ± 2.9</td>
<td>weak</td>
<td>( \Xi_c^+ )</td>
</tr>
<tr>
<td>( \Xi_c(2645)^{+} )</td>
<td>( \frac{3}{2}^+ )</td>
<td>1</td>
<td>0</td>
<td>1</td>
<td>2646.6 ± 1.4</td>
<td>3.1</td>
<td>( \Xi^0 \pi )</td>
</tr>
<tr>
<td>( \Xi_c(2645)^{0} )</td>
<td>( \frac{3}{2}^+ )</td>
<td>1</td>
<td>0</td>
<td>1</td>
<td>2646.1 ± 1.2</td>
<td>&lt; 5.5</td>
<td>( \Xi_c \pi )</td>
</tr>
<tr>
<td>( \Xi_c(2790)^{+} )</td>
<td>( \frac{1}{2}^+ )</td>
<td>0</td>
<td>1</td>
<td>1</td>
<td>2789.2 ± 3.2</td>
<td>&lt; 15</td>
<td>( \Xi_c^+ )</td>
</tr>
<tr>
<td>( \Xi_c(2790)^{0} )</td>
<td>( \frac{1}{2}^+ )</td>
<td>0</td>
<td>1</td>
<td>1</td>
<td>2791.9 ± 3.3</td>
<td>&lt; 12</td>
<td>( \Xi_c^0 )</td>
</tr>
<tr>
<td>( \Xi_c(2815)^{+} )</td>
<td>( \frac{3}{2}^+ )</td>
<td>0</td>
<td>1</td>
<td>1</td>
<td>2816.5 ± 1.2</td>
<td>&lt; 3.5</td>
<td>( \Xi^0 \pi ), ( \Xi_1 \pi ), ( \pi \pi )</td>
</tr>
<tr>
<td>( \Xi_c(2815)^{0} )</td>
<td>( \frac{3}{2}^+ )</td>
<td>0</td>
<td>1</td>
<td>1</td>
<td>2818.2 ± 2.1</td>
<td>&lt; 6.5</td>
<td>( \Xi_c^+ ), ( \Xi_1 \pi ), ( \Xi_c \pi )</td>
</tr>
<tr>
<td>( \Xi_c(2890)^{0} )</td>
<td>( \frac{1}{2}^+ )</td>
<td>0</td>
<td>0</td>
<td>0</td>
<td>2971.1 ± 1.7</td>
<td>25.2 ± 3.0</td>
<td>see Table 7 of [2]</td>
</tr>
<tr>
<td>( \Xi_c(3055)^{+} )</td>
<td>( \frac{3}{2}^+ )</td>
<td>0</td>
<td>0</td>
<td>0</td>
<td>2971.7 ± 9.5</td>
<td>43.5</td>
<td>see Table 7 of [2]</td>
</tr>
<tr>
<td>( \Xi_c(3080)^{+} )</td>
<td>( \frac{3}{2}^+ )</td>
<td>0</td>
<td>0</td>
<td>0</td>
<td>3054.2 ± 1.3</td>
<td>17 ± 13</td>
<td>( \Lambda_c^0 K^0 ), ( \Lambda_c^+ K^+ )</td>
</tr>
<tr>
<td>( \Xi_c(3080)^{0} )</td>
<td>( \frac{3}{2}^+ )</td>
<td>0</td>
<td>0</td>
<td>0</td>
<td>3076.5 ± 0.6</td>
<td>6.2 ± 1.1</td>
<td>see Table 7 of [2]</td>
</tr>
<tr>
<td>( \Xi_c(3080)^{0} )</td>
<td>( \frac{3}{2}^+ )</td>
<td>0</td>
<td>0</td>
<td>0</td>
<td>3082.8 ± 2.3</td>
<td>5.2 ± 3.6</td>
<td>see Table 7 of [2]</td>
</tr>
<tr>
<td>( \Xi_c(3123)^{+} )</td>
<td>( \frac{3}{2}^+ )</td>
<td>0</td>
<td>0</td>
<td>0</td>
<td>3122.9 ± 1.3</td>
<td>4.4 ± 3.8</td>
<td>( \Lambda_c^0 K^0 ), ( \Lambda_c^+ K^+ )</td>
</tr>
<tr>
<td>( \Omega_c^0 )</td>
<td>( \frac{3}{2}^+ )</td>
<td>1</td>
<td>0</td>
<td>1</td>
<td>2697.5 ± 2.6</td>
<td>weak</td>
<td>( \Omega_c )</td>
</tr>
<tr>
<td>( \Omega_c(2770)^{0} )</td>
<td>( \frac{3}{2}^+ )</td>
<td>1</td>
<td>0</td>
<td>1</td>
<td>2768.3 ± 3.0</td>
<td>weak</td>
<td>( \Omega_c )</td>
</tr>
</tbody>
</table>

\[
\frac{\Gamma (\Lambda_c(2595)^{0+} \rightarrow \Sigma_c^+ \pi^0)}{\Gamma (\Lambda_c(2625)^{0+} \rightarrow \Sigma_c \pi)} = 0.06,
\]

\[
\frac{\Gamma (\hat{\Lambda}_c(2595)^{0+} \rightarrow \Sigma^+_c \pi)}{\Gamma (\hat{\Lambda}_c(2625)^{0+} \rightarrow \Sigma^+_c \pi)} = 78.3.
\]  

Both symmetric states \( \Lambda_{c2} \) and \( \hat{\Lambda}_{c2} \) are thus ruled out as the predicted ratio \( R \) is either too small or too big compared to experiment. However, the assignment of \( \Lambda_c^{0+}(5^{+}) \) for \( \Lambda_c(2880) \) has an issue with the spectrum: The quark model indicates a \( \Lambda_{c2}(5^{+},J^P=5/2^+) \) state around 2910 MeV which is close to the mass of \( \Lambda_c(2880) \), while the mass of \( \Lambda_c^{0+}(5^{+}) \) is higher [1].

It is interesting to notice that, based on the diquark idea, the quantum numbers \( J^P = 5/2^+ \) have been correctly predicted in [13] for the \( \Lambda_c(2880) \) before the Belle experiment.

The highest \( \Lambda_c(2940)^+ \) was first discovered by BaBar in the \( D^0 p \) decay mode [10] and confirmed by

Belle in the decays $\Sigma^0 \pi^+, \Sigma^+_c \pi^-$ which subsequently decay into $\Lambda_c^+ \pi^+ \pi^-$ [11]. Since the mass of $\Lambda_c(2940)^+$ is barely below the threshold of $D^{*0}p$, this observation has motivated the authors of [14] to suggest an exotic molecular state of $D^{*0}p$ with a binding energy of order 6 MeV and $J^P = \frac{3}{2}^-$ for $\Lambda_c(2940)^+$. The quark potential model predicts a $\frac{3}{2}^-$ $\Lambda_c$ state at 2900 MeV and a $\frac{3}{2}^+ \Lambda_c$ state at 2910 MeV [1]. A similar result of 2906 MeV for $\frac{3}{2}^+ \Lambda_c$ is also obtained in the relativistic quark model [15].

Given the uncertainty of order 50 MeV for the quark model calculation, this suggests that the possible allowed $J^P$ numbers of the highest $\Lambda_c(2940)^+$ are $\frac{5}{2}^-$ and $\frac{7}{2}^+$. Hence, the potential candidates are $\check{\Lambda}_c \Sigma(\frac{5}{2}^-)$, $\check{\Lambda}_c \Sigma(\frac{7}{2}^+)$, $\check{\Lambda}'_c (\frac{5}{2}^-)$, $\check{\Lambda}'_c (\frac{7}{2}^+)$, $\check{\Lambda}''_c (\frac{3}{2}^-)$ and $\check{\Lambda}''_c (\frac{5}{2}^+)$. Since the predicted ratios differ significantly for different $J^P$ quantum numbers, the measurements of the ratio of $\Sigma_c^+\pi/\Sigma_c \pi$ will enable us to discriminate the $J^P$ assignments for $\Lambda_c(2940)$ [2]. Note that it has been argued in [8] that $\Lambda_c(2940)$ is the first radial excitation of $\Sigma_c$ (not $\Lambda_c$!) with $J^P = 3/2^+$.

2.2. $\Sigma_c$

The highest isotriplet charmed baryons $\Sigma_c(2800)^{+++,+0}$ decaying to $\Lambda_c^+ \pi$ were first measured by Belle [14]. They are most likely to be the $J^P = \frac{3}{2}^+ \Sigma_c(2800)^{+++,+0}$ states because the $\Sigma_c(2760)^{++}$ baryons decays principally into the $\Lambda_c \pi$ system in a $P$-wave, while $\Sigma_c(2780)^{++,0}$ decays mainly to the two pion system $\Lambda_c \pi \pi$ in a $P$-wave. The state $\Sigma_c(2800)^{++,0}$ can decay into $\Lambda_c \pi$ in an $S$-wave, but it is very broad with width of order 406 MeV. Therefore, $\Sigma_c(2800)^{++,0}$ are likely to be $\Sigma_c(2800)^{++}$ with a possible small mixing with $\Sigma_c(2780)^{++}$.

2.3. $\Xi_c$

The states $\Xi_c(2790)$ and $\Xi_c(2815)$ form a doublet $\Xi_c(\frac{1}{2}^-, \frac{3}{2}^-)$. Since the diquark transition $1^- \rightarrow 0^+ + \pi$ is prohibited, $\Xi_c(\frac{1}{2}^-, \frac{3}{2}^-)$ cannot decay to $\Xi_c \pi$. The dominant decay mode is $[\Xi_c^0 \pi]_S$ for $\Xi_c(\frac{1}{2}^-)$ and $[\Xi_c^+ \pi]_S$ for $\Xi_c(\frac{3}{2}^-)$, where $\Xi_c^0$ stands for $\Xi_c(2645)$.

The new charmed strange baryons $\Xi_c(2980)^+$ and $\Xi_c(3080)^+$ that decay into $\Lambda_c^+ K^- \pi^+$ were first observed by Belle [17] and confirmed by BaBar [18]. For the charmed states $\Xi_c(2980)$ and $\Xi_c(3080)$, they could be the first positive-parity excitations of $\Xi_c$ in viewing of their large masses. Since the mass difference between the antitriplets $\Lambda_c$ and $\Xi_c$ for $J^P = \frac{1}{2}^+, \frac{3}{2}^-$ is of order 180 $\sim$ 200 MeV, it is conceivable that $\Xi_c(2980)$ and $\Xi_c(3080)$ are the counterparts of $\Lambda_c(2765)$ and $\Lambda_c(2880)$, respectively, in the strange charmed baryon sector. As noted in passing, the state $\Lambda_c(2765)^+$ could be an even-parity orbital excitation or a radial excitation and $\Lambda_c(2880)$ has the quantum numbers $J^P = \frac{5}{2}^+$, it is thus tempting to assign $J^P = 1^+_2$ for $\Xi_c(2980)$ and $\frac{5}{2}^+$ for $\Xi_c(3080)$. The possible strong decays of the first positive-parity excitations of the $\Xi_c$ states are summarized in Table VII of [2]. Since the two-body modes $\Xi_c^0 \pi$, $\Lambda_c K$, $\Xi_c^+ \pi$ and $\Sigma_c K$ are in $P (F)$ waves and the three-body modes $\Xi_c^+ \pi \pi$ and $\Lambda_c K \pi$ are in $S (D)$ waves in the decays of $\frac{1}{2}^+_2 (\frac{5}{2}^+_2)$, this explains why $\Xi_c(2980)$ is broader than $\Xi_c(3080)$. Since both $\Xi_c(2980)$ and $\Xi_c(3080)$ are above the $DA$ threshold, it is important to search for them in the $DAC$ spectrum as well.

Two new $\Xi_c$ resonances $\Xi_c(3055)$ and $\Xi_c(3123)$ were recently reported by BaBar [19] with masses and widths shown in Table III.

2.4. $\Omega_c$

At last, the $J^P = \frac{3}{2}^+ \Omega_c(2770)$ charmed baryon was recently observed by BaBar in the decay $\Omega_c(2770)^0 \rightarrow \Omega_c^0 \gamma$ [20]. With this new observation, the $\frac{3}{2}^+$ sextet is finally completed. However, it will be very difficult to measure the electromagnetic decay rate because the width of $\Omega_c^0$, which is predicted to be of order 0.9 keV [21], is too narrow to be experimentally resolvable.

The possible spin-parity quantum numbers of the newly discovered charmed baryon resonances that have been suggested in the literature are summarized in Table III. Some of the predictions are already ruled out by experiment. For example, $\Lambda_c(2880)$ has $J^P = \frac{3}{2}^+$ as seen by Belle. Certainly, more experimental studies are needed in order to pin down the quantum numbers.

3. Strong decays

Due to the rich mass spectrum and the relatively narrow widths of the excited states, the charmed baryon system offers an excellent ground for testing the ideas and predictions of heavy quark symmetry and light flavor SU(3) symmetry. The pseudoscalar mesons involved in the strong decays of charmed baryons such as $\Sigma_c \rightarrow \Lambda_c \pi$ are soft. Therefore, heavy quark symmetry of the heavy quark and chiral symmetry of the light quarks will have interesting implications for the low-energy dynamics of heavy baryons interacting with the Goldstone bosons.

The strong decays of charmed baryons are most conveniently described by the heavy hadron chiral Lagrangians in which heavy quark symmetry and chiral symmetry are incorporated [22, 23]. The Lagrangian...
involves two coupling constants $g_1$ and $g_2$ for $P$-wave transitions between s-wave and s-wave baryons [22], six couplings $h_3 - h_7$ for the S-wave transitions between s-wave and p-wave baryons, and eight couplings $h_8 - h_{15}$ for the D-wave transitions between s-wave and p-wave baryons [24].

### 3.1. Strong decays of s-wave charmed baryons

In principle, the coupling $g_1$ can be determined from the decay $\Sigma^+ \rightarrow \Sigma^0 \pi^+$, unfortunately, this strong decay is kinematically prohibited since the mass difference between $\Sigma^+$ and $\Sigma^0$ is only of order $65$ MeV. Consequently, the coupling $g_1$ cannot be extracted directly from the strong decays of heavy baryons. As for the coupling $g_2$, one can use the measured rates of $\Sigma^{++} \rightarrow \Lambda^+_c \pi^+$, $\Sigma^{+0} \rightarrow \Lambda^+_c \pi^0$ and $\Sigma^{00} \rightarrow \Lambda^+_c \pi^-$ as inputs to obtain

$$ |g_2| = 0.605^{+0.039}_{-0.043}, \quad 0.57 \pm 0.04, \quad 0.60 \pm 0.04, \quad (6) $$

respectively, where we have neglected the tiny contributions from electromagnetic decays. Hence, the averaged $g_2$ is

$$ |g_2| = 0.591 \pm 0.023. \quad (7) $$

Using this value of $g_2$, the predicted total width of $\Xi^+_c \pi^+$ is found to be in the vicinity of the current limit $\Gamma(\Xi^+_c \pi^+) < 3.1$ MeV [22].

It is clear from Table [IV] that the strong decay width of $\Sigma_c$ is smaller than that of $\Sigma^+_c$ by a factor of $\sim 7$, although they will become the same in the limit of heavy quark symmetry. This is ascribed to the fact that the c.m. momentum of the pion is around $90$ MeV in the decay $\Sigma_c \rightarrow \Lambda_c \pi$, while it is two times bigger in $\Sigma^+_c \rightarrow \Lambda_c \pi$. Since $\Sigma_c$ states are significantly narrower than their spin-1 counterparts, this explains why the measurement of their widths came out much later.

### 3.2. Strong decays of p-wave charmed baryons

Some of the $S$-wave and $D$-wave couplings of p-wave baryons to s-wave baryons can be determined. In principle, the coupling $h_2$ is readily extracted from $\Lambda_c(2595)^+ \rightarrow \Sigma^0 \pi^+$ with $\Lambda_c(2595)$ being identified as $\Lambda_{c1}(1^{++})$. However, since $\Lambda_c(2595)^+ \rightarrow \Sigma_c \pi$ is kinematically barely allowed, the finite width effects of the intermediate resonant states could become important [26]. Before proceeding to a more precise determination of $h_2$, we make several remarks on the partial widths of $\Lambda_c(2595)^+$ decays. (i) PDG [8] has assumed the isospin relation, namely, $\Gamma(\Lambda^+_c \pi^+ \pi^-) = 2\Gamma(\Lambda^+_c \pi^0 \pi^0)$ to extract the branching ratios for $\Sigma_c \pi$ modes. However, the decay $\Lambda_c(2595) \rightarrow \Lambda_c \pi \pi$ occurs very close to the threshold as $m_{\Lambda_c(2595)} - m_{\Lambda_c} = 308.9 \pm 0.6$ MeV. Hence, the phase space is very sensitive to the small isospin-violating mass differences between members of pions and charmed Sigma baryon multiplets. Since the neutral pion is slightly lighter than the charged one, it turns out that both $\Lambda^+_c \pi^0 \pi^0$ and $\Lambda^+_c \pi^0 \pi^0$ have very similar rates. (ii) Taking $B(\Lambda_c(2595)^+ \rightarrow \Sigma^+_c \pi^+ \pi^-) \approx 0.5$ and using the measured ratios of $\Lambda_c(2595)^+ \rightarrow \Sigma^+_c \pi^+$ and $\Sigma^0 \pi^+$ rela-

\[\begin{array}{|c|c|c|c|c|c|c|}
\hline
\text{Decay} & \text{Expt.} & \text{HHChPT} \\
\hline
\Sigma^+ c^+ \rightarrow \Lambda^+_c \pi^+ & 2.23 \pm 0.30 & \text{input} \\
\Sigma^+ c^+ \rightarrow \Lambda^+_c \pi^0 & < 4.6 & 2.5 \pm 0.2 \\
\Sigma^0 c^0 \rightarrow \Lambda^+_c \pi^- & 2.2 \pm 0.4 & \text{input} \\
\Sigma^+ (2520)^+ \rightarrow \Lambda^+_c \pi^+ & 14.9 \pm 1.9 & \text{input} \\
\Sigma^+ (2520)^+ \rightarrow \Lambda^+_c \pi^0 & < 17 & 16.6 \pm 1.3 \\
\Sigma^0 (2520)^0 \rightarrow \Lambda^+_c \pi^- & 16.1 \pm 2.1 & \text{input} \\
\Xi^+ (2645)^0 \rightarrow \Xi^0 \pi^+ \pi^0 & < 3.1 & 2.7 \pm 0.2 \\
\Xi^+ (2645)^0 \rightarrow \Xi^0 \pi^0 \pi^0 & < 5.5 & 2.8 \pm 0.2 \\
\hline
\end{array} \]
Table V Same as Table IV except for p-wave charmed baryons.

<table>
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<tr>
<th>Decay</th>
<th>Expt.</th>
<th>HHCChPT</th>
</tr>
</thead>
<tbody>
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<td>input</td>
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<tr>
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<td>0.72^{+0.43}_{-0.30}</td>
</tr>
<tr>
<td>$\Lambda_c(2595)^+ \to (\Sigma_c^+ \pi^-)$</td>
<td>0.67^{+0.41}_{-0.31}</td>
<td>0.77^{+0.46}_{-0.32}</td>
</tr>
<tr>
<td>$\Lambda_c(2625)^+ \to (\Sigma_c^0 \pi^-)$</td>
<td>1.5^{+0.93}_{-0.65}</td>
<td></td>
</tr>
<tr>
<td>$\Lambda_c(2625)^+ \to (\Sigma_c^+ \pi^-)$</td>
<td>&lt; 0.10</td>
<td>&lt; 0.029</td>
</tr>
<tr>
<td>$\Lambda_c(2625)^+ \to (\Sigma_c^+ \pi^-)$</td>
<td>&lt; 0.09</td>
<td>&lt; 0.029</td>
</tr>
<tr>
<td>$\Lambda_c(2625)^+ \to (\Sigma_c^+ \pi^-)$</td>
<td>&lt; 0.19</td>
<td>&lt; 0.21</td>
</tr>
<tr>
<td>$\Sigma_c(2800)^0 \to \Lambda_c^- \pi^0$</td>
<td>75^{+22}_{-17}</td>
<td>input</td>
</tr>
<tr>
<td>$\Sigma_c(2800)^0 \to \Lambda_c^- \pi^-$</td>
<td>62^{+0.09}_{-0.08}</td>
<td>input</td>
</tr>
<tr>
<td>$\Xi_c(2800)^0 \to \Xi_c^- \pi^0$</td>
<td>61^{+0.18}_{-0.17}</td>
<td>input</td>
</tr>
<tr>
<td>$\Xi_c(2800)^0 \to \Xi_c^- \pi^-$</td>
<td>&lt; 15</td>
<td>8.0^{+4.7}_{-3.4}</td>
</tr>
<tr>
<td>$\Xi_c(2800)^0 \to \Xi_c^- \pi^-$</td>
<td>&lt; 12</td>
<td>8.5^{+5.0}_{-3.5}</td>
</tr>
<tr>
<td>$\Xi_c(2815)^0 \to \Xi_c^- \pi^0$</td>
<td>&lt; 3.5</td>
<td>3.4^{+2.0}_{-1.4}</td>
</tr>
<tr>
<td>$\Xi_c(2815)^0 \to \Xi_c^- \pi^-$</td>
<td>&lt; 6.5</td>
<td>3.6^{+2.1}_{-1.5}</td>
</tr>
</tbody>
</table>

Assuming the coupling $h_2$ obtained from (10) and the experimental observation that the $\Xi_c^\pi\pi$ mode in $\Xi_c(2815)$ decays is consistent with being entirely via $\Xi_c^\pi\pi$, the predicted $\Xi_c(2790)$ and $\Xi_c(2815)$ widths are shown in Table V. The predictions are consistent with the current experimental limits.

Some information on the coupling $h_3$ can be inferred from the strong decays of $\Sigma_c(2800)$. Assuming the widths of the states $\Sigma_c(2800)^0, \Sigma_c^0$ and $\Sigma_c^+$ are dominated by the two-body $D$-wave modes $\Lambda_c^- \pi$, $\Sigma_c^- \pi$ and $\Sigma_c^0 \pi$, and applying the quark model relation $|h_3| = |h_1| |h_2|$, we then have

$$|h_3| \leq (0.86^{+0.08}_{-0.10}) \times 10^{-3} \text{MeV}^{-1},$$

which improves the previous limit (10) by a factor of 4.

Acknowledgments

I wish to thank Chun-Khiang Chua for collaboration on this interesting subject and the organizers for organizing this very stimulating workshop.

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Hints of a New Spectroscopy

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There are several reasons to believe that some of the new particles observed at B-factories have no-ordinary quark composition. We briefly illustrate the diquark-antidiquark model and the recent experimental discoveries which confirm some of its most striking predictions.

I. INTRODUCTION

The field of hadron spectroscopy has been revitalized by the discovery of a number of new particles at B-factories. The first of the series is a 1$^{+\pm}$ state, the X(3872), decaying to J/$\psi\pi\pi$, found by Belle and later confirmed by CDF, D0 and BaBar [1]. The difficulty in interpreting this particle as a standard charmonium has opened the way to alternative interpretations. The diquark-antidiquark, [cq][\bar{c}q] picture [2], at the moment, seems the most promising. There are at least two reasons for this:

- The X(3872) decays with the same rate in J/$\psi\rho$ and J/$\psi\omega$, therefore maximally violating isospin. In the diquark picture, two states are needed to explain this decay pattern, namely an $X_u$ and an $X_d$, with a difference in mass of the order of $m_u - m_d$. Recently Belle and BaBar have shown the existence of a second X at a mass of 3875 MeV, confirming the diquark model prediction [3].

- The stringiest prediction of the diquark model was the existence of charged particles decaying to J/$\psi(\pi^+\pi^-\rho^0)$ [2]. Last summer a state of this kind has been discovered by Belle: the Z(4430) decaying to J/$\psi\pi^+$ [4].

After the discovery of the X(3872), BaBar has found a new state produced in ISR, the Y(4260) [4]. Again a charmonium interpretation was not tenable for this particle, despite its decay to J/$\psi\pi\pi$. The diquark model explains the Y(4260) as a 1$^{-}\pm$ state made up of two diquarks in P-wave which decays to $f_0(980)$, a strong candidate for a light four-quark state, and J/$\psi$. The increasing number of X, Y, Z states being found, casts serious doubts on the fact that all of them can be explained, separately, as effects (threshold effects, cusp effects...) or loosely bound molecules [4] of open charm mesons. If other ways of aggregation of quarks in matter are indeed possible, a unified explanation of all these particles, and of the coming ones (hopefully!), will emerge clearly from data.

II. FROM LIGHT SCALARS TO XYZ

The strongest theoretical motivation for the diquark-antidiquark picture lies in its consistent description of the scalar mesons below 1 GeV, namely $f_0$, $a_0$, $\kappa$, $\sigma$. These are likely bound states of a spin zero diquark

$$q_i\bar{q}_a = \epsilon_{ijk}\epsilon_{ab\gamma}q_i^{[\alpha}_{(\gamma}\gamma^k\kappa^\gamma, \hspace{1cm} (1)$$

where latin indices label flavor and greek letters are for color, and an anti-diquark $\bar{q}^0_\alpha$. The color is saturated as in a standard $q\bar{q}$ meson: $q^a_i\bar{q}_a$. Therefore, since a spin zero diquark is in a 3-flavor representation, nonets of $qq$ states are allowed (crypto-exotic states). The sub-GeV scalar mesons represent most likely in the lowest tetraquark nonet.

The $qq$ model of light-scalars is very effective at explaining the most striking feature of these particles: the inverted pattern in the mass-versus-I$_3$ diagram [6], with respect to what observed for ordinary $q\bar{q}$ mesons. This aspect is not explicable using a $q\bar{q}$ model. For example, in the $qq$ model, the $f_0(980)$ should be an $s\bar{s}$ state [6] while the $I = 1$, $a_0(980)$, should be a $u\bar{u} + d\bar{d}$ state. If this were the case, the degeneracy of the two particles would appear quite mysterious.

Beyond a correct description of the mass-I$_3$ pattern, the tetraquark model offers the possibility to explain the decay rates of scalars at a level never reached in standard $q\bar{q}$ descriptions. The decay Lagrangian into two pseudoscalar mesons, e.g. $\sigma \rightarrow \pi\pi$, is:

$$\mathcal{L}_{exch.} = c_fS_i^j\epsilon^{jtu}\epsilon_{irs}\Pi_i^r\Pi_u^s, \hspace{1cm} (2)$$

where $i,j$ are the flavor labels of $q^i$ and $\bar{q}^j$, while $r,s,t,u$ are the flavor labels of the quarks $q^r$, $q^s$ and $\bar{q}^t$, $\bar{q}^u$. $c_f$ is the effective coupling weighting this interaction term and $S,\Pi$ are the scalar and pseudoscalar matrices. This Lagrangian describes the quantum exchange amplitude for the quarks to tunnel out of their diquark shells to form ordinary mesons [6]. Such mechanism is the alternative to the color string breaking $qq\rightarrow B\bar{B}$, i.e., a baryon-antibaryon decay, phase-space forbidden to sub-GeV scalar mesons.

The problem with (2) is that it is not able to describe the decay $f_0 \rightarrow \pi\pi$ because $f_0 = (q^i\bar{q}^i + q^j\bar{q}^j)/\sqrt{2}$, being 1, 2, 3 the $u, d, s$ flavors so that, see
III. HEAVY-LIGHT DIQUARKS

The successful theoretical interpretation of the light scalar mesons in terms of diquarks suggest that such structures could exist also at higher mass scales. An heavy-light diquark is still bound in color $3_c$, but: 1. We cannot state anymore that there is a Fermi statistics that forces the diquarks to be in $3_f$ or $6_f$ depending if the diquark spin is zero or one (a charm quark by no way can be considered identical to a light $q$ with respect to strong interactions); 2. The spin-spin interaction is weakened by 1/mc. According to lattice studies, light diquarks are preferably formed with spin= 0. On the other hand, heavy-light diquarks could appear equally well in spin 0 and 1. Therefore our Ansatz for heavy-light diquarks is:

$$Q_{1\alpha}^i = \epsilon_{\alpha\beta\gamma} Q_{C}^\beta \gamma q_i^\gamma \text{ spin } 1^- \quad (3)$$
$$Q_{0\alpha} = \epsilon_{\alpha\beta\gamma} Q_{C}^\beta \gamma q_i^\gamma \text{ spin } 0^+ \quad (4)$$

The flavor $i$ is carried by the light quark, while $Q$ = c for all $X,Y,Z$. Such spin 1 diquarks are likely the building blocks of the new particles, and since the flavor of Q is the flavor of the light quark, we can still accommodate particles like the partners in $(X,3872)$ and its partners in $SU(3)$ nonets. In our notation:

$$X(3872) = X_d = Q_2^0 Q_2^0 + Q_2^1 Q_1^1 \quad (5)$$
$$X(3875) = X_u = Q_1^0 Q_1^0 + Q_1^1 Q_1^1 \quad (6)$$
$$Y(4260) = (Q_1^0 Q_0^0)_p\text{-wave} \quad (7)$$
$$Z(4433) = (Q_1^0 Q_1^0 + Q_1^1 Q_1^1)_{2S\text{-wave}} \quad (8)$$

We have shown in [12] that, assuming:

$$X_d \to J/\psi\pi^+\pi^- \quad (9)$$
$$X_u \to D^0 \bar{D}^0 \pi^0, \quad (10)$$

we obtain a simple rule for the ratios of branching ratios of $B$ decays. With an obvious notation:

$$\frac{B^0}{B^+}_{KJ/\psi\pi} = \left( \frac{B^0}{B^+}_{KDD\pi} \right)^{-1} \quad (11)$$

This is to be confronted with the most recent experimental data, giving:

$$0.94 \pm 0.24 \pm 0.10 = \frac{1}{1.33 \pm 0.69 \pm 0.52} \quad (12)$$

Such quite reasonable agreement gives credit to our assignments for $X(3872)$ and $X(3875)$. Clearly enough, the situation of the remaining $1^{++}$ charged partners has to be clarified. We have reasons to believe that they can be very broad and this could cast some doubts on their actual visibility.

TABLE I: Some $L = 0$ neutral tetraquark wave functions constructible from diquarks. Consider that $C|ff\rangle = (1)^{L+S} |ff\rangle$.

<table>
<thead>
<tr>
<th>$J^{PC}$</th>
<th>Wave Funct.</th>
</tr>
</thead>
<tbody>
<tr>
<td>0++</td>
<td>$Q_0^0 Q_0^0 \vee (Q_1^1 Q_1^1)_{J=0}$</td>
</tr>
<tr>
<td>1++</td>
<td>$Q_0^0 Q_0^0 \vee (Q_1^1 Q_1^1)_{J=1}$</td>
</tr>
<tr>
<td>1−+</td>
<td>$Q_1^0 Q_1^0 \vee (Q_1^1 Q_1^1)_{J=2}$</td>
</tr>
<tr>
<td>2++</td>
<td>$(Q_1^1 Q_1^1)_{J=3}$</td>
</tr>
</tbody>
</table>

IV. $Z(4433)$

As mentioned in the introduction, the $Z(4433)$, most likely a $1^{+-}$ state, is the first charged particle observed that plausibly fits very well a diquark-antidiquark interpretation. Due to its decay to $\psi(2S)\pi^{+}$ we think that $Z(4433)$ is itself a radial excitation of one of the lowest lying $1^{+-}$ states predicted by the tetraquark model to be at a mass of $\sim 3880$ MeV [2]. In particular it is striking to observe that:

$$M(\psi(2S)) - M(J/\psi) \approx 590 \text{ MeV} \approx 4433 - 3880 \text{ MeV}$$

So the search of $X(1^{+-};1S)$ states, neutral and charged, is particularly urgent.

Scanning higher energy regions one approaches baryon-baryon thresholds, like $2M_{\Lambda} = 4572$ MeV and $M_{\Lambda} + M_{\Sigma_c} = 4379$ MeV. As mentioned above, the baryon-antibaryon decay channel is expected to be quite natural for a diquark-antidiquark system. This would mean that, above a certain threshold, all charmed tetraquark states are expected to be very broad.

V. 1/N AND TETRAQUARKS

In the large number of colors 1/N expansion, four-quark states of the kind $q_i q^a \bar{q} \bar{q}^a (= \mathcal{O}(x))$ are suppressed, in the sense that they do not appear as leading terms in the 1/N expansion of the two-point correlation function $\langle \mathcal{O}(x)\mathcal{O}(0) \rangle$. Indeed, the leading term of such a correlator would represent a disconnected graph (see the left diagram in Fig. 1).

But if we replace quarks with diquarks: $q^a \to q^\alpha$, there are no disconnected parts and the leading color diagram looks like the one depicted in Fig. 2. In other
words, if the diquarks are assumed to be the building blocks, we evade the $1/N$ difficulty with standard tetraquarks.

This reminds of the Corrigan-Ramond (CR) [13] large $N$ limit where quarks and ‘larks’ are introduced, transforming as the $N$ and $N(N-1)/2$ representations of $SU(N)$. Larks, $\ell$, are therefore antisymmetric objects, $\ell_{\alpha\beta} = -\ell_{\beta\alpha}$, coinciding with antiquarks if $N = 3$. A theory of only larks is equivalent to QCD. In the CR expansion, it is therefore possible to consider baryons at large $N$, if the baryon is represented by a color saturated $qq\ell\bar{\ell}$ state (for three quarks one can neutralize the color with an $\epsilon_{\alpha\beta\gamma}$ if the colors are three: there are no color singlets made up of three quarks for $N > 3$). Since larks have two color indices (like gluons in the ‘t Hooft limit [14], but with the arrow pointing in the same directions - gluon lines $A_\beta^\gamma$ have opposite arrows - ), the large $N$ power counting works differently (see Fig. 3). A diquark-antidiquark state is then like a lark-antilark state $\ell\bar{\ell}$ (see Fig. 2).

VI. CONCLUSIONS: BEYOND SPECTROSCOPY

The interest for such problems has implications beyond spectroscopy itself. This is not merely a chemistry exercise, aimed at a classification of new ‘elements’ whose nature has very little impact to QCD fundamental issues. QCD is extremely predictive only in a narrow range of very high energy phenomena, while the study of hadron structures and interactions remains a very difficult field which is mainly approached with the use of Effective Theories and Lattice studies. Being able to assess that new forms of aggregation of quarks are possible, such as diquarks, opens a window on a territory poorly known. Moreover, diquarks are essential to the theory of color superconductivity [12], which is at the basis of the comprehension of an entirely new sector of the QCD phase diagram.

The skepticism of the community about the tetraquarks is mainly driven by the shock following the ‘discovery’ and the disappearance of the baryonic pentaquarks. It is quite possible that multiquark baryons might exist at higher masses. Anyway, there is no clear logical connection between the four-quark mesons we are discussing here and the pentaquark baryons; one should also remind that the case for four-quark mesons is based on the phenomenology of light scalars since a very long time. Some other recent investigations on diquark based tetraquark charmed mesons can be found in [16].

We believe that a strong experimental effort aimed at searching the new predicted particles and clearly discriminate between models is of great importance for a progress in the comprehension of key aspects of non-perturbative QCD.

Acknowledgments

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Charm Meson Spectroscopy at BaBar and CLEO-c

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In this mini-review we report on the most recent progress in charm meson spectroscopy. We discuss the precision measurements performed by the BaBar and CLEO-c experiments in the non strange charm meson part and we present the newly discovered strange charm meson excited states.

1. Introduction

During the last few years many new D, Ds, charmonium, and charmed baryon excited states have been discovered. Some of these states were not expected theoretically: their masses, widths, quantum numbers, and decay modes did not fit the existing spectroscopic classification, which was based mostly on potential model calculations. The theoretical models had to be improved and new approaches have been developed to explain the data; the possibility of a non-quark-antiquark interpretation of these states has also been widely discussed. Charmonium, and charmed baryon excited states results are discussed elsewhere in these proceedings. In this report an overview of recent results on non strange charm mesons production is presented. Then, recent results on excited DsJ meson production will be presented and their behavior will be discussed.

2. Non Strange Charm Mesons

2.1. Measurement of the Absolute Branching Fractions $B \rightarrow D\pi, D^*\pi, D^{**}\pi$ with a Missing Mass Method

Our understanding of hadronic $B$-meson decays has improved considerably during the past few years with the development of models based on the Heavy Quark Effective Theory (HQET), where collinear [1, 2] or $k_T$ [3, 4] factorization theorems are considered. Models such as the QCD-improved Factorization (QCDF) [5, 6] and the Soft Collinear Effective Theory (SCET) [1, 7] use the collinear factorization, while the perturbative QCD (pQCD) approach [8, 9] uses the $k_T$ factorization. In these models the amplitude of the $B \rightarrow D^{(*)}\pi$ two-body decay carries information about the difference $\delta$ between the strong-interaction phases of the two isospin amplitudes $A_{1/2}$ and $A_{3/2}$ that contribute [10, 11]. A non-zero value of $\delta$ provides a measure of the departure from the heavy-quark limit and the importance of the final-state interactions in the $D^{(*)}\pi$ system. With the measurements by the BaBar [12] and BELLE [13] experiments of the color-suppressed $B$ decay $\bar{B}^0 \rightarrow D^{(*)0}\pi^0$ providing evidence for a sizeable value of $\delta$, an improved measurement of the color-favored decay amplitudes ($B^- \rightarrow D^{(*)0}\pi^-$ and $\bar{B}^0 \rightarrow D^{(*)+}\pi^-$) is of renewed interest. In addition, the study of $B$ decays into $D$, $D^*$, and $D^{**}$ mesons will allow tests of the spin symmetry [14–17] imbedded in HQET and of non-factorizable corrections [18] that have been assumed to be negligible in the case of the excited states $D^{**}$ [19].

A measurement of the branching fractions is presented for the decays $B^- \rightarrow D^0\pi^-$, $D^{*0}\pi^-$, $D^{**0}\pi^-$ and $\bar{B}^0 \rightarrow D^+\pi^-$, $D^{*+}\pi^-$, $D^{**+}\pi^-$ [20] with a missing mass method, based on a sample of 231 million $\Upsilon(4S) \rightarrow B\overline{B}$ pairs collected by the BaBar detector at the PEP-II $e^+e^-$ collider. One of the B mesons is fully reconstructed and the other one decays to a reconstructed $\pi$ and a companion charmed meson identified by its recoil mass, inferred by the kinematics of the two body $B$ decay. This method, compared to the previous exclusive measurements [21], does not imply that the $\Upsilon(4S)$ decays into $B^+$ and $B^0$ with equal rates, nor rely on the $D$, $D^*$, or $D^{**}$ decay branching fractions. The number of fully reconstructed $B$ mesons $B_{rec}^{d\bar{d}}$ is extracted from a fit to its mass distribution. In the decay $\Upsilon(4S) \rightarrow B_{rec}^{d\bar{d}} \overline{B}_{X\pi}$, where $\overline{B}_{X\pi}$ is the recoiling $\overline{B}$ which decays into $\pi^- X$, the invariant mass of the $X$ system is derived from the missing 4-momentum $p_X$ applying the energy-momentum conservation:

\[
px = p_{\Upsilon(4S)} - p_{B_{rec}^{d\bar{d}}} - p_{\pi^-}.
\]

The 4-momentum of the $\Upsilon(4S)$, $p_{\Upsilon(4S)}$, is computed from the beam energies and $p_{\pi^-}$ and $p_{B_{rec}^{d\bar{d}}}$ are the measured 4-momenta of the pion and of the reconstructed $B_{rec}^{d\bar{d}}$, respectively. The $B_{rec}^{d\bar{d}}$ energy is constrained by the beam energies. The $\overline{B} \rightarrow D\pi^-$, $\overline{B} \rightarrow D^*\pi^-$, or $\overline{B} \rightarrow D^{**}\pi^-$ signal yields peak at the $D$, $D^*$, and $D^{**}$ masses in the missing mass spectrum, respectively. The signal yield of the different modes, is extracted from the missing mass spectra. The $D\pi$ and $D^*\pi$ signal yields are extracted by a $\chi^2$ fit to the background subtracted missing mass distribution in the range $1.65 - 2.20$ GeV/$c^2$. The $D^{**}$ yield is obtained by counting the candidates in excess in the missing mass range $2.2 - 2.8$ GeV/$c^2$. This range is chosen in order to keep most of the excess and no
2.2. Precision Measurement of $D^0$ mass by CLEO-c

The $D^0$ ($c\bar{u}$) and $D^\pm$ ($c\bar{d}, c\bar{d}$) mesons form the ground states of the open charm system. The knowledge of their masses is important for its own sake, but a precision determination of the $D^0$ mass has become more important because of the recent discovery of a narrow state known as X(3872) [24–27]. Many different theoretical models have been proposed [28–31] to explain the nature of this state, whose present average of measured masses is $M(X) = 3871.2 \pm 0.5$ MeV [32]. A provocative and challenging theoretical suggestion is that X(3872) is a loosely bound molecule of $D^0$ and $D^{*0}$ mesons [31]. This suggestion arises mainly from the closeness of $M[X(3872)]$ to $M(D^{*0}) + M(D^0) = 2M(D^0) + [M(D^{*0}) - M(D^0)] = 2(1864.1 \pm 1.0) + (142.12 \pm 0.07)$ MeV = 3870.32 $\pm$ 2.0 MeV based on the PDG [32] average value of the measured $D^0$ mass, $M(D^0) = 1864.1 \pm 1.0$ MeV. This gives the binding energy of the proposed molecule, $E_b[X(3872)] = M(D^{*0}) + M(D^0) - M[X(3872)] = -0.9 \pm 2.1$ MeV. Although the negative value of the binding energy would indicate that X(3872) is not a bound state of $D^0$ and $D^{*0}$, its $\pm 2.1$ MeV error does not preclude this possibility. It is necessary to measure the masses of both $D^0$ and X(3872) with much improved precision to reach a firm conclusion. Recently, CLEO-c reported a precision measurement of the $D^0$ mass, and provided a more constrained value of the binding energy of X(3872) as a molecule. Several earlier measurements of the $D^0$ mass exist. The PDG [32] resulting average $D^0$ mass is based on the measured $D^0$ masses as $M(D^0)_{AVG} = 1864.1 \pm 1.0$ MeV. They also list a fitted mass, $M(D^0)_{FIT} = 1864.5 \pm 0.4$ MeV, based on the updated results of measurements of $D^\pm$, $D_s^\pm$, $D^{*\pm}$, $D^{*0}$, and $D_s^{*\pm}$ masses and mass differences. In its recent measurement, CLEO-c analyzes 281 pb$^{-1}$ of $e^+e^-$ annihilation data taken at the $\Psi(3770)$ resonance at the Cornell Electron Storage Ring (CESR) with the CLEO-c detector to measure the $D^0$ mass using the reaction $\Psi(3770) \rightarrow D^0 \overline{D}^0$, with $D^0 \rightarrow K_S^0\Phi$, $K_S^0 \rightarrow \pi^-\pi^+$ and $\Phi \rightarrow K^+K^-$. The choice of the $D^0 \rightarrow K_S^0\Phi$ mode is motivated by the determination of the $D^0$ mass not depending on the precision of the determination of the beam energy. Since $M(\Phi) + M(K_S^0) = 1517$ MeV is a substantial fraction of $M(D^0)$, the final state particles have small momenta and the uncertainty in their measurement makes a small contribution to the total uncertainty in $M(D^0)$. This consideration favors $D^0 \rightarrow K_S^0\Phi$ over the more prolific decays $D^0 \rightarrow K^-\pi^+$ and $D^0 \rightarrow K^-\pi^+\pi^-\pi^+$ in which the decay particles have considerably larger momenta and therefore greater sensitivity to the measurement uncertainties. An additional advantage of the $D^0 \rightarrow K_S^0\Phi$ reaction is that in fitting for $M(D^0)$ the mass of $K_S^0$ can be constrained to its value which is known with precision [32]. The final result of this measurement is $M(D^0) = 1864.847 \pm 0.150$ (stat) $\pm$ 0.095 (syst) MeV. Adding the errors in quadrature, $M(D^0) = 1864.847 \pm 0.178$ MeV is obtained. This is significantly more precise than the current PDG average [32]. This result for $M(D^0)$ leads to $M(D^{*0}) = 3871.81 \pm 0.36$ MeV. Thus, the binding energy of X(3872) as a $D^0 \overline{D}^{*0}$ molecule is $E_b = (3871.81 \pm 0.36) - (3871.2 \pm 0.5) = +0.6 \pm 0.6$ MeV. This result provides a strong constraint for the theoretical predictions for the decays of X(3872) if it is a $D^0 \overline{D}^{*0}$ molecule [31]. The error in the binding energy is now dominated by the error in the X(3872) mass measurement, which will hopefully improve as the results from the analysis of larger luminosity data from various experiments become available. This analysis is published [33].

3. Strange charm mesons

Much of the theoretical work on the $c\bar{s}$ system has been performed in the limit of heavy $c$ quark mass using potential models [34–37] that treat the $c\bar{s}$ system much like a hydrogen atom. Prior to the discovery of the $D_{sJ}(2317)^+$ meson, such models were successful at explaining the masses of all known $D$ and $D_s$ states and even predicting, to good accuracy, the masses of many $D$ mesons (including the $D_{s1}(2536)^+$ and $D_{s2}(2575)^+$) before they were observed (see Fig. 1). Several of the predicted $D_s$ states were not confirmed experimentally, notably the lowest mass $J^P = 0^+$ state (at around 2.48 GeV/$c^2$) and the second lowest mass $J^P = 1^+$ state (at around 2.58 GeV/$c^2$). Since the predicted widths of these two states were large, they would be hard to observe, and thus the lack of
established by BELLE. The family at all but is instead some type of exotic particle. The low mass (160 MeV meson state, the narrow width could be explained by DK to D isospin-violating DK low the measured to be around 2 GeV/c², is below the DK threshold. Thus, this particle is forced to decay either electromagnetically, of which there is no experimental evidence, or through the observed isospin-violating D⁺π⁰ strong decay. The intrinsic width is small enough that only upper limits have been measured (the best limit previous to this analysis being Γ < 4.6 MeV at 95% CL as established by BELLE [40]). If the D_{sJ}⁺(2317) is the missing 0⁺ c̅s meson state, the narrow width could be explained by the lack of an isospin-conserving strong decay channel. The low mass (160 MeV/c² below expectations) is more surprising and has led to the speculation that the D_{sJ}⁺(2317) does not belong to the D_{sJ}⁺ meson family at all but is instead some type of exotic particle, such as a four-quark state [43].

The D_{sJ}(2460)⁺ meson has been observed decaying to D⁺π⁺γ [38–42], D⁺π⁺π⁻ [40], and D⁺γ [40–42]. The intrinsic width is small enough that only upper limits have been measured (the best limit previous to this analysis being Γ < 5.5 MeV at 95% CL as established by BELLE [40]). The D⁺γ decay implies a spin of at least one, and so it is natural to assume that the D_{sJ}(2460)⁺ is the missing 1⁺ c̅s meson state. Like the D_{sJ}⁺(2317), the D_{sJ}(2460)⁺ is substantially lower in mass than predicted for the normal c̅s meson. This suggests that a similar mechanism is deflecting the masses of both mesons, or that both the states belong to the same family of exotic particles.

The spin-parity of the D_{sJ}⁺(2317) and D_{sJ}(2460)⁺ mesons has not been firmly established. The decay mode of the D_{sJ}⁺(2317) alone implies a spin-parity assignment from the natural J⁺P series {0⁺, 1⁺, 2⁺,...}, assuming parity conservation. Because of the low mass, the assignment J⁺P = 0⁺ seems most reasonable, although experimental data have not ruled out higher spin. It is not clear whether electromagnetic decays such as D_{sJ}⁺(2112)⁺γ can compete with the strong decay to D⁺π⁰, even with isospin violation. Thus, the absence of experimental evidence for radiative decays such as D_{sJ}⁺(2317) → D_{sJ}⁺(2112)⁺γ is not conclusive.

Experimental evidence for the spin-parity of the D_{sJ}(2460)⁺ meson is somewhat stronger. The observation of the decay to D⁺γ alone rules out J = 0. Decay distribution studies in B → D_{sJ}(2460)⁺D⁺γ [41, 42] favor the assignment J = 1. Decays to either D⁺π⁰, D⁰K⁺, or D⁺K⁺ would be favored if they were allowed. Since these decay channels are not observed, this suggests, when combined with the other observations, the assignment J⁺P = 1⁺. In this case, the decay to D_{sJ}(2317)⁺γ is allowed, but it may be small in comparison to the D_{sJ}(2112)⁺γ decay mode.

An updated analysis of the D_{sJ}(2317)⁺ and D_{sJ}(2460)⁺ mesons using 232 fb⁻¹ of e⁺e⁻ → c̅c data is presented here. Established signals from the decay D_{sJ}(2317)⁺ → D_{sJ}⁺π⁰ and D_{sJ}(2460)⁺ → D_{sJ}⁺π⁰γ, D_{sJ}(2460)⁺ → D_{sJ}⁺π⁺γ, and D_{sJ}⁺π⁺π⁻ are confirmed. A detailed analysis of invariant mass distributions of these final states including consideration of the background introduced by reflections of other c̅c decays produces the following mass values:

\[ m(D_{sJ}(2317)⁺) = (2319.6 ± 0.2 ± 1.4) \text{ MeV/c}^2 \]
\[ m(D_{sJ}(2460)⁺) = (2460.1 ± 0.2 ± 0.8) \text{ MeV/c}^2 \]

where the first error is statistical and the second systematic. Upper 95% CL limits of Γ < 3.8 MeV and Γ < 3.5 MeV are calculated for the intrinsic D_{sJ}(2317)⁺ and D_{sJ}(2460)⁺ widths. All results are consistent with previous measurements.

The following final states are investigated: D_{sJ}⁺π⁰, D_{sJ}⁺γ, D_{sJ}(2112)⁺π⁰, D_{sJ}(2112)⁺γ, D_{sJ}(2460)⁺γ, D_{sJ}(2460)⁺π⁻, and D_{sJ}(2460)⁺π⁺π⁻. No statistically significant evidence of new decay modes is observed. The following branching ratios are measured:

\[ \frac{B(D_{sJ}(2460)⁺ → D_{sJ}⁺γ)}{B(D_{sJ}(2460)⁺ → D_{sJ}⁺π⁺γ)} = 0.337 ± 0.036 ± 0.038 \]
\[ \frac{B(D_{sJ}(2460)⁺ → D_{sJ}⁺π⁺π⁻)}{B(D_{sJ}(2460)⁺ → D_{sJ}⁺π⁰γ)} = 0.077 ± 0.013 ± 0.008, \]

where the first error is statistical and the second systematic. The data are consistent with the decay
$D_{sJ}(2460)^+ \rightarrow D_s^+ \pi^0 \gamma$ proceeding entirely through $D_s^*(2112)^+\pi^0$.

Since the results presented here are consistent with $J^P = 0^+$ and $J^P = 1^+$ spin-parity assignments for the $D_{sJ}^*(2317)^+$ and $D_{sJ}(2460)^+$ mesons, these two states remain viable candidates for the lowest lying $p$-wave $c\bar{s}$ mesons. The lack of evidence for some radiative decays, in particular $D_{sJ}^*(2317)^+ \rightarrow D_s^*(2112)^+\gamma$ and $D_{sJ}(2460)^+ \rightarrow D_s^*(2112)^+\gamma$, are in contradiction with this hypothesis according to some calculations, but large theoretical uncertainties remain. No state near the $D_{sJ}^*(2317)^+$ mass is observed decaying to $D_s^+\pi^\pm$. If charged or neutral partners to the $D_{sJ}^*(2317)^+$ exist (as would be expected if the $D_{sJ}^*(2317)^+$ is a four-quark state), some mechanism is required to suppress their production in $e^+e^-$ collisions. This analysis is realized in inclusive $c\bar{c}$ production using $232$ fb$^{-1}$ of data collected by the BaBar experiment near $\sqrt{s} = 10.6$ GeV and is published in [45].

3.2. The $D_{s1}(2536)^+$ Case

For a complete understanding of the charmed strange meson spectrum, a comprehensive knowledge of the parameters of all known $D_{sJ}^+$ mesons is mandatory. In this part of the presentation, a precision measurement of the mass and the decay width of the meson $D_{s1}(2536)^+$ is presented. The mass is currently reported by the PDG with a precision of 0.6 MeV/$c^2$, while only an upper limit of 2.3 MeV/$c^2$ is given for the decay width [32]. These values are based on measurements with 20 times fewer reconstructed $D_{s1}^+$ candidates compared to this one. The BaBar experiment, in addition to its excellent tracking and vertexing capabilities, provides a rich source of charmed hadrons, enabling an analysis of the $D_{s1}^+$ with high statistics and small errors.

Since the uncertainty of the $D^{*+}$ mass is large (0.4 MeV/$c^2$ [32]), a measurement of the mass difference defined by

$$\Delta m(D_{s1}^+) = m(D_{s1}^+) - m(D^{*+}) - m(K^0_s),$$

is performed. Additionally, due to the correlation between the masses, the $D_{s1}^+$ signal in the mass difference spectrum is much more narrow than the one from the $D_{s1}^+$ mass spectrum alone leading to a high precision measurement of the mass and the decay width of the meson $D_{s1}(2536)^+$ using the decay mode $D_{s1}^+ \rightarrow D^{*+}K^0_s$. The mass difference between $D_{s1}^+$ and $D^{*+}K^0_s$ for the two reconstructed decay modes is measured to be

$$\Delta \mu(D_{s1}^+)_{K4\pi} = 27.290 \pm 0.028 \pm 0.031 \text{ MeV}/c^2,$$

$$\Delta \mu(D_{s1}^+)_{K6\pi} = 27.180 \pm 0.023 \pm 0.043 \text{ MeV}/c^2,$$

with the first error denoting the statistical uncertainty and the second one the systematic uncertainty. These results correspond to a relative error of 0.15% for the mass difference. This lies within the range of precision achievable with the Babar detector: the $J/\psi$ mass has been reconstructed with a relative error of 0.05% [46].

Combining the results, while taking the systematic errors including the uncertainties of the $D^{*+}$ mass ($\pm$0.4 MeV/$c^2$) and of the $K^0_s$ mass ($\pm$0.022 MeV/$c^2$) into account, yields a final value for the $D_{s1}^+$ mass of

$$m(D_{s1}^+) = 2534.85 \pm 0.02 \pm 0.40 \text{ MeV}/c^2,$$

while the PDG value for the mass is given as $2535.35 \pm 0.34 \pm 0.50 \text{ MeV}/c^2$. The error on the measured $D_{s1}^+$ mass is dominated by the uncertainty of the $D^{*+}$ mass. The mass difference between the $D_{s1}^+$ and the $D^{*+}$ follows from these results as

$$\Delta m = m(D_{s1}^+) - m(D^{*+}) = 524.85 \pm 0.02 \pm 0.04 \text{ MeV}/c^2.$$ 

The decay width is measured to be

$$\Gamma(D_{s1}^+)_{K4\pi} = 1.112 \pm 0.068 \pm 0.131 \text{ MeV}/c^2,$$

$$\Gamma(D_{s1}^+)_{K6\pi} = 0.990 \pm 0.059 \pm 0.119 \text{ MeV}/c^2.$$ 

The final combined value for decay width is

$$\Gamma(D_{s1}^+) = 1.03 \pm 0.05 \pm 0.12 \text{ MeV}/c^2.$$ 

The result for the mass difference $\Delta m = m(D_{s1}^+) - m(D^{*+})$ represents an improvement in precision by a factor of 14 compared with the current PDG value of $525.3 \pm 0.6 \pm 0.1 \text{ MeV}/c^2$. It deviates by $1\sigma$ from the larger PDG value. The precision achieved is comparable with other recent high precision analyses performed at Babar like the $\Lambda_c$ mass measurement ($m(\Lambda_c) = 2286.46 \pm 0.04 \pm 0.14 \text{ MeV}/c^2$) [47]. Furthermore, this analysis presents for the first time a direct measurement of the $D_{s1}^+$ decay width with small errors rather than just an upper limit, which is currently stated by the PDG as $2.3 \text{ MeV}/c^2$. This analysis is also realized in inclusive $c\bar{c}$ production using $232$ fb$^{-1}$ of data collected by the Babar experiment near $\sqrt{s} = 10.6$ GeV and is detailed in [48].

3.3. $D_{s2}(2573)^+$ and New Strange Charmed Mesons

Here, a new $c\bar{s}$ state and a broad structure observed in the decay channels $D^0K^+$ and $D^+K^0_S$ are reported. This analysis is based on a $240$ fb$^{-1}$ inclusive $c\bar{c}$ data sample recorded near the $\Upsilon(4S)$ resonance by the Babar detector at the PEP-II asymmetric-energy $e^+e^-$ storage rings.

Three inclusive processes [20] are reconstructed:

$$e^+e^- \rightarrow D^0K^+X, D^0 \rightarrow K^-\pi^+ \quad (1)$$

$$e^+e^- \rightarrow D^0K^+X, D^0 \rightarrow K^-\pi^+\pi^0 \quad (2)$$

$$e^+e^- \rightarrow D^+K^0_SX, D^+ \rightarrow K^-\pi^+\pi^+, K^0_S \rightarrow \pi^+\pi^0 \quad (3)$$

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$$e^+e^- \rightarrow D^+K^0_SX, D^+ \rightarrow K^-\pi^+\pi^+, K^0_S \rightarrow \pi^+\pi^0 \quad (3)$$

By
Selecting events in the D signal regions, Fig. 2 shows the $D^0 K^+$ invariant mass distributions for channels (1) and (2), and the $D^+ K_S^0$ invariant mass distribution for channel (3). To improve mass resolution, the nominal $D$ mass and the reconstructed 3-momentum are used to calculate the $D$ energy for channels (1) and (3). Since channel (2) has a poorer $D^0$ resolution, each $K^-\pi^+\pi^0$ candidate is kinematically fit with a $D^0$ mass constraint and a $\chi^2$ probability greater than 0.1% is required.

The fraction of events having more than one $DK$ combination per event is 0.9% for channels (1) and (3) and 3.4% for channel (2). In the following, the term reflection will be used to describe enhancements produced by two or three body decays of narrow resonances where one of the decay products is missed.

The three mass spectra in Fig. 2 present similar features:

- A single bin peak at 2.4 GeV/c^2 due to a reflection from the decays of the $D_{s1}(2536)^+$ to $D^{*0}K^+$ or $D^{*+}K_S^0$ in which the $\pi^0$ or $\gamma$ from the $D^*$ decay is missed. This state, if $J^P = 1^+$, cannot decay to $DK$.
- A prominent narrow signal due to the $D_{s1}(2573)^+$.
- A broad structure peaking at a mass of approximately 2.7 GeV/c^2.
- An enhancement around 2.86 GeV/c^2. This can be seen better in the expanded views shown in the insets of Fig. 2.

Different background sources are examined: combinatorial, possible reflections from $D^*$ decays, and particle misidentification.

Backgrounds come both from events in which the candidate $D$ meson is correctly identified and from events in which it is not. The first case can be studied combining a reconstructed $D$ meson with a kaon from another $D$ meson in the same event, using data with fully reconstructed $DD$ pairs or Monte Carlo simulations. No signal near 2.7 or 2.86 GeV/c^2 is seen in the $DK$ mass plots for these events. The second case can be studied using the $D$ mass sidebands. The shaded regions in Fig. 2 show the $DK$ mass spectra for events in the $D$ sideband regions normalized to the estimated background in the signal region. No prominent structure is visible in the sideband mass spectra. The dotted histogram in (a) is from $e^+e^- \rightarrow c\bar{c}$ Monte Carlo simulations incorporating previously known $D_s$ states with an arbitrary normalization.

The possibility that the features at 2.7 and 2.86 GeV/c^2 could be a reflection from $D^*$ or other higher mass resonances is considered. Candidate $DK$ pairs where the $D$ is a $D^*$-decay product are identified by forming $D\pi$ and $D\gamma$ combinations and requiring the invariant-mass difference between one of those combinations and the $D$ to be within $\pm2\sigma$ of the known $D^*-D$ mass difference. No signal near 2.7 or 2.86 GeV/c^2 is seen in the $DK$ mass plots for these events. Events belonging to these possible reflections (except for the $D^{*0} \rightarrow D^0\gamma$ events, which could not be isolated cleanly) have been removed from the mass distributions shown in Fig. 2 (corresponding to $\approx8\%$.}

Figure 2: The $DK$ invariant mass distributions for (a) $D_{s1}^{*0} K^{-+} K^{+}$, (b) $D_{s1}^{*0} K^{-+} \pi^0 K^{+}$ and (c) $D_{s1}^{*0} K^{-+} \pi^0 K^{+}$. The shaded histograms are for the $D$-mass sideband regions. The dotted histogram in (a) is from $e^+e^- \rightarrow c\bar{c}$ Monte Carlo simulations incorporating previously known $D_s$ states with an arbitrary normalization. The insets show an expanded view of the 2.86 GeV/c^2 region. The solid curves are the fitted background threshold functions from the three separate fits.
of the final sample).

The presence of resonant structures can be visually enhanced by subtracting the fitted background threshold function from the data. Fig. 3 shows the background-subtracted \( D_{sJ}(2860)^+ \), \( D_{sJ}(2573)^+ \), and \( D_{sJ}(2860)^+ \) invariant mass distributions in the 2.86 GeV/c² mass region. Fig. 3(d) shows the sum of the three mass spectra.

In the following, the structure in the 2.86 GeV/c² mass region is labelled \( D_{sJ}(2860)^+ \) and the one in the 2.7 GeV/c² mass region is labelled \( X(2690)^+ \). The three \( DK \) mass spectra shown in Fig. 2 from 2.42 GeV/c² to 3.1 GeV/c² (excluding the \( D_{sJ}(2536)^+ \) reflection) are first fitted separately using a binned relativistic Breit-Wigner lineshapes where spin-2 is activistic Breit-Wigner for the \( \chi^3_{2\%} \) using this background expression and one spin-2 relativistic Breit-Wigner for the \( \chi^3_{2\%} \) for the \( D_{sJ}(2536)^+ \) reflection. The fits give consistent values for the parameters of the three structures.

When the three mass distributions of Fig. 2 are fitted simultaneously, the resulting resonance parameters are found consistent with those obtained with previous separate fits. For \( D_{sJ}(2537)^+ \) resonance, mass and width are:

\[
m(\text{D}_{sJ}(2537)^+) = (2572.2 \pm 0.3 \pm 1.0) \text{ MeV/c}^2
\]

\[
\Gamma(\text{D}_{sJ}(2537)^+) = (27.1 \pm 0.6 \pm 5.6) \text{ MeV/c}^2,
\]

where the first errors are statistical and the second systematic. For the new states, the following values were extracted:

\[
m(\text{D}_{sJ}(2860)^+) = (2856.6 \pm 1.5 \pm 5.0) \text{ MeV/c}^2
\]

\[
\Gamma(\text{D}_{sJ}(2860)^+) = (47 \pm 7 \pm 10) \text{ MeV/c}^2.
\]

\[
m(\text{X}(2690)^+) = (2688 \pm 4 \pm 3) \text{ MeV/c}^2
\]

\[
\Gamma(\text{X}(2690)^+) = (112 \pm 7 \pm 36) \text{ MeV/c}^2.
\]

In summary, in 240 fb⁻¹ of data collected by the BaBar experiment, a new \( D^+ \) state is observed in the inclusive \( D^+ \) mass distribution near 2.86 GeV/c² in three independent channels. The decay to two pseudoscalar mesons implies a natural spin-parity for this state: \( J^P = 0^+, 1^- \). It has been suggested that this new state could be a radial excitation of \( D_{sJ}(2317) \) [49] although other possibilities cannot be ruled out. In the same mass distributions a broad enhancement around 2.69 GeV/c² is also observed, it is not possible to associate it to any known reflection or background. This analysis is published [50].

Another BaBar analysis [51], has searched for resonances in \( B \rightarrow D^{(*)} D^{(*)} K \) decays in 22 decay modes using 347 fb⁻¹ data sample recorded at the \( T(4S) \) resonance. The \( D K \) and \( D^* K \) invariant mass distributions are built with 8 decay modes each. Both distributions show a resonant enhancement around 2700 MeV/c². However, due to an unknown structure at low mass in the \( D K \) invariant mass distribution and to the possible additional resonances in the signal region in the \( D^* K \) invariant mass distribution, a full Dalitz analysis in necessary and is ongoing in order to extract the \( D_{sJ}(2700)^+ \) parameters.

4. Conclusion

Although the nature of the newly discovered charm resonances is not yet fully understood, the resonances are interpreted as molecular or hybrid states in most theoretical papers. It will be interesting to see if these interpretations are confirmed by future measurements and analyses.

References

Figure 3: Fitted background-subtracted $DK$ invariant mass distributions for (a) $D^0_{K^+\pi^-}\pi^+$, (b) $D^0_{K^+\pi^-\pi^+\pi^0}K^+$, (c) $D^+_{K^-\pi^-\pi^+}K^0_s$, and (d) the sum of all modes in the 2.86 GeV/c$^2$ mass region. The curves are the fitted functions described in the text.

[20] Charge-conjugate reactions are implied throughout.
[22] Here $D^{**}$ refers collectively to all the charmed meson excited states with masses in the 2.2 – 2.8 GeV/c$^2$ range.
[48] B. Aubert et al. (BABAR Collaboration), hep-ex/0608044.

[51] B. Aubert et al. (BABAR Collaboration), preliminary, see Ph. Grenier’s talk at 42nd REN-CONTRES DE MORIOND, La Thuile, Italy, 10 - 24 March 2007, QCD PROCEEDINGS.
The BES-III experiment at the high luminosity Tau-Charm factory

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Interesting results from BES-II and other experiments raised actually many new questions which shall be answered by its upgrade program, BEPCII and BES-III. The design and current status of BEPCII and BES-III are reported.

I. PHYSICS MOTIVATION

In early 80’s, Chinese government decided to build an $e^+e^-$ collider running at the tau-charm energy region, called BEPC, which is completed in 1989. The only detector at the machine is called Beijing spectrometer (BES). In mid 90’s, there has been a minor upgrade of the detector, which is then called BES-II. Since then, hundreds of papers have been published on the international journals, some with significant impacts to the community. The upgrade of BEPC was decided at the beginning of this century, called BEPCII, which has a designed luminosity of $10^{33}$ cm$^{-2}$s$^{-1}$, an increase of a factor of 100. The corresponding detector, called BES-III, adopted latest detector technology to minimize systematic errors in order to match the unprecedented statistics.

The physics program of the BES-III experiment includes light hadron spectroscopy, charmonium, electroweak physics from charmed mesons, QCD and hadron physics, tau physics and search for new physics. Due to its huge luminosity and small energy spread, the expected event rate per year is historical, as listed in table I.

**TABLE I:** $\tau$-Charm productions at BEPCII in one year’s running($10^7$s).

<table>
<thead>
<tr>
<th>Data Sample</th>
<th>CoM energy (MeV)</th>
<th>Luminosity ($10^{33}$cm$^{-2}$s$^{-1}$)</th>
<th>#Events</th>
</tr>
</thead>
<tbody>
<tr>
<td>$J/\psi$</td>
<td>3097</td>
<td>0.6</td>
<td>$10 \times 10^9$</td>
</tr>
<tr>
<td>$\tau^+\tau^-$</td>
<td>3670</td>
<td>1.0</td>
<td>$12 \times 10^6$</td>
</tr>
<tr>
<td>$\psi(2S)$</td>
<td>3686</td>
<td>1.0</td>
<td>$3.0 \times 10^9$</td>
</tr>
<tr>
<td>$D^0\bar{D}^0$</td>
<td>3770</td>
<td>1.0</td>
<td>$18 \times 10^6$</td>
</tr>
<tr>
<td>$D^+D^-$</td>
<td>3770</td>
<td>1.0</td>
<td>$14 \times 10^6$</td>
</tr>
<tr>
<td>$D_s^+D_s^-$</td>
<td>4030</td>
<td>0.6</td>
<td>$1.0 \times 10^6$</td>
</tr>
<tr>
<td>$D_s^0D_s^-$</td>
<td>4170</td>
<td>0.6</td>
<td>$2.0 \times 10^6$</td>
</tr>
</tbody>
</table>

It is well known that $J/\psi$ and $\psi'$ decays is an ideal laboratory for light hadron studies since it has a huge production cross section with a very clean and gluon reach environment. We plan to study the meson and baryon spectroscopy, search for glueballs and other exotics such as hybrids and multi-quark states. Recently, BES-II found several new structures and threshold enhancements in various decays channels [1], which leads to a number of speculations. It is clear that more data is needed to understand these results, study their decay properties and establish a theoretical framework to accommodate them. Fig. 1 shows a comparison of the X(1835) signal at BES-II and the corresponding expectation at BES-III. (a) 2.5 days of BES-III, (b) two years of BES-II

![Fig. 1: A comparison of the X(1835) signal at BES-II and the corresponding expectation at BES-III.](image)

Recently, a lot of new $XYZ$ resonance have been found above or below the open charm threshold from B decays. BES-III should be able to study the direct production of $1^-$ states and some of the low mass states, such as $Y(2175)$. We will systematically study charmonium transitions and their decays, search for rare decays and new phenomena, and calibrate lattice QCD calculations.

For Charm physics, the precision of CKM matrix elements can be significantly improved by measuring the leptonic and semi-leptonic decays of charmed mesons, and test the unitarity of CKM matrix. The $D\bar{D}$-mixing can be measured at the level of $10^{-4}$ and the CP violation will be searched for at the level of $10^{-3}$. Rare and forbidden decays can be searched for at a typical level of $10^{-8}$.

The tau mass measurement can be improved by a factor of two over the BES-II results with a new beam.
energy calibration based on the Compton scattering technique.

A summary of BES-III physics program, called yellow book, is under preparation. It will be published at the beginning of next year.

II. STATUS OF THE COLLIDER

The new BEPCII has two storage rings with a circumference of 224 m, one for electron and one positron, each with 93 bunches spaced by 8 ns. The total current of the beam is designed to be 0.93 A, the crossing angle of two beams 22 mrad. The luminosity is expected to be $10^{33}$ cm$^{-2}$s$^{-1}$ at the beam energy of 1.89 GeV, the bunch length is estimated to be 1.5 cm and the energy spread 5.16 x 10$^{-4}$.

At this moment, all the LINAC equipments have been installed and successfully tested. Parameters such as the beam current, emittance and energy spread etc. for both electron and positron beams, have been measured. All the design specifications have been satisfied.

The storage rings have been installed in two phases. In the first phase, conventional magnets at the interaction region were installed and tested. Both electron and positron beams have been stored up to 500 mA with a reasonable life time. Synchrotron radiation beams have been delivered to user for about three months in total. Tests of $e^+e^-$ collision have been performed with an estimated peak luminosity of about $10^{31}$ cm$^{-2}$s$^{-1}$. In the second phase, superconducting quadrupoles for the final beam focusing at the interaction region have been installed and beams have successfully stored. Collision of $e^+e^-$ beams have been observed and synchrotron radiation run will start soon.

We plan to move the BES-III detector into the interaction region early next year after the luminosity is more than $3 \times 10^{31}$ cm$^{-2}$s$^{-1}$ and the beam background is under control.

III. STATUS OF THE BES-III CONSTRUCTION

The BES-III detector, as shown in Fig. 2, consists of a drift chamber in a small cell structure filled with a helium-based gas, an electromagnetic calorimeter made of CsI(Tl) crystals, Time-of-Flight(TOF) counters for particle identification made of plastic scintillators, a muon system made of Resistive Plate Chambers(RPC), and a super-conducting magnet providing a field of 1T. Current status of the construction is summarized in the following.

The drift chamber has a cylindrical shape with two chambers jointed at the end flange: an inner chamber without outer wall and an outer chamber without inner wall. There are a total of 6 stepped end flanges made of 18 mm Al plates in order to give space for the focusing magnets. The inner radius of the chamber is 63 mm and the outer radius is 810 mm, with a length of 2400 mm. Both the inner and outer cylinder of the chamber are made of carbon fiber with a thickness of 1 mm and 10 mm respectively. A total of 6300 gold-plated tungsten wires(3% Rhenium) with a diameter of 25 um are arranged in 43 layers, together with a total of 22000 gold-plated Al wires for field shaping. The designed single wire spatial resolution and dE/dx resolution are 130 µm and 6%, respectively.

All the wiring have been completed with a very high quality, the wire tension and the leakage current are well under control. The assembly of the chamber has been completed together with all preamplifiers and related electronics. The whole chamber has been tested using cosmic-rays for three months. The obtained single wire resolution is about 120 µm, as shown in Fig. 3, well satisfying our design goal. The chamber has now been installed successfully into the BES-III detector.

The CsI(Tl) crystal electromagnetic calorimeter consists of 6240 crystals, 5280 in the barrel, and 960 in two endcaps. Each crystal is 28 cm long, with a front face of about 5.2 x 5.2 cm$^2$, and a rear face of about 6.4 x 6.4 cm$^2$. All crystals are tiled by 1.5° in the azimuth angle and 1-3° in the polar angle, respectively, and point to a position off from the interaction...
point by a few centimeters. The designed energy and position resolution are 2.5% and 6 mm at 1 GeV, respectively.

All the crystals have been produced and shipped, been tested, and assembled. Fig. 4 shows the test results of the light yield, uniformity and radiation hardness. All the barrel crystals have been installed into the mechanical structure, which has been installed into the BES-III detector as well.

The readout electronics of crystals, including preamplifiers, main amplifiers and charge measurement modules are tested at the IHEP E3 beam line together with a crystal array and photodiodes. Results from the beam test shows that the energy resolution of the crystal array reached the design goal of 2.5% at 1 GeV and the total noise achieved the level of less than 1000 equivalent electrons, corresponding to an energy of 220 KeV.

The particle identification at BES-III is based on the momentum and dE/dx measurements by the drift chamber, and the TOF measurement by plastic scintillators. The barrel scintillator bar is 2.4 m long, 5 cm thick and 6 cm wide. A total of 176 such scintillator bars constitute two cylinders, to have a good efficiency and time resolution. For the endcap, a total of 48 fan-shaped scintillators form a single layer. A 2-inch fine-mesh phototube is directly attached to each end of the scintillator to collect the light. The intrinsic time resolution is designed to be 90 ps including contributions from electronics and the common start/stop time. Such a time resolution, together with contributions from the beam size, momentum uncertainty, etc. can distinguish charged $\pi$ from K mesons for a momentum up to 0.9 GeV at the $2\sigma$ level.

Beam tests show that the intrinsic time resolution can be better than 90 ps and 75 ps for the barrel and the endcap TOF counters, respectively. Currently, all the PMTs and scintillators have been delivered and tested. The average attenuation length of all barrel scintillators is 4.8 m and the relative light yield exceeds our specification. All barrel scintillator have been assembled outside of the MDC and installed into the BES-III detector successfully.

The BES-III muon chamber is made of Resistive Plate Chambers (RPC) interleaved in the magnet yoke. There are a total of 9 layers in the barrel and 8 layers in the endcap, with a total area of about 2000 m$^2$. The readout strip is 4 cm wide, alternated between layers in x and y directions. The RPC is made of bakelite with a special surface treatment without
linseed oil [4]. Such a simple technique for the RPC production shows a good quality and stability at a low cost. All RPCs have been manufactured, tested, assembled and installed with satisfaction.

The BES-III super-conducting magnet has a radius of 1.48 m and a length of 3.52 m. It uses the Al stabilized NbTi/Cu conductor with a total of 920 turns, making a 1.0T magnetic field at a current of 3400 amp. The total cold mass is 3.6 t with a material thickness of about 1.92 $X_0$. In collaboration with WANG NMR of California, the magnet is designed and manufactured at IHEP.

The magnet was successfully installed into the iron yoke of the BES-III, together with the valve box. A stable magnetic field of 1.0T at a current of 3368 A was achieved. The field mapping together with superconducting quadrupole magnets for final focusing of the beam has been completed, and results are shown in Fig. 5. The uniformity of the magnetic field is satisfactory.

The BES-III offline software consists of a framework based on GAUDI, a Monte Carlo simulation based on GEANT4, an event reconstruction package, a calibration package and a database package using MySQL. Currently all codes are working as a complete system, and tests using cosmic-ray data and beam test data are underway. Analysis tools such as the particle identification, secondary vertex finding, kinematic fitting, event generator and partial wave analysis are still under development. Data challenge of the whole system is planned.

In summary, the BEPCII and BES-III construction went on smoothly. Currently all the mass production of detector components have completed, most of the assembly and installation of the detector are finished. We plan to take data in 2008.


The PANDA Experiment at FAIR

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The physics program of the future FAIR facility covers a wide range of topics that address central issues of strong interactions and QCD. The antiproton beam of unprecedented quality in the momentum range from 1 GeV/c to 15 GeV/c will allow the PANDA experiment to make high precision, high statistics measurements, which include charmonium and open charm spectroscopy, the search for exotic hadrons and the study of in-medium modifications of hadron masses.

1. Introduction

One of the most challenging and fascinating goals of modern physics is the achievement of a fully quantitative understanding of the strong interaction, which is the subject of hadron physics. Significant progress has been achieved over the past few years thanks to considerable advances in experiment and theory. New experimental results have stimulated a very intense theoretical activity and a refinement of the theoretical tools.

Still there are many fundamental questions which remain basically unanswered. Phenomena such as the confinement of quarks, the existence of glueballs and hybrids, the origin of the masses of hadrons in the context of the breaking of chiral symmetry are longstanding puzzles and represent the intellectual challenge in our attempt to understand the nature of the strong interaction and of hadronic matter.

Experimentally, studies of hadron structure can be performed with different probes such as electrons, pions, kaons, protons or antiprotons. In antiproton-proton annihilation particles with gluonic degrees of freedom as well as particle-antiparticle pairs are copiously produced, allowing spectroscopic studies with very high statistics and precision. Therefore, antiprotons are an excellent tool to address the open problems.

The recently approved FAIR facility (Facility for Antiproton and Ion Research), which will be built as a major upgrade of the existing GSI laboratory in Germany, will provide antiproton beams of the highest quality in terms of intensity and resolution, which will provide an excellent tool to answer these fundamental questions.

The PANDA experiment (Pbar ANnihilations at DArmstadt) will use the antiproton beam from the High-Energy Storage Ring (HESR) colliding with an internal proton target and a general purpose spectrometer to carry out a rich and diversified hadron physics program, which includes charmonium and open charm spectroscopy, the search for exotic hadrons and the study of in-medium modifications of hadron masses.

This paper is organized as follows: in section 2 we will give an overview of the FAIR facility and the HESR; in section 3 we will discuss some of the most significant items of the PANDA experimental program; in section 4 we will give a brief description of the PANDA detector. Finally in section 5 we will present our conclusions.

2. The FAIR facility

The planned FAIR complex is shown in Fig. 1. The heart of the system consists of two synchrotron rings, called SIS100 and SIS300, housed in the same tunnel, which will provide proton and ion beams of unprecedented quality. The SIS100, a 100 t·m proton ring, will feed the radioactive ion and antiproton beam lines for experiments to be carried out in the High-Energy Storage Ring (HESR), the Collector and Cooler rings (CR) and the New Experimental Storage Ring (NESR). The SIS300 will deliver high energy ion beams for the study of ultra relativistic heavy ion collisions.

The accelerators of FAIR will feature significant improvements in system parameters over existing facilities:

- **beam intensity**: increased by a factor of 100 to 1000 for primary and 10000 for secondary beams;
- **beam energy** will increase by a factor 30 for
heavy ions;

- **beam variety**: FAIR will offer a variety of beam lines, from antiprotons to protons, to uranium and radioactive ions;

- **beam precision**: availability of cooled antiproton and ion beams (stochastic and electron cooling);

- **parallel operation**: full accelerator performance for up to four different, independent experiments and experimental programs.

These features will make FAIR a first rate facility for experiments in particle, nuclear, atomic, plasma and applied physics.

### 2.1. The High-Energy Storage Ring

The antiproton beam will be produced by a primary proton beam from the SIS100. The \( \bar{p} \) production rate will be of approximately \( 2 \times 10^7 / s \). After \( 5 \times 10^5 \bar{p} \) have been produced they will be transferred to the HESR, where internal experiments in the \( \bar{p} \) momentum range from 1 GeV/c to 15 GeV/c can be performed.

The layout of the HESR is shown in Fig. 2. It is a racetrack ring, 574 meters in length, with two straight sections which will host the electron cooling and, respectively, the PANDA experiment. Two modes of operation are foreseen: in the high-luminosity mode peak luminosities of \( 2 \times 10^{32} \text{cm}^{-2}\text{s}^{-1} \) will be reached with a beam momentum spread \( \delta p/p = 10^{-4} \), achieved by means of stochastic cooling; in the high-resolution mode for beam momenta below 8 GeV/c electron cooling will yield a smaller beam momentum spread \( \delta p/p = 10^{-5} \) at a reduced luminosity of \( 10^{31} \text{cm}^{-2}\text{s}^{-1} \). The high-resolution mode will allow to measure directly the total width of very narrow (below 1 MeV) resonances.

### 3. The PANDA Physics Program

The PANDA experiment has a rich experimental program whose ultimate aim is to improve our knowledge of the strong interaction and of hadron structure. The experiment is designed to fully exploit the extraordinary physics potential arising from the availability of high-intensity, cooled antiproton beams. Significant progress beyond the present understanding of the field is expected thanks to improvements in statistics and precision of the data.

Many experiments are foreseen in PANDA. In this paper we will discuss the following:

- charmonium spectroscopy;
- search for gluonic excitations (hybrids and glue-balls);
- study of hadrons in nuclear matter;
- open charm spectroscopy.

### 3.1. Charmonium Spectroscopy

Ever since its discovery in 1974 [1] charmonium has been a powerful tool for the understanding of the strong interaction. The high mass of the \( c \) quark (\( m_c \approx 1.5 \text{ GeV/c}^2 \)) makes it plausible to attempt a description of the dynamical properties of the \( (c\bar{c}) \) system in terms of non-relativistic potential models, in which the functional form of the potential is chosen to reproduce the asymptotic properties of the strong interaction. The free parameters in these models are to be determined from a comparison with the experimental data.

Now, more than thirty years after the \( J/\psi \) discovery, charmonium physics continues to be an exciting and interesting field of research. The recent discoveries of new states (\( \eta_c \), X(3872)), and the exploitation of the B factories as rich sources of charmonium states have given rise to renewed interest in heavy quarkonia, and stimulated a lot of experimental and theoretical activities. Over the past few years a significant progress has been achieved by Lattice Gauge Theory calculations, which have become increasingly more capable of dealing quantitatively with non-perturbative dynamics in all its aspects, starting from the first principles of QCD.
3.1.1. Experimental Study of Charmonium

Experimentally charmonium has been studied mainly in $e^+e^-$ and $\bar{p}p$ experiments.

In $e^+e^-$ annihilations direct charmonium formation is possible only for states with the quantum numbers of the photon $J^{PC} = 1^{--}$, namely the $J/\psi$, $\psi'$ and $\psi(3770)$ resonances. Precise measurements of the masses and widths of these states can be obtained from the energy of the electron and positron beams, which are known with good accuracy. All other states can be reached by means of other production mechanisms, such as photon-photon fusion, initial state radiation, B-meson decay and double charmonium.

On the other hand all $\sigma c$ states can be directly formed in $\bar{p}p$ annihilations, through the coherent annihilation of the three quarks in the proton with the three antiquarks in the antiproton. This technique, originally proposed by P. Dalpiaz in 1979 [2], could be successfully employed a few years later at CERN and Fermilab thanks to the development of stochastic cooling. With this method the masses and widths of all charmonium states can be measured with excellent accuracy, determined by the very precise knowledge of the initial $\bar{p}p$ state and not limited by the resolution of the detector.

The cross section for the process:

$$\bar{p}p \to (\tau c) \to \text{final state}$$

is given (in units $\hbar = c = 1$) by the well known Breit-Wigner formula:

$$\sigma_{BW}(E) = \frac{2J + 1}{4} \frac{\pi B_{in} B_{out} \Gamma_R^2}{k^2 (E - M_R)^2 + \Gamma_R^2/4}$$

where $E$ and $k$ are the center-of-mass (c.m.) energy and momentum; $J$, $M_R$ and $\Gamma_R$ are the resonance spin, mass and total width and $B_{in}$ and $B_{out}$ are the branching ratios into the initial ($\bar{p}p$) and final states. Due to the finite energy spread of the beam, the measured cross section is a convolution of the Breit-Wigner cross section, eq. (2), and the beam energy distribution function $f(E, \Delta E_B)$; the effective production rate $\nu$ is given by:

$$\nu = L_0 \left\{ \epsilon \int dE f(E, \Delta E_B) \sigma_{BW}(E) + \sigma_b \right\}$$

where $L_0$ is the instantaneous luminosity, $\epsilon$ an overall efficiency × acceptance factor and $\sigma_b$ a background term.

The parameters of a given resonance can be extracted by measuring the formation rate for that resonance as a function of the c.m. energy $E_{cm}$. The accurate determination of masses and widths depends crucially on the precise knowledge of the absolute energy scale and on the beam energy spectrum.

The technique is illustrated in Fig. 3 which shows a scan of the $\chi_{c1}$ resonance carried out at the Fermilab antiproton accumulator by the E835 experiment [3] using the process $\bar{p}p \to \chi_{c1} \to J/\psi \gamma$. For each point of the scan the horizontal error bar in (a) corresponds to the width of the beam energy distribution. The actual beam energy distribution is shown in (b). This scan allowed the E835 experiment to carry out the most precise measurement of the mass ($3510.719 \pm 0.051 \pm 0.019$ MeV/$c^2$) and total width ($0.876 \pm 0.045 \pm 0.026$ MeV) of this resonance.

3.1.2. The charmonium spectrum

The spectrum of charmonium states is shown in Fig. 4. It consists of eight narrow states below the open charm threshold (3.73 GeV) and several tens of states above the threshold.

All eight states below $D \bar{D}$ threshold are well established, but whereas the triplet states are measured with very good accuracy, the same cannot be said for the singlet states.

The $\eta_c$ was discovered almost thirty years ago and many measurements of its mass and total width exist, with six new measurements in the last four years. Still the situation is far from satisfactory. The Particle Data Group (PDG) [4] value of the mass is $2980.4 \pm 1.2$ MeV/$c^2$, an average of eight measurements with an internal confidence level of 0.026: the error on the $\eta_c$ mass is still as large as 1.2 MeV/$c^2$, to be compared with few tens of KeV/$c^2$ for the $J/\psi$ and $\psi'$ and few hundreds of KeV/$c^2$ for the $\chi_{cJ}$ states. The situation is even worse for the total width: the PDG average is $25.5 \pm 3.4$ MeV, with an overall confidence level of only 0.001 and individual measurements ranging from 7 MeV to 34.3 MeV. The most recent measurements have shown that the $\eta_c$ width is larger than was previously believed, with values which are difficult to ac-
comodate in quark models. This situation points to the need for new high-precision measurements of the $\eta_c$ parameters.

The first experimental evidence of the $\eta_c(2S)$ was reported by the Crystal Ball collaboration [5], but this finding was not confirmed in subsequent searches in $p\bar{p}$ or $e^+e^-$ experiments. The $\eta_c(2S)$ was finally discovered by the Belle collaboration [6] in the hadronic decay of the $B$ meson $B \to K + \eta_c(2S) \to K + (K_s K^- \pi^+) \psi$ with a mass which was incompatible with the Crystal Ball candidate. The Belle finding was then confirmed by CLEO [7] and BaBar [8] which observed this state in two-photon fusion. The PDG value of the mass is $3638 \pm 4$ MeV/$c^2$, corresponding to a surprisingly small hyperfine splitting of $48 \pm 4$ MeV/$c^2$, whereas the total width is only measured with an accuracy of $50\%$. The study of this state has just started and all its properties need to be measured with good accuracy.

The $^3P_1$ state of charmonium ($h_c$) is of particular importance in the determination of the spin-dependent component of the $q\bar{q}$ confinement potential. The Fermilab experiment E760 reported an $h_c$ candidate in the decay channel $J/\psi \pi^0$ [9], with a mass of $3526.2 \pm 0.15 \pm 0.2$ MeV/$c^2$. This finding was not confirmed by the successor experiment E835, which however observed an enhancement in the $\eta_c\gamma$ [10] final state at a mass of $3525.8 \pm 0.2 \pm 0.2$ MeV/$c^2$. The $h_c$ was finally observed by the CLEO collaboration [11] in the process $e^+e^- \to \psi' \to h_c + \pi^0$ with $h_c \to \eta_c + \gamma$, in which the $\eta_c$ was identified via its hadronic decays. They found a value for the mass of $3524.4 \pm 0.6 \pm 0.4$ MeV/$c^2$. It is clear that the study of this state has just started and that many more measurements will be needed to determine its properties, in particular the width.

The region above $D\bar{D}$ threshold is rich in interesting new physics. In this region, close to the $D\bar{D}$ threshold, one expects to find the four $1D$ states. Of these only the $1^3D_1$, identified with the $\psi(3770)$ resonance, has been found. The $J = 2$ states ($1^3D_2$ and $1^3D_2$) are predicted to be narrow, because parity conservation forbids their decay to $D\bar{D}$. In addition to the $D$ states, the radial excitations of the $S$ and $P$ states are predicted to occur above the open charm threshold. None of these states have been positively identified.

The experimental knowledge of this energy region comes from data taken at the early $e^+e^-$ experiments at SLAC and DESY and, more recently, at the $B$-factories, CLEO-c and BES. The structures and the higher vector states observed by the early $e^+e^-$ experiments have not all been confirmed by the latest much more accurate measurements by BES [12]. A lot of new states have recently been discovered at the $B$-factories, mainly in the hadronic decays of the meson: these new states ($X, Y, Z, ...$) are associated with charmonium because they decay predominantly into charmonium states such as the $J/\psi$ or the $\psi'$, but their interpretation is far from obvious. The situation can be roughly summarized as follows:

- the $Z(3931)$ [13], observed in two-photon fusion and decaying predominantly into $D\bar{D}$, is tentatively identified with the $2(2S)$;
- the $X(3940)$ [14], observed in double charmonium events, is tentatively identified with the $\eta_c(3S)$;
- for all other new states ($X(3872), Y(3940), Y(4260), Y(4320) and so on$) the interpretation is not at all clear, with speculations ranging from the missing $c\bar{c}$ states, to molecules, tetraquark states, and hybrids. It is obvious that further measurements are needed to determine the nature of these new resonances.

The main challenge of the next years will be thus to understand what these new states are and to match these experimental findings to the theoretical expectations for charmonium above threshold.

### 3.1.3. Charmonium in PANDA

Charmonium spectroscopy is one of the main items in the experimental program of PANDA, and the design of the detector and of the accelerator are optimized to be well suited for this kind of physics. PANDA will represent a substantial improvement over the Fermilab experiments E760 and E835.
• up to ten times higher instantaneous luminosity ($\mathcal{L} = 2 \times 10^{32}\text{cm}^{-2}\text{s}^{-1}$ in high-luminosity mode, compared to $2 \times 10^{31}\text{cm}^{-2}\text{s}^{-1}$ at Fermilab);
• better beam momentum resolution ($\Delta p/p = 10^{-5}$ in high-resolution mode, compared with $10^{-4}$ at Fermilab);
• a better detector (higher angular coverage, magnetic field, ability to detect the hadronic decay modes).

At full luminosity PANDA will be able to collect several thousand $\sigma$ states per day. By means of fine scans it will be possible to measure masses with accuracies of the order of 100 KeV and widths to 10% or better. The entire energy region below and above open charm threshold will be explored.

3.2. Gluonic Excitations

One of the main challenges of hadron physics, and an important item in the PANDA physics program, is the search for gluonic excitations, i.e. hadrons in which the gluons can act as principal components. These gluonic hadrons fall into two main categories: glueballs, i.e. states of pure glue, and hybrids, which consist of a $q\bar{q}$ pair and excited glue. The additional degrees of freedom carried by gluons allow these hybrids and glueballs to have $J^{PC}$ exotic quantum numbers: in this case mixing effects with nearby $q\bar{q}$ states are excluded and this makes their experimental identification easier. The properties of glueballs and hybrids are determined by the long-distance features of QCD and their study will yield fundamental insight into the structure of the QCD vacuum.

Antiproton-proton annihilations provide a very favourable environment in which to look for gluonic hadrons. Two particles, first seen in $\pi N$ scattering $^{13}$ with exotic quantum numbers $J^{PC} = 1^{-+}$, $\pi_1(1400)$ $^{16}$ and $\pi_1(1600)$ $^{17}$, are clearly seen in $\overline{p}p$ annihilation at rest. On the other hand a narrow state at 1500 MeV/c$^2$ discovered in $\overline{p}p$ annihilations by the Crystal Barrel experiment $^{18}$ is considered the best candidate for the glueball ground state ($J^{PC} = 0^{++}$), even though the mixing with nearby $q\bar{q}$ states makes this interpretation difficult.

So far the experimental search for glueballs and hybrids has been mainly carried out in the mass region below 2.2 MeV/c$^2$. PANDA will extend the search to higher masses and in particular to the charmonium mass region, where light quark states form a structureless continuum and heavy quark states are far fewer in number. Therefore exotic hadrons in this mass region could be resolved and identified unambiguously.

3.2.1. Charmonium Hybrids

The spectrum of charmonium hybrid mesons can be calculated within the framework of various theoretical models, such as the bag model, the flux tube model, the constituent quark model and recently, with increasing precision, from Lattice QCD (LQCD). For these calculations the parameters are fixed according to the properties of the known $q\bar{q}$ states. All model predictions and LQCD calculations agree that the masses of the lowest lying charmonium hybrids are between 4.2 GeV/c$^2$ and 4.5 GeV/c$^2$. Three of these states are expected to have $J^{PC}$ exotic quantum numbers ($0^{-+}$, $1^{-+}$, $2^{-+}$), making their experimental identification easier since they will not mix with nearby $q\bar{q}$ states. These states are expected to be narrower than conventional charmonium, because their decay to open charm will be suppressed or forbidden below the $D\overline{D}$ threshold. The cross sections for the formation and production of charmonium hybrids are estimated to be similar to those of normal charmonium states, which are within experimental reach. Formation experiments will generate only non-exotic charmonium hybrids, whereas production experiments will yield both exotic and non-exotic states. This feature can be exploited experimentally: the observation of a state in production but not in formation will be, in itself, a strong hint of exotic behavior.

3.2.2. Glueballs

The glueball spectrum can be calculated within the framework of LQCD in the quenched approximation $^{19}$. In the mass range accessible to PANDA as many as 15 glueball states are predicted, some with exotic quantum numbers ($oddballs$). As with hybrids, exotic glueballs are easier to identify experimentally since they do not mix with conventional mesons. The complications arising from mixing with normal $q\bar{q}$ states is well illustrated by the case of the $f_0(1500)$. As mentioned above, this narrow state, observed at LEAR by the Crystal Barrel $^{18}$ and Obelix $^{20}$ experiments, is considered the best candidate for the ground state glueball. However this interpretation is not unique, and relies on the combined analysis of the complete set of two-body decays of the $f_0(1500)$ and two other scalar states, the $f_0(1370)$ and the $f_0(1710)$. This analysis yields the following mixing picture $^{21}$:

$$|f_0(1710)| = 0.39|gg > + 0.14|N\overline{N} > + 0.14|N\overline{N} >$$

(4)

$$|f_0(1500)| = - 0.69|gg > + 0.37|N\overline{N} > - 0.62|N\overline{N} >$$

(5)

$$|f_0(1370)| = 0.60|gg > - 0.13|N\overline{N} > - 0.79|N\overline{N} >$$

(6)

where $|N\overline{N} > = (|u\overline{u} > + |d\overline{d} >)/\sqrt{2}$. Other scenarios for the scalar meson nonet not involving a glueball have been proposed and this makes the interpretation of the $f_0(1500)$ as the ground state glueball ambiguous. This example highlights the need to extend the glueball search to higher mass regions, which are free of the problem of mixing with conventional $q\bar{q}$ states.
3.3. Hadrons in Nuclear Matter

The study of medium modifications of hadrons embedded in hadronic matter is aimed at understanding the origin of hadron masses in the context of spontaneous chiral symmetry breaking in QCD and its partial restoration in a hadronic environment. So far experiments have been focussed on the light quark sector: evidence of mass changes for pions and kaons have been deduced by the study of deeply bound pionic atoms [22] and of K meson production in proton-nucleus and heavy-ion collisions [23].

The high-intensity \( \bar{p} \) beam of up to 15 GeV/c will allow an extension of this program to the charm sector both for hadrons with hidden and open charm. The in-medium masses of these states are expected to be affected primarily by the gluon condensate. Recent theoretical calculations predict small mass shifts (5-10 MeV/c\(^2\)) for the low-lying charmonium states [24] and more consistent effects for the \( \chi_{cJ} \) (40 MeV/c\(^2\)), \( \psi' \) (100 MeV/c\(^2\)) and \( \psi(3770) \) (140 MeV/c\(^2\)) [25].

D mesons, on the other hand, offer the unique opportunity to study the in-medium dynamics of a system with a single light quark. Recent theoretical calculations agree in the prediction of a mass splitting for D mesons in nuclear matter but, unfortunately, they disagree in sign and size of the effect.

Experimentally the in-medium masses of charmonium states can be reconstructed from their decay into di-leptons and photons, which are not affected by final state interaction. D meson masses, on the other hand, need to be reconstructed by their weak decays into pions and kaons which makes the direct measurement of mass modifications difficult. Therefore other signals have been proposed for the detection of in-medium mass shifts of D mesons: in particular it has been speculated that a lowering of the D\( \bar{D} \) threshold would result in an increased D and \( \bar{D} \) production in \( \bar{p} \)-nucleus annihilations [26] or in an increase in width of the charmonium states lying close to the threshold [27].

Another study which can be carried out in PANDA is the measurement of \( J/\psi \) and D meson production cross sections in \( \bar{p} \) annihilation on a series of nuclear targets. The comparison of the resonant \( J/\psi \) yield obtained from \( \bar{p} \) annihilation on protons and different nuclear targets allows to deduce the \( J/\psi \)-nucleus dissociation cross section, a fundamental parameter to understand \( J/\psi \) suppression in relativistic heavy ion collisions interpreted as a signal for quark-gluon plasma formation.

4. Open Charm Physics

The HESR running at full luminosity and at \( \bar{p} \) momenta larger than 6.4 GeV/c would produce a large number of D meson pairs. The high yield (e.g. 100 charm pairs per second around the \( \psi(4040) \)) and the well defined production kinematics of D meson pairs would allow to carry out a significant charmed meson spectroscopy program which would include, for example, the rich D and \( D_s \) meson spectra.

The B-factory experiments have discovered several new resonances in the D and \( D_s \) sectors, where two are extremely narrow: the \( D_{sJ}^*(2317) \) [28] and the \( D_{sJ}^*(2317) \) [29]. These new states appear at unexpected locations, since their masses are more than 140 MeV/c\(^2\) lower than expected from potential models, as shown in Fig. 5. This has given rise to speculations about their nature. It is important to verify these findings by means of new measurements. Threshold pair production can be employed for precision measurements of the mass and the width of the narrow excited D states.

5. The PANDA Detector

In order to carry out the physics program discussed above the PANDA detector must fulfil a number of requirements: it must provide (nearly) full solid angle coverage, it must be able to handle high rates (2 \( \times \) 10\(^7\) annihilations/s) with good particle identification and momentum resolution for \( \gamma \), e, \( \mu \), \( \pi \), K and p. Additional requirements include vertex reconstruction capability and, for charmonium, a pointlike interaction region, efficient lepton identification and excellent calorimetry (both in terms of resolution and of sensitivity to low-energy showers).

A schematic view of the PANDA detector is shown in Fig. 6. The antiprotons circulating in the HESR
hit an internal hydrogen target (either pellet or cluster jet), while for the nuclear part of the experimental program wire or fiber targets will be used. The apparatus consists of a central detector, called Target Spectrometer (TS) and a Forward Spectrometer (FS).

The TS, for the measurement of particles emitted at laboratory angles larger than $5^\circ$, will be located inside a solenoidal magnet which provides a field of 2 T. Its main components will be a microvertex silicon detector, a central tracker (either a straw tube detector or a time projection chamber), an inner time-of-flight telescope, a cylindrical DIRC (Detector of Internally Reflected Light) for particle identification, an electromagnetic calorimeter consisting of PbWO$_4$ crystals, a set of muon counters and of multiwire drift chambers.

The FS will detect particles emitted at polar angles below $10^\circ$ in the horizontal and $5^\circ$ in the vertical direction. It will consist of a 2 T-m dipole magnet, with tracking detectors (straw tubes or multiwire chambers) before and after for charged particle tracking. Particle identification will be achieved by means of Čerenkov and time-of-flight detectors. Other components of the FS are an electromagnetic and a hadron calorimeter.

All detector components are currently being developed within a very active R&D program. This continued development implies that the choice has not yet been finalized for all detector elements.

6. Conclusions

The availability of high-intensity, cooled antiproton beams at FAIR will make it possible to perform a very rich experimental program.

The PANDA experiment will perform high-precision hadron spectroscopy from $\sqrt{s} = 2.25$ GeV to $\sqrt{s} = 5.5$ GeV and produce a wealth of new results:

- precision measurement of the parameters of all charmonium states, both below and above open charm threshold, with the possible discovery of the missing states (e.g. the D-wave states), which will lead to a full understanding of the charmonium spectrum;
- the observation/discovery of glueballs and hybrids, particularly in the mass range between 3 and 5 GeV/c$^2$, yielding new insights into the structure of the QCD vacuum;
- the measurement of mass shifts of charmonium and open charm mesons in nuclear matter, related to the partial restoration of QCD chiral symmetry in a dense nuclear medium;
- open charm spectroscopy ($D$ and $D_s$ spectra).

All these new measurements will make it possible to achieve a very significant progress in our understanding of QCD and the strong interaction. We are looking forward to many years of exciting hadron physics at FAIR.
References


LHCb status and charm physics program

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LHCb is a dedicated flavor physics experiment that will observe the 14 TeV proton–proton collisions at CERN’s Large Hadron Collider (LHC). Construction of the LHCb detector is near completion, commissioning of the detector is well underway, and LHCb will be fully operational and ready to take data in advance of the projected May 2008 turn-on date for the LHC. The LHCb software trigger will feature a dedicated channel for events containing $D^*$ mesons that will dramatically enhance the statistical reach of LHCb in many charm physics measurements. The LHCb charm physics program is initially focused on mixing and CP violation measurements in two body decay modes of $D^0$. A much broader program is possible and will be explored as manpower allows. We intend to use both promptly produced charm and secondary charm from $B$ meson decays in measurements. Initial studies have focused on using secondary $D^{∗+}$ mesons for mixing measurements in two body decays. Preliminary Monte Carlo studies indicate that LHCb may obtain a statistical precision of $\sigma_{\text{stat}}(x^2) = \pm 0.064 \times 10^{-3}$ and $\sigma_{\text{stat}}(y') = \pm 0.87 \times 10^{-3}$ from a time dependent mixing analysis of wrong sign two body $D^0 \to \pi^- K^+$ decays and a statistical precision of $\sigma_{\text{stat}}(y_{CP}) = \pm 0.5 \times 10^{-3}$ from a ratio of the lifetimes of $D^0$ decays to the final states $K^- K^+$ and $K^- \pi^+$ in 10 fb$^{-1}$ of data.

1. LHCb status

As the dedicated flavor experiment at CERN’s Large Hadron Collider (LHC), LHCb is designed to optimally exploit the large $b\bar{b}$ production cross-section in the LHC 14 TeV proton–proton collisions for precision measurements of $b$ hadron properties. Figure 1 shows the layout of the LHCb detector. In high energy hadronic collisions that produce $b\bar{b}$ pairs, the $b$- and $\bar{b}$-hadrons are predominantly produced into the same forward cone—a fact that led to LHCb’s single-arm spectrometer design [2]. The angular acceptance of the detector extends from approximately 10 mrad around the beam axis to 300 mrad in the magnetic bending plane and to 250 mrad in the non-bending plane.

Many of the features that make LHCb an excellent $B$ physics laboratory also make LHCb well-suited for many charm physics studies at unprecedented levels of precision. The silicon Vertex Locator (VELO) will provide the excellent vertex resolutions necessary for time dependent measurements—an estimated 45 fs proper time resolution for $D^0 \to K^- \pi^+$ decays where the $D^0$ mesons are produced in $b$-hadron decays. The LHCb tracking system will supply precise momentum measurements—an estimated 6 MeV mass resolution for two body decays of $D^0$ mesons. The LHCb Ring Imaging Cherenkov (RICH) detectors will provide excellent $K^- \pi$ discrimination over a wide momentum range from 2 GeV/c to 100 GeV/c. Finally, the LHCb trigger system will have a high statistics charm stream, described in Section 2, so that the large charm production in LHC collisions can be exploited for precision measurements.

As of September 2007, construction of the LHCb detector is well advanced with commissioning activities underway for all sub-detectors. LHCb is on-schedule to be complete and ready for data taking by the projected LHC turn-on date in May 2008.

2. LHCb trigger and $D^*$ stream

LHCb will have a two stage trigger: a fast hardware trigger called the Level 0 Trigger (L0) followed by a software High Level Trigger (HLT). Although the triggers are designed to favor $b\bar{b}$ events, the HLT will feature a dedicated $D^*$ stream for selecting charm events at a high rate.

At design operation, LHCb will observe bunch crossings at 40 MHz with a luminosity of $2 \times 10^{32}$ cm$^{-2}$ s$^{-1}$. The L0 trigger is designed to reduce the 40 MHz input rate to approximately 1 MHz while efficiently favoring $b\bar{b}$ events. Using the fact that the decay products of $b$-hadrons typically have significant transverse momentum, the L0 trigger reads data from the calorimeters and the muon detectors to identify quickly individual candidate hadrons, electrons, photons, and muons that have a few GeV of transverse energy or momentum. The
Table I: Estimated yield of reconstructed secondary $D^{*+} \to \pi^{+}\pi^{0\pi}$(h$^{-}h^{+}$) candidates in 2 fb$^{-1}$ of LHCb data passing the LHCb trigger sequence.

<table>
<thead>
<tr>
<th>Two body $D^{0}$ mode</th>
<th>HLT Yield in 2 fb$^{-1}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$D^{0} \to K^{+}\pi^{+}$</td>
<td>$50 \times 10^{6}$</td>
</tr>
<tr>
<td>$D^{0} \to K^{+}\pi^{-}$</td>
<td>$5 \times 10^{6}$</td>
</tr>
<tr>
<td>$D^{0} \to \pi^{+}\pi^{+}$</td>
<td>$2 \times 10^{6}$</td>
</tr>
<tr>
<td>$D^{0} \to \pi^{-}K^{+}$</td>
<td>$0.2 \times 10^{6}$</td>
</tr>
</tbody>
</table>

L0 $E_{T}$ thresholds must be tuned on real data, but, to illustrate the expected scale, current detector studies indicate a trigger threshold for hadrons of $E_{T} > 3.5$ GeV.

The HLT software trigger performs an event reconstruction to identify events of specific physics interest, reducing the 1 MHz L0 output/HLT input rate to approximately 2 kHz. The HLT has access to all of the detector information and uses it to reconstruct final state particle candidates and the locations of the primary proton-proton interactions. These objects are then used in several parallel channels to identify events of specific physical interest. Although the configuration of the HLT channels has yet to be finalized, the HLT will contain a high yield stream for charm events. In the current configuration, 300 Hz, 15% of the recorded 2 kHz of LHCb events, are allocated to an inclusive $D^{*}$ stream. Table I shows potential yields of some key charm decays in a preliminary configuration the $D^{*}$ stream. These estimates are based on the performance of the HLT on fully simulated LHCb events. LHCb will record at least $50 \times 10^{6}$ fully reconstructed $D^{*+} \to \pi^{+}\pi^{0\pi}$(h$^{-}h^{+}$) candidates per 2 fb$^{-1}$ (a nominal year of LHCb data at design luminosity), where the $D^{*+}$ originates in a b-hadron decay and $h, h^{*} \in \{K, \pi\}$. The subscript on the $\pi^{+}$ labels it as the tagging ‘slow’ pion. Studies indicate that this HLT configuration also yields a similar number of reconstructed prompt $D^{*+}$ candidates in each mode.

### 3. $D^{*+}$ event selection

The charm physics program at LHCb is initially focused on mixing and CP violation measurements in two body decays of $D^{0}$. LHCb note 4 details a preliminary selection of wrong sign (WS) $D^{0} \to K^{-} \pi^{+}$ decays from $D^{*+} \to D^{0}\pi^{+}$, where the $D^{*+}$ originates from a B meson decay. The intent of the selection is to provide a sample of candidates suitable for a time dependent mixing analysis (see Section 4). The performance of the selection on fully simulated LHCb data predicts a yield of approximately 230,000 true WS decays in 10 fb$^{-1}$ of LHCb data (5 years of nominal LHCb data taking) with a background-to-signal ratio of 1.07 < $B/S < 5.28$ at the 90% confidence level.

### 3.1. Selection backgrounds

The primary backgrounds accepted by this selection result from combinatoric coincidences and can be divided into two classes: random slow pion backgrounds where properly reconstructed right sign $D^{0} \to K^{-} \pi^{+}$ decays are combined with random pions produced somewhere else in the event to mimic a $D^{*\pm}$ decay, and pure combinatoric coincidences where the candidate $D^{0}$ decay products come from different decays in the event. These backgrounds will be separated from the signal with the typical method of fitting the reconstructed $D^{0}$ mass ($m_{D^{0}}$) and $D^{*+}$-$D^{0}$ mass difference ($\Delta m$) distributions.

### 3.2. $D^{*+}$ decay vertex

In order to perform a time dependent mixing measurement, both the creation and decay vertices of the $D^{0}$ must be precisely determined. The $D^{0}$ decay vertex can be measured from its decay products very precisely in the VELO. Table II shows the results of vertex resolution studies in fully simulated LHCb events. The coordinate system in the table is that defined in Figure 1 with the primary proton-proton collisions along the $z$ axis. The $D^{0}$ decay vertex can be determined with a resolution of 257 $\mu$m along the beam axis. To provide some context for this value, the mean laboratory flight distance for a 60 GeV/c $D^{0}$ (the mean momentum of $D^{0}$ from $D^{*+}$ from B mesons) is approximately 4 mm.

In contrast, the $D^{*+}$ decay vertex ($D^{0}$ creation vertex) is poorly determined from its decay products. The small mass difference between the $D^{*+}$ and its decay products leads to a narrow laboratory frame angle between the $D^{0}$ and $\pi^{0}$ momenta. The $D^{+}$ column in Table II shows that the resolution of the $D^{*+}$ decay vertex estimated only from its decay products is 4322 $\mu$m, the same size as the mean laboratory flight distance of $D^{0}$ mesons. The precision of the $D^{*+}$ vertex must be improved by including additional tracks from particles created with the $D^{*+}$. For $D^{*+}$ from B decays, this means finding additional charged particles created at the B decay vertex.

In studies of fully simulated events, 63% of $B \to D^{*+}X$ decays in triggered events produce an additional reconstructed charged particle at the B meson decay vertex that can be used to improve the precision of the estimated $D^{0}$ birth vertex. As shown in the $B_{\text{part}}$ column of Table II, using such additional tracks dramatically improves the precision of the estimated $D^{0}$ production vertex, and, consequently, the measured $D^{0}$ proper time. The subscript on $B_{\text{part}}$ signifies that the parent $B$ is partially reconstructed. Figure 2 shows the dramatic improvement in measured proper time obtained by using the $B_{\text{part}}$ decay vertex as the $D^{0}$ production vertex. In these plots, the reconstructed proper time is signed to represent
whether the $D^0$ momentum and flight direction are aligned (positive proper time) or anti-aligned (negative proper time). When the $D^{*+}$ decay vertex is used in calculating the $D^0$ proper time, its resolution dominates the exponential decay distribution as shown in Figure 2a. When the $B_{\text{part}}$ decay vertex is used, as in Figure 2b, the proper time resolution is relatively narrow and the reconstructed proper time distribution closely reproduces the generated proper time. Preliminary work in [4] has demonstrated that this partial $B$ reconstruction can be done with an 80% efficiency, and that LHCb can use $D^{*+}$ mesons from $B$ decays for time dependent analyses.

4. Charm mixing measurements at LHCb

Mixing observables are commonly expressed in terms of two dimensionless parameters: the mass difference parameter, $x$, and the full width difference parameter, $y$, defined by

$$x = \frac{2(m_1 - m_2)}{\Gamma_1 + \Gamma_2}, \quad y = \frac{\Gamma_1 - \Gamma_2}{\Gamma_1 + \Gamma_2},$$

where the subscripts denote the mass eigenstates of the $D^0$-$\overline{D}^{0}$ system. Various measurements are sensitive to different combinations of these variables, and each of these should be explored with the highest possible precision to gain a full understanding of the charm mixing phenomenon. Preliminary work at LHCb has focused on the measurement of mixing parameters in a time dependent analysis of two body WS decays (Section 4.1 below) and in an analysis of the ratios of two body lifetimes (Section 4.2 below). However, future plans include multi-body mixing measurements (Section 4.3).

4.1. Time dependent WS $D^0 \rightarrow \pi^- K^+$

The time dependent analysis of WS $D^0 \rightarrow \pi^- K^+$ decays is one of the long established methods of searching for $D^0-\overline{D}^0$ mixing [5-8]. If $D^0$ and $\overline{D}^0$ mix, a meson created as a $D^0$ may decay to the WS final state $\pi^- K^+$ either directly, by a doubly Cabibbo suppressed (DCS) decay, or indirectly, by mixing into a $\overline{D}^0$ meson that undergoes a Cabibbo favored (CF) decay. Interference between the two processes leads to a decay time dependence that can be expanded to leading order in the small parameters $x$ and $y$ (in the absence of CP violation) as

$$r_{\text{WS}}(t) \propto e^{-\Gamma t} \left( R_D + \sqrt{R_D y'} (\Gamma t) + \frac{1}{2} R_M (\Gamma t)^2 \right),$$

where $R_D$ is the ratio of the DCS decay rate to the CF decay rate, $R_M = (x'^2 + y'^2)/2 = (x'^2 + y'^2)/2$ is the mixing rate, and $x'$ and $y'$ are rotated with respect to the parameters $x$ and $y$ by the relative strong phase between the CF and DCS decays, $\delta$:

$$x' \equiv x \cos \delta + y \sin \delta,$$

$$y' \equiv y \cos \delta - x \sin \delta.$$
Hence, a detailed analysis of the WS $D^0$ proper time distribution is sensitive to $x^2$ and $y'$. The strong phase $\delta$ that relates these quantities to the mixing parameters $x$ and $y$ must be measured independently. Since $x^2$ enters the decay time distribution only in $R_M = (x^2 + y'^2)/2$, the values of $x^2$ measured by this method are highly anti-correlated to the values of $y'$.

Because of the small values of $x$ and $y$, it is only recently that this method has yielded values over 3$\sigma$ from $x^2 = 0$, $y' = 0$ [5]. The large charm statistics at LHCb should enable LHCb to improve significantly on searching for CP violation in two body decays. The selection described in Section 3, with appropriate modifications of the final state particle identification criteria, yields approximately $8 \times 10^6 D^{*+}$ tagged $D^0 \to K^-\pi^+$ decays and $3 \times 10^6 D^{*+}$ tagged $D^0 \to \pi^-\pi^+$ decays originating from $b$-hadron decays in 10 fb$^{-1}$ of LHCb data. LHCb expects to use its statistical advantage over current experiments to improve the precision of the lifetime ratio measurement of $y_{CP}$.

The LHCb statistical sensitivity to $y_{CP}$ has been estimated with a toy Monte Carlo study similar to the WS mixing toy study described in Section 4.1. The selection described in Section 3, with appropriate modifications of the final state particle identification criteria, yields approximately $8 \times 10^6 D^{*+}$ tagged $D^0 \to K^-\pi^+$ decays and $3 \times 10^6 D^{*+}$ tagged $D^0 \to \pi^-\pi^+$ decays originating from $b$-hadron decays in 10 fb$^{-1}$ of LHCb data. The signal to background ratios of the selection are $S/B = 4.8$ for the $K^-\pi^+$ mode and $S/B = 2.6$ for the $\pi^-\pi^+$ mode. Using the estimated $D^0 \to K^-\pi^+$ yield and signal to background ratio with proper time resolution and acceptance functions determined from fully simulated LHCb events, the estimated statistical precision of $y_{CP}$ is $\sigma_{stat}(y_{CP}) = \pm 0.5 \times 10^{-3}$.

5. Searches for CP violation at LHCb

The Standard Model (SM) predicts any CP violation in charm interactions to be very small. Observable CP violation at the level of 1% would be an unambiguous sign of new physics [14]. Each of the mixing measurements performed at LHCb will be analyzed in charge conjugate subsets to measure possible CP violating effects. In addition, LHCb will perform time integrated CP violation searches in as many charm decays as are possible. Initial studies have focused on searching for CP violation in two body decays of $D^{*+}$ tagged $D^0$ mesons, in particular the CP eigenstate decays $D^0 \to K^-\pi^+$ and $D^0 \to \pi^-\pi^+$. These singly Cabibbo suppressed decays, in which a small CP violation is predicted by the SM, are particularly...
sensitive to CP violation enhanced by well-motivated new physics scenarios [15]. Experimental measurements in this channel have steadily reduced the upper limit of CP violation with increasing data set sizes and improved treatments of systematic uncertainties [12, 16, 17, 18, 19]. However, CP violation at the order of 1% has not been ruled out.

With $8 \times 10^6$ tagged $D^0 \to K^-K^+$ decays and $3 \times 10^6$ tagged $D^0 \to \pi^-\pi^+$ decays (Section 4.2), LHCb will have the statistical power to search for CP asymmetries to order $O(0.0004)$ or below, provided the systematic uncertainties can be controlled to this level.

Charm meson production asymmetries and final state particle detector asymmetries, particularly the detector asymmetries associated with the tagging slow pion, are expected to be the primary sources of systematic uncertainties in CP asymmetry measurements. Methods of measuring the production and detections asymmetries precisely from data are under development. Also, advantage may be gained by comparing the asymmetries of related decays. For example, the decays $D^0 \to K^-K^+$ and $D^0 \to \pi^-\pi^+$ are subject to the same production and slow pion detection asymmetries, so the difference of their measured asymmetries will have a much smaller systematic uncertainty than either asymmetry measured separately. Although this difference will be small, it is an observable that can be measured very precisely, and, if found to be significantly different from zero, can provide evidence of direct CP violation in at least one of the two decay channels.

6. Multi-body channels

LHCb will also investigate the use of $D$ meson decays to three or more final state hadrons in mixing and CP violation measurements. However, development of multi-hadron charm analyses is less advanced than the two body $D^0$ program. For example, a time dependent amplitude analysis of the three body decay $D^0 \to K_S^0\pi^+\pi^-$ is directly sensitive to the mixing parameters $x$ and $y$. This decay mode should be efficiently reconstructible at LHCb. Preliminary work is also under way to develop selections for $D^0$ decays to four hadrons, both for HLT triggering and for analysis. Further development in four body decays will investigate the feasibility of time dependent amplitude analyses for mixing measurements. The technology of four body amplitude analyses is already quite advanced [20].

In four body decays, CP violation searches will include analyses of quantities that are odd under the time reversal operation in addition to complete amplitude analyses of the decays. Although studies are still in their earliest stages, LHCb should be able to reconstruct with acceptable signal-to-background ratios

![Figure 4](image)

Figure 4: Contours representing the $1 \sigma$, $2 \sigma$, and $3 \sigma$ regions for the toy WS mixing study (ellipses) and for the toy lifetime ratio study (bands). The central values are $\sqrt{x^2} = 8.4 \times 10^{-3}$ and $y = 6.4 \times 10^{-3}$, the very preliminary averages of [21]. For purposes of this plot, it is assumed that $x'^2 = x^2$ and $y' = y_{CP} = y$.

the decays $D^0 \to K_S^0h^+h^-$ and three body decays of charged $D^+$ containing at least one kaon. Amplitude analyses of these modes will expand the scope for CP violation searches in charm decays.

7. Conclusions

The LHCb trigger will provide LHCb with charm physics data sets of unprecedented statistics. Current analysis in charm physics has focused on mixing measurements and CP violation searches in two body decays $D^0 \to h^-h^+$, but a broader charm physics program is envisaged. Toy Monte Carlo studies indicate that with 10 fb$^{-1}$ LHCb can achieve a statistical precision of $\sigma_{stat}(x'^2) = \pm 0.064 \times 10^{-3}$ and $\sigma_{stat}(y') = \pm 0.87 \times 10^{-3}$ with a two body wrong sign mixing analysis, and a statistical precision of $\sigma_{stat}(y_{CP}) = \pm 0.5 \times 10^{-3}$ with a two body lifetime ratio measurement. Figure 4 summarizes these precisions by showing the intersection of the $1 \sigma$ ellipse in $(x'^2, y')$, also scaled to $2 \sigma$ and $3 \sigma$, and the toy lifetime ratio $1 \sigma$ band in $y_{CP}$, also scaled to $2 \sigma$ and $3 \sigma$.

References


[3] The use of charge conjugate modes is implied unless otherwise indicated.


Charm Physics Opportunities at a Super Flavor Factory

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The primary physics goals of a high luminosity $e^+e^-$ flavor factory are discussed, including the possibilities to perform detailed studies of the CKM mechanism of quark mixing, and constrain virtual Higgs and non-standard model particle contributions to the dynamics of rare $B_{u,d,s}$ decays. The large samples of $D$ mesons and tau leptons produced at a flavor factory will result in improved sensitivities to rare $D$ processes - mixing, CP violation and rare decays - and lepton flavor violation searches, respectively. Recent developments in accelerator physics have demonstrated the feasibility to build an accelerator that can achieve luminosities of $O(10^{36})$ cm$^{-2}$ s$^{-1}$ at $\sqrt{s} = 10$ GeV. The capability to run at $\sqrt{s} = 3.770$ GeV with luminosity of $10^{35}$ cm$^{-2}$s$^{-1}$ is included in the initial design. This report emphasizes the charm physics that can be probed at a Super Flavor Factory.

These proceedings aim to present a brief overview of the SuperB effort with a special emphasis on the charm physics program of such a facility. In the interest of completeness (and time) some passages from the SuperB Conceptual Design Report[1] are reproduced here.

1. Introduction

Elementary particle physics in the next decade will be focused on the investigation of the origin of electroweak symmetry breaking and the search for extensions of the Standard Model (SM) at the TeV scale. The discovery of New Physics will likely produce a period of excitement and progress recalling the years following the discovery of the $J/\psi$. In this new world, attention will be riveted on the detailed elucidation of new phenomena uncovered at the LHC; these discoveries will also provide strong motivation for the construction of the ILC. High statistics studies of heavy quarks and leptons will have a crucial role to play in this new world.

The two asymmetric $B$ Factories, PEP-II [2] and KEKB[3], and their companion detectors, BaBar [4] and Belle[5], have over the last seven years produced a wealth of flavour physics results, subjecting the quark and lepton sectors of the Standard Model to a series of stringent tests, all of which have been passed. With the much larger data sample made possible by a Super $B$ Factory, qualitatively new studies will be possible. These studies will provide a uniquely important source of information about the details of the New Physics uncovered at hadron colliders in the coming decade.

The continued detailed studies of heavy quark and heavy lepton (henceforth heavy flavour) physics will not only be pertinent in the next decade; they will be central to understanding the flavour sector of New Physics phenomena. A Super Flavour Factory such as SuperB will be a partner with LHC, and eventually, ILC, experiments, in ascertaining exactly what kind of New Physics has been found. The capabilities of SuperB in measuring CP-violating asymmetries in very rare $b$ and $c$ quark decays, accessing branching fractions of heavy quark and heavy lepton decays in processes that are either extremely rare or forbidden in the Standard Model, and making detailed investigations of complex kinematic distributions will provide unique and important constraints in, for example, ascertaining the type of supersymmetry breaking or the kind of extra dimension model behind the new phenomena that many expect to be manifest at the LHC.

The SuperB Conceptual Design Report[1] is the founding document of a nascent international enterprise aimed at the construction of a very high luminosity asymmetric $e^+e^-$ flavour factory. A possible location for SuperB is the campus of the University of Rome “Tor Vergata”, near the INFN National Laboratory of Frascati.

The exciting physics program that can be accomplished with a very large sample of heavy quark and heavy lepton decays produced in the very clean environment of an $e^+e^-$ collider; with a peak luminosity in excess of $10^{36}$ cm$^{-2}$ s$^{-1}$ at the $\Upsilon(4S)$ resonance is described in Ref.[1] and summarized below. This is program complementary to that of an experiment such as LHCb at a hadronic machine. The physics reach of LHCb and SuperB in the $b$-sector are compared in Figure 1. Luminosities of $10^{35}$ cm$^{-2}$s$^{-1}$ at the $\psi(3770)$ are expected. This report focuses on the charm physics that can be probed both near the $\Upsilon(4S)$ resonance and near charm production threshold.

The conceptual design of a new type of $e^+e^-$ collider that produces a nearly two-order-of-magnitude increase in luminosity over the current generation of asymmetric $B$ Factories is described in Ref.[1]. The key idea is the use of low emittance beams produced in an accelerator lattice derived from the ILC Damping Ring Design, together with a new collision region, again with roots in the ILC final focus design, but with important new concepts developed in this design effort. Remarkably, SuperB produces this very large improvement in luminosity with circulating currents and wallplug power similar to those of the current $B$ Factories. The design of an appropriate detector,

Figure 1: Comparison of SuperB with 50 ab$^{-1}$ and and upgraded LHCb 100 fb$^{-1}$. Design luminosity for SuperB is 15 ab$^{-1}$/year. Design luminosity for LHCb is 2 fb$^{-1}$/year. This comparison assumes that SuperB does not integrate luminosity at the $\Upsilon(5S)$. An upgraded LHCb could integrate luminosity at a 10 times greater rate than LHCb.

1.1. The Physics Case for SuperB

By measuring mixing-dependent $CP$-violating asymmetries in the $B$ meson system for the first time, PEP-II/BaBar and KEKB/Belle have shown that the CKM phase accounts for all observed $CP$-violating phenomena in $b$ decays. The Unitarity Triangle construction provides a set of unique overconstrained precision tests of the self-consistency of the three generation Standard Model. Figure 2 shows the current status of the Unitarity Triangle construction, incorporating measurements from BaBar and Belle, as well as the $B_s$ mixing measurement of CDF; the addition of $CP$ asymmetry measurements, together with the improvement in the precision of $CP$-conserving measurements, has made this uniquely precise set of Standard Model tests possible.

The fact that the CKM phase has now been shown to be consistent with all observed $CP$-violating phenomena is both a triumph and an opportunity. In completing the experimentally-verified Standard Model ansatz (except, of course, for the Higgs), it intensifies the mystery of the creation of the baryon-antibaryon asymmetry of the universe: the observed $CP$-violation is too small for the Standard Model to account for electroweak baryogenesis. This intriguing result opens the door to two possibilities: the matter antimatter asymmetry is produced by another mechanism, such as leptogenesis, or baryogenesis proceeds through the additional $CP$-violating phases that naturally arise in many extensions of the Standard Model. These extra phases produce measurable effects in the weak decays of heavy flavour particles. The detailed pattern of these effects, as well as of rare decay branching fractions and kinematic distributions, is, in fact, diagnostic of the characteristics of New Physics at or below the TeV scale.

By the end of this decade, the two $B$ Factories will have accumulated a total of $\sim 2$ ab$^{-1}$. Even at this level, most important measurements pertinent to the Unitarity Triangle construction will still be statistics limited: an even larger data sample would provide increasingly stringent tests of three-generation CKM unitarity. There are two main thrusts here. The first is the substantial remaining improvement that can still be made in the Unitarity Triangle construction. Here measurements in $B$, $D$ and $\tau$ decay play an important role, as do improvements in lattice QCD calculations of hadronic matrix elements. This important physics goal is NOT, however, the sole, or even the primary, motivation for a Super $B$ Factory. The precision of our knowledge of the Unitarity Triangle will perforce improve to the limit allowed by theoretical uncertainties as we pursue the primary goal: improving the precision of the measurement of $CP$.
asymmetries, rare decay branching fractions, and rare decay kinematic distributions in penguin-dominated $b \to s$ transitions, to a level where there is substantial sensitivity to New Physics effects. This requires data samples substantially larger than the current $B$ Factories will provide. Some of these measurements are accessible at the LHC [6], but the most promising approach to this physics is Super$B$, a very high luminosity asymmetric $B$ Factory, which is also, of course, a Super Flavour Factory, providing large samples of $b$ and $c$ quark and $\tau$ lepton decays.

Super$B$, having an initial luminosity of $10^{36}$ cm$^{-2}$s$^{-1}$, will collect 15 ab$^{-1}$ in a New Snowmass Year [7], or 75 ab$^{-1}$ in five years. A data sample this large will make the Unitarity Triangle tests, in their manifold versions, the ultimate precision test of the flavour sector of the Standard Model, and open up the world of New Physics effects in very rare $B$, $D$, and $\tau$ decays.

A primary tool for isolating new physics is the time-dependent $CP$ asymmetry in decay channels that proceed through penguin diagrams, such as the $b \to s\bar{s}s$ processes $B^0_d \to \phi K^0_s$ and $B^0_s \to (K\bar{K})_s CP K^0$ or similar transitions such as $B^0_d \to \eta' F^0$, $B^0_s \to f_0 K^0$, $B^0_d \to \pi^0 K^0$, $B^0_s \to \rho^0 K^0$, $B^0_d \to \omega K^0$, and $B^0_s \to \pi^0 \pi^0 K^0$. The dominant contribution to these decays is the combination of CKM elements $V_{tb} V_{ts}^*$; these amplitudes have the same phase as the charmonium channels $b \to c\bar{s}s$, up to a small phase shift of $V_{ts}$ with respect to $V_{tb}$. New heavy particles contribute new loop amplitudes, with new phases that can contribute to the $CP$ asymmetry and the $S$ coefficient of the time-dependent analysis, so that the measured $CP$ violation parameter could be substantially different from $\sin 2\beta$.

Physics beyond the Standard Model can affect rare $B$ decay modes, through observables such as branching fractions, $CP$-violating asymmetries and kinematic distributions. These decays do not typically occur at tree level, and thus their rates are strongly suppressed in the Standard Model. Substantial enhancements in the rates and/or variations in angular distributions of final state particles could result from the presence of new heavy particles in loop diagrams, resulting in clear evidence of New Physics. Moreover, because the pattern of observable effects is highly model-dependent, measurements of several rare decay modes can provide information regarding the source of the New Physics. An extended run at the $T(5S)$ is also contemplated; such a run would yield a wealth of interesting new $B^0_d$ decay results.

The Super$B$ data sample will also contain unprecedented numbers of charm quark and $\tau$ lepton decays. This data is also of great interest, both for its capacity to improve the precision of existing measurements and for its sensitivity to New Physics. This interest extends beyond weak decays; the detailed exploration of new charmonium states is also an important objective. Limits on rare $\tau$ decays, particularly lepton-flavour-violating decays, already provide important constraints on New Physics models. Super$B$ may have the sensitivity to actually observe such decays. The accelerator design will allow for longitudinal polarization of the $e^-$ beam, making possible uniquely sensitive searches for a $\tau$ electric dipole moment, as well as for $CP$-violating $\tau$ decays.

Some measurements in charm and $\tau$ physics are best done near threshold. Super$B$ also has the capability of running in the 4 GeV region. Short runs at specific center-of-mass energies in this region, representing perhaps 10% of data taking time, would produce data samples substantially larger than those currently envisioned to exist in the next decade. The many advantages of analysis at threshold are enumerated in Section 2.1.

1.2. The Super$B$ Design

Given the strong physics motivation, there has been a great deal of activity over the past few years aimed at designing an $e^+e^-$ $B$ Factory that can produce samples of $B$ mesons 50 to 100 times larger than will exist when the current $B$ Factory programs end. Several approaches were tried before the design[1] described briefly here was developed.

Upgrades of PEP-II [8] and KEKB [9] to Super $B$ Factories that accomplish this goal have been proposed at SLAC and at KEK. These machines are extrapolations of the existing $B$ Factories, with higher currents, more bunches, and smaller $\beta$ functions (1.5 to 3 mm). They also use a great deal of power ($\geq 100$ MW), and the high currents (as much as 10 A) pose significant challenges for detectors. To minimize the substantial wallplug power, the SuperPEP-II design doubled the current RF frequency, to 958 MHz. In the case of SuperKEKB, a factor of two increase in luminosity is assumed for the use of crab crossing, which will soon be tested at KEKB.

SLAC has no current plans for an on-site accelerator-based high energy physics program, so the SuperPEP-II proposal is moribund. As of this writing, no decision has been made on SuperKEKB. In the interim, the problematic power consumption and background issues associated with the SLAC and KEK-based Super $B$ Factory designs stimulated a new approach, using low emittance beams, to constructing a Super $B$ Factory with a luminosity of $10^{36}$, but with reduced power consumption [10].

The current machine concept, which has roots in ILC R&D: a very low emittance storage ring, based on the ILC damping ring minimum emittance growth lattice and final focus, that incorporates several novel accelerator concepts and appears capable of meeting all design criteria, while reducing the power consumption, which dominates the operating costs of the facility, to a level similar to that of the current $B$ Factories.
Due to similarities in the design of the low emittance rings and the final focus, operation of SuperB can serve as a system test for these linear collider components.

By utilizing concepts developed for the ILC damping rings and final focus in the design of the SuperB collider, it is possible to produce a two-order-of-magnitude increase in luminosity with beam currents that are comparable to those in the existing asymmetric $B$ Factories. Background rates and radiation levels associated with the circulating currents are comparable to current values; luminosity-related backgrounds such as those due to radiative Bhabhas, increase substantially. With careful design of the interaction region, including appropriate local shielding, and straightforward revisions of detector components, upgraded detectors based on BaBar or Belle are a good match to the machine environment: in this discussion, we use BaBar as a specific example. Required detector upgrades include: reduction of the radius of the beam pipe, allowing a first measurement of track position closer to the vertex and improving the vertex resolution (this allows the energy asymmetry of the collider to be reduced to 7 on 4 GeV); replacement of the drift chamber, as the current chamber will have exceeded its design lifetime; replacement of the endcap calorimeter, with faster crystals having a smaller Molière radius, since there is a large increase in Bhabha electrons in this region.

The SuperB design has been undertaken subject to two important constraints: 1) the lattice is closely related to the ILC Damping Ring lattice, and 2) as many PEP-II components as possible have been incorporated into the design. A large number of PEP-II components can, in fact, be reused: The majority of the HER and LER magnets, the magnet power supplies, the RF system, the digital feedback system, and many vacuum components. This will reduce the cost and engineering effort needed to bring the project to fruition.

The SuperB concept is a breakthrough in collider design. The invention of the “crabbed waist” final focus can, in fact, have impact even on the current generation of colliders. A test of the crabbed waist concept is planned to take place at Frascati in 2007; a positive result of this test would be an important milestone as the SuperB design progresses. The low emittance lattice, fundamental as well to the ILC damping ring design, allow high luminosity with modest power consumption and demands on the detector.

SuperB appears to be the most promising approach to producing the very high luminosity asymmetric $B$ Factory that is required to observe and explore the contributions of physics beyond the Standard Model to heavy quark and $\tau$ decays.

2. Charm Physics at SuperB

It is a truth universally accepted that charm studies played a seminal role in the evolution and acceptance of the Standard Model. Yet the continuing importance of this sector is not widely appreciated, since the Standard Model electroweak phenomenology for charm decays is on the dull side: the CKM parameters are known, $D^0\bar{D}^0$ oscillations are slow, $CP$ asymmetries are small or absent and loop-driven decays are extremely rare.

Yet on closer examination, a strong case emerges in two respects, both of which derive from this apparent dullness:

- Detailed and comprehensive analyses of charm transitions will continue to provide us with new insights into QCD’s nonperturbative dynamics, and advance us significantly towards establishing theoretical control over them. Beyond the intrinsic value of such lessons, they will also calibrate our theoretical tools for $B$ studies; this will be essential to saturate the discovery potential for New Physics in $B$ transitions.

- Charm decays constitute a novel window into New Physics.

Lessons from the first item will have an obvious impact on the tasks listed under the second. They might actually be of great value even beyond QCD, if the New Physics anticipated for the TeV scale is of the strongly interacting variety.

The capabilities of a Super Flavour Factory are well matched to these goals. It allows uniquely clean determinations of CKM parameters, with six of the nine matrix elements impacted by charm measurements. New Physics signals can easily exceed Standard Model predictions by considerable factors such that there will be no ambiguity in interpreting them, yet they are unlikely to be large; this again requires the clean environment and huge statistics that a Super Flavour Factory can provide.

A number of other facilities either currently running or soon to commence operation provide competition in the area of charm physics. The current $B$ Factory program is expected to produce a sample of about $10^{10}$ charm hadrons from operation at or near the $\Upsilon(4S)$ resonance. The CLEO-c experiment at CESR is operating in the charm threshold region, and anticipates collecting a total of $5 \times 10^8 D^0\bar{D}^0$ pairs and about $7 \times 10^5 D^+_s D^-_s + D^+_s D^-_s$ through coherent production. The BESIII experiment at BEPCII expects first $e^+e^-$ collisions in 2008, and will collect large charmonium samples, in addition to exceeding the CLEO-c data set in open charm production. Although there will be no successors to the Fermilab fixed target charm production experiments, the LHC will produce copious quantities of charm (notably, charm physics
forms a part of the LHCb physics program); these are expected to result in very large samples of charmed hadrons in final states with reconstructible topologies.

Most of the benchmark charm measurements will still be statistics-limited after the CLEO-c, BESIII and B Factory projects, and many will not be achievable in a hadronic environment. SuperB is the ideal machine with which to pursue these measurements to their ultimate precision. Operation near the \( \Upsilon(4S) \) will provide enormous samples of charm hadrons, in a clean environment and with a detector well-suited for charm studies. The charm physics program would benefit further from the ability to operate in the threshold region, in order to exploit the quantum correlations associated with coherent production. The expected lower luminosity at threshold would be partly compensated by the higher production cross-section, resulting in a comparable charm production rate. To estimate the reach of SuperB from operation at the charm threshold, we have assumed a simple dependence of the luminosity on the center-of-mass energy: \( \mathcal{L}_{\text{peak}} \propto s \). Thus, we expect that SuperB (which will integrate \( \sim 15 \, \text{ab}^{-1} \) per year operating at the \( \Upsilon(4S) \)) can accumulate \( \sim 150 \, \text{fb}^{-1} \) per month when operated at the \( \psi(3770) \).

### 2.1. Advantages of Threshold Production

The production rate of charm during threshold running at a SuperB and \( \Upsilon(4S) \) running is comparable. Although the luminosity for charm threshold running is expected to be an order of magnitude lower, the production cross section is 3 times higher than at \( \sqrt{s} = 10.58 \, \text{GeV} \). Charm threshold data has distinct and powerful advantages over continuum and \( b \rightarrow c \) charm production data accumulated above \( B \) production threshold.

**Charm Events at Threshold are Extremely Clean:** The charged and neutral multiplicities in \( \psi(3770) \) events are only 5.0 and 2.4 - approximately 1/2 the multiplicity of continuum charm production at \( \sqrt{s} = 10.58 \, \text{GeV} \).

**Charm Events at Threshold are pure \( D\bar{D} \):** No additional fragmentation particles are produced. The same is true for \( \sqrt{s} = 4170 \, \text{MeV} \) production of \( D\bar{D}^*, D^+\bar{D}^0, D^+\bar{D}^- \), and for threshold production of \( \Lambda_c\bar{\Lambda}_c \). This allows use of kinematic constraints, such as total candidate energy and beam constrained mass, and permits effective use of missing mass methods and neutrino reconstruction. The crisp definition of the initial state is a uniquely powerful advantage of threshold production that is absent in continuum charm production.

**Double Tag Studies are Pristine:** The pure production of \( D\bar{D} \) states, together with low multiplicity and large branching ratios characteristic of many \( D \) decays permits effective use of double-tag studies in which one \( D \) meson is fully reconstructed and the rest of the event is examined without bias but with substantial kinematic knowledge. The techniques pioneered by Mark III and extended by CLEO-c\[13, 14\] allow precise absolute branching fraction determination. Backgrounds under these conditions are heavily suppressed which minimizes both statistical errors and systematic uncertainties.

**Signal/Background is Optimum at Threshold:** The cross section for the signal \( \psi(3770) \rightarrow D\bar{D} \) is about 1/2 the cross section for the underlying continuum \( e^+e^- \rightarrow \text{hadrons} \). By contrast, for \( c\bar{c} \) production at \( \sqrt{s} = 10.58 \, \text{GeV} \) the signal is only 1/4 of the total hadronic cross section.

**Neutrino Reconstruction:** The undetected energy and momentum is interpreted as the neutrino four-vector. For leptonic and semileptonic charm decays the signal is observed in missing mass squared distributions and for double-tagged events these measurements have low backgrounds. The missing mass resolution is about one pion mass. For semileptonic decays the \( q^2 \) resolution is excellent, about 3 times better than in continuum charm reconstruction at \( \sqrt{s} = 10.58 \, \text{GeV} \). Neutrino reconstruction at threshold is clean.

**Quantum Coherence:** The production of \( D \) and \( \bar{D} \) in a coherent \( C = -1 \) state from \( \psi(3770) \) decay is of central importance for the subsequent evolution and decay of these particles. The same is true for \( D\bar{D}(n)\pi^0(m)\gamma \) produced at \( \sqrt{s} \sim 4 \, \text{GeV} \) where \( C = -1 \) for even \( m \) and \( C = +1 \) for odd \( m \). The coherence of the two initial state \( D \) mesons allows both simple and sophisticated methods to measure \( D\bar{D} \) mixing parameters, strong phases, \( CP \) eigenstate branching fractions, and \( CP \) violation\[15–19\].

### 2.2. Lessons on Strong Dynamics

Detailed analyses of (semi)leptonic decays of charm hadrons provide a challenging test bed for validating lattice QCD (LQCD), which is the only known framework with realistic promise for a truly quantitative treatment of charm hadrons that can be systematically improved. Such studies form the core of the ongoing CLEO-c and the nascent BESIII programs; they are also pursued very profitably at the B Factories. Central goals are measuring the decay constants \( f_{D^+} \) and \( f_{D_s} \) and going beyond total rates for semileptonic \( D^+ \), \( D^0 \) and \( D_s^+ \) decays, on the Cabibbo allowed and forbidden level by extracting the form factors etc. It is essential to analyze lepton spectra and perform “meaningful” Dalitz plot studies. To quantify “meaningful” we can compare to analyses on \( K_{e4} \) decays. With a sample size of 30,000 events as it became available in 1977 one was able to begin extracting dynamical information. Precise measurements are possible now with NA48/2 and E685 each having ac-

cumulated 400,000 events. For charm we are nowhere near that level yet: CLEO-c will have about 10,000 semileptonic charm decays – comparable to kaon studies in the late 1970s. Since for charm the phase space is larger (actually a good thing, since it opens up more domains of interest) it seems reasonable to aim for sample sizes of $10^6$ events. Again, this is well beyond the reach of CLEO-c and most probably of BESIII as well. Such high quality studies will greatly improve our understanding of hadronization and provide an even richer test bed for LQCD with the lessons to be learned of crucial importance for extracting $V_{ub}$ from semileptonic $B$ decays. Our knowledge of charm baryon decays is also rather limited; e.g., no precision data on absolute branching ratios or semileptonic decay distributions exist. CLEO-c will not run above the charm baryon threshold, and BESIII cannot.

2.2.1. Leptonic Charm Studies

In the Standard Model the leptonic decay width is given by [20]:

$$
\Gamma(D^+ \to \ell^+ \nu) = \frac{G_F^2}{8\pi} f_{D^+}^2 m_c^2 M_{D^+} \left(1 - \frac{m_{\ell}^2}{M_{D^+}^2}\right)^2 |V_{cd}|^2
$$

$$
\Gamma(D^0 \to \ell^+ \nu) = \frac{G_F^2}{8\pi} f_{D^0}^2 m_c^2 M_{D^0} \left(1 - \frac{m_{\ell}^2}{M_{D^0}^2}\right)^2 |V_{cd}|^2
$$

Taking $|V_{cd}|$ and $|V_{cb}|$ from elsewhere, one uses Eq.(1) to extract $f_{D^+}$ and $f_{D^+}$. The ratio $R_f$ of the leptonic decay rates of the $D^+_s$ and the $D^+$ is proportional to $(f_{D^+_s}/f_{D^+})^2$, for which the lattice calculation is substantially more precise. A significant deviation from its predicted value would be a clear sign of New Physics, probably in the form of a charged Higgs exchange [21]. On the other hand, the ratio of the rates of tauonic and muonic decays for either $D^+$ or $D^+_s$ is independent of both form factors and CKM elements, and serves as a useful cross-check in this context.

CLEO-c has published a measurement of $f_{D^+}$ [22–24], and several measurements of $f_{D^+_s}$ [25–27]. These measurements have benefited from a “double-tag” method uniquely possible at threshold, where a $D^+/(s)$-$D^+_s/(s)$ pair is produced with no extra particles. The latest results are

$$
f_{D^+} = (222.6 \pm 16.7^{+2.8}_{-3.4}) \text{ MeV} \ .
$$

$$
f_{D^0} = (275 \pm 10 \pm 5) \text{ MeV} \ .
$$

$$
f_{D^+_s}/f_{D^+} = 1.24 \pm 0.10 \pm 0.03 \ .
$$

Babar has also measured $f_{D^0} = (283 \pm 17 \pm 7 \pm 14) \text{ MeV}[28]$. The central values for $f_{D^+_s}$ and $f_{D^+_s}/f_{D^+}$ are slightly above, but consistent with, the present LQCD calculations. It is important to note that the desired 1–3% accuracy level has not yet been reached on either the experimental or theoretical side. While LQCD practitioners expect to reach this level over the next decade, the experimental precision is likely to fall significantly short of this goal, even after BESIII. Since larger statistics will certainly allow reduction of the systematic errors in the current results, it is clear that data accumulated by SuperB from a relatively short run (~ 1 month) at charm threshold would allow the desired improvement of the experimental precision (see discussion below, and Table 1). Validating LQCD on the $O(1\%)$ level will have important consequences for $B_d$ and $B_s$ oscillations, since it would give us demonstrated confidence in the theoretical extrapolation to $f_{B_d}$ and $f_{B_s}$. 

2.2.2. Semileptonic Charm Studies

In the area of semileptonic decays, CLEO-c has made the most accurate measurements for the inclusive $D^0$ and $D^+$ semileptonic branching fractions – $B(D^0 \to X \ell \nu) = (6.46 \pm 0.17 \pm 0.13\%)$ and $B(D^+ \to X \ell \nu) = (16.13 \pm 0.20 \pm 0.33\%)$ [29] – and expects to do the same for $D^+_s$. Such data provide important “engineering input" for other $D$ and $B$ decay studies. However, a central goal must be to go beyond the total rates for these decays and to extract the form factors etc. In order to do so, it is essential to analyze lepton spectra and perform “meaningful" Dalitz plot studies. To quantify “meaningful", it is instructive to compare to analyses on $K_{e4}$ decays. With a sample size of 30,000 events which became available in 1977, one was able to begin extracting dynamical information. Precise measurements are now possible, with NA48/2 and E685 each having accumulated 400,000 events [30, 31]. For charm we are nowhere near that level: CLEO-c will have about 10,000 semileptonic charm decays – comparable to kaon studies in the late 1970s. Since for charm the phase space is larger, thereby opening more domains of interest, a reasonable target sample size is $10^6$ events, which is far beyond the reach of CLEO-c, and most probably, of BESIII.

Three-family unitarity constraints on the CKM matrix yield rather precise values for $|V_{cs}|$ and $|V_{cd}|$. Using these, one can extract the form factors from analyses of exclusive semileptonic charm decays. Both the normalization and $q^2$ dependence must be measured. Existing LQCD studies do not allow us to determine the latter from first principles; instead a parametrization originally proposed by Becirevic and Kaidalov (BK) is used [32]. Recent and forthcoming results from CLEO-c, Babar and Belle [33, 34] are expected to be statistics limited, and will not reach the desired 1–3% level.

The current status can be characterized by comparing the measured value of the ratio $R_{sl}$, which is independent of $|V_{cd}|$, to that inferred from a recent
LQCD calculation [35]:

\[
R_{sl} = \sqrt{\frac{\Gamma(D^+ \to \mu^+\nu)\Gamma(D \to \pi\nu\bar{\nu})}{\Gamma(D^+ \to \pi\nu\bar{\nu})}} = \left\{ \begin{array}{c} 0.237 \pm 0.019 \text{ (exp)} \\ 0.212 \pm 0.028 \text{ (theo)} \end{array} \right. \tag{5}
\]

The values are nicely consistent, yet both are still far from the required level of precision.

While operation in the \( \Upsilon \) region will produce large quantities of charm hadrons, there are significant backgrounds and one pays a price in statistics when using kinematic constraints to infer neutrino momenta, etc.. On the other hand, even a limited run at charm threshold will generate the statistics required to study (semi)leptonic decays with the desired accuracy. Assuming that systematic uncertainties in tracking and muon identification will provide a limit to the precision at the 0.5% level, we estimate the integrated luminosity from threshold running required to achieve a similar statistical uncertainty. As shown in Table I we expect to be able to measure \( f_{D^+} \) and \( f_D \) and their ratio with better than 0.5% statistical uncertainty with integrated luminosities of at least 100 fb\(^{-1}\).

Table I Statistics required to obtain 0.5% statistical uncertainties on corresponding branching fractions using double-tagged events, when running at threshold.

<table>
<thead>
<tr>
<th>Channel</th>
<th>Integrated luminosity (fb(^{-1}))</th>
</tr>
</thead>
<tbody>
<tr>
<td>( D^+ \to \mu^+\nu )</td>
<td>500</td>
</tr>
<tr>
<td>( D_s^+ \to \mu^+\nu )</td>
<td>100</td>
</tr>
</tbody>
</table>

For semileptonic decays, a case-by-case study is necessary. One also has to distinguish between merely determining the branching ratio and performing a “meaningful” Dalitz plot analysis, as discussed above. The required integrated luminosities are given in Table II. It is clear that the \( \sim 150 \text{ fb}^{-1} \) anticipated from one month of running in the threshold region would provide the desired statistics for most measurements. Note that while \( D_s \) mesons are not produced at the \( \psi(3770) \), short runs at other energies are possible.

2.3. Precision CKM Measurements

Studies of leptonic decay constants and semileptonic form factors will yield a set of measurements, including \( |V_{cd}| \) and \( |V_{cb}| \), at the few percent level. These measurements will constrain theoretical calculations, and those that survive will be validated for use in a variety of areas in which interesting physics cannot be extracted without theoretical input. This broader impact of charm measurements extends beyond those measurements that can be performed directly at charm threshold, and has a large impact on the precision determination of CKM matrix elements.

The determination of \( |V_{cd}| \) and \( |V_{cb}| \) is limited by ignorance of \( f_B \sqrt{B_{Bs}} \) and \( f_B \sqrt{B_{Bs}} \); improved determinations of \( f_B \) and \( f_{Bs} \) are required. Precision measurements of \( f_D \) and \( f_{Ds} \) can validate the theoretical treatment of the analogous quantities for \( B \) mesons. Similarly, improved form factor calculations in the decays \( D \to \pi\ell\nu \) and \( D \to \rho\ell\nu \) and inclusive semileptonic charm decays enable improved precision in \( |V_{ub}| \) and \( |V_{cb}| \).

The precision measurement of the UT angle \( \gamma \) depends on decays of \( B \) mesons to final states containing neutral \( D \) mesons. A variety of charm measurements impact these analyses, including: improved constraints on charm mixing amplitudes, – important for the GLW method [36, 37], measurements of relative rates and strong phases between Cabibbo-favoured and -suppressed decays measurement of the relative rate and relative strong phase \( \delta \) between \( D^0 \) and \( \bar{D}^0 \) decay to \( K^+\pi^- \) – important for ADS method[38, 39], and studies of charm Dalitz plots tagged by hadronic flavor or \( CP \) content [40–42]. Note that the latter two measurements can only be performed with data from charm threshold.

2.3.1. Overconstraining the Unitarity Triangle

At present three-family unitarity constraints yield more precise values for \( |V_{us}| \) and \( |V_{cd}| \) than direct mea-

Table II Statistics required to obtain 0.5% statistical uncertainties on corresponding branching fractions (column 2) or one million signal events (column 3) using double-tagged events, when running at threshold.

<table>
<thead>
<tr>
<th>Channel</th>
<th>Integrated luminosity (fb(^{-1}))</th>
<th>Integrated luminosity (fb(^{-1}))</th>
</tr>
</thead>
<tbody>
<tr>
<td>( D^0 \to K^-e^+\nu )</td>
<td>1.3</td>
<td>33</td>
</tr>
<tr>
<td>( D^0 \to K^{*-}e^+\nu )</td>
<td>17</td>
<td>425</td>
</tr>
<tr>
<td>( D^0 \to \pi^-e^+\nu )</td>
<td>20</td>
<td>500</td>
</tr>
<tr>
<td>( D^0 \to \rho^-e^+\nu )</td>
<td>45</td>
<td>1125</td>
</tr>
<tr>
<td>( D^+ \to K^0e^+\nu )</td>
<td>9</td>
<td>225</td>
</tr>
<tr>
<td>( D^+ \to K^{\circ}e^+\nu )</td>
<td>9</td>
<td>225</td>
</tr>
<tr>
<td>( D^+ \to \pi^0e^+\nu )</td>
<td>75</td>
<td>1900</td>
</tr>
<tr>
<td>( D^+ \to \rho^0e^+\nu )</td>
<td>110</td>
<td>2750</td>
</tr>
<tr>
<td>( D_s^+ \to \phi e^+\nu )</td>
<td>85</td>
<td>2200</td>
</tr>
<tr>
<td>( D_s^+ \to K^0e^+\nu )</td>
<td>1300</td>
<td>33000</td>
</tr>
<tr>
<td>( D_s^+ \to K^{\circ}e^+\nu )</td>
<td>1300</td>
<td>33000</td>
</tr>
</tbody>
</table>
measurements. Since it is conceivable that a fourth family exists (with neutrinos so heavy that the $Z^0$ could not decay into them), one would like to obtain more accurate direct determinations. This should be possible if LQCD is indeed validated at the $O(1\%)$ level through its predictions on form factors and their ratios.

From four-family unitarity, and using current experimental constraints [43] we can infer for a fourth quark doublet $(t', b')$:

$$|V_{cb'}| = \sqrt{1 - |V_{cb}|^2 - |V_{cs}|^2 - |V_{cb}|^2} \lesssim 0.5 , \quad (6)$$

$$|V_{ts'}| = \sqrt{1 - |V_{ub}|^2 - |V_{cs}|^2 - |V_{ts}|^2} \lesssim 0.5 . \quad (7)$$

These loose bounds are largely due to the 10% error on $|V_{cs}|$.

2.4. Charm as a Window to New Physics

While significant progress can be guaranteed for the Standard Model studies outlined above, the situation is much less certain concerning the search for New Physics. No sign of it has yet been seen, but we have only begun to approach the regime of experimental sensitivity in which a signal for New Physics could realistically emerge in the data. The interesting region of sensitivity extends several orders of magnitude beyond the current status.

New Physics scenarios in general induce flavor-changing neutral currents that a priori have no reason to be as strongly suppressed as in the Standard Model. More specifically, they could be substantially stronger for up-type than for down-type quarks; this can occur in particular in models that reduce strangeness-changing neutral currents below phenomenologically acceptable levels through an alignment mechanism.

In such scenarios, charm plays a unique role among the up-type quarks $u$, $c$ and $t$; for only charm allows the full range of probes for New Physics. Since top quarks do not hadronize [44], there can be no $T^0\bar{T}^0$ oscillations (recall that hadronization, while hard to bring under theoretical control, enhances the observability of $CP$ violation). As far as $u$ quarks are concerned, $\pi^0$, $\eta$ and $\eta'$ do not oscillate, and decay electromagnetically, not weakly. $CP$ asymmetries are mostly ruled out by $CPT$ invariance. Our basic contention can then be formulated as follows: charm transitions provide a unique portal for a novel access to flavor dynamics with the experimental situation being a priori quite favourable. The aim is to go beyond “merely” establishing the existence of New Physics around the TeV scale – we want to identify the salient features of this New Physics as well. This requires a comprehensive study, i.e., that we also search in unconventional areas such as charm decays.

2.4.1. On New Physics scenarios

In a scenario in which the LHC discovers direct evidence of SUSY via observation of sleptons or squarks, the Super Flavour Factory program becomes even more important. The sfermion mass matrices are a new potential source of flavor mixing and $CP$ violation and contain information about the SUSY-breaking mechanism. Direct measurements of the masses can only constrain the diagonal elements of this matrix. However, off-diagonal elements can be measured through the study of loop-mediated heavy flavor processes. As a specific example, a minimal flavor violation scenario such as mSUGRA with moderate tan $\beta$, could result in a SUSY partner mass spectrum that is essentially indistinguishable from an SU(5) GUT model with right-handed neutrinos. However the mSUGRA scenario would be expected to yield no observable effects in the heavy flavor sector, whereas the SU(5) model is expected to produce measurable effects in time-dependent $CP$ violation in penguin-mediated hadronic and radiative decays.

While there is no compelling scenario that would generate observable effects in charm, but not in beauty and strange decays, it is nevertheless reassuring that such scenarios do exist. One should keep in mind that New Physics signals in charm $CP$ asymmetries are particularly clean, since the Standard Model background (which often exists in $B$ decays) is largely absent. The consequence is that New Physics could produce signals that exceed Standard Model predictions by an order of magnitude or more – something that is of great help in interpreting the signals. We will focus on the most promising areas; more details can be found in several recent reviews [17, 45, 46].

2.4.2. $D^0\bar{D}^0$ oscillations

Oscillations of neutral $D$ mesons driven by the two quantities $x_D = \Delta M_D/\Gamma_D$ and $y_D = \Delta \Gamma_D/2\Gamma_D$ lead to an effective violation of the Standard Model $\Delta C = \Delta Q$ and $\Delta C = \Delta S$ rules in semileptonic and nonleptonic channels. The status of the Standard Model prediction can be summarized as [17]: while one predicts $x_D \sim O(10^{-3}) \sim y_D$, at present one cannot rule out $x_D, y_D \sim 0.01$.

Many different charm decay modes can be used to search for charm mixing. The appearance of “wrong-sign” kaons in semileptonic decays would provide direct evidence for $D^0\bar{D}^0$ oscillations (or another process with origin beyond the Standard Model). The wrong-sign hadronic decay $D^0 \to K^+\pi^-$ is sensitive to linear combinations of the mass and lifetime differences, denoted $x^p_D$ and $y_D$. The relation of these parameters to $x_D$ and $y_D$ is controlled by a strong phase difference. Direct measurements of $x_D$ and $y_D$ independent of unknown strong interaction phases, can also be made using time-independent studies of amplitudes present in multi-body decays of the $D^0$, for example, $D^0 \to K^0_{\pi^0\pi^0\pi^-}$. Direct evidence for $y_D \neq 0$ can also appear through lifetime differences between decays to $CP$ eigenstates. The measured quantity in this case,
$y_{CP}$, is equivalent to $y_D$ in the absence of $CP$ violation. Another approach is to study quantum correlations near threshold [17–19] in $e^+e^- \rightarrow D^0\overline{D}^{0}(\pi^0)$ and in $e^+e^- \rightarrow D^0\overline{D}^{0}\gamma$, which yield C-odd and C-even $D^0\overline{D}^{0}$ pairs, respectively.

Very recently, several new results have suggested that charm mixing may be at the upper end of the range of Standard Model predictions. BaBar finds evidence for oscillations in $D^0 \rightarrow K^+\pi^-$ with 3.9σ significance [47], while Belle sees a 3.2σ effect in $D^0 \rightarrow K^+K^-$, with results using $D^0 \rightarrow K_S^0\pi^+\pi^-$ supporting the claim [48]. These results are consistent with previous measurements, some of which had hinted at a mixing effect [49–53]. The results are not systematic limited, and further improvements are anticipated.

The charm decays subgroup of the Heavy Flavor Averaging Group [54] is preparing world averages of all the charm mixing measurements, taking into account correlations between the measured quantities. A preliminary average is available, giving:

$$x_D = (8.7\pm3.0) \times 10^{-3} \quad \text{and} \quad y_D = (6.6^{+2.1}_{-2.0}) \times 10^{-3}.$$  

Contours in the $(x_D, y_D)$ plane are shown in Fig. 3. The significance of the oscillation effect in the preliminary world averages exceeds 5σ.

At present no clear signal has emerged. Since no single measurement exceeds 5σ significance, it is too early to consider charm oscillations as definitively established. Nonetheless, even if one accepts the central value, the interpretation of these new results in terms of New Physics is inconclusive. For one thing, it is not yet clear whether the effect is caused by $x_D \neq 0$ or $y_D \neq 0$ or both, though the latter is favored and this point may be clarified soon. As shown in Table III, SuperB will be able to observe both lifetime and mass differences in the $D^0$ system, if they lie in the range of Standard Model predictions. It should be noted that the full benefit of measurements in the $D^0 \rightarrow K^+\pi^-$ system (and indeed for other hadronic decays) can only be obtained if the strong phases are measured. This can be achieved with a short ($\sim 1$ month) period of data taking at charm threshold.

A serious limitation in the interpretation of charm oscillations in terms of New Physics is the theoretical uncertainty on the Standard Model prediction. Nonetheless, if oscillations indeed occur at the level suggested by the latest results, this will open the window to searches for $CP$ asymmetries that do provide unequivocal New Physics signals.

Table III Summary of the expected precision on charm mixing parameters. For comparison we put the reach of the $B$ Factories at $2 \text{ ab}^{-1}$. The estimates for SuperB assume that systematic uncertainties can be kept under control.

<table>
<thead>
<tr>
<th>Mode</th>
<th>$B$ Factories</th>
<th>SuperB</th>
</tr>
</thead>
<tbody>
<tr>
<td>$D^0 \rightarrow K^+K^-$</td>
<td>$y_{CP}$</td>
<td>$2.3 \times 10^{-3}$</td>
</tr>
<tr>
<td></td>
<td>$y_D$</td>
<td>$2.3 \times 10^{-3}$</td>
</tr>
<tr>
<td></td>
<td>$x_D^2$</td>
<td>$1.2 \times 10^{-4}$</td>
</tr>
<tr>
<td>$D^0 \rightarrow K_S^0\pi^+\pi^-$</td>
<td>$y_D$</td>
<td>$2.3 \times 10^{-3}$</td>
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<tr>
<td></td>
<td>$x_D$</td>
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<tr>
<td>Average</td>
<td>$y_D$</td>
<td>$1.2 \times 10^{-3}$</td>
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<tr>
<td></td>
<td>$x_D$</td>
<td>$2.3 \times 10^{-3}$</td>
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2.4.3. $CP$ Violation With and Without Oscillations

Several factors favor dedicated searches for $CP$ violation in charm transitions:

- Within the Standard Model, the effective weak phase is highly diluted, namely $\sim O(\lambda^4)$, and can arise only in singly-Cabibbo-suppressed transitions, where one expects asymmetries to reach the $O(0.1\%)$ level; significantly larger values would signal New Physics. Any asymmetry in Cabibbo-allowed or -doubly suppressed channels requires the intervention of New Physics – except for $D^\pm \rightarrow K_0^\pm\pi^\pm$ [17] where the $CP$ impurity in $K^0_S$ induces an asymmetry of $3.3 \times 10^{-3}$, CLEO-c measures $A_{CP} = \pm (0.6 \pm 1.0 \pm 0.3)\%$[14]. One should keep in mind that in going from Cabibbo-allowed to Cabibbo-singly and -doubly suppressed channels, the Standard Model rate is suppressed by factors of about twenty and four hundred, respectively. One would expect that this suppression will enhance the visibility of New Physics.

Figure 3: Likelihood contours in the $(x_D, y_D)$ plane from HFAG [54]. These preliminary world averages use all available charm mixing results.
• Strong phase shifts required for direct $CP$ violation to emerge are, in general, large, as are the branching ratios into relevant modes. Although large final state interactions complicate the interpretation of an observed signal in terms of the microscopic parameters of the underlying dynamics, they enhance its observability.

• With the Standard Model providing one amplitude, observable $CP$ asymmetries can be linear in New Physics amplitudes—unlike the case for rare decays—thus increasing the sensitivity.

• Decays to multibody final states contain more dynamical information than given by their widths; their decay distributions as described by Dalitz plots or $T$-odd moments can exhibit asymmetry, observable ability of the underlying dynamics, they enhance its observability. Final state interactions complicate the interpretation of an observable effect two conditions have to be satisfied simultaneously: a transition must respect animal symmetries down to the 10$^{-4}$ level or better.

As already mentioned CKM dynamics does not support any $CP$ violation in Cabibbo allowed and doubly suppressed channels due to the absence of a second weak amplitude. In singly Cabibbo suppressed transitions one expects $CP$ asymmetries, albeit highly diluted ones of order $\lambda^4 \sim 10^{-3}$ or less [56].

**CP asymmetries involving oscillations**

For final states that are common to $D^0$ and $\bar{D}^0$ decays one can search for $CP$ violation manifesting itself with the help of $D^0\bar{D}^0$ oscillations in qualitative—though certainly not quantitative—analogy to $B_d \to J/\psi K_S^0$. Such common states can be $CP$ eigenstates—like $D^0 \to K^+K^-/\pi^+\pi^-/K^0_S\rho^0$—but do not have to be: two very promising candidates are $D^0 \to K^0_S\pi^+\pi^-$, where one can bring the full Dalitz plot machinery to bear, and $D^0 \to K^+\pi^- \text{vs.} \ D^0 \to K^-\pi^+$, since its Standard Model amplitude is doubly Cabibbo suppressed. Undertaking time-dependent Dalitz plot studies requires a higher initial overhead, yet in the long run this should pay handsome dividends exactly since Dalitz analyses can invoke many internal correlations that in turn serve to control systematic uncertainties.

**2.4.4. Experimental Status and Future Benchmarks**

Time-integrated $CP$ asymmetries have been searched for and sensitivities of order 1% [several %] have been achieved for Cabibbo-allowed and -singly suppressed modes with two [three] body final states [58]. A Dalitz-plot analysis of time-integrated $CP$ asymmetries provides constraints $O(10^{-3})$ [59].

Time-dependent $CP$ asymmetries (i.e., those involving $D^0\bar{D}^0$ oscillations) are still largely terra incognita.

Since the primary goal is to establish the intervention of New Physics, one “merely” needs a sensitivity level above the reach of the Standard Model: “merely” does not mean this can easily be achieved. As far as direct $CP$ violation is concerned, this means asymmetries down to the $10^{-3}$ or $10^{-4}$ level in Cabibbo-allowed channels and down to the 1% level or better in doubly Cabibbo-suppressed modes. In Cabibbo-singly-suppressed decays one wants to reach the $10^{-3}$ range (although CKM dynamics can produce effects of that order, future advances might sharpen the Standard Model predictions). For time-dependent asymmetries in $D^0 \to K^0_S\pi^+\pi^-$, $K^+K^-$, $\pi^+\pi^-$ etc., and in $D^0 \to K^+\pi^-$, one should strive for the $O(10^{-4})$ and $O(10^{-3})$ levels, respectively.

When striving to measure asymmetries below the 1% level, one has to minimize systematic uncertain-
ties. There are at least three powerful weapons in this struggle: i) resolving the time evolution of asymmetries that are controlled by $x_D$ and $y_D$, which requires excellent vertex detectors; ii) Dalitz plot consistency checks; iii) quantum statistics constraints on distributions, $T$-odd moments, etc. [18].

### 2.4.5. Experimental reach of New Physics searches

In this section we briefly summarize the experimental reach of Super $B$ for New Physics sensitive channels in the charm sector. Table IV shows the expected 90% confidence level upper limits that may be obtained on various important rare $D$ decays, including suppressed flavor-changing neutral currents, lepton flavor-violating and lepton number-violating channels, from one month of running at the $\psi(3770)$. It is expected that the results from running at the $T(4S)$ will be systematics limited before reaching this precision.

For studies of $D^0\bar{D}^0$ mixing, running in the $T$ region appears preferable, and, if the true values of the mixing parameters are unobservably small, the upper limits on both $x_D$ and $y_D$ can be driven to below 0.1% in several channels ($D^0 \rightarrow K^+\pi^-$, $K^0\bar{K}^0$, $\pi^0\pi^+\pi^-$, etc.) Therefore, Super $B$ can study charm mixing if $x_D$ and $y_D$ lie within the ranges predicted by the Standard Model, and recently observed. The sensitivity to mixing-induced $CP$ violation effects obviously depends strongly on the size of the mixing parameters. If one or both of $x_D$ and $y_D$ are $O(1\%)$, as indicated by the most recent results, Super $B$ will be able to make stringent tests of New Physics effects in this sector.

The situation for searches of direct $CP$ violation is clearer: the Super $B$ statistics will be sufficient to observe the Standard Model effect of $\sim 3 \times 10^{-3}$ in $D^+ \rightarrow K^0\pi^+$ [17], and other channels can be pursued to a similar level. Within three body modes, uncertainties in the Dalitz model are likely to become the limiting factor. However, model-independent $T$-odd moments can be constructed in multibody channels, and limits in the $10^{-4}$ region appear obtainable.

### 2.5. Summary: Charm Physics at Super $B$

One does not have to be an incorrigible optimist to argue that the best might still be ahead of us in the exploration of the weak decays of charm hadrons. Detailed studies of leptonic and semileptonic charm decays will allow experimental verification of improvements in lattice QCD calculations, down to the required $O(1\%)$ level of precision. This will result in significant improvements in the precision of CKM matrix elements. The possibility to operate with $e^+e^-$ collision energies in the charm threshold region further extends the physics reach and the charm program of the Super Flavour Factory.

### Table IV

Expected 90% confidence level upper limits that may be obtained on various important rare $D$ decays, from 1 month of Super $B$ running at the $\psi(3770)$.

<table>
<thead>
<tr>
<th>Channel</th>
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<tbody>
<tr>
<td>$D^0 \rightarrow e^+e^-$, $D^0 \rightarrow \mu^+\mu^-$</td>
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<td>$2 \times 10^{-8}$</td>
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<tr>
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<td>$3 \times 10^{-8}$</td>
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While no evidence for New Physics has yet been found in charm decays, the searches have only recently entered a domain where one could realistically hope for an effect. New Physics typically induces flavor-changing neutral currents. Those could be considered for an effect. New Physics typically induces flavor-changing neutral currents, lepton flavor-violating and lepton number-violating channels, from one month of running at the $\psi(3770)$. It is expected that the results from running at the $T(4S)$ will be systematics limited before reaching this precision.

For studies of $D^0\bar{D}^0$ mixing, running in the $T$ region appears preferable, and, if the true values of the mixing parameters are unobservably small, the upper limits on both $x_D$ and $y_D$ can be driven to below 0.1% in several channels ($D^0 \rightarrow K^+\pi^-$, $K^0\bar{K}^0$, $\pi^0\pi^+\pi^-$, etc.) Therefore, Super $B$ can study charm mixing if $x_D$ and $y_D$ lie within the ranges predicted by the Standard Model, and recently observed. The sensitivity to mixing-induced $CP$ violation effects obviously depends strongly on the size of the mixing parameters. If one or both of $x_D$ and $y_D$ are $O(1\%)$, as indicated by the most recent results, Super $B$ will be able to make stringent tests of New Physics effects in this sector.

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<td>$1 \times 10^{-8}$</td>
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</table>
A Super Flavour Factory would allow conclusive measurements. SuperB, with data taken at the $T(4S)$ and near threshold, will complete the charm program down to the Standard Model level.

References


[7] The New Snowmass Year is an updating of the convention that multiplying peak luminosity by a “year” containing $10^7$ seconds provides a good measure of actual running time, the effects of accelerator and detector down time, dead time effects and the difference between peak and average luminosity. PEP-II/BaBar experience has shown that a “New Snowmass Year” with $1.5 \times 10^7$ seconds is a more precise estimator of actual performance at a B Factory.


[10] The starting point of this effort was an attempt to leverage the active development effort in support of a high energy linear collider that has been going on for the past two decades. The idea, which has antecedents dating to the mid 1980’s [11], was to achieve the high luminosity by using very low emittance beams with high disruption, and to recapture at least the positron beam and recirculate it, to minimize power consumption. It is, however, a substantial challenge to produce luminosities of the order of $10^{36} \text{cm}^{-2}\text{s}^{-1}$ while having center-of-mass energy spread less than 10 MeV and keeping the power consumption to a tolerable level. This proved to be a difficult problem [12].


The future of charm physics: a discussion

B.D. Yabsley (ed.)
School of Physics, University of Sydney. NSW 2006, AUSTRALIA.

We closed the CHARM 2007 workshop with a lively panel discussion on the future of the field. This document presents a summary of the key points, and a lightly edited transcript of the discussion itself.

1. Introduction

The CHARM 2007 organisers (and advisory committee) deliberately chose to close the workshop not with a summary talk, but with a panel discussion. This document, accordingly, tries to respect the idiosyncrasy of that discussion, rather than seeking to be comprehensive or to be too “on-message”.

The panel discussion followed four talks on future facilities — the relevant speakers were all panel members — and could helpfully be read together with those presentations. The raw material for the discussion was a set of questions from workshop attendees, and some absent committee members.

2. Key points

- Charm physics is intellectually alive as a field, and we plan to keep it that way.

Mixing and CPV

- We now have good evidence for mixing, and may even breach the $5\sigma$ threshold in individual channels, with the $2\text{ ab}^{-1}$ that’s currently within reach. With twice that dataset, the chance will be good. (And CDF will help!)

- There’s no obstacle (in principle) to calculating potential beyond-the-Standard-Model contributions to mixing on the lattice. But current methods won’t cope with the nonlocal operators that we believe dominate the SM contribution.

- Finding and studying CP violation in the charm sector is (we hope) the next big thing. CPV-in-mixing has long been a goal, but is only now within reach, with evidence for mixing established; controlling systematics at the $10^{-3} - 10^{-4}$ level may be challenging. And let’s not forget direct CPV studies ...

QCD, spectroscopy, lattice ...

- The charm sector is a good laboratory for QCD; we’re still learning from it, and continued to do so at this workshop.

- Not every bump is a new state: it’s possible to get complicated structure from a single pole. Beyond the mere “discovery” of “new states”, experimental determination of quantum numbers and decay properties is important to provide enough information for theoretical analysis. Interpretation can be difficult even so.

- Very large data samples, and improvement in quality (e.g. BES-II → BES-III), will help.

- A really good precision test of lattice QCD requires the use of CKM-free quantities, e.g. the ratio of semileptonic to leptonic decay rates. Very large data samples will help here, too.

- Experimentalists and non-lattice theorists should interpret any lattice calculation in an informed critical spirit. Is the calculation unquenched? With $2+1$ sea quarks? How thorough is the systematic error analysis? If the result is to be compared to experiment, there should be a positive answer on all three fronts.

Facilities in the future

- There are plans at BaBar to maintain an analysis effort (including priority-setting) after data-taking stops. Still, it may be tough.

- There should be active experiments in the flavour sector, to complement the LHC.

- There is strong support for a so-called super-B or super-flavour factory in our community, in addition to BES-III. While the B-physics programme is its selling point, it will be highly capable in charm (& charmonium & ISR & ...)”

- The way forward in hadronic $p\pi$ physics is to take a global approach, with a full-capability detector: PANDA. (Earlier-generation experiments were already at their limits.)

- It would be very helpful to form a “shopping list” of measurements we want LHCb to make in the charm sector. It could have real influence. Analysis will be limited by manpower and interest as much as the trigger … and note that the upgrade will likely have a software trigger, to the benefit of charm studies.
3. Dramatis personae

Panel members:

[DA] David Asner  
[DB] Diego Bettoni  
[AE] Aida El-Khadra  
[EG] Eugene Golowich  
[PS] Patrick Spradlin  
[YW] Yifang Wang  
and Bruce Yabsley  
(chair: comments listed in italics)

Contributors from the floor:

[EE] Estia Eichten  
[BM] Brian Meadows  
[KP] Klaus Peters  
[AP] Alexey Petrov  
[JR] Jonas Rademaker  
[AS] Alan Schwartz  
[KS] Kamal Seth  
[??] at least one significant contributor is currently anonymous

4. The discussion

Thank you everyone for joining us for this last session. You know everyone on the panel: they don’t need any further introduction. Let me just show you what we’re going to do. You can see the outline here, divided into

- current issues:
  - mixing and CP violation, and
  - everything else;

- the future of the field:
  - in general, and
  - concerning future facilities.

We’ll try to give about 15 minutes on each of those categories, picking out questions out as we go. But to begin, I’m going to throw the microphone to Gene: we’ve heard four talks from experimentalists, I thought we’d get a few remarks from the other side of the divide; and then we’ll head to questions.

4.1. Overview

[EG] Bruce asked me to be provocative and nasty, but it’s really not in my nature. Actually over lunch we were talking about the future of the field, and I was drifting off, and ended up in a fantasy world where things were done the right way. And in this world the LHC was in fact built and came on the air, and found the Higgs, and found many new events that we couldn’t explain with the Standard Model. And people had realised that in order to interpret these possible signals of new physics, we would also have to have flavour physics studies of rare phenomena, so that we could start to see patterns emerging … and working symbiotically together, the LHC and the flavor sector would get to the root of what was happening, something that would be very difficult if not impossible to do with the LHC alone.

But then I woke up. And I thought about a colloquium I’d given recently, where one of the chief experimentalists there took me into his office and shut the door and said to my face, “Flavor physics is dead!” and apparently he’s not the only one who said it: some pretty important people have said it. And when something like that is said over and over it begins to have a truth of itself.

So, anyone who’s been at this workshop realises that that’s not the case: that not only are we getting wonderful results, there are all sorts of ideas — arguments even — about what to do for the future, about what’s going on, and that’s the sign of an intellectually alive effort. And part of what we’re doing on this panel, I think, is trying to see whether or not we can realise the potential of the field.

I know that in my own shopping-list I tend to divide it between electroweak physics and QCD, and I’d like to see efforts in both halves; I’m interested in both. My number one wish, as far as the electroweak sector is involved, is that we find and are able to study patterns of CP violation in charm physics. As regards the QCD part of things, I really find that charm has turned out to be a terrific laboratory for QCD and its workings, we’re still trying to figure out what’s going on; all the data’s not in yet, and I think we’re going to need every bit of data to understand what’s going on, and that’s been a terrific part of this workshop.

So the completion and growth of those two areas would be what I’d hope most for, but by no means saturates what I want to see. Now Bruce has handed out to us a bunch of questions that we’ve gathered from various folks; so
let me throw it back to Bruce now — and make sure you have a live microphone.

4.2. Mixing: \((x, y)\), the lattice, and CPV

The first question is a reasonably simple one: whether we can firmly establish that mixing occurs — that \(x\) and \(y\) are not equal to zero — with (say) 2\(ab^{-1}\) of data. And going further than that, whether we can determine \((x, y)\) themselves with twice as much data. Anyone want to take that?

[DA] Sure. If the current central values hold up, then 2\(ab^{-1}\) is close to sufficient for 5\(σ\) signals in two or three channels. If the current central values don’t hold up, one might hope that at least 4\(ab^{-1}\) is . . . so yes, that’s certainly the right ballpark.

Brian, you don’t disagree with that? Either of the Brians [Petersen and Meadows], since you’re sitting together? No?

[JR] Just one comment: CDF will contribute to these measurements very soon as well. So there will be some help even if the situation is marginal with 2\(ab^{-1}\): contributions from hadron colliders [will help resolve \(x\) and \(y\)].

Alright, so the next question may or may not be easy because I genuinely don’t know the answer to it — I had a pretty fair idea of the answer of that first one — and that’s the following: Calculations of New Physics contributions to mixing rely on determinations of four-fermion matrix elements: what can we expect the lattice to contribute to that? So that’s a question for Aida.

[AE] So the methods for calculating 4-fermion operators that contribute to \(\bar{B}\bar{B}\) mixing are quite mature and ongoing — there’s a result by the HPQCD collaboration using NRQCD heavy quarks — and in principle there’s no barrier towards repeating these calculations for charm. HPQCD would use the HISQ action, and our collaboration [Fermilab] also has an analysis in progress for \(\bar{B}\bar{B}\) mixing, and we can certainly do this for the charm sector too.

I should add a note of caution however: what we can do with current methods are calculations of matrix elements of local operators. So we cannot say whether the observed \(\bar{D}\bar{D}\) mixing is Standard Model or not, because that involves non-local operators — because you have light particles propagating — and that’s something that we can’t really calculate with current methods. But if one wants to know the expectation values of local operators that would come from beyond the Standard Model: those can be calculated.

OK so let me see if I understand the caveat: What about the extent to which the observed pattern of mixing comes from the interference of SM and NP contributions? Does your caveat about long-distance operators mean that one has a grey area there? Or do they not interfere because they are short and long distance?\(^1\)

[AP] In New Physics, \(ΔM\) or mass-difference comes from heavy intermediate states, and even those intermediate states are not present in the SM: squarks, gluinos, whatever. So regardless of what happens, at the end of the day you get a bunch of matrix elements at the charm scale . . . and there are only 8 of them. So regardless of the New Physics model, you will get those 8 operators. And so the matrix elements of those . . .

[AE] And they can be calculated . . .

[AP] Yes.

It may not all be this easy.

We often say that CP-violation-in-mixing is the real New Physics signal, and this is part of the propaganda, but up until now the focus has essentially been on \(x\) and \(y\). Do the experimental plans that the existing and future collaborations have, reflect the need to measure CPV in mixing? Comments?

[DA] The short answer is of course yes. Every mixing paper in the last five years has had CP-violating limits although they’ve been trivial; and one of the most active efforts now of the HFAG-charm mixing group is to come up with meaningful CP-violating limits which are actually posted on the site, if you choose to go look at it . . .

So yes.

But that’s a nominal yes, isn’t it? I understand that (say) for the \(D^0 \to K^+K^-\) channel we included a CP-violation measurement, but that was essentially trivial, in that it came almost for free with our \(γCP\) measurement . . . aren’t there more difficult CPV studies?

[DA] Well it comes for free with a \(D^*\) tag as well, for other measurements . . .

Well, at an \(e^+e^-\) machine it does.

\(^1\)The discussion moved on from this question, and it wasn’t followed up at the time. For the record, the answer is that in addition to the problem of hadronic uncertainties in the SM predictions, “we do not know the relative phase between the SM contribution and that from any NP model, so that \(x_0\) will lie between the extreme limiting cases of constructive and destructive interference.” [1]
Yes. Well, perhaps the fact that the LHCb mixing studies are currently CP-conserving might indicate that the experimentalists haven’t caught up with this . . . but it’s certainly on the agenda.

Our plans have recently changed focused, to look more at CP violation. What I was presenting was some of the more complete studies, incomplete as they were. And our initial efforts were towards mixing measurements. But there’s been a fundamental shift in what we’ve been studying, towards CP violation, and CP-violation-in-mixing.

The statement is always in the papers but the emphasis isn’t always in the slides, if you see what I mean, or in people’s spoken remarks. So I guess that’s what was behind the question.

Just related to that, to do the CP violation, at least the direct one, you need to get asymmetries also from detector effects and so on, so how well are LHCb and SuperB geared towards handling this? (And I guess you have similar issues in B-physics, right?)

This is actually the next question, or part of the next question, which is about systematics: It’s one thing to do a percent-level measurement of mixing, but it’s quite another to do a $10^{-3}$ or $10^{-4}$ measurement of CPV. That’s a whole lot more demanding, and it’s going to be systematics that are the problem. Has LHCb thought about this?

Yes, but unfortunately we’ve only begun thinking about systematics. We have a lot of ideas on how to evaluate them in the data itself. But that’s a work in progress.

Alan wanted to speak to that.

Actually I just had a comment on [the earlier point]. You have to keep in mind that to see any CP-asymmetry effects, you need a nonzero $x$ and $y$. But it’s only been since March, since March, that we’ve had such evidence in hand [2, 3]. So that, so when you say do experimental plans reflect going after [the CPV phase]; since March yes, but before March, even in Belle when we measured those things, as you said, they were not the headline measurements, because we didn’t know what $x$ and $y$ were.

Well this is true if you stick to CPV-in-mixing. In both of those cases. I was actually wondering why the person who asked the question wanted to know about time-dependent CP asymmetries, I mean of course you have CP-violating asymmetries that are independent of $x$ and $y$ — direct CPV effects — and you know, those are easy to study, although experimental issues are similar; I don’t know, if you look at $D^+$ decaying to a state that contains charged kaons, detector effects for kaons of different charges are significant . . .

But the point is that you can have CP violating signals that don’t have [time-dependence] . . .

I just had something . . . I think the last mixing paper that I read — experimental paper — that I read that didn’t say anything about CP-violation in mixing must have been at least 10 years ago.

4.3. Rare decay measurements

We’re going to hop forward to some stuff on rare decays now: something not unrelated to CPV. This actually cropped up in one of the talks earlier: that there hasn’t been much work at the B-factories on rare decays. Despite our propaganda, and despite the fact that we put it in all of our proposals and statements, and all the rest of it — and I certainly did, for the charm programme at Belle — so the question is what the prospects are for rare charm decays in the rest of the B-factory era, or perhaps for a super-flavor factory. Why aren’t there more rare decay measurements?

Everybody complains about manpower, right?

[Well I’ve seen rare decay measurements done, and then not turned into an ongoing programme, because the judgement was that other work would better advance people’s careers. And that judgement is probably correct.] So, you know, “manpower” . . . there are manpower issues but they’re not always just the sheer number of people.

To be fair: until rare decays can probe SM sensitivities they are just not that interesting. We have a long history of not finding new physics, and if you want to sign up a grad student to work on something where we know he’s not going to find new physics . . . well we can do that, [and there are many ways of doing it, but there’s little motivation for it.] But when we do have sensitivity [at the level of the SM predictions], I suspect there’ll be rather a large line of people who want to perform the measurements.

But I wonder about this argument, because we’ve been busy doing charm physics at B-factories, and I’ve had guys bowl up to me at conferences and say “The beauty sector is the most rich in physics at the moment (which is true), and offers the best prospects for
finding physics beyond the SM at the moment (which is also true), and therefore it’s the most exciting thing, so why would you work on anything else? And in particular, why are you working on charm?” And in the terms in which that question is posed, I don’t know that there is an answer to it. And isn’t [the statement about rare decays] a species of that same argument?

[DA] Well I’d rather address your first part first: Rare decays are much more interesting in charm than they are in B, because rare decays actually include CP violation and mixing, which are also rare in charm but not in B; so from the get-go rare charm processes are a broader field. And I’m not actually certain how what I said translated to the analogy [with charm-versus-beauty].

[KS] As long as you have this up, I have a question for David Asner. When you were comparing what could be done at LHCb and the super-B-factories, most of the time the things were sort-of comparable, but there were a couple of instances in which you said you really need SuperB for that . . . Now I have to ask you a question whose flavour is completely different from [that of] your talk. Does one, for those rare cases for which SuperB is essential, spend $500 million?

[DA] That’s a good question: it depends what part of the programme of SuperB you think is most important. If you are a tau physics enthusiast, there is no comparison between the tau physics programme at SuperB and the tau physics programme at LHCb. [And for] anything that requires a primary vertex (including, probably, charm semileptonic decays) the performance will simply be better at SuperB. To be fair, I’ve probably represented LHCb in a rather favourable light. After all, it’s going to be an existing experiment.

Will the world end if we don’t have the charm physics programme from SuperB? No. We’re not trying to sell SuperB on its charm physics programme. But if you can motivate SuperB based on the way that the B-physics interacts with the LHC programme, then you get the charm and the tau and the charmonium and the ISR programmes etc. for free, assuming that you can assemble the workforce.

[KS] But if there’s one thing I have learned, it’s that when someone is blessed with huge luminosities, they have the advantage of making cuts and throwing away so much data that they can [do almost anything]. I always used to say that the B-factories, if they wanted to get at whatever I might be doing, they can cut everything out and still beat us on total statistical and other kinds of errors.

And that is what makes me wonder about comparison between LHCb and the SuperB. LHCb will have so much production of whatever-you-want, that you wonder whether they have the ability to throw away a lot of it, and become, you know . . .

[DA] Well they do throw away a lot of events: it’s called the trigger, right?

[KS] Yes, the trigger, and whatever else you want . . . and still reach the statistical precision [of SuperB].

There are also issues of systematics, remember: things that a hadron experiment by its nature doesn’t do as well. Alan?

[AS] There are many final states that would be very challenging at a hadron collider: π⁰ and η and charged ρ, you know, those will be very difficult at a hadron machine. You will not measure B⁰ → π⁰π⁰ at LHCb — I don’t think — and I don’t think you’ll see B⁰ → ρ⁺ρ⁻ and many of the η and η' decays will be very hard. I mean, obviously they can be simulated, but there are many many final states where I don’t know how it’s going to happen in a hadron environment.

[BM] Just one thing to add to what Alan was saying: One thing you probably want to do in looking for CP-violation in mixing is, you want to see if there’s a difference in x and y for D⁰ and D⁻. And then you want to see if they’re the same in different decay modes, in all the different decay modes. And I think that at SuperB that will be much easier to do.

4.4. Do we claim too many new states?

I’m going to move on, to something that’s going to induce a different kind of argument, but this is also about perceptions of what is exciting and what’s not. One of our questioners wanted to know whether we are claiming too many new states, either quarkonium-like or in general. And what about dynamics? Particularly near threshold, since a lot of these states do seem to be near threshold, . . . An awful lot of new states are being claimed, some of them on our experiments’ websites. [I think Yifang has his hand up to answer this.]

[YW] I think, in fact, we do see a lot of resonances, or particles, or structures — whatever you want to call them — near the threshold, and I think there’s a very simple experimental reason, in that near the threshold the background is much
lower, so wide resonances are much easier to see . . .

**OK, but let me pick you up on the words that you used there:** yes it’s easy to see “it”, and you said “resonances”, “new particles”, or “structures”, but just because you see nontrivial structure, it doesn’t mean that there’s a new pole sitting [on the complex plane].

[YW] I agree: I said all these possibilities, because now, particularly at BES [where there are so many things that we see], we have a lot of arguments about whether new particles or structures or enhancements (or whatever) exist, and I think shows that we need more statistics, and we need a much better detector.

No, but hang on: don’t we also need better interpretation?

[YW] Of course, later on: I think the important thing, right now, if we see a signature of (say) 5σ, 4σ sometimes, 6σ, then many people believe it, some people don’t believe it, and for both experimentalists and theorists, [the confusion] makes it hard to work together. Now if you have a much better detector, and significantly improved statistics, if you see a signal at a level of say 20σ, or 50σ, then there’s no question of whether these things exist or not, and it makes it easier for both experimentalists and theorists.

Perhaps it will clarify the point [if we consider another example. There have been many claims of new states at Belle, and one of the most interesting pieces of recent work concerns the extra enhancement below the J/ψ, in the e+e− → π+π−ψ cross-section. Yes? I think everyone has seen this. Now in the paper [4] we were very careful not to say “this is a new state”. But in arXiv postings and on the web, it is being given a name, a “Y” and a number, with it being implied or stated that it’s a new state.] And that reflects the fact that for (I think) 80% of the people in the community saying that you see a significant enhancement — and it is an enhancement, and it is significant, OK, that part’s not in question: a better detector won’t change that — but the first thing we always lean to is, “it’s a new state”. And that’s been the first interpretation rather than the last one, for lots of things from Belle, from BaBar, from BES, and what this question is asking is, **Is that really a good thing — and if it’s not, how do we get past it?** And Estia wants to speak to that.

[EE] I mean, when you have wide structures, it’s always rather difficult to figure out what’s going on, and you can certainly show examples where a single pole will give rise to complicated-looking structure.

So there’s two separate questions. For the theorists, they have to figure out how to interpret the structures that are seen in experiment, in terms of what’s underlying them. I don’t know how the experiments can do that, but what experiment can certainly do is describe the properties of those structures — make sure you understand its quantum numbers, if you interpret it as a state, its decay properties etc.: these things will certainly make it a lot easier for theorists to try to disentangle wide structures into what actually is underlying them [and whether] there are additional states.

Every bump is not a new state. You can show examples when it is clearly not . . .

And was it you that showed that point about the π+π−ψ and the π+π−ψ′ [cross sections, in ISR data from the B-factories], and that maybe . . .

[EE] Yes, it may very well be that there’s one state with different decay modes at different energies, whose branching fraction varies as a function of energy. Now, I can’t calculate that one, because I don’t know what it is — if it’s a hybrid of some sort, then I don’t know how to do the calculation of how much it’ll decay into those two modes — but that’s (in a sense) a theorist’s or a phenomenologist’s job to try to figure that all out, but experimentalists can certainly try to disentangle any other decay modes that are in the same region, for example, besides the ψ or the ψ′, maybe there’s another mode in between that can fill us in on what’s going on. But that’s all an experimentalist should [be expected to do]. The interpretation is difficult.

[EG] This goes back as long as I can remember in hadron physics: the σ(600) is something that people call different things . . . I don’t know what to call it, but if I do know how to use it, you can call it whatever you want.

[AP] There’s actually a little counter-example to that: You all know the story of the pentaquark, which was very well predicted by several groups, and very well observed by several different experiments, and yet it’s gone. And so it’s natural that you see so many states, well not states, but structures. Time passes, and you sort out which of those are real states, and which of those are just threshold things . . .

[YW] I think I haven’t finished what I want to say, so let me finish, and then we can go forward. I think it’s very important to have more data. If you have 4σ, 5σ, how do you disentangle whether it’s a real particle, structure, a fake, or so on? But once you have significantly improved
statistics, you can certainly establish whether this is a real structure, and you have the possibility to understand its decay properties, dynamics and so on, and then you can have a much better understanding.

Put it this way: up to now, we have studied a lot of electroweak physics, but we understand QCD much less well. And we know that we should see a lot more resonant states, hybrids or whatever, glueballs, beyond the conventional $q\bar{q}$ or baryon states: certainly more than what we see.

So with significantly improved statistics ... put it this way, we now see many $4\sigma$ signals. Over time, BES-III will increase statistics by a factor of a hundred, and on top of that the detector is much better — corresponding to an improvement, for the same statistics, by a factor of two or four — so this is an improvement of almost a thousand in total. So you can imagine: with one thousand more better signals, I think it will be revolutionary, and for many unresolved problems or questions, [it will give us] a much better understanding. So I think it’s important to have more data. And by that time, once we have firmly established particles and their properties, theorists can work together, and figure out all these details. And [regarding] QCD: by that time, we’ll have a much better understanding.

4.5. Comparing the lattice to experiment

We’re going to move on. This is a question I think you’ll even like: Suppose experimental and lattice values for decay constants disagree by $2\sigma$. Is that a discrepancy, and should we be worried about it? And if $2\sigma$ is not, then what about $3\sigma$ or $4\sigma$? [To AE:] Do you have a a view from your side?

[AE] I would say that at present, the situation is not yet completely clear, whether or not there is indeed a discrepancy. I’d also like to point out that the experimental values for the decay constants have to assume a value for $V_{cs}$, which is not very well determined at this point. It’s hard to imagine right now that $V_{cs}$ could change too much, but you can have at present a few small effects which would bring experiment and theory back together. I think once the error bars decrease more, then we would be more worried.

To really do a precision test of lattice QCD without any assumptions about CKM angles, we need to compare the CKM-free quantities like the ratio of the semileptonic decay rate to the leptonic decay rate, where the CKM angle does not contribute. So [at the level of precision reached by HPQCD] I don’t think that just looking at the decay constant itself is really [enough]: you want to look at the ratio of semileptonic decays to leptonic decays; you want to look at the shapes of the form factors; you want to see if you extract $V_{cs}$ from the semileptonic, does that agree with the result from the leptonic case, to see if there really is a discrepancy. So I think we need more information than just the decay constant. And I think that’s part of the programme.

And that’s a point where Yifang’s argument about overwhelming statistics is important, because with overwhelming statistics you have the luxury of doing those assumption-free tests that you’re talking about.

[AE] That’s certainly true. I mean, two sigma statistical fluctuations are not unheard of; and while Peter Lepage might vehemently disagree with me for saying this, it’s not absolutely impossible that (say) the HPQCD result might have slightly larger errors than they [claim], in which case the discrepancy is not as large as it currently is. So I think it’s certainly important to monitor and keep in mind, but ah . . .

Does anyone else want to speak to that? It’s a fairly complete answer, but . . .

[AP] I’d like to ask a question about the errors that lattice calculations are assigning to their numbers. I mean, I used to see nice lattice predictions with very small error bars that were completely overtaken by new lattice predictions, and people would say that they’re inconsistent, but “That’s OK, because the previous error bar did not include quenching errors.” Now, how can one claim an error bar if there is an error bar sitting on top of that, which you don’t even know what it is? Now I understand that with the new techniques this [is no longer the case]. But when you now do things at the level of a percent, then none of these errors, they’re not independent, and you have to take into account correlations between them.

[AE] And we do. And we do. We have learned a lot in the last ten years on this. I want to make this point very clearly, and I think it requires non-lattice theorists and experimentalists maybe to have some knowledge, but if you read a paper and it says we are doing a quenched calculation

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2The reference is to the disagreement in $f_{D_s}$, between the new lattice calculation ($241 \pm 3$) MeV [5] and the CLEO-c average ($275 \pm 10 \pm 5$) MeV presented in Steven Blusk’s talk at the workshop.
and we get this result with this error, just, don’t even . . . they did an error analysis . . .

But it’s quenched.

[AE] But it’s quenched. And even if they have some estimate for the quenching error, but if it’s not anywhere close to 20% or so, you just shouldn’t believe it. The same thing with calculations that have two flavours of sea quarks, that’s not enough: there are three in the real world. So, you know, I don’t think it makes a whole lot of sense to compare quenched lattice calculations to experiment. I also don’t think it makes a lot of sense to compare lattice calculations based on two flavours of light sea quarks to experiment. [These comparisons are not interesting to people outside of lattice field theory.] And that’s the reason I only showed results based on 2 + 1 sea quarks.

And the second point to your question, it’s a valid point: we have a long history as a community in lattice field theory, certainly [more than] 10 years ago people were not worried about systematic errors and as we got more sophisticated we became more worried. And not everyone in the field understands that they need to be very careful with the systematic error analysis. But I think that people who are claiming to do serious calculations need to be able to show you plots: Do you understand the light quark mass dependence of your result? Have you looked at the lattice spacing dependence? I mean you need to have a very sophisticated systematic error analysis, and if you don’t, you shouldn’t believe the result. [Certainly not at the few percent level.]

And it’s always good to have more than one person, because these analyses are complicated. I mean more than one group — one person never does these analyses by themselves — but more than one group do these calculations with different methodologies and see that you have that consistency. So, you know, we have the HPQCD result, we have our result, which is a factor of two less precise than their result, but within our error, the two results are certainly very consistent.

[DA] Can I jump in real quick?

Quickly . . .

[DA] Speaking to the experimental determination of the CKM-independent quantities, at best I can tell from the two reports I mentioned earlier, it looks like the measurement is at about the 10% level. Part of the BES-III programme will get that to, you know, a few percent, and SuperB will be required to get 7%, but these are part of . . .

We might come back to that later.

4.6. BaBar data after datataking ends

Before Gene goes, I just thought we’d ask the question you see here: whether there will be an analysis effort on BaBar beyond 2009. And what the plan is for doing such an obviously difficult thing? I mean, we have some BaBarians or ex-BaBarians here in the room: anyone want to speak to it?

[??] As far as I understand, there will be three years of . . . at least the infrastructure for doing analysis for three years after datataking. What does “infrastructure” mean?

[??] All the computing, so you can process all the data, do the simulations, and put new decay modes in.

Right, I mean, if you want to do ten decay modes, will you have a supply of slaves to do it, or will you have to do it all yourself?

[??] Clearly, there is a very limited number of new graduate students coming on, already now. So, . . . they will be graduating in the three years . . .

So what’s being done to prioritise measurements?

[??] We have a core list of results which were scheduled to come out . . . somewhere next year or the year after that. And they are being prioritised within B-physics and charm physics, so for instance, the mixing analysis will be updated with the full dataset. So there is a list of people, with their names assigned to core analyses . . .

I don’t know, does someone who’s lived through this part of the life-cycle of a big experiment want to speak to this question?

[DB] In addition to that, there is also serious consideration being given to conserving the data for future use well beyond 2011, because it is recognised . . .

This goes to the second question posted here actually, what to do when the data becomes a heritage matter.

[DB] I’m not aware that there’s a solution to that yet, but certainly that question is being considered. My question is: when it’s several years out from the running of the experiment and you’ve got one or two dedicated groups at a slow burn working on this data, how [do] you maintain quality? How do you stop people with agendas and nothing better to do, and a lot of BaBar data . . .
You don’t avoid that.

That just happens?

It’s natural. Things die, and there’s a rate of dying, and then there is that not that close supervision, and the quality suffers.

This might open a can of worms, but are there any plans to make the data publicly available on the internet, just like astronomy?

But that then multiplies the issue that I was just raising, doesn’t it. I mean, if some high-school student can run a job can you imagine the number of new states we’re going to be seeing . . . [laughter]

But what’s wrong with that? It’s not going to be signed by BaBar . . .

. . . but the answer is yes.

Gene leaves

4.7. Fermilab: charm, and antiprotons

We’ve perforce been skipping around here, so I’ll just continue on this slide for a moment, [can we hear] what the plans are at Fermilab for charm studies — at D0 and CDF?

I already knew the answer to that question, but what about non-mixing topics?

I have no good overview, except for charm production . . . [from elsewhere: “and the rare decays”].

. . . and the rare decays. I guess the question is whether there will be an update or successor to that analysis, or whether that’s the last thing we’ll see. [to AP: “You have a lot of comments!”]

No, I just happen to be at an institution that participates in CDF, and which actually does rare-decay charm analysis. So, you know, the first paper from CDF-II was on $D^0 \rightarrow \mu^+\mu^-$ [6] . . . and right now they’re doing an update of that.

There was a second question posted here — it’s outside my awareness — about what the support in the community is for the antiproton source proposal at Fermilab. Does anyone want to comment on that?

I can say my personal opinion — I cannot speak for the community, of course — because I have been following a little the proposal for the antiproton source, to make a dedicated experiment at the pbar accumulator — my personal impression is that the next step, the next generation of hadronic $p_T$ physics will have to be a global approach like PANDA, with a detector like PANDA. A detector like we had in E835 has more-or-less said whatever it could say, and I don’t know whether Kam would agree, but we were discussing the other day: already at the end we were sort of at the limit of what could be done in this kind of experiment.

So I don’t know what was the point of the question: my impression is that there is very little one can do in the limited space which is available there, unless one can dig a bigger hole, and really do a general-purpose detector.

As a former worker in the field: I can tell you that there is just not enough constituency and it doesn’t rank high enough in the priority, so at least my feeling is that it is very sad that the only place which has currently the ability to produce antiprotons is not planning [anything for the future]. But that’s the way life goes: GSI is coming, [the current facility has] outlived its capabilities, and I don’t see much future for that at Fermilab. That is my personal opinion, but I make that opinion after having heard the arguments and the things that have been proposed.

OK, I’ll take that as an answer. I’m going to skip to the second part ...

About SuperB?

Yep. [i.e. What is the interest and the support of our community for the Super-B project?]

OK I’ll start with a question to Alan — or maybe to you, Bruce — as to how many names are on the Belle letter of intent for SuperBelle?

At least half [of the collaboration].

Because SuperB has three-hundred-and-twenty signatories, and there’s some overlap between the two, but it’s not overwhelming [i.e. the overlap is not the dominant part], but we’re talking about hundreds, five or six hundred people in our field, theorists and experimentalists, who are excited to do work, who have produced these documents . . . so I think there is considerable support for a super-B-factory.
I’d like to say that although we are not part of the super-B [project, we strongly support it] . . . we think it’s very good for the field, and very good for the community. Without super-B . . . at least for us, we at BES, five years from now, the charm field would only be at BES, and um, we don’t want that.

4.8. What to measure at LHCb

There was a pessimist’s question on exactly this point, which you saw, which I’m not going to put up, because we’re running out of time. [Instead I want to consider] a couple of LHC-related questions — essentially LHCb — the first one was answered by Patrick’s talk; so the second one is whether we should be making a shopping-list of the measurements that we want LHCb to make — the people in this room, this community — how much [would that help?]

[PS] I think that that’s an excellent idea. People complain about the manpower issue at the B-factories: at LHCb it’s even worse right now. So having a list of things to look at — and a prioritised list — would be helpful for our planning. As scientists come off of the building phase of the experiment they’re going to be looking for analyses. And if we have an enumerated list of charm physics [topics] we might be able to woo them into this field.

OK, so a specific question, as to whether there’s a charmed baryon programme at LHCb.


Will there be?

[PS] I don’t know. It would take more than just me to be doing it.

OK, but realistically: would the people in this room saying “We want such-and-such to be measured” make a difference to that, or is it just the people in the collaboration?

[PS] I think it’s more getting individual scientists and groups in the collaboration interested in that work.

[JF] I just wanted to second that. I think that at the moment, suggesting exciting analyses might actually still have an influence, for the very reason that many people who have worked very hard on building the detector are moving into physics at this moment, and we can have influence over [the choices that they make].

Maybe we should look into the list idea.

4.9. On being a successful sideline:

4.9.1. ISR work at a “super-B-factory”?

Going to another sideline point about the future, this is one for SuperB: the ISR programme at the B-factories has become quite an important part of the work lately, and through accumulated accidents that kind of dropped out of this meeting. [But there’s good physics being done using ISR at present.] So, what ISR programme is foreseen for the super-B or super-flavor factories?

[DA] There are a couple of pages in the SuperB CDR precisely on this physics, and I can tell you that at every SuperB meeting that is about the physics (not the accelerator) there are talks on this subject. So it’s on the radar. The data will get read out: it will be there on tape. But it will require people who are interested in this physics from the current experiments showing up and doing that physics . . . I don’t think the ability to do this physics, as interesting as it is, will drive whether or not a super-B machine gets built . . .

No, I think that’s right . . . Because it took us a long time to warm up to doing ISR physics. It was in the BaBar book — we didn’t have a book, but it was in the BaBar book — but it even took them a while to warm up to it, so I guess that’s what’s behind the question.

4.9.2. Charm studies at LHCb?

Some general questions about the future now: This is actually a question that I care about as well, as to how easy it’s going to be doing charm studies parasitically at facilities that aren’t dedicated to it. This is really an LHCb question: we managed it fine at Belle and BaBar, but Belle and BaBar run with open triggers. Now . . .

[PS] Yes. We are in some sense competing for bandwidth. And we may have to make a case for every exclusive charm channel that we put in, and certainly it needs to be optimised so that it doesn’t eat into the total bandwidth too much. But I think that anything that someone is genuinely interested in doing can be put into the trigger at this stage. It’s more of a manpower issue, even at this stage. And if I may make a comment about the poor neglected LHCb upgrade — I’m sorry I didn’t mention anything about that — the trigger is going to have to be completely overhauled for an LHCb upgrade in order to accommodate the increased luminosity, and although it’s by no means finalised, the prevailing idea for the trigger there is to go with a full software trigger. And once you’ve done
that, it opens up the whole field: if you can efficiently create a charm trigger, you don’t have these $p_T$ cut limitations that are applied at the hardware level.

So, yes. It’s a matter of manpower and interest and influence, I think.

Anyone else on that one? No?

4.10. Charm & the heavy ion community

OK this is a question that I wasn’t expecting, but it’s a potentially useful one, and it’s whether a closer collaboration with the heavy ion community would be beneficial, either for charm or quarkonium physics. Anyone with experience of heavy ion work? Or friends in heavy ion work?

[AP] There is already close collaboration. First of all, they are already going to share an accelerator: the LHC. And you know, those guys need fragmentation functions, things like that: they get it from us.

[KS] Already at Brookhaven, people are beginning to do spectroscopy, which was not their [original] goal. So they are of necessity coming into the field, and they often ask questions: it is not quite a collaboration yet, but yes, each is beginning to get interested in the other. I think it’s a developing thing.

Right, I assume that’s the thought behind the question: that the communities are somewhat disjoint.

[KS] At RHIC it’s beginning. At LHC it remains to be seen, because it is still in the future.

[AP] There are already a couple of groups: I mean, at our university, Michigan State, the groups are both . . . there are groups on both CDF and D0, CLEO, and STAR and PHENIX and so people are . . . maybe it needs to be formalised.

[KP] Also at GSI with the FAIR project we have two distinct communities . . . [but] it’s exactly ten metres from my office to [that of] my colleague who leads the heavy ion community there, and our intention is to place an application for a virtual institute . . . to manifest [our cooperation] a little more: to really concentrate on charm and the various issues with the nucleus . . . things are moving together to some extent.

4.11. On charmed baryon measurements

OK. We are . . . running out of time. I thought I’d just throw one slightly odd question in. There was a comment along the way about the measurements that are being made in baryons, and that there appear to be some kind of imbalance in it: that you can have some of these quite sophisticated studies that measure the spin and parity of some of the baryons, or search for new states or whatever, but there are quite fundamental things that we know poorly, particularly about the $\Omega_c$. And some of those measurements aren’t being made. Is there a way around that? Or do we need a super-motivated group to just do the work at the B-factories? I mean, CLEO-c was going to run at $\Lambda_c$ threshold at one point, and famously . . .

[DA] SuperB — the INFN proposal — has plans to spend a month at the $\Lambda_c$ threshold. But it’s at $\Lambda_c$ threshold.

Right: they talk about it, but CLEO-c did as well. Is it going to be the first thing that gets dropped, in the same way that . . .

[DA] I think the answer is, “of course”.

Right. OK. That’s not necessarily a criticism . . .

[DA] The entire 4 GeV programme would probably be the first thing to get dropped if the 10 GeV luminosity wasn’t up to snuff.

[KP] To give you one example, our group started the idea to measure . . . the $D_{s1}(2536)$ which we have seen this morning [7]; and it really has reached excellent resolution . . . But to give you an idea of the timescales, this PhD student has spent three years on it, to get really all the systematics nailed down to that point where you can reach these 100 keV systematic errors . . . and that’s exactly what kills a lot of these mass measurements. If there is not really an important issue, you just drop it because these systematic studies . . . there are so many other things you could do, which are more geared to the mainstream.

Less mainstream and more difficult.

[KP] Yes.

4.12. Thanks

OK look, I . . . don’t want to keep you here any longer, and I think we’ve covered a fair amount of territory on a range of things. So thanks to the people who served on the panel, and to all of the speakers I guess. And um, I would certainly like to thank the committee, for what’s been a good workshop.
References


