SLAC – PUB – 4514
December 1987
(T/E)

THE PHYSICS OF HEAVY FLAVORS*

FREDERICK J. GILMAN
Stanford Linear Accelerator Center,
Stanford University, Stanford, California 94305

ABSTRACT

We review the physics of heavy quark flavors, including weak decays, onium, tau leptons, mixing, the Kobayashi–Maskawa matrix, and CP violation in B decay.

INTRODUCTION

The study of heavy flavors is a possible path to uncovering physics beyond the standard model. Such is the case especially with respect to the study of rare decays and of CP violation. Within the standard model, the physics of heavy flavors entails the study of the strong interactions or of the electroweak interactions in the presence of the strong interactions. With heavy quarks, the situation is often both cleaner and theoretically simpler to analyze than for systems involving only light quarks. As a result we can sometimes pinpoint the underlying dynamical mechanisms, as well as extract the values of key parameters.

This is a field that is reaching maturity. Many of the questions that remain open are quantitative rather than qualitative ones. To complement this kind of question, there are high statistics data samples available from electron-positron annihilation and, more recently, from the use of vertex detectors in fixed target experiments at hadron machines. For the tau lepton and for states containing heavy quark flavors, this has meant thousands of events in major decay channels. Correspondingly, we are beginning to probe some of the rarer decays, or to establish significant limits thereon. What follows is a rather quick, personal view of the status of the physics of heavy flavors.

---

*Work supported by the Department of Energy, contract DE–AC03–76SF00515.

Invited talk presented at the Eighth International Conference on High Energy Physics at Vanderbilt, Nashville, Tennessee, October 8–10, 1987

WEAK DECAYS OF HEAVY QUARKS

We have a solid general framework within which to calculate the weak decays of heavy quarks. Starting with the electroweak interactions and their gauge group, $SU(2) \times U(1)$, we add the corrections due to the strong interactions through the use of the renormalization group equations for the coefficients of the operators, with anomalous dimensions computed from QCD.[1]

These calculations are carried out at the quark level. A first stage in their application to actual hadrons is to consider weak decays inclusively and to simply neglect any other constituent of the decaying hadron aside from the heavy quark. In such a spectator model, as it is called, one directly carries over the quark level calculation as the hadron level result; spectator quarks and gluons are assumed to arrange themselves into final state particles, together with the quarks or leptons coming from the heavy quark, at no cost or benefit in the overall rate. This simple spectator model in fact gives a qualitative, if not semiquantitative, account of the known data.

It is clear, however, that there are corrections to this picture since the lifetimes of different species of charmed particles, which are all the same in the spectator model, differ by a factor of two or so. The data on charmed particle lifetimes has recently undergone a qualitative improvement. The Fermilab photoproduction experiment E691, using silicon strip vertex detection, provides clean data with high statistics; the result is precise lifetime measurements for different charmed species.[2] Figure 1 shows the data that lead to the measured D_s lifetime from observation of decay distributions for the $D_s \to \pi\phi$ and $D_s \to K\bar{K}^*$ modes. The results on charm lifetimes from E691 are shown in Table 1.

So, to do better than the factor of two level of agreement, we need to go beyond the spectator model. Final state interactions, annihilation diagrams, interference between different amplitudes, and color (mis)matching have all entered the discussion. All have roles to play. The general situation is reviewed elsewhere;[3] we discuss here only a few recent developments.

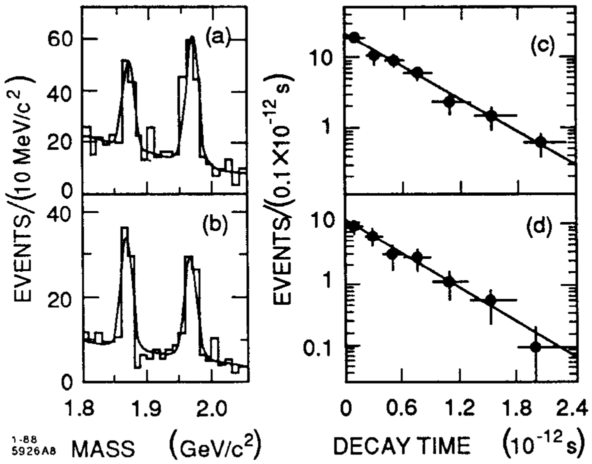

Fig. 1. Data for measurement of the lifetime of the D_s from (a) the $\phi\pi$ and (b) the $K\bar{K}^*$ decay modes, from Ref. 2.

Table 1. Charmed Particle Lifetime Measurements from Ref. 2.

Mode	Signal	Background	Lifetime (psec)
D^+	2992	1354	$1.090 \pm 0.030 \pm 0.025$
D^0	4212	975	$0.422 \pm 0.008 \pm 0.010$
D_s^+	228	75	$0.47 \pm 0.04 \pm 0.02$
Λ_c	93	85	$0.22 \pm 0.03 \pm 0.02$

- Final State Interactions

Final state interactions must be present. The question is their importance. Direct evidence of their magnitude is provided by the Mark III data on $D \to \bar{K}\pi$ modes: $D^0 \to K^-\pi^+$, $D^0 \to \bar{K}^0\pi^0$, and $D^+ \to \bar{K}^0\pi^+$. The $K\pi$ final state can occur with isospin 1/2 and 3/2, and so there is one triangle relation between the three amplitudes. Without final state interactions the amplitudes should be relatively real. The experimentally measured branching fractions demand that the isospin 1/2 and isospin 3/2 amplitudes have a large phase between them.[4] Similar comments hold for the $D \to \bar{K}^*\pi$ channel, but, on the contrary, $D \to \bar{K}\rho$ shows only a small phase difference. Thus, at least in some cases, final state interactions are very important.

We note in passing that the absolute D meson nonleptonic branching ratios from Mark III have been revised downward[4] by 19% to 25%, thereby removing the "charm deficit" in B decays.[5] The newest Mark III branching ratio for $D^0 \to K^-\pi^+$ of $4.2 \pm 0.4 \pm 0.4\%$ has also been confirmed in an independent manner by the HRS, which obtains[6] $4.0 \pm 0.6\ ^{+0.7}_{-0.6}\%$.

- Annihilation Diagrams

Quark-antiquark annihilation must occur in the decay of a meson if only leptons are present in the final state. In particular, the decay $D_s^+ \to \tau^+\nu_\tau$ is expected with a roughly 2% branching ratio. Similarly, the new bound[7]

$$B(D^+ \to \mu^+\nu_\mu) < 6.2 \times 10^{-4}$$

from the Mark III may be used to assert that $f_D < 290$ MeV, where the probability of the c and $\bar{d}$ quarks in the D^+ to annihilate has been summarized in the pseudoscalar decay constant, f_D.

The question which remains outstanding is again a quantitative one of the magnitude quark-antiquark annihilation or of W exchange in nonleptonic decays. The observation[8] of the decay $D^0 \to \phi\bar{K}^0$ at the 1% level in branching ratio looks, on the face of it, to be evidence for W exchange (One needs to get rid of the $\bar{u}$ quark in the initial state since it does not appear in the final hadrons.) It can be argued, however, that such a final state can be generated without W exchange by final state interactions.[9] As this one decay mode is not decisive, we seek more evidence. This comes from observing that if W exchange is important in D^0 decay, annihilation should play at least as important a role in D_s^+ decay, and final state interactions should be different there as well. However, the results of recent experiments[2] indicate that the lifetime of the D_s^+ appears to be at least as long as the D^0 (see Table 1), and the search for the specific mode $D_s \to \rho\pi$ has turned up only upper limits. The best of these limits is derived from the data shown in Figure 2, from which

$$B(D_s^+ \to \rho^0\pi^+)/B(D_s^+ \to \phi\pi^+) < 0.08$$

is established. As of now, the conclusion is that while annihilation and W exchange are surely present, they are not very important.

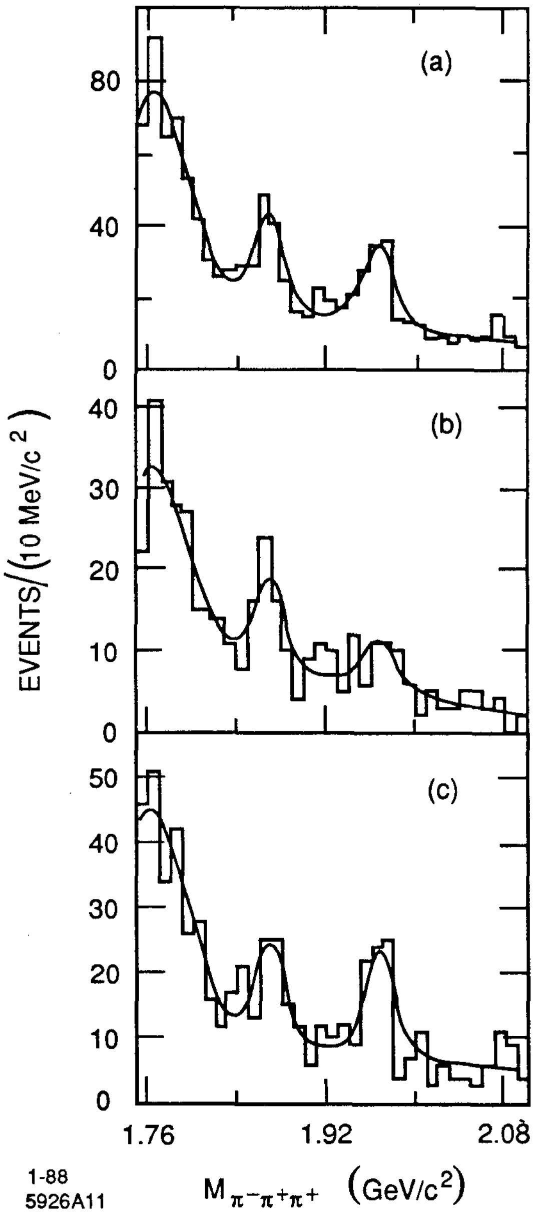

Fig. 2. Data showing peaks for D^+ and D_s^+ decay into (a) 3π, (b) $\rho\pi$, and (c) non-$\rho\pi$ modes, yielding a limit on $B(D_s \to \rho\pi)$, from Ref. 2.

That, however, leaves us with another question: Where have all the strange quarks gone in D_s decay? For if the $\bar{s}$ quark in a D_s^+ does not annihilate, it must appear in the final state together with an s quark from (Cabibbo-allowed) charm decay. However, the modes that satisfy this criterion, like $D_s \to \phi\pi$, $D_s \to K\bar{K}$, and $D_s \to K\bar{K}^*$, do not appear to have branching ratios that would allow us to account for the majority of D_s decays. An answer to the question has now come from the Mark II and Mark III collaborations, which observe[10,11] very substantial modes involving eta's and eta prime's (which contain $s\bar{s}$ valence quarks). Figure 3 shows the evidence[11] from the Mark III for $D_s \rightarrow \eta\pi$, while Figures 4 and 5 show the Mark II invariant mass plots indicating[10] signals for both $D_s \rightarrow \eta\pi$ and $D_s \rightarrow \eta'\pi$, respectively. These results correspond to[11] $B(D_s \rightarrow \eta\pi)/B(D_s \rightarrow \phi\pi) \approx 2.5$ and to[10] $B(D_s \rightarrow \eta'\pi)/B(D_s \rightarrow \eta\pi) \geq 1$ Modes involving eta's and eta prime's are large; this problem may well be solved.

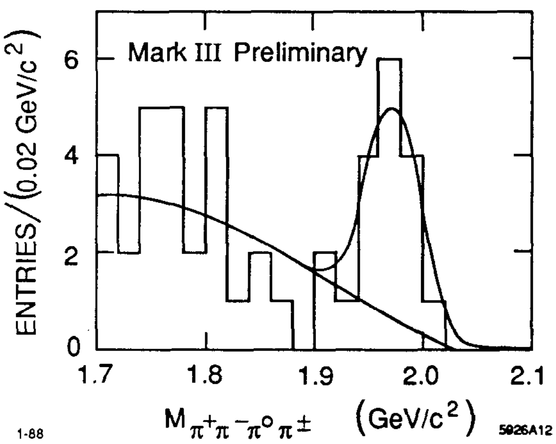

Fig. 3. Signal for $D_s^+ \rightarrow \eta\pi^+$ from Ref. 11.

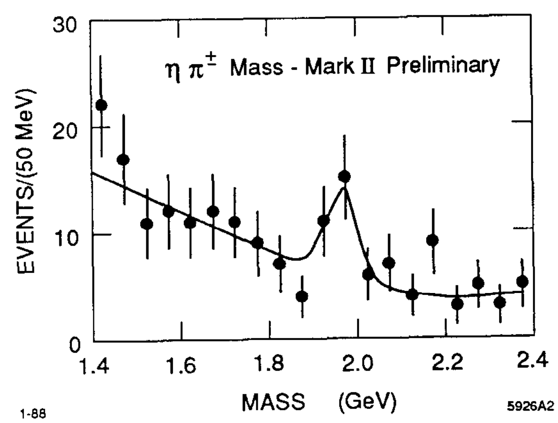

Fig. 4. Signal for $D_s^+ \rightarrow \eta\pi^+$ from Ref. 10.

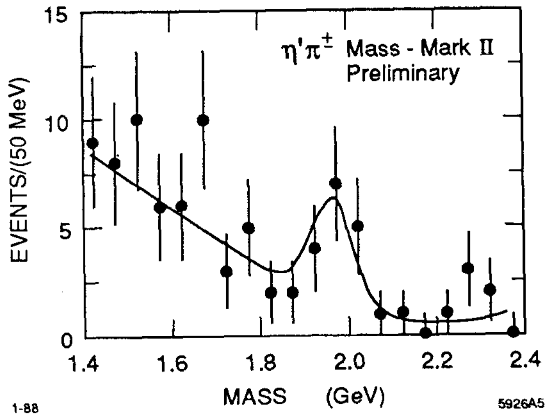

Fig. 5. Signal for $D_s^+ \to \eta'\pi^+$ from Ref. 10.

ONIUM

Bound states of heavy quarks and heavy antiquarks provide us with the showcase of our understanding of strong interactions spectroscopy. The same flavor-independent potential can fit the spectra of both charmonium and bottomonium.[11] In this area also we are asking detailed, quantitative questions, and have much beautiful data to supply us with answers.

An example is provided by the three χ_b' (*i.e.*, $2^3P_{0,1,2}$) states which were recently clearly separated by the CUSB collaboration,[12] as shown in Figure 6. Now that we have both χ_b and χ_b' states and their mass splittings, it becomes interesting to ask if we understand this theoretically.

The splitting of these states, which is due to spin-orbit and tensor terms in the nonrelativistic potential, can be expressed in terms of one absolute mass difference and one ratio,

$$R_\chi = \frac{M(^3P_2) - M(^3P_1)}{M(^3P_1) - M(^3P_0)} = \frac{2a - \frac{12}{5}b}{a + 6b} \quad , \tag{1}$$

where a and b are the matrix elements of the spin-orbit and tensor terms, respectively. In a picture where one thinks of Lorentz vector and scalar exchanges as giving rise to the effective potential between the heavy quark and antiquark, an expansion in powers of v^2/c^2 gives the spin-independent potential $v(r) + s(r)$, plus various spin-dependent pieces which yield[13]

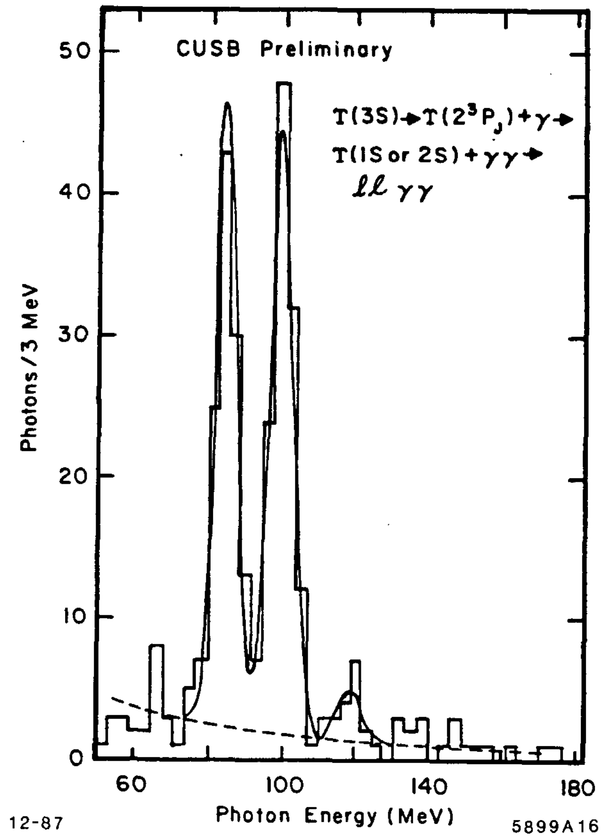

Fig. 6. Observation of the χ_b' states in radiative Υ'' decays from Ref. 12.

$$a = \frac{1}{2m^2}\left\langle -\frac{ds}{rdr} + 3\frac{dv}{rdr} \right\rangle \quad , \tag{2a}$$

$$b = \frac{1}{12m^2}\left\langle \frac{dv}{rdr} - \frac{d^2v}{dr^2} \right\rangle \quad , \tag{2b}$$

where m is the mass of the heavy quark. If only a Coulomb-like vector part of the potential, $v(r) \propto 1/r$, is present, $R_\chi = 0.8$. As the strength of the scalar term, $s(r)$, is increased, there is more cancellation between the two terms on the right-hand side of Eq. (2a); the matrix element a decreases, and R_χ drops below 0.8.

The most recent experimental results[12] are $R_\chi = 0.67 \pm 0.06$ and $R_{\chi'} = 0.69 \pm 0.05$ for bottomonium. A rather simple model accounts for these numbers.[14]

We take the Cornell potential,[15]

$$V(r) = \frac{-\beta}{r} + kr = \frac{-0.52}{r} + \frac{r}{(2.34 \text{ GeV}^{-1})^2} \quad , \tag{3}$$

where the two coefficients having been adjusted to fit the charmonium spectrum (although the model does a quite adequate job in describing bottomonium as well). The Schrodinger equation can be put in dimensionless form by using the variables $\rho = \mu\beta r = r/r_B$ and $K = k/(\beta^3\mu^2)$, where μ is the reduced mass and $r_B = 1/\beta\mu$ is the Bohr radius of the corresponding purely Coulomb potential problem. The assumption is then made that the $-\beta/r$ piece of the potential is a Lorentz four-vector, and the kr piece is a Lorentz scalar for all values of r.

Figure 7 then shows[14] the ratio of mass splittings R_χ for the 1P and 2P levels of the Cornell potential as a function of the scaled variable K. The arrows indicate the values of K corresponding to charmonium and bottomonium. The agreement with experiment is quite good for both charmonium (where experimentally, $R_\chi = 0.48 \pm 0.01$) and bottomonium, considering that nothing about spin-dependent effects was used as an input in the choice of parameters. For charmonium, however, the absolute magnitude of the χ_c splittings is about a factor of two smaller than experiment.

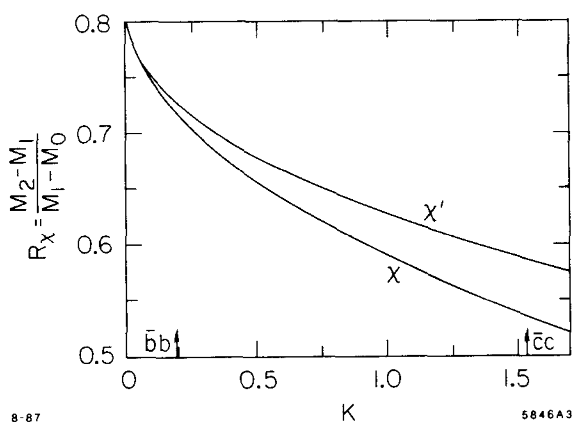

Fig. 7. The ratio R_χ for the Cornell potential as a function of the scaled, dimensionless variable K for the χ and χ' states, respectively. The arrows indicate the values of K corresponding to charmonium and bottomonium with $m_c = 1.84$ GeV and $m_b = 5.17$ GeV, respectively.

For quark masses above ~ 13 GeV, which corresponds to $K = 1/34$, $R_\chi > R_{\chi'}$, opposite to the situation for bottomonium. Even with very high mass quarks, very high radial excitations (which "live" primarily in the confining part of the potential) revert back to the situation for bottomonium: One gets larger values of R_χ as we go up in principal quantum number. More generally, the behavior of R_χ as we go from the lowest P states to their radial excitations is sensitive to the radial dependence of the Lorentz character of the effective interaction between heavy quarks, and can be used as a tool to understand this more detailed feature of the potential.

THE TAU LEPTON

The tau and the tau neutrino seem to fit nicely into the standard assignment of third generation leptons.[16] The limits on the mass of the tau neutrino keep going down year by year, with the latest ARGUS limit[17] being $m_{\nu_\tau} < 35$ MeV.

During the past year there was a flurry of activity after the HRS collaboration claimed[18] a branching ratio at the 5% level for the mode, $\tau^- \to \nu_\tau \pi^- \eta$, that is not expected in the standard model. The $\eta\pi$ system, which is G odd, has natural spin-parity and in the standard model it must come from the vector current, which is G even; we have by definition a process that involves a second class current. A succession of results have come out since then contradicting this claim, with the latest and best limit coming from the Crystal Ball data[19] shown in Figure 8:

$$B(\tau^- \to \nu_\tau \pi^- \eta) < 0.3\% \quad .$$

That brings us to the one problem that is still outstanding in tau physics, the "missing" one-prongs. The decays $\tau \to \nu_\tau e \bar{\nu}_e$, $\tau \to \nu_\tau \mu \bar{\nu}_\mu$, $\tau \to \nu_\tau \pi$, $\tau \to \nu_\tau 2\pi$, $\tau \to \nu_\tau 3\pi$, and $\tau \to \nu_\tau 4\pi$, plus a number of smaller modes occur at the expected rates[20] and the sum of their exclusive branching ratios accounts for about 90% or so of tau decays. Moreover, there seem to be no other modes of consequence which haven't been included. The difficulty of getting to 100% centers on accounting for all the one-prong decays of the tau.

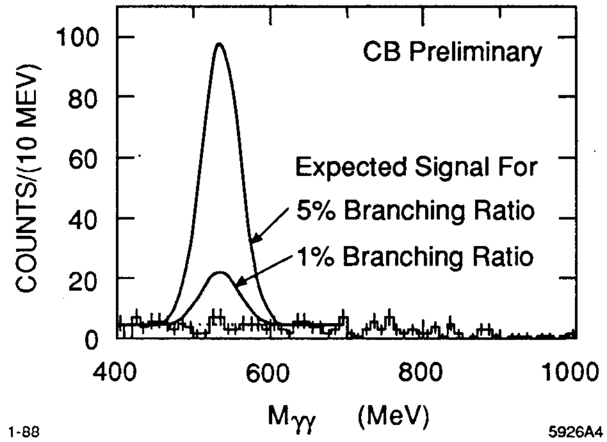

Fig. 8. Upper limit on an η peak from the decay $\tau \rightarrow \nu_\tau \eta \pi$, from Ref. 19.

In more detail, the problem arises as follows. Consider first three-prong decays, which are shown in Table 2. All the theoretical calculations of decay branching ratios[20] are normalized to that for $\tau \rightarrow \nu_\tau e \bar{\nu}_e$, for which we take[21] the world average value of 17.9%. Where theory and experiment can be compared, they are in excellent agreement. Moreover, the sum of all the exclusive measurements[16] is in agreement with the inclusive three-prong branching ratio.

Table 2. Three charged prong decays of the tau.

Decay Mode	Branching Ratio (%)	
	Theory[20]	Experiment[16]
$\tau^- \rightarrow \nu_\tau 2\pi^- \pi^+$		6.7 ± 0.4
$\tau^- \rightarrow \nu_\tau 2\pi^- \pi^+ \pi^0$	4.9	5.0 ± 0.5
$\tau^- \rightarrow \nu_\tau (K\pi)^-$	0.3	0.4 ± 0.1
$\tau^- \rightarrow \nu_\tau K^- \pi^- \pi^+ (\pi^0)$		0.22 ± 0.14
$\tau^- \rightarrow \nu_\tau K^- K^+ \pi^-$		0.22 ± 0.14
$\tau^- \rightarrow \nu_\tau 2\pi^- \pi^+ 3\pi^0$	< 0.4	
Total Exclusive:		12.5 ± 0.6
Total Inclusive:		13.3 ± 0.3

The agreement between theory (where there is a prediction of some accuracy) and experiment (where there is a definite measurement) is also very good in the case of tau decays involving one charged prong, as shown in Table 3. Note in particular that aside from the leptonic decays (whose branching ratio is used as an input), there is excellent agreement between theory and experiment for the semihadronic decays $\tau \to \nu_\tau \pi$, $\tau \to \nu_\tau K$, and $\tau^- \to \nu_\tau \pi^- \pi^0$.

Table 3. One charged prong decays of the tau.

Decay Mode	Branching Ratio (%)	
	Theory[20]	Experiment[16]
$\tau^- \to \nu_\tau e^- \bar{\nu}_e$	17.9 (Input)	17.7 ± 0.4
$\tau^- \to \nu_\tau \mu^- \bar{\nu}_\mu$	17.4	17.7 ± 0.4
$\tau^- \to \nu_\tau \pi^-$	10.9	10.9 ± 0.6
$\tau^- \to \nu_\tau \pi^- \pi^0$	22	22.8 ± 1.0
$\tau^- \to \nu_\tau \pi^- 2\pi^0$	≤ 7.1	7.5 ± 0.9
$\tau^- \to \nu_\tau \pi^- 3\pi^0$	1.0	$0.54 \pm 0.28 \pm 1.06$
$\tau^- \to \nu_\tau \pi^- 4\pi^0$	< 0.1	
$\tau^- \to \nu_\tau \pi^- 5\pi^0$	< 0.1	
$\tau^- \to \nu_\tau (K\bar{K})^-$	< 0.26	
$\tau^- \to \nu_\tau (K\bar{K}\pi)^-$	< 0.5	
$\tau^- \to \nu_\tau \eta \pi^- \pi^0$	0.15	< 0.9
$\tau^- \to \nu_\tau K^-$	0.7	0.6 ± 0.2
$\tau^- \to \nu_\tau (K\pi)^-$ 0.9	0.9	1.2 ± 0.3
Total Exclusive:		78.9 ± 1.6
Total Inclusive:		86.6 ± 0.3

The branching ratio for $\tau^- \to \nu_\tau \pi^- 2\pi^0$ that appears in Table 3 agrees with a new measurement[19] of the Crystal Ball from the data shown in Figure 9 of $7.4 \pm 0.6 \pm 1.3\%$, which also supplies a branching ratio for $\tau^- \to \nu_\tau \pi^- 3\pi^0$ of $0.54 \pm 0.28 \pm 1.06\%$ and the upper limit on $\tau^- \to \nu_\tau \eta \pi^- \pi^0$. We have accounted for all the purely leptonic modes and all the modes of the form $\tau^- \to \nu_\tau(n\pi)^-$ of any substance, as well as Cabibbo-suppressed modes. Although now with smaller experimental error bars, this is the situation that was already noted over two years ago.[20] Where are the remaining one-prong decays?

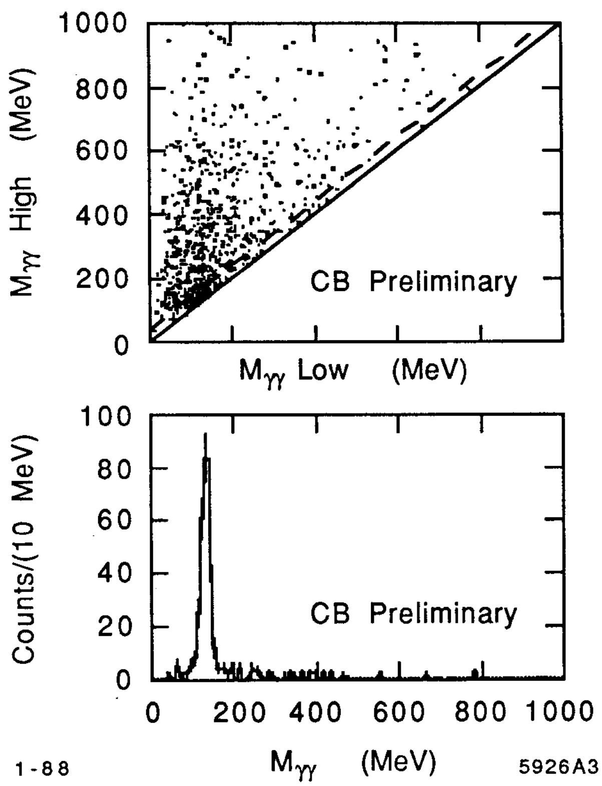

Fig. 9. Scatter plot of $M_{\gamma\gamma}$(high) *versus* $M_{\gamma\gamma}$(low) and the diagonal projection for the measurement of the mode $\tau \to \nu_\tau\pi^-\pi^\circ\pi^\circ$, from Ref. 19.

With other conventional (or even unconventional) modes that contribute to one-prong tau decays severely limited, there are two possible ways out of the problem. The first is that the branching ratio for $\tau \to \nu_\tau e\bar{\nu}_e$, which we took to be 17.9% (and to which we normalized all our theoretical predictions), should be $\approx$ 19%. This would scale up all the predicted branching ratios by $\approx$ 6% and make the sum for theory agree with the measurement of the one-prong inclusive branching ratio. Of course it is one thing to scale up the theory by a common factor, as all the predictions are normalized to the single mode $\tau^- \to \nu_\tau e^-\bar{\nu}_e$, and another to get all the individually measured experimental branching ratios to change, all in the same direction! Nevertheless this is what the present world average value of the tau lifetime indicates (by about 2σ) is the solution, and it is the easiest (purely from the point of view of not straining the standard model) way out.

Second, there is the possibility of something new. This is the direction in which the directly measured branching ratios, and particularly that for $\tau \to \pi^- +$ multineutrals, points. But what? It is hard to find a scenario for this situation which is not very contrived (if it is not to be in conflict with other existing experiments).

The experiments necessary to decide between these two possibilities are hard as well. I suspect that although the puzzle has been sharpened considerably in the past year, it will take some additional measurements over the next several years by detectors well-instrumented for detection of neutrals to resolve it.

MIXING

As in the neutral K system, the neutral D and neutral B systems are capable of exhibiting mixing between, for example, an initial $B_d^0(\bar{b}d)$ and its charge conjugate state, $\bar{B}_d^0(\bar{d}b)$. A typical signature (although hardly the only one) arises from the ensuing semileptonic decay involving a negatively charged lepton instead of the positively charged one which would come from a B_d^0. Calling the eigenstates of the mass matrix B_1 and B_2, with $\Delta M = M_1 - M_2$ and $\Delta\Gamma = \Gamma_1 - \Gamma_2$, the relationship to experiment is made through the quantity

$$r = \frac{(\Delta M)^2 + (\Delta\Gamma/2)^2}{2\Gamma^2 + (\Delta M)^2 - (\Delta\Gamma/2)^2} \approx \frac{(\Delta M/\Gamma)^2}{2 + (\Delta M/\Gamma)^2} \quad , \tag{4}$$

where the last approximation follows when $\Delta\Gamma \leq \Delta M$, as should be the case for the $B - \bar{B}$ system. When the initial B is tagged as to being a B^0 rather than $\bar{B}^0$, $r = \ell^-/\ell^+$, the number of "wrong" to "right" sign leptons in its semileptonic decay. For uncorrelated $B^0 + \bar{B}^0$ pairs it follows that $2r/(1+r^2) = \ell^\pm\ell^\pm/\ell^+\ell^-$, but for correlated pairs produced at the $\Upsilon(4S)$, $r = \ell^\pm\ell^\pm/\ell^+\ell^-$.

For the $D^0 - \bar{D}^0$ system we expect that r is of order 10^{-3} or so.[22] The tightest upper limit,[23]

$$r < 5 \times 10^{-3}$$

comes from E691. On the other hand, the Mark III has three events from operating at the ψ'' which have Kaons of the same sign in nonleptonic decays of the final pair of D mesons.[23] These events also could arise from the doubly Cabibbo-suppressed decay of one of the D mesons and the Cabibbo-allowed decay of the other, and only "look like" mixing. This is also expected[22] at the level of a few times 10^{-3}, while the observed events, if real correspond to a signal at the 10^{-2} level. It will take more experiments to decide what is the level of mixing and of doubly Cabibbo-suppressed decays in the D system.

As already noted, we expect mixing in the neutral B system to be due to ΔM. Before the fact, a theoretical guestimate on the high end for $(\Delta m/\Gamma)_{B_d} = x_d$ was ~ 0.2. This past year the ARGUS collaboration found[24] $x_d = 0.73 \pm 0.18$; the mixing time is not so different from the lifetime. For theorists this has meant an upward adjustment in the combination of a hadronic matrix element, a Kobayashi–Maskawa (K–M) matrix element, and, most importantly, in the value of m_t. For experimentalists, this together with the b lifetime means that in some situations not only will B_d mesons live long enough to leave a measurable gap, but that in this time there is a nonnegligible chance that they will oscillate into the corresponding antiparticle state. The B_s meson must have large mixing in the three-generation standard model, which has important consequences for observing CP violation for the B_s system as we will see in the last Section.

The Kobayashi–Maskawa Matrix

In the standard model, the left-handed quarks are assigned to the weak isospin doublets:

$$\begin{pmatrix} u \\ d' \end{pmatrix}_L \qquad \begin{pmatrix} c \\ s' \end{pmatrix}_L \qquad \begin{pmatrix} t \\ b' \end{pmatrix}_L \quad ,$$

where the upper components are chosen to be the mass eigenstates u, c, and t, and the essential complication that the weak and mass eigenstates are not the same is entirely represented in terms of a matrix transformation[25] operating on the lower components. It takes us from the mass eigenstates (d, s, and b) to the weak eigenstates (d', s', and b'), and is represented by the unitary (K–M) matrix:

$$\begin{pmatrix} d' \\ s' \\ b' \end{pmatrix} = \begin{pmatrix} V_{ud} & V_{us} & V_{ub} \\ V_{cd} & V_{cs} & V_{cb} \\ V_{td} & V_{ts} & V_{tb} \end{pmatrix} \begin{pmatrix} d \\ s \\ b \end{pmatrix} \quad . \tag{5}$$

As V is a 3×3 unitary matrix, it is specified by nine parameters. When we take into account that each of the six quark fields can be changed by a phase with no change in the physics, but that a common phase change of all quark fields would not change the matrix V, we are left with $9 - (6 - 1) = 4$ parameters. These can be considered as three mixing angles and one phase, with a non-zero value of this phase inducing CP violation. More on this in a moment.

At the present time the three angles and one phase of the three-generation K–M matrix are limited by direct measurements of the *magnitudes* of the K–M matrix elements $V_{ud}, V_{us}, V_{cd}, V_{cs}, V_{cb}$, and bounds on the magnitude of V_{ub}. This determines two of the angles (or combinations of the angles) fairly well, and bounds a third one. The key experimental restrictions can be stated as[26]

$$|V_{us}| = 0.221 \pm 0.002 \tag{6}$$

from strange particle decays, and[27]

$$|V_{cb}| = 0.046 \pm 0.010 \tag{7}$$

from the b lifetime, and

$$0.07 \leq |V_{ub}/V_{cb}| \leq 0.23 \quad , \tag{8}$$

where the upper bound comes from the absence of a signal for $b \to u + \ell\bar{\nu}_\ell$ in semileptonic B decay and the lower bound from the ARGUS observation[5] of exclusive baryonic B decays, shown in Figure 10, which involve $b \to u + d\bar{u}$ at

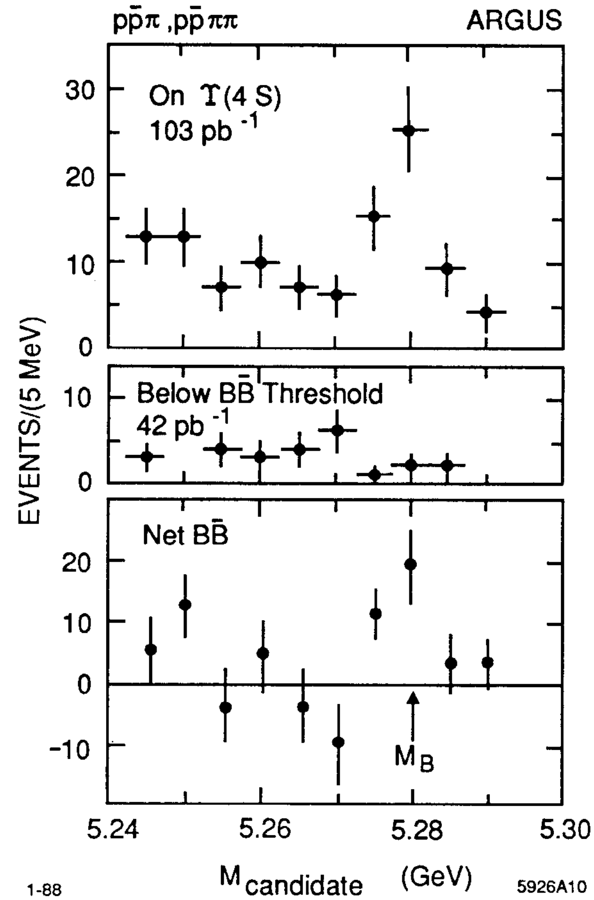

Fig. 10. Mass distribution of ARGUS candidates for $B^- \rightarrow p\bar{p}\pi^-$ and $\bar{B}^0 \rightarrow p\bar{p}\pi^+\pi^-$ and their charge conjugates from Ref. 5 in running at the $\Upsilon(4S)$, below threshold, and the bin-by-bin net yield, by subtraction, coming from $B\bar{B}$.

the quark level. The present results[28] from CLEO are not conclusive on the existence of these modes, although there is an apparent signal in the $p\bar{p}\pi$ mode (see Figure 11) which is comparable to that of ARGUS when interpreted as coming from B decay, if similar data selection cuts are made.

As noted at the beginning of this Section, the standard model allows for CP violation in the form of phases in the quark mixing matrix, the K–M matrix.[25] When there are three generations of quarks and leptons, there is precisely one CP violating phase, δ. The theoretical expressions for all CP violating quantities then contain a factor of $\sin\delta$.

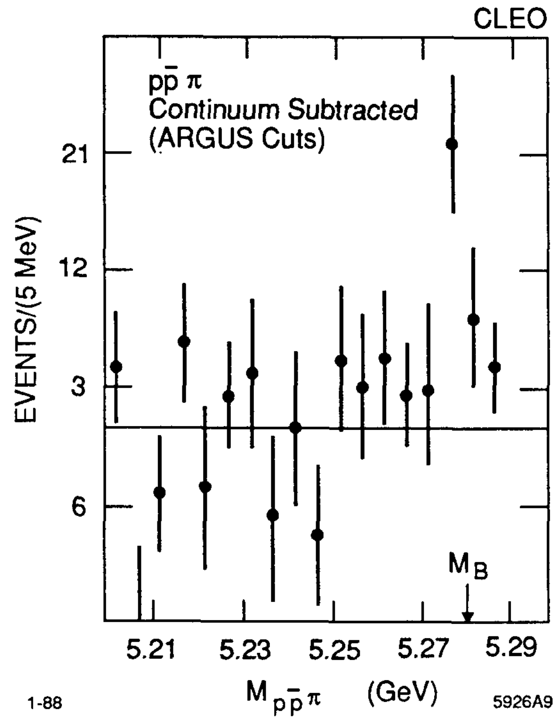

Fig. 11. Mass distribution of CLEO candidates for $B^- \to p\bar{p}\pi^-$ and charge conjugate after bin-by-bin subtraction, from Ref. 29.

CP violation has still only been observed in the neutral K system. There it is conveniently parametrized in terms of ϵ, which characterizes CP violation in the Kaon mass matrix, and ϵ', which is non-zero only due to CP violation in the neutral Kaon decay amplitude. Over the last ten years the expectations for ϵ and for ϵ' have been refined again and again as new experimental and theoretical information became available on hadronic matrix elements, K–M matrix elements, m_t, *etc.*[29]

As the present year has proceeded the standard model "explanation" of CP violation has looked better and better. In particular, there have been two important new experimental results for ϵ'/ϵ. First came the preliminary result from a test run of the Fermilab experiment:[30]

$$\epsilon'/\epsilon = 3.5 \pm 3.0 \pm 2.0 \ \times 10^{-3} \quad ,$$

and then this past summer, the preliminary result from the CERN experiment[31]

$$\epsilon'/\epsilon = 3.5 \pm 0.7 \pm 0.4 \pm 1.2 \ \times 10^{-3} \quad .$$

Both experiments have the capability of eventually decreasing both their statistical and systematic error bars below the 10^{-3} level. If $\epsilon'/\epsilon \sim 3.5 \times 10^{-3}$, then CP violating effects from heavy quark loops is a likely interpretation and, especially if m_t is large, the result is not wildly different than expectations. It would seem that the wind is blowing in the direction of the standard model and the explanation of CP violation in terms of the K–M phase.

CP VIOLATION IN B DECAY

The possibilities for observation of CP violation in B decays are much richer than for the neutral K system. The situation is even reversed, in that for the B system the variety and size of CP violating asymmetries in decay amplitudes far overshadows that in the mass matrix.[32]

To start with the familiar, however, it is useful to consider the phenomenon of CP violation in the mass matrix of the neutral B system. Here, in analogy with the neutral K system, one defines a parameter ϵ_B. It is related to p and q, the coefficients of the B° and $\bar{B}^\circ$, respectively, in the combination which is a mass matrix eigenstate by

$$\frac{q}{p} = \frac{1-\epsilon_B}{1+\epsilon_B} \quad .$$

The charge asymmetry in $B^\circ\bar{B}^\circ \to \ell^\pm\ell^\pm + X$ is given by[33]

$$\frac{\sigma(B^\circ\bar{B}^\circ \to \ell^+\ell^+ + X) - \sigma(B^\circ\bar{B}^\circ \to \ell^-\ell^- + X)}{\sigma(B^\circ\bar{B}^\circ \to \ell^+\ell^+ + X) + \sigma(B^\circ\bar{B}^\circ \to \ell^-\ell^- + X)} = \frac{|\frac{p}{q}|^2 - |\frac{q}{p}|^2}{|\frac{p}{q}|^2 + |\frac{q}{p}|^2} \tag{9}$$

$$= \frac{Im(\Gamma_{12}/M_{12})}{1+\frac{1}{4}|\Gamma_{12}/M_{12}|^2} \tag{10}$$

where we define $< B^\circ|H|\bar{B}^\circ >= M_{12} - \frac{i}{2}\Gamma_{12}$. The quantity $|M_{12}|$ is measured in $B-\bar{B}$ mixing and we may estimate Γ_{12} by noting that it gets contributions from B° decay channels which are common to both B° and $\bar{B}^\circ$, *i.e.*, K–M-suppressed decay modes. This causes the charge asymmetry for dileptons most likely to be in the ballpark of a few times 10^{-3}, and at best 10^{-2}. For the foreseeable future, we might as well forget it experimentally.

Turning now to CP violation in decay amplitudes, in principle this can occur whenever there is more than one path to a common final state. For example, let us consider decay to a CP eigenstate, f, like ψK_s°. Since there is substantial $B^\circ - \bar{B}^\circ$ mixing, one can consider two decay chains of an initial B° meson:

$$\begin{array}{ll} B^\circ \to B^\circ & \searrow \\ & \quad f \quad , \\ B^\circ \to \bar{B}^\circ & \nearrow \end{array}$$

where f is a CP eigenstate. The second path differs in its phase because of the mixing of $B^\circ \to \bar{B}^\circ$, and because the decay of a $\bar{B}$ involves the complex conjugate of the K–M factors involved in B decay. The strong interactions, being CP invariant, give the same phases for the two paths. The amplitudes for these decay chains can interfere and generate non-zero asymmetries between $\Gamma(B^\circ(t) \to f)$ and $\Gamma(\bar{B}^\circ(t) \to f)$. Specifically,

$$\Gamma(\bar{B}^\circ(t) \to f) \sim e^{-\Gamma t} \left(1 - \sin[\Delta m\ t] Im\left(\frac{p}{q}\rho\right)\right) \tag{11a}$$

and

$$\Gamma(B^\circ(t) \to f) \sim e^{-\Gamma t} \left(1 + \sin[\Delta m\ t] Im\left(\frac{p}{q}\rho\right)\right) \quad . \tag{11b}$$

Here we have neglected any lifetime difference between the mass matrix eigenstates (thought to be very small) and set $\Delta m = m_1 - m_2$, the difference of the eigenstate masses, and $\rho = A(B \to f)/A(\bar{B} \to f)$, the ratio of the amplitudes, and we have used the fact that $|\rho| = 1$ when f is a CP eigenstate in writing Eqs. (11a) and (11b). From this we can form the asymmetry:

$$A_{\text{CP Violation}} = \frac{\Gamma(B) - \Gamma(\bar{B})}{\Gamma(B) + \Gamma(\bar{B})} = \sin[\Delta m\ t] Im\left(\frac{p}{q}\rho\right) \quad . \tag{12}$$

In the particular case of decay to a CP eigenstate, the quantity $Im\left(\frac{p}{q}\rho\right)$ is given entirely by the K–M matrix and is independent of hadronic amplitudes. To measure the asymmetry experimentally, however, one must know if one starts with an initial B° or $\bar{B}^\circ$, *i.e.*, one must "tag."

We can also form asymmetries where the final state f is not a CP eigenstate. Examples are $B_d \to D\pi$ compared to $\bar{B}_d \to \bar{D}\bar{\pi}$; $B_d \to \bar{D}\pi$ compared to $\bar{B}_d \to D\bar{\pi}$; or $B_s \to D_s^+ K^-$ compared to $\bar{B}_s \to \bar{D}_s^- K^+$. These is a decided disadvantage here in theoretical interpretation, in that the quantity $Im\left(\frac{p}{q}\rho\right)$ is now dependent on hadron dynamics.

It is instructive to look not just at the time-integrated asymmetry between rates for a given decay process and its CP conjugate, but to follow the time dependence,[34] as given in Eqs. (11a) and (11b). As an example, Figure 12 shows[35] the time dependence for the quark level process $\bar{b} \to \bar{c}c\bar{s}$ (solid curve) in comparison to that for $b \to c\bar{c}s$ (dashed curve). At the hadron level this could be, for example, B_d in comparison to $\bar{B}_d$ decaying to the same, (CP self-conjugate) final state, ψK_s°. As discussed before, $|\rho| = 1$ in this case.

The three parts of Figure 12 show the situation for $\Delta m/\Gamma = 0.2$ (at the high end of theoretical prejudice before the ARGUS result[24] for B_d mixing), $\Delta m/\Gamma = \pi/4$ (near the central value from ARGUS), and $\Delta m/\Gamma = 5$ (roughly the minimum value expected for the B_s in the three-generation standard model, given the central value of ARGUS for B_d). The advantages of having $\Delta m/\Gamma$ for the B_d° system as suggested by ARGUS (Figure 12b) rather than previous theoretical estimates (Figure 12a) are very apparent. When we go to mixing parameters expected for the B_s° system (Figure 12c), the effects are truly spectacular. In fact, in this last case, the time average asymmetry is washed out by the many oscillations in one lifetime and a study of the time dependence of the asymmetry is a necessity.

A second path to the same final state could arise in several other ways besides through mixing. For example, one could have two cascade decays that end up with the same final state, such as:

$$B_u^- \to D^\circ K^- \to K_s^\circ \pi^\circ K^-$$

and

$$B_u^- \to \bar{D}^\circ K^- \to K_s^\circ \pi^\circ K^- \quad .$$

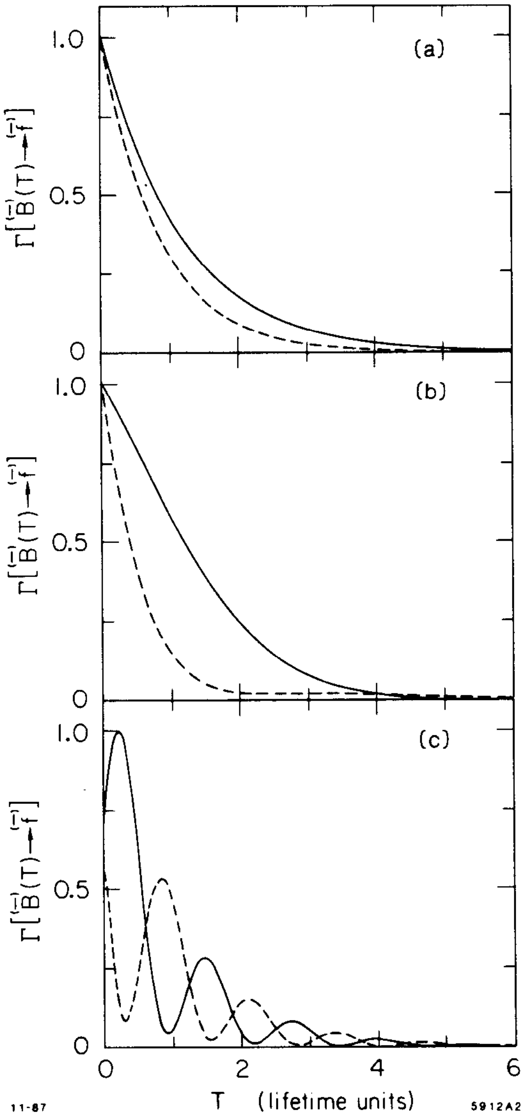

Fig. 12. The time dependence for the quark level process $\bar{b} \rightarrow \bar{c}c\bar{s}$ (solid curve) in comparison to that for $b \rightarrow c\bar{c}s$ (dashed curve). At the hadron level this could be, for example, $B_d \rightarrow \psi K_s^\circ$ (dashed curve) in comparison to $\bar{B}_d \rightarrow \psi K_s^\circ$ (solid curve). (The curves are interchanged for the ψK_s° final state because it is odd under CP.) The three subgraphs correspond to (a) $\Delta m/\Gamma = 0.2$, (b) $\Delta m/\Gamma = \pi/4$, and (c) $\Delta m/\Gamma = 5$.

Another possibility is to have spectator and annihilation graphs contribute to the same process.[36] Still another is to have spectator and "penguin" diagrams interfere. This latter possibility is the analogue of the origin of the parameter ϵ' in neutral K decay, but as discussed previously, there is no reason to generally expect a small asymmetry here. Indeed, with a careful choice of the decay process, large CP violating asymmetries are expected.

Note that not only do these routes to obtaining a CP violating asymmetry in decay rates not involve mixing, but they do not require one to know whether one started with a B or $\bar{B}$, *i.e.*, they do not require "tagging." These decay modes are in fact "self-tagging" in that the properties of the decay products (through their electric charges or flavors) themselves fix the nature of the parent B or $\bar{B}$.

Even with potentially large asymmetries, the experimental task of detecting these effects is a monumental one. When the numbers for branching ratios, efficiencies, *etc.* are put in, it appears that 10^7 to 10^8 produced B mesons are required to end up with a significant asymmetry (say, 3σ), depending on the decay mode chosen.[32] This is beyond the samples available today (of order a few times 10^5) or in the near future ($\sim 10^6$). On the other hand, it is possible to envision such samples at new electron-positron colliders, fixed target experiments and, at hadron colliders, especially the SSC.[32] A great deal of experimental work needs to be done to explore both technique and physics to achieve the goal of observing CP violation in the B system. A good start has already been made. With the excitement within the experimental community that has been growing over the past few years, it begins to seem likely that in the next five years we will see the experimental situation develop to the point that this physics is capable of being attacked.

REFERENCES

1. See, for example, the lectures of F. J. Gilman, *Proceedings of the Fourteenth SLAC Summer Institute on Particle Physics,* edited by E. C. Brennan (SLAC, Stanford, 1987), p. 191.
2. Witherell, M., invited talk at the International Symposium on Production and Decay of Heavy Flavors, Stanford, California, September 1–5, 1987 (unpublished); Raab, J. R. *et al.,* Fermilab preprint FERMILAB-PUB-87/144-E, 1987 (unpublished).
3. Some recent reviews from various viewpoints are found in R. Ruckl, *Proceedings of the XXIII International Conference on High Energy Physics,* edited by S. C. Loken (World Scientific, Singapore, 1987), p. 797; Stech, B., invited talk at the International Europhysics Conference, Uppsala, Sweden, June 25–July 1, 1987 and Heidelberg preprint HD-THEP-87-19, 1987 (unpublished); Bigi, I. I. Y., invited talk at the International Symposium on the Production and Decay of Heavy Flavors, Stanford, California, September 1–5, 1987 and SLAC preprint SLAC-PUB-4455, 1987 (unpublished).
4. Hitlin, D., lectures presented at the Charm Physics Symposium/Workshop, Beijing, China, June 4–16, 1987 and Caltech preprint CALT-68-1463, 1987 (unpublished).
5. Schmidt-Parzefall, W., invited talk at the 1987 International Symposium on Lepton and Photon Interactions at High Energies, Hamburg, Germany, July 27–31, 1987.
6. Snyder, A., invited talk at the International Symposium on Production and Decay of Heavy Flavors, Stanford, California, September 1–5, 1987 and SLAC-PUB-4466, 1987 (unpublished).
7. Stockdale, I. E., talk at the International Symposium on Production and Decay of Heavy Flavors, Stanford, California, September 1–5, 1987 and SLAC-PUB-4466, 1987 (unpublished).
8. Albrecht, H. *et al.,* Phys. Lett. **158B**, 525 (1985); Bebek, C. *et al.,* Phys. Rev.

Lett. **56,** 1983 (1986); Baltrusaitis, R. M. *et al.,* Phys. Rev. Lett. **56,** 2136 (1986).

9. Donoghue, J. F., Phys. Rev. **D33,** 1516 (1986).

10. Wormser, G., talk at the International Symposium on Production and Decay of Heavy Flavors, Stanford, California, September 1–5, 1987 and SLAC–PUB–4466, 1987 (unpublished).

11. Stockdale, I. E., Ref. 7.

12. Lee–Franzini, J., invited talk at the XVth SLAC Summer Institute on Particle Physics, August 10–21, 1987 (unpublished).

13. Rosner, J. L., *Proceedings of the 1985 Symposium on Lepton and Photon Interactions at High Energies,* edited by M. Konuma and K. Takahashi (Kyoto University, Kyoto, 1986), p. 448. See also H. J. Schnitzer, Phys. Lett. **134B,** 253 (1984) and Phys. Rev. Lett. **35,** 1540 (1975) for the original use of the ratio R_χ.

14. Dib, C. O., Franzini, P. J. and Gilman, F. J., SLAC preprint SLAC–PUB–4402, 1987 and Phys. Rev., in press.

15. Eichten, E. *et al.,* Phys. Rev. **D17,** 3090 (1978) and **D21,** 203 (1980).

16. See the review by K. K. Gan and M. L. Perl, SLAC–PUB–4331, 1987 (unpublished).

17. Albrecht, H. *et al.,* DESY preprint DESY 87–148, 1987 (unpublished).

18. Derrick, M. *et al.,* Phys. Lett. **189B,** 260 (1987).

19. Lowe, S. T., invited talk at the International Symposium on Production and Decay of Heavy Flavors, Stanford, California, September 1–5, 1987 and SLAC–PUB–4449, 1987 (unpublished).

20. Gilman, F. J. and Rhie, S. H., Phys. Rev. **D31,** 1066 (1985); Gilman, F. J., lectures given at the Charm Physics Symposium/Workshop, Beijing, China, June 4–16, 1987 and SLAC–PUB–4352, 1987 (unpublished) and references to previous work therein.

21. This number is computed using the data on both the modes $\tau^- \rightarrow \nu_\tau e^- \bar{\nu}_e$ and $\tau^- \rightarrow \nu_\tau \mu^- \bar{\nu}_\mu$ together with $e - \mu$ universality to make a fit to both their branching ratios. For a separate determination see Ref. 16.

22. See, for example, I. I. Y. Bigi, lectures at the Theoretical Advanced Study Institute, Santa Fe, New Mexico, July 5–August 1, 1987 and SLAC preprint SLAC–PUB–4439, 1987 (unpublished), and references therein.

23. See the review by G. Gladding, invited talk at the International Symposium on Production and Decay of Heavy Flavors, Stanford, California, September 1–5, 1987 (unpublished).

24. Albrecht, H. *et al.,* Phys. Lett. **192B,** 245 (1987).

25. Kobayashi, M. and Maskawa, T., Prog. Theor. Phys. **49,** 652 (1973).

26. Particle Data Group, Phys. Lett. **170B,** 74 (1986), and references therein.

27. Most of the error quoted in Eq. (7) is not from the experimental uncertainty in the value of the b lifetime, but in the theoretical uncertainties in choosing a value of m_b and in the use of the quark model to represent inclusively semileptonic decays which, at least for the B, are dominated by a few exclusive channels. We have made the error bars larger than they are often stated to reflect these uncertainties.

28. Berkelman, K., invited talk at the International Symposium on the Production and Decay of Heavy Flavors, Stanford, California, September 1–5, 1987 and Cornell preprint CLNS 97/102, 1987 (unpublished).

29. See F. J. Gilman, invited talk at the International Symposium on the Production and Decay of Heavy Flavors, Stanford, California, September 1–5, 1987 and SLAC preprint SLAC–PUB–4469, 1987 (unpublished) for a brief, recent review of the situation.

30. Winstein, B., invited talk at the Division of Particles and Fields Meeting, Salt Lake City, Utah, January 14–17, 1987 (unpublished).

31. Mannelli, I., invited talk at the International Symposium on Lepton and Photon Interactions at High Energies, Hamburg, Germany, July 27–31,

1987 and CERN preprint CERN–EP/87–177, 1987 (unpublished).

32. Foley, K. J. *et al.*, to be published in the Proceedings of the Workshop on Experiments, Detectors, and Experimental Areas for the Supercollider, Berkeley, California, July 7–17, 1987, and SLAC preprint SLAC–PUB–4426, 1987 (unpublished) review the current status of CP violation in B decay and give references to previous work.

33. Pais, A. and Treiman, S. B., Phys. Rev. **D12**, 2744 (1975); Okun, L. B. *et al.*, Nuovo Cim. Lett. **13**, 218 (1975).

34. The importance of this has been particularly emphasized by I. Dunietz and J. L. Rosner, Phys. Rev. **D34**, 1404 (1986).

35. These graphs were constructed by R. Kauffman, in accord with the paper of Dunietz and Rosner, Ref. 34, but with somewhat different parameters: $s_1 = 0.22, s_2 = 0.09, s_3 = 0.05$ and $\delta = 150°$ and the values of $\Delta M/\Gamma$ given in the text and figure captions.

36. This possibility has been particularly emphasized by L. L. Chau and H. Y. Cheng, Phys. Lett. **165B**, 429 (1985).