STUDIES OF THE STRONG AND ELECTROWEAK INTERACTIONS AT THE Z^0 POLE*

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Abstract

This thesis presents studies of the strong and electroweak forces, two of the fundamental interactions that govern the behaviour of matter at high energies. We have used the hadronic decays of Z^0 bosons produced with the unique experimental apparatus of the e^+e^- Linear Collider at the Stanford Linear Accelerator Center (SLAC) and the SLAC Large Detector (SLD) for these measurements. Employing the precision tracking capabilities of the SLD, we isolated samples of Z^0 events containing primarily the decays of the Z^0 to a chosen quark type. With an inclusive selection technique, we have tested the flavor independence of the strong coupling, α_{s_1} by measuring the rates of multi-jet production in isolated samples of light (uds), c, and b quark events. We find: $\alpha_{*}^{uds}/\alpha_{*}^{all} = 0.987 \pm 0.027(stat) \pm 0.022(syst) \pm 0.022(theory)$, $\alpha_s^c/\alpha_s^{all} = 1.012 \pm 0.104(stat) \pm 0.102(syst) \pm 0.096(theory), \alpha_s^b/\alpha_s^{all} = 1.026 \pm 0.006(theory)$ $0.041(stat) \pm 0.041(syst) \pm 0.030(theory)$, which implies that the strong interaction is independent of quark flavor within our present experimental sensitivity. We have also measured the extent of parity-violation in the $Z^0 c\bar{c}$ coupling, given by the parameter A_{c}^{0} using a sample of fully and partially reconstructed D^{\star} and D^{+} meson decays and the longitudinal polarization of the SLC electron beam. This sample of charm quark events was derived with selection techniques based on their kinematic properties and decay topologies. We find $A_c^0 = 0.73 \pm 0.22(stat) \pm 0.10(syst)$. This value is consistent with that expected in the electroweak standard model of particle interactions.

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^{*}Ph.D. Thesis

This thesis is dedicated to:

my grandfather

Stanley S. Kotarski (d. 1.988)

for his constant support and encouragement May this make him proud...

and to

my son Andrew. D. Hildreth (b. 1994)

for reestablishing my wonder at the miracle of the universe

I wish they could have known each other

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Ah, a chance to reflect on the halcyon days of graduate school... Believe it or not, studying physics here at Stanford has been more fun that just about anything else I've done. I guess that's why I've been called the "teflon graduate student" – none of the prototypically "bad" experiences of grad school life seemed to bother me. I even (sort-of) enjoyed the opportunity to study for the qual. Well, enough about pain and torture...

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Chapter 1

Overview

This thesis presents two studies of the physics encompassed in the Standard Model of particle physics. The first study is a test of the flavor independence of strong interactions, a fundamental assumption of the theory of strong interactions, Quantum Chromo-Dynamics (QCD). The second analysis represents the first direct measurement of parity violation in the coupling of the Z^0 boson to charm quarks. The data used to perform these investigations was acquired at the Stanford Linear Accelerator Center (SLAC) using the SLAC Large Detector (SLD) at the SLAC Linear Collider (SLC) during the calendar year 1993. All of the data was produced at a center-of-mass energy equal to the Z^0 boson mass.

The thesis begins with introductory chapters which present an overview of the strong and electroweak interactions that make up the Standard Model. Chapters follow which contain descriptions of the detector and accelerator apparatus and its performance which are relevant to both analyses. At this point, the thesis is divided into two parts, the first containing a description of the analysis comprising the test of QCD in its entirety, the second containing the measurement of the parity violation in the coupling of the Z^0 boson to charm quarks. Supporting appendices follow, adding complementary information to that presented in the preceding chapters. In particular, Appendix C, presents an overview of the types of accelerator-induced backgrounds that can be encountered at a linear collider and various techniques for their reduction.

The following section gives a brief layman's overview of the components of what

has become known as the Standard Model of particle physics. Those portions relevant to this thesis will be introduced in much greater detail in the following chapters.

1.1 The Standard Model

In standard usage, the term "Standard Model" refers to the quantum field theories of the electroweak' and strong interactions that describe the interactions between the fundamental particles that make up the universe.

These fundamental particles, all of which are fermions $(spin \frac{1}{2})$, are divided into quarks and leptons. Quarks carry electroweak and color charges, and thus interact via both the electroweak and strong interactions; they combine to form composite particles called *hadrons* like the proton and neutron. The leptons consist of particles like the electron which carry electroweak charges, and neutrinos which only carry the weak charge.

The fundamental forces of nature are represented in these theories by vector (spin 1) bosons which mediate the interactions between the fundamental fermions. The electroweak force contains four such bosons: the massless photon (γ), which mediates the familiar electromagnetic force, the neutral Z^0 ($m_Z \sim 91 \text{ GeV}/c^2$), and the charged W^+ and W^- ($m_W \sim 80 \text{ GeV}/c^2$), all of which are responsible for the weak force. The strong interactions are mediated by massless particles called gluons.

The one ingredient missing from the above description is some means by which the masses of the fundamental fermions and bosons can be generated. This addition to the theoretical framework is commonly referred to **as** the **Higgs** mechanism. Its effect will be discussed in detail in the next 'chapter.

As a summary, the fermions of the theory are listed in Table 1.1.

^{&#}x27;The electroweak force contains a description of the electromagnetic and weak forces in a common field-theoretical language – these two seemingly unrelated forces are different aspects of one underlying interaction.

Table 1.1: The fundamental fermions of the Standard Model, with their electric charges and approximate masses in units of GeV/c^2 .

Quarks				
Charge	Name (Mass)			
$+\frac{2}{3}e$	$u (\sim 0.1)$ $c (\sim 1.5)$ $t (\sim 175)$			
$-\frac{1}{3}e$	$d ~(\sim 0.1)$	s (~ 0.3)	b (~ 5)	
Leptons				
Charge Name (Mass)				
-1e	e (~ 0.05)	$\mu~(\sim 0.1)$	$ au~(\sim 1.7)$	
0	$\nu_e~(<\sim 10^{-6})$	$\nu_{\mu} \; (< \sim 0.003)$	$ u_{ au} \ (< \sim 0.03)$	

Chapter 2

The Electroweak Interaction: An Introduction

Bccquerel's conjecture' that Roentgen's X rays[2] may have some relation to phosphoresence in uranic salts has turned out to be true, but not in any way he could have imagined at the time. Becquerel's 1896 discovery of what is now called radioactivity[3] while pursuing this question marks the starting point of the study of subatomic physics as well as the first experimental manifestation of the weak nuclear force. Since this result predates even the characterization of the electron as a fundamental particle[4] by a few years[†], it is almost certain that the true implications of this discovery were incomprehensible at the time.

The discovery and subsequent characterization of radioactivity, however, have functioned as major catalysts for the enormous growth in our understanding of the subatomic world that has occurred over the past century. Consideration of β decay spectra in conjunction with the unsolved problem of nuclear structure led **Pauli** to propose the existence of the electron neutrino[6] in 1930; Fermi provided the first "correct" description of the low-energy weak interaction couched in the language of quantum field theory only three years later[7]. Out of all of the developments that unfolded over the next two decades, the most striking is the discovery in 1957 that

^{&#}x27;Ref. [I] has been invaluable to the preparation of this thesis.

the discovery of the nucleus[5] was still 15 years hence

the weak interaction maximally violates parity[8], which led (after some confusion) to the V – A theory of the weak interaction[9] with the four-fermion interaction based on mediating charged vector bosons, denoted W^* .

Of course, the development of weak interaction theory did not proceed independently of the growing population of elementary particles. Studies of μ decay in the 1940's indicated that the weak coupling was similar in magnitude and structure to that observed in β decay, which led many to postulate that the weak interaction was universal. As the quest for unification continued, the electromagnetic interaction also came under scrutiny, and the search for a combined theory of the weak and electromagnetic interactions was begun. The earlier introduction by Yang and Mills of a non-abelian theory of particle interactions[10] provided the basis, and after a decade or more of casting about for the correct formulation of the theory, Weinberg[11] and Salam[12] proposed a theory of the electroweak interaction which contained the underlying symmetry $SU(2) \ge U(1)[13]$. T is unification of the two theories necessitated the introduction of a new gauge boson, the Z^0 , which is the weak interaction analogue to the photon. The Higgs method [14] of spontaneously breaking the symmetry of the theory [15] was used to transform the massless gauge bosons which result from the Yang-Mills theory into the massless photon and the massive Z^0 and W^{\pm} . One essential difficulty remained, however: this theory seemed not to be renormalizable. This impasse was solved by 't Hooft[16] in 1971, and the theory was essentially complete.

It remained for the experimentalists to find some evidence of the weak neutral current or the massive gauge bosons themselves. This first came in studies of neutrino scattering, where the elastic scattering[17] $\nu_{\mu}e \rightarrow \nu_{\mu}e$ and the inelastic scattering[18] $\nu_{\mu}N \rightarrow \nu_{\mu} + X$ were seen. Corroborating evidence for the doublet structure of $SU(2) \times U(1)$ was provided in the following year by the discovery of the charm quark[19]. The apparent symmetry between the generations of quarks and leptons was upset shortly thereafter by the observation of another heavy lepton, the τ [20]. The discovery of the Υ resonances in $\mu^+\mu^-$ production[21] provided the first quark of the third generation, and reestablished the symmetry of the electroweak theory. From this point, the search was on for the elusive top quark, the last member of the third generation. Eventually, the W^{\pm} and Z^0 bosons were observed directly in $p\bar{p}$

collisions[22] in 1982 and 1983, providing the final piece of material evidence that the Weinberg-Salam model of the electroweak interaction properly describes Nature. The recent observation of the long-sought top quark[23] provides the final piece of the particle complement puzzle.

As it stands, what has come to be called the Standard Model of electroweak interactions is a remarkably successful theory. Since its initial formulation, virtually all of its predictions have been verified to ever-increasing precision. Attention has now turned from confirmation of the model's validity to the questions underlying the assumptions involved in its construction: why are there three generations of fermions? What is the relationship between quarks and leptons, given that both of them are necessary for the theory to be renormalizable[24]? What is the nature of CP violation? Just what is the true Higgs mechanism? The current program of precision measurements of electroweak parameters was undertaken to provide clues to the answers of these questions. Hopefully, the answers will be forthcoming.

2.1 The Electroweak Standard Model

In this section, we present an overview of the components of the electroweak theory described in the Standard Model*.

The unification of the electromagnetic and weak forces is accomplished by modifying the eletromagnetic current

$$-iej_{\mu}^{em}A^{\mu} = -ie\left(\bar{\psi}\gamma_{\mu}\mathcal{Q}\psi\right)A^{\mu} \tag{2.1}$$

to include the weak interaction. Here, A^{μ} is the four-potential of the electromagnetic field, satisfying $\partial^2 A = 0$, ψ is the fermion spinor wavefunction, γ_{μ} is the usual spinor matrix, and Q is the charge operator, which has eigenvalue -1 for the electron. To complete this unification, the electromagnetic current is replaced by two new currents: an $SU(2)_L$ isotriplet of weak currents J,, which couple to three vector bosons \mathbf{W}^{μ} , given by

$$-ig\mathbf{J}_{\mu}\cdot\mathbf{W}^{\mu} = -ig\bar{\chi}_{L}\gamma_{\mu}\mathbf{T}\cdot\mathbf{W}^{\mu}\chi_{L}$$
(2.2)

[‡]More detailed expositions of this type can be found in, e.g., Ref. [25] or Ref. [26].

Table 2.1: The L isodoublets and R isosinglets of $SU(2)_L \ge U(1)_Y$

Quarks

$$\binom{u}{d}_{L}$$
 $(u)_{R}$ or $(d)_{R}$ $\binom{c}{b}_{L}$ $(c)_{R}$ or $(b)_{R}$ $\binom{t}{b}_{L}$ $(t)_{R}$ or $(b)_{R}$

Leptons

$$\begin{pmatrix} \nu_e \\ e \end{pmatrix} (e)_R \qquad \begin{pmatrix} \nu_\mu \\ \mu \end{pmatrix}_L \qquad (\mu)_R \qquad \begin{pmatrix} \nu_\tau \\ \tau \end{pmatrix}_L \qquad (\tau)_R$$

and a U(1) weak hypercharge current j^{Y}_{μ} coupled to a fourth vector boson B^{μ} , given by

$$-i\frac{g'}{2}j^{Y}_{\mu}B^{\mu} = -ig'\bar{\psi}\gamma_{\mu}\frac{Y}{2}\psi B^{\mu} .$$
 (2.3)

The operators **T** and Y are the generators of the $SU(2)_L$ and $U(1)_Y$ gauge groups; the operators **T** satisfy the commutation relation

$$\left[T^{i}, T^{j}\right] = i_{\tilde{i}_{j}ik}T^{k}$$
(2.4)

The new coupling constants of the theory are g and g'. The fundamental fermions have been arranged in left-handed isodoublets χ_L and right-handed isosinglets ψ_R , which transform as

$$\chi_L \rightarrow \chi'_L = e^{i\boldsymbol{\alpha}(x)\cdot\mathbf{T}+i\boldsymbol{\beta}(x)Y}\chi_L$$

$$\psi_R \rightarrow \psi'_R = e^{i\boldsymbol{\beta}(x)Y}\psi_R \qquad (2.5)$$

under the infinitesimal gauge transformations given by $\alpha(x)$ and $\beta(x)$. The grouping of the fermions is shown in Table 2.1. Given a field corresponding to any fermion f, the left-handed and right-handed components can be projected out by means of the operator γ_5 :

$$f_{R} = \frac{1}{2}(1 - \gamma_5)f$$
, $f_{L} = \frac{1}{2}(1 + \gamma_5)f$. (2.6)

The general electroweak interaction

$$-ig(J^{i})^{\mu}W^{i}_{\mu} - \frac{g'}{2}(j^{Y})^{\mu}B_{\mu}$$
(2.7)

can be rewritten in terms of the electric-charge raising and lowering bosons

$$W^{\pm}_{\mu} = \sqrt{\frac{1}{2}} \left(W^{1}_{\mu} \mp W^{2}_{\mu} \right)$$
(2.8)

which only couple to the left-handed fermion doublets, and a linear combination of two neutral bosons A^{μ} and Z^{μ}

$$-ig(J^{3})_{\mu}(W^{3})^{\mu} - \frac{'g}{2}(j^{Y})^{\mu}B_{\mu} = -i\left(g\sin\theta_{W}J_{\mu}^{3} + g'\cos\theta_{W}\frac{y_{\mu}}{2}\right)A^{\mu} \\ -i\left(g\cos\theta_{W}J_{\mu}^{3} - g'\sin\theta_{W}\frac{y_{\mu}^{Y}}{2}\right)Z^{\mu}.$$
 (2.9)

This linear combination is just the fields W^3 and B rotated by an angle θ_W in the (W^3, B) plane. The value of the so-called Weinberg angle θ_W is not predicted by the theory and must be determined from experiment. Relations between the couplings g and g' and θ_W can be derived by noting that the generators of the two groups also satisfy the relation

$$Q = T^3 + \frac{Y}{2} , \qquad (2.10)$$

where Q is the electric charge. This leads to an expression of the electromagnetic current in terms of the two newly introduced weak currents:

$$j_{\mu}^{em} = J_{\mu}^3 + \frac{1}{2} j_{\mu}^Y$$
 (2.11)

This **choice** of currents, with the electric charge given by Eq. 2.10, guarantees that the gauge boson A_{μ} which mediates the electromagnetic interaction will remain massless after the spontaneous symmetry breaking of the Higgs mechanism (to be discussed in the next section) that gives masses to the other gauge bosons.

If we identify the first term on the right side of Eq. 2.9 with the electromagnetic current, we see immediately that

$$g \sin \theta_W = g' \cos \theta_W = e$$
 (2.12)

In addition, the relationship

$$\tan \theta_W = \frac{g'}{g} \tag{2.13}$$

also holds. We can now rewrite the second term on the right hand side of Eq. 2.9 as

$$-i\frac{g}{\cos\theta_W}\left(J^3_\mu - \sin^2\theta_W j^{em}_\mu\right)Z^\mu \equiv -i\frac{g}{\cos\theta_W}J^{NC}_\mu Z^\mu.$$
 (2.14)

The definition

$$J_{\mu}^{NC} = J_{\mu}^{3} - \sin^{2}\theta_{W}j_{\mu}^{em}$$
(2.15)

completes the specification of the two observed neutral interactions in terms of the new currents J and j^{Y} that were introduced. Thus, the $SU(2)_{L} \ge U(1)_{Y}$ formulation of the electroweak theory does contain both the electromagnetic and weak forces with their appropriate properties. The only missing element is the means to arrive at the huge mass splitting between the photon and the weak vector bosons. This is provided by the Higgs mechanism, the topic of the next section.

2.1.1 Spontaneous Symmetry Breaking:

The Higgs Mechanism

As the demands of gauge invariance and renormalizability prohibit the introduction of such terms as $M_W^2 W_\mu W^\mu$ into the theory, another mechanism must be invoked that preserves the desired aspects of the theory. This is provided by introducing a term in the electroweak lagrangian of the form

$$\mathcal{L} = \left| i \left(\partial_{\mu} - g \mathbf{T} \cdot \mathbf{W}_{\mu} - \frac{g'}{2} B_{\mu} \right) \phi \right|_{\mathbf{I}}^{2} - V(\phi) , \qquad (2.16)$$

where there are four scalar fields ϕ_i and the potential $V(\phi)$ is given by

$$V(\phi) = \mu^2 \phi^{\dagger} \phi + \lambda (\phi^{\dagger} \phi)^2 . \qquad (2.17)$$

To preserve the gauge invariance of the lagrangian, the fields ϕ must be members of $SU(2) \ge U(1)$ multiplets. The least complicated choice is that of a Y = 1 isodoublet, which is given by

$$\phi = \begin{pmatrix} (\phi_1 + i\phi_2)/\sqrt{2} \\ (\phi_3 + i\phi_4)/\sqrt{2} \end{pmatrix} .$$
 (2.18)

If we choose the parameters $\mu^2 < 0$ and $\lambda > 0$, the potential $V(\phi)$ possesses an SU(2)-invariant manifold of minima whose locus is given by

$$\phi^{\dagger}\phi \equiv \frac{1}{2}\left(\phi_1^2 + \phi_2^2 + \phi_3^2 + \phi_4^2\right) = -\frac{\mu^2}{2\lambda} \equiv \frac{1}{2}v^2 .$$
 (2.19)

We can now "spontaneously break" the SU(2) symmetry by assuming that the physical vacuum is given by some specific choice of the minimum of $V(\phi)$, e.g.

$$\phi_{1} = \phi_{2} = \phi_{4} = 0, \qquad \phi_{3}^{2} = -\frac{\mu^{2}}{\lambda} \equiv v^{2}.$$
 (2.20)

The vacuum ϕ_0 is then given (from Eq. 2.18) by

$$\phi_0 = \sqrt{\frac{1}{2}} \begin{pmatrix} 0\\ v \end{pmatrix} \quad . \tag{2.21}$$

This choice of vacuum is not SU(2)Invariant, but the lagrangian of Eq. 2.16 still is, as it does not care where the minimum in the vacuum is located.

To examine the consequences of this choice of vacuum, we can consider small oscillations of the fields about this minimum. These fluctuations can be parametrized by the four real fields θ_1 , θ_2 , θ_3 , and h using the form

$$\phi(x) = \frac{e^{i\vec{T}\cdot\vec{\theta}(x)/\nu}}{\sqrt{2}} \begin{pmatrix} 0\\ \nu + h(x) \end{pmatrix}$$
(2.22)

Substitution of this form into Eq. 2.16 results in three massless fields[§] θ_i a term proportional to $2\mu^2 h^2$, and a number of cross terms proportional to $(W^i)_{\mu}\partial^{\mu}\theta_i$. We are free to exploit the SU(2) gauge symmetry by performing a gauge transformation

$$\phi \to \phi' = e^{-i\vec{T} \cdot \vec{\theta}(x)/\nu} \phi , \qquad (2.23)$$

which leaves us with the field

$$\phi(x) = \begin{pmatrix} 0\\ \frac{\nu+h(x)}{\sqrt{2}} \end{pmatrix} .$$
 (2.24)

[§]These are the three Goldstone bosons that result from the spontaneous symmetry breaking[27].

Upon inserting this parametrization into the lagrangian (2.16), we find that that the field h, which we now call the Higgs particle, has acquired a mass $-\frac{\lambda v^2}{2}$, and the gauge bosons W' have acquired masses via the term

$$\left| \left(-ig\mathbf{T} \ \mathbf{W}_{\mu} - i\frac{g'}{2}B_{\mu} \right) \phi \right|^{2} \tag{2.25}$$

which, after squaring and some algebra, yields

$$\left(\frac{1}{2}vg\right)^2 W^+_{\mu} W^{-\mu} + \frac{1}{8}v^2 \left[gW^3_{\mu} - g'B_{\mu}\right]^2 + 0 \left[g'W^3_{\mu} + gB_{\mu}\right]^2 .$$
(2.26)

The masses of the charged and neutral bosons can be read off from this expression by identifying the coefficients of the various terms with $M_W^2 W^+ W^-$, $\frac{1}{2} M_Z^2 Z^2$ and $\frac{1}{2} M_A^2 A^2$, respectively. This gives

$$M_W = \frac{1}{2}vg \tag{2.27}$$

and, upon normalization of the fields,

$$A_{\mu} = \frac{g' W_{\mu}^3 + g B_{\mu}}{\sqrt{g^2 + {g'}^2}} \quad \text{with} \quad M_A = 0 \tag{2.28}$$

$$Z_{\mu} = \frac{gW_{\mu}^3 - g'B_{\mu}}{\sqrt{g^2 + {g'}^2}}$$
 with $M_Z = \frac{1}{2}\sqrt{g^2 + {g'}^2}$. (2.29)

Thus, we have arrived at a theory which contains three massive and one massless vector boson. Using the information from Eq. 2.12, we see that the Higgs mechanism has given us a new relationship between the W and Z masses:

$$\frac{M_W}{M_Z} = \frac{g}{\sqrt{g^2 + {g'}^2}} = \cos\theta_W \;. \tag{2.30}$$

This prediction can be tested experimentally to verify that the Higgs mechanism is responsible for the symmetry breaking of the electroweak theory.

2.2 Specifying the Standard Model

As we have just seen, the three parameters g (the SU(2) coupling constant), g' (the U(1) coupling constant), and $\langle phi_0 \rangle$ or equivalently, v (the Higgs vacuum expectation

value) suffice to specify all of the tree level processes that occur in the standard model. However, none of these parameters are experimentally accessible. Instead, we can define three combinations of them which can be measured experimentally. First, from Eq. 2.12, we have

$$\alpha \equiv \frac{e^2}{4\pi} = \frac{gg'^2}{4\pi\sqrt{g^2 + g'^2}} \,. \tag{2.31}$$

Next, we take from the low energy behaviour of the weak interaction the Fermi constant G_F , which is given by the charged weak current at $q^2 \rightarrow 0$. In this case, we can rewrite the lagrangiaii as

$$\mathcal{L}_{CC} = -i\frac{g}{\sqrt{2}} \left(J^{\mu}W^{+}_{\mu} + J^{\mu\dagger}W^{-}_{\mu} \right) . \qquad (2.32)$$

In the low q^2 limit, this leads to the Fermi four-point amplitude

$$\mathcal{M}_{CC} = \left(\frac{g}{\sqrt{2}}J_{\mu}\right) \left(\frac{1}{M_{W}^{2}}\right) \left(\frac{g}{\sqrt{2}}J^{\mu\dagger}\right) \equiv \frac{4G_{F}}{\sqrt{2}}J_{\mu}J^{\mu\dagger} , \qquad (2.33)$$

whence we can extract the relation

$$G_F = \frac{g^2}{4\sqrt{2}M_W^2} = \frac{1}{2\sqrt{2}\langle\phi_0\rangle^2} .$$
 (2.34)

Finally, we have

$$M_Z = \frac{1}{\sqrt{2}} \langle \phi_0 \rangle \sqrt{g^2 + {g'}^2} . \qquad (2.35)$$

The values of these parameters are shown in Table 2.2 Any other measurable combination of these three parameters, such as

$$\sin^2 \theta_W = \frac{g^2}{g^2 + {g'}^2} , \qquad (2.36)$$

can be compared against its predicted value using the above three parameters to test the Standard Model.

At higher order, of course, the unspecified fermion masses and Higgs mass and couplings enter into the theoretical predictions in the form of radiative corrections. This has the effect of modifying the observables in such a way as to allow experimental access, given the Standard Model calculations as input, to the as-yet-unobserved Higgs, and allows the verification of the higher-order behaviour of the theory (or, hopefully, to some violation of the theoretical predictions!).

Table 2.2: The parameters which specify the electroweak Standard Model at tree level

Parameter	Measured Value	Precision
$lpha G_F \ M_Z$		4.5 x 10 ⁻⁸ 1.7 x 10 ⁻⁵ 7.7 x 10 ⁻⁵

2.3 Z^0 Production and Decay

Since 1989, we have had direct access to the weak interactions by studying the production and decay of Z^0 bosons in the pristine environment of e^+e^- collisions. Since the Z^0 couples to all known fermions, the detailed vertex structure of the electroweak interaction can be studied through measuring the Z^0 coupling to the initial and final state particles.

To aid in our discussion of physics at the Z^0 pole, we can rewrite the weak neutral current

$$-i\frac{g}{\cos\theta_W}\left(J^3_\mu - \sin^2\theta_W j^{em}_\mu\right)Z^\mu \tag{2.37}$$

in terms of the quantum numbers of the electroweak interaction, T^3 and Q (cf. Eq. 2.10)

$$-i\frac{g}{\cos\theta_W}\bar{\psi}_f\gamma^{\mu}\left[\frac{1}{2}(1-\gamma^5)T_f^3-\sin^2\theta_W Q_{f_{\rm I}}\psi_f Z_{\mu}\right]. \tag{2.38}$$

We can express the vertex factor in terms of the vector (v_f) and axial vector (a_i) couplings to the Z^0

$$-i\frac{g}{\cos\theta_W}\gamma^{\mu}\frac{1}{2}\left(v_f - a_f\gamma^5\right),\qquad(2.39)$$

where

$$v_f = T_f^3 - 2\sin^2\theta_W Q_f$$

 $a_f = T_f^3$. (2.40)

Here, Q_f and T_f^3 are the charge and the third component of weak isospin for the fermion f. For completeness, the values of the vector and axial vector couplings for each of the fermions is listed in Table 2.3.

Table 2.3: The vector and axial vector couplings for the fermions to the Z^0 .

Fermion	a_f	v_f
ν_e, ν_μ, ν_τ	$\frac{1}{2}$	12
e, μ , $ au$	$-\frac{1}{2}$	$-\frac{1}{2}$ + $2\sin^2 \theta_W$
$u,\ c,\ t$	$\frac{1}{2}$	$\frac{1}{2} - \frac{4}{3}\sin^2 heta_W$
d, s,	$b -\frac{1}{2}$	$-\frac{1}{2}+\frac{2}{3}\sin^2\theta_W$

In Figure 2.1, we shown the total cross section for $e^+e^- \rightarrow \mu^+\mu^-$. The large resonance around the Z^0 pole is quite obvious. Since it dominates the photon-mediated processes by a factor of ~ 800 , we can safely neglect the effects of the electromagnetic interaction when discussing processes that occur at the Z^0 pole. In addition, at the pole the effects of $\gamma - Z^0$ interference vanish or become negligible, so WC are left to consider the electroweak current as consisting of pure Z^0 exchange. As will be discussed in Chapter 4, the SLC is capable of producing a longitudinally polarized electron beam, which, as we shall see momentarily, allows us to study the chiral structure of the electroweak interaction in more detail than previously possible.

The differential cross section for $e^+e^- \rightarrow f\bar{f}$ at the Z^0 pole can be expressed as[31]

$$\frac{d\sigma}{d\Omega}\Big|_{Z} = \frac{\alpha^{2}}{4\sin^{4}2\theta_{W}(s-M_{Z}^{2})^{2}+\Gamma_{Z}^{2}s^{2}/M_{Z}^{2}} \cdot \\ \left[(1+\cos^{2}\theta)\left(v_{e}^{2}+a_{e}^{2}\right)-8\cos\theta v_{e}a_{e}v_{f}a_{\pi}\right. \\ \left.-P_{z}^{-}\left[2\left(1\,\cos^{2}\theta\right)v_{e}a_{e}\left(v_{f}^{2}\,a_{f}^{2}\right),4\cos\theta\left(v_{e}^{2}\,a_{e}^{2}\right)v_{f}a_{f}\right](2.41) \right]$$

Here, P_z^- is the longitudinal polarization of the electron beam, where $P_z^- = 1(-1)$ corresponds to right- (left-) handed electrons, s is the square of the center-of-mass collision energy, and the width of the Z^0 resonance is specified by Γ_z . The angle θ is the polar angle of the outgoing fermion with respect to the electron beam direction. We have neglected in this representation extra terms that. could arise if the positron beam is polarized, and terms that are due to transverse beam polarization, since neither apply in our particular case. Upon inspection of this formula, it can be



Figure 2.1: The total cross section for $e^+e^- \rightarrow \mu^+\mu^-$ vs. center-of-mass energy

seen that both the polarization dependent and the polarization-independent parts of this expression contain polar-angle anti-symmetric terms, which implies that all Z^0 decays will exhibit fermion production asymmetries. As we shall see below, these are enhanced for the polarized case.

In general, the effects of $\gamma - Z$ interference can be included by adding the supplementary cross section

$$\frac{d\sigma}{d\Omega}\Big|_{\gamma Z} = -2Q_f \left(1 - \frac{M_Z^2}{s}\right) \frac{s}{(s - M_Z^2)^2 + \Gamma_Z^2 s^2/M_Z^2} \cdot (2.42)$$
$$\left\{ \left[(1 \cdot \cos^2\theta) v_e v_f \cdot 2\cos\theta a_e a_f \right] - P_z^- \left[(1 \cdot \cos^2\theta) a_e v_f \cdot 2\cos\theta v_e \epsilon \right] \right\} \cdot (2.42)$$

As mentioned above, this contribution vanishes at the pole, where $M_Z^2 = s$. Since the energy dependence of this cross section is different from that of the pure Z^0 exchange, however, the total fermion production cross section and any fermion production asymmetries depend on the center-of-mass energy. In considering these effects on the final state fermion production asymmetries, we will see that they are small.

2.4 Renormalization of the Electroweak Coupling

In this section, we consider the radiative corrections to the electroweak coupling by examining the higher-order diagrams that modify the tree-level coupling. A general discussion of electroweak radiative corrections can be found in Ref. [32]. We divide the discussion of the virtual corrections as in Ref. [32] into "soft" corrections, which can be considered as responsible for "smearing" the Z^0 line shape, and "hard" corrections which determine the higher-order "unsmeared" resonance through the renormalization of the clectroweak parameters. Since a proper exposition of the renormalization of the full electroweak theory would be too lengthy, we present merely a sketch of the procedure here.

2.4.1 Soft Radiative Corrections

The predominant soft radiative corrections occur to $\mathcal{O}(\alpha)$ and are due primarily to initial state radiation (ISR). The diagrams that contribute to ISR at the Z^0 pole are shown in Figure 2.2. Their effect is to modify the shape of the Z^0 resonance as a function of center-of-mass energy. The relative sign of this effect is easily understood: since ISR is a relatively probably occurance, some fraction of collisions that occur when the colliding beams are at the Z^0 pole will take place at a reduced center-ofmass energy, where the Z^0 production cross section is lower. Thus, the peak cross section is reduced by initial state radiation. In addition, those collisions which occur



Figure 2.2: The leading-order initial state QED corrections to the Z^0 resonance shape.



Figure 2.3: Variation of the Z^0 resonance shape from ISR (see text).

at a center-of-mass energy above the Z^0 pole can radiate down to the pole, where the cross-section is larger. This results in an overall upward shift of the observed peak cross section, as well as an asymmetry in the shape of the resonance, with the high energy side being larger than the low energy side[¶]. A comparison of the treelevel cross section with two higher-order calculations is shown *in* Figure 2.3. The Bonneau-Martin calculation[33] is a fixed-order calculation performed at $\mathcal{O}(\alpha)$. The Fadin-Kuraev result[34] is obtained by evolving the radiating beam particles down to a lower energy by a process similar to the Altarelli-Parisi equations of QCD^{||} and summing the contributions of the leading terms in the splitting kernel to $\mathcal{O}(\alpha^2)$. One can see that the reduction in the peak cross section is quite large, ~ 30%; consideration of these effects is crucial in making precision electroweak measurements at the Z^0 pole. CHAPTER 2. THE ELECTROWEAK INTERACTION: INTRODUCTION 18



Figure 2.4: A schematic view of the different electroweak radiative corrections, including (a) initial and final state bremsstrahhmg, (b) vertex corrections, (c) box diagrams, and (d) self-energy ("oblique") corrections.

2.4.2 Hard Radiative Corrections

A complete consideration of the radiative corrections to the electroweak lagrangian should include the effects of all of the diagrams shown schematically in Figure 2.4. The important effects of the initial-state bremsstrahlung graphs of Fig. 2.4a were discussed in the previous section. We turn now to an overview of the other corrections which modify the electroweak parameters. The most important of these to our present discussion are the oblique corrections of Fig. 2.4d, which provide the primary contributions to the renormalized couplings **and** are common to all electroweak interactions. The vertex corrections (Fig. 2.4b) are process-specific in the sense that they depend on the masses and flavors of the external fermions. The box corrections (Fig. 2.4c) are in most cases negligible[35] and will not be considered further. We will follow the notation of Hollik[36] and Bardin[37] while keeping in mind the organizational simplicity of Kennedy and Lynn[38].

Renormalization of the Electromagnetic Coupling

As an example of the renormalization procedure, we consider first the renormalization of the electromagnetic coupling cr. By gauge invariance*, the only contribution to the renormalization of the electric charge comes from the photon self-energy diagrams,

The e^+e^- pair can always radiate down to the Z^0 peak. See Section 3.1.4.

^{&#}x27;See Section 3.4.1 for a discussion of the Ward-Takahashi identity.



Figure 2.5: A schematic view of (a) the photon self-energy correction and the definition of the vertex factors and (b) the infinite series of one-particle irreducible diagrams that contribute to the renormalization of the electric charge.

shown schematically in Figure 2.5. It is convenient to define the function

$$\Pi'_{QQ}(q^2) \equiv \frac{\Pi_{QQ}}{q^2} , \qquad (2.43)$$

where Π_{QQ} is the coefficient of $g^{\mu\nu}$ in the photon self-energy in Fig. 2.5a. If the oneparticle irreducible diagrams of Fig. 2.5b are summed to all orders[39], one obtains

$$\frac{-ie_0^2}{q^2} \left(1 + e_0^2 \Pi_{QQ} \frac{1}{q^2} + \dots \right) g^{\mu\nu} = \frac{-ig^{\mu\nu}}{q^2} \cdot \frac{e_0^2}{1 - e_0^2 \Pi'_{QQ}}$$
(2.44)

as the expression for the complete photon propagator, where e_0 is the bare electric charge. This suggests the definition of an effective charge $e_{eff} \equiv$ e given by

$$e(q^2) = \frac{e_0^2}{1 - e_0^2 \Pi'_{QQ}} , \qquad (2.45)$$

where now the coupling e depends on the q^2 of the interaction, i.e., it "runs". The value of e at $q^2 = 0$ is chosen as the reference coupling; this can be determined from the g - 2 of the electron or from other measurements such as the Josephson effect. This allows us to define the familiar quantity α by $4\pi\alpha = e^2(q^2 = 0)$, where the value of α from this definition is given in Table 2.2. From this definition, the running of a is given by

$$\frac{1}{4\pi\alpha(q^2)} = \frac{1}{4\pi\alpha} - \left[\Pi'_{QQ}(q^2) - \Pi'_{QQ}(0)\right]$$
(2.46)

Though the quantity Π'_{QQ} diverges, the difference $\Pi'_{QQ}(q^2) - \Pi'_{QQ}(0)$ is finite. Of course, all of the difficulties involved in renormalizing the couplings are involved in actually calculating the quantity Π'_{QQ} in a convenient gauge. We will omit a discussion of that here, and merely mention that the value of $\alpha(M_Z^2)$ is necessary for the comparison of experimental results obtained at the Z^0 pole with theoretical predictions. This quantity can be obtained from existing measurements of $e^+e^- \rightarrow f\bar{f}$ by a theoretical argument[40] using the optical theorem to relate the hadronic corrections to forward Bhabha scattering to the total cross section for $e^+e^- \rightarrow$ hadrons. A recent result of this procedure is[41]

$$\alpha^{-1}(M_Z^2) = 129.08 \pm 0.10$$
, (2.47)

where the dominant uncertainty is from the imprecise measurements of the total cross section for $e^+e^- \rightarrow$ hadrons at Iow-energy (2 GeV < E_{CM} < 3.6 GeV).

Renormalization of the Weak Couplings

Because of the number of different diagrams involved, a discussion of the renormalization of the weak couplings becomes quite complicated. As a summary of the types of diagrams involved, we show in Figures 2.G ant! 2.7 some of the contributing processes. Their effect on the couplings is most easily expressed in terms of the induced modification to the form of the neutral current vertex[†]. If we rewrite the weak neutral current in terms the parameters chosen to specify the theory, we obtain

$$-iM_{Z}\left[4\sqrt{2}G_{F}\right]^{\frac{1}{2}}\gamma^{\mu}\left[\frac{1}{2}(1-\gamma^{5})T_{f}^{3}-\sin^{2}\theta_{W}\Big|_{S}Q_{f}\right],\qquad(2.48)$$

where, for notational convenience, we have used the Sirlin definition of $\sin^2 \theta_W[42]$:

$$\sin^2 \theta_W \Big|_S \equiv 1 - \frac{M_W^2}{M_Z^2} \tag{2.49}$$

which elevates the tree-level relation of Eq. 2.30 to a definition which will be taken as valid to all orders in perturbation theory.

 $^{^{\}dagger}As$ it is not relevant to our purposes, we will omit here any discussion of the charged weak current.

$$\gamma \cdots \gamma = i \quad e^2 \Pi_{QQ} g^{\mu\nu} + \cdots$$

$$Z \cdots \gamma = i \quad \frac{e^2}{cs} (\Pi_{3Q} - s^2 \Pi_{QQ}) g^{\mu\nu} + \cdots$$

$$Z \cdots \gamma Z = i \quad \frac{\partial}{c^2 s^{22}} (\Pi_{33} 2 s^2 \Pi_{3Q} + s^4 \Pi_{QQ}) g^{\mu\nu} + \cdots$$

$$W \cdots \Psi = i \quad \frac{e^2}{s^2} \Pi_{11} g^{\mu\nu} + \cdots$$

Figure 2.6: A schematic view of self-energy diagrams that occur in calculating radiative corrections to the weak interaction and the definition of the vertex factors involved. Here, $s^2(c^2) = \sin^2 \theta_W(\cos^2 \theta_W)$.



Figure 2.7: A schematic view of self-energy diagrams that occur in calculating radiative corrections to the weak interaction and the definition of the vertex factors involved. Here, $s^2(c^2) = \sin^2 \theta_W(\cos^2 \theta_W)$.

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When the effects of the one-loop corrections are included, the form of the current from Eq. 2.48 can be retained with the addition of an overall normalization factor ρ_f and an additional vertex form factor κ_f , where the current is now given by[‡]

$$-iM_Z \left[4\sqrt{2}G_F \rho_f\right]^{\frac{1}{2}} \gamma^{\mu} \left[\frac{1}{2} \left(1-\gamma^5\right) T_f^3 - \kappa_f \sin^2 \theta_W \Big|_S Q_f\right] . \tag{2.50}$$

Note that, as defined, the form factors ρ_f and κ_f are fermion-flavor dependent. They can, however, be separated into universal parts (due to the boson self-energies) and non-universal parts (from the vertex corrections and self-energies of the external fermion lines) as follows:

$$\rho_f = 1 + (\Delta \rho)_{univ} + (\Delta \rho)_{non-univ}$$

$$\kappa_f = 1 + (\Delta \kappa)_{univ} + (\Delta \kappa)_{non-univ}.$$
(2.51)

The universal parts of the form factors contain essentially all of the the dependence of the couplings on the top quark mass (m_t) and Higgs mass, except in the case of b quarks, where large vertex corrections are important. In general, with the preceding exception, the universal portions of the form factors are much larger than the non-universal contributions. To lowest order, we have

$$1 + (\Delta \rho)_{univ} = \frac{1}{1 - \Delta \bar{\rho}} + \dots$$
$$(\Delta \kappa)_{univ} = \frac{\cos^2 \theta_W |_S}{\sin^2 \theta_W |_S} \Delta \bar{\rho} , \qquad (2.52)$$

where [43]

$$\Delta \bar{\rho} = N \frac{G}{8\pi^2 \sqrt{2}} \left[1 + \frac{G_F m_t^2}{8\pi^2 \sqrt{2}} (19 - 2\pi^2) \right]$$
(2.53)

is the result of calculating the correction including 2-loop irreducible diagrams.

The form of Equation 2.50 suggests that we define an "effective" renormalized version of sir? θ_W given by

$$\sin^2 \theta_W^{eff} \equiv \kappa_f \sin^2 \theta_W |_S \equiv \sin^2 \theta_W \Big|_{univ} + (\Delta \kappa)_{non-univ} \sin^2 \theta_W \Big|_S \,. \tag{2.54}$$

 $^{^{\}ddagger}$ this factorization of the corrections is only valid in the case where the **box** diagrams of Fig. 2.4c arc neglected.



Figure 2.8: The dependence of $\sin^2 \theta_W^{eff}$ on m_t and m_{Higgs} .

With this definition, we can rewrite Eq. 2.50 as

$$-iM_{Z}\left[4\sqrt{2}G_{F}\rho_{f}\right]^{\frac{1}{2}}\gamma^{\mu}\left[\frac{1}{2}\left(1-\gamma^{5}\right)T_{f}^{3}-\sin^{2}\theta_{W}^{eff}Q_{f}\right]$$
(2.55)

Figure 2.8 shows the dependence of $\sin^2 \theta_W^{eff}$ on the top and Higgs masses[§]. We note in passing that the "universal" weak mixing angle defined in Eq. 2.54 is the same as the s_*^2 of Kennedy and Lynn[38], which is constructed to be independent of the vertex corrections.

Now, we are able to continue our discussion of electroweak physics at the Z^0 pole, including the leading order corrections as we proceed by replacing the couplings in the tree-level expressions by their renormalized counterparts.

2.5 Electroweak Asymmetries: Testing the Standard Model

As mentioned above, the measurement of a different combination of the parameters q, q', and $\langle \phi_0 \rangle$ (or equivalently α , G_F , and M_Z) which specify the weak interaction can serve as a test of the validity of the theory. Different combinations of these parameters have different sensitivity to the various radiative corrections introduced by physical processes beyond the Standard Model. Thus, if a difference between prediction and experiment is seen, a comparison of the various observables may yield some indication of its origin. Since the goal of these measurements is precision in determining the electroweak parameters, every attempt has been made to eliminate sources of experimental error. One powerful technique to accomplish this is to form ratios of cross sections where the parameter of interest has been isolated by the angular- or polarization-dependence of the differential cross section. This has the advantage that the errors inherent in determining an absolute cross section, such as those due to luminosity normalization and absolute detection efficiency, cancel in the asymmetry which is formed. Using this method allows the isolation and examination of various combinations of the electroweak coupling. Several of these asymmetries will be discussed below.

Before we proceed, however, it will be useful to rewrite the Born cross section for fermion production at the Z^0 pole (Eq. 2.41) as follows:

$$\frac{d\sigma}{d\Omega}\Big|_{Z} = K\Big(v_e^2, a_e^2\Big)\Big(v_f^2, a_f^2\Big) \times \\ (1 - A_e P_e)(1, \cos^2\theta) + 2A_f(A_e - P_e)\cos\theta, \qquad (2.56)$$

with

$$K = \frac{\alpha^2}{4\sin^4 2\theta_W (s - M_Z^2)^2 + \Gamma_Z^2 s^2 / M_Z^2}$$
(2.57)

where we have factored out the term $(v_e^2 + a_e^2) (v_f^2 + a_f^2)$ from the previous expression and have defined the quantity

$$A_f = \frac{2v_f a_f}{v_f^2 + a_f^2}$$
(2.58)

[§]This dependence was calculated using the program ZSHAPE[44], without the value of $\alpha^{-1}(M_Z^2)$ given in Eq. 2.47.

for each species of fermion f. The difference of A_f from zero is a measure of the parity violation in the Zff coupling. The signed longitudinal polarization of the electron beam will henceforth be referred to as P_e .

2.5.1 The Left-Right Asymmetry

Perhaps the simplest example of an asymmetry which can be formed in this manner is the Left-Right Asymmetry, denoted $A_{LR}[45]$. At the SLC, the longitudinal polarization of the electron beam allows us to examine the parity violation in Z^0 production directly by forming the asymmetry

$$A_{LR} = \frac{\sigma(e_R^+ e_L^- \to Z^0 \to f\bar{f}) - \sigma(e_L^+ e_R^- \to Z^0 \to f\bar{f})}{\sigma(e_R^+ e_L^- \to Z^0 \to f\bar{f}) + \sigma(e_L^+ e_R^- \to Z^0 \to f\bar{f})},$$
(2.59)

where f is any outgoing fermion. More specifically, from the expression from Eqs. 2.56 and 2.40 with $P_e = 1$,

$$A_{LR} = A, = \frac{2v_e a_e}{v_e^{2} + a_e^{2}} = \frac{2\left[1 - 4\sin^2\theta_W^{eff}(M_Z^2)\right]}{1 + \left[1 - 4\sin^2\theta_W^{eff}(M_Z^2)\right]^2}.$$
 (2.60)

Here, the dependence on the renormalized electrowcak parameters has been introduced by replacing $\sin^2 \theta_W$ in Eq. 2.40 with $\sin^2 \theta_W^{eff}$. For arbitrary beam polarization with $P_e < 1$,

$$A_{LR} = \frac{1}{P_e} A_m = \frac{1}{P_e} \frac{N_L - N_R}{N_L + N_R}$$
(2.61)

where the quantity $A_{,,,}$ is the measured asymmetry, given by the difference between the number of Z^0 events produced with a left-handctl electron beam (N_L) and the number produced with a right-handed electron beam (N_R) over the total. As can be seen from this expression, the higher the beam polarization, the larger the asymmetry. For $P_e \sim 1$, the expected value of A_{LR} is about 0.16.

The quantity A_{LR} is uniquely suited for a precision determination of the value of $\sin^2 \ell_W^{eff}$ due to several important properties. First, as can be seen from Eq. 2.60, there is no dependence on the final state fermion couplings. This implies that all final state fermions can be used[¶] to measure A_{LR} , which greatly increases the statistical

power of the measurement. This also implies that, if symmetric acceptance limits are chosen, A_{LR} is insensitive to any final state radiative corrections[46], such as those due to QCD. In addition, as long as the local detector efficiency for observing fermions and antifermions is the same, A_{LR} is independent of detector acceptance.

We give the sensitivity of A_{LR} to $\sin^2 \theta_W^{e'f}$ here as a benchmark against which we can compare the other asymmetries:

$$\frac{dA_{LR}}{d\sin^2\theta_W} = -7.84 \tag{2.62}$$

which implies that

$$\delta \sin^2 \theta_W^{eff} = \delta A_{LR} / 7.84 . \qquad (2.63)$$

The sensitivities of all of the electroweak asymmetries are given in Table 2.4.

The SLD Collaboration has performed two measurements of A_{LR} based on running during 1992 and 1993. The combined results for A_{LR} and for $\sin^2 \theta_W^{eff}$ after radiative corrections have been applied are[47]:

$$A_{LR}^0 = 0.1637 \pm 0.0075$$

 $\sin^2 \theta_W^{eff} = 0.2294 \pm 0.0010$.

2.5.2 The au Polarization Asymmetry

The angular correlations from the decays of the final state fermions can also be analyzed to provide information on the electroweak couplings. This measurement is most easily made in τ decays, as it is the only fermion which can be "spin-analyzed" in a convenient manner, i.e., its charge is readily identified and it often decays two a two-body final state whose helicity components can be measured. Since the natural polarization of the final state fermion is used in this analysis, longitudinal beam polarization is not necessary. We can form the final state polarization asymmetry between the production of left-handed (τ_L) and right-handed (τ_R) taus:

$$\mathcal{P}_{\tau} \equiv \frac{\frac{d\sigma}{d\Omega}(\tau_L) - \frac{d\sigma}{d\Omega}(\tau_R)}{\frac{d\sigma}{d\Omega}(\tau_L) + \frac{d\sigma}{d\Omega}(\tau_R)}, \qquad (2.64)$$

[¶]Because of contributions from t-channel Bhabha scattering, the e^+e^- final state is usually excluded.

where the τ helicity is obtained by examining its decay. If we insert the unpolarized fermion production cross section, we find

$$\mathcal{P}_{\tau} = \frac{2A_e \cos\theta + A_f (1 + \cos^2\theta)}{(1 + \cos^2\theta) + 2A_e A_f \cos\theta} \tag{2.65}$$

This asymmetry has two interesting properties. First, if it is integrated over a $\cos \theta$ -symmetric acceptance, the dependence on the initial state coupling drops out, leaving

$$\langle \mathcal{P}_{\tau} \rangle = A_{\tau} \tag{2.66}$$

Since the electroweak quantum numbers of the τ and the electron are identical, this quantity has the same sensitivity to $\sin^2 \theta_W^{eff}$ as A_{LR} , namely

$$\frac{d\langle \mathcal{P}_{\tau} \rangle}{d\sin^2 \theta_W} = -7.84 \ . \tag{2.67}$$

Second, by measuring the forward-backward asymmetry in \mathcal{P}_{τ} , one can isolate $A_{,\tau}$, which provides yet another way to access $\sin \frac{eff_{I}}{eW}$ Thus the τ polarization asymmetry is a remarkable tool with which to probe the electroweak sector. Unfortunately, the fraction of Z^0 decays which produce a $\tau^+\tau^-$ pair is only 3%, and less than half of these events contain a τ decay mode suitable for spin analysis. In addition, the analyzing power of the helicity analysis is not unity, which further dilutes the measured asymmetry. All told, each Z^0 event has approximately 200 times less statistical weight for the τ polarization asymmetry than is does for A_{LR} . Comparisons of the statistical weight for each of the asymmetries are also summarized in Table 2.4.

2.5.3 The Forward-Backward Asymmetry

The fermion forward-backward asymmetry, which one can measure at any e^+e^- experiment, provides information on the initial and final state couplings of the electroweak interaction. This method has been used at lower center-of-mass energies[48] to measure the extent of $\gamma - Z$ interference. At the Z^0 pole, it allows direct access to the fermion-Z couplings through the parity-violation present in any weak interaction. We define the forward-backward asymmetry for a fermion of type f as

$$A_{FB}^{f} = -\frac{\sigma^{f}(\cos\theta > 0) - \sigma^{f}(\cos\theta < 0)}{\sigma^{f}(\cos\theta > 0) - \sigma^{f}(\cos\theta < 0)} = \frac{3}{4} \frac{2v_{e}a_{e}}{v_{e}^{2} + a_{e}^{2}} \frac{2v_{f}a_{f}}{v_{f}^{2} + a_{f}^{2}} = \frac{3}{4} A_{e}A_{f}.$$
 (2.68)

Table 2.4: A comparison of the different observable fermion production asymmetries that can be measured at the Z^0 pole. The relative statistical sensitivities of each measurement (last column) are calculated by assuming reasonable selection efficiencies for each of the analyses. Those measurements using *b* quarks are assumed to have been performed with a lepton tag. The magnitude of the longitudinal beam polarization is assumed to be 63%. The size of the measured asymmetry has not been taken into account.

Parameter	Couplings	Raw Asymmetry	$dA/d\sin^2\theta_W$	Rel. Stat. Sens.
A_{LR}	A_e	$ P_e A_e \sim 0.1$	-7.84	1
\mathcal{P}_{r}	$A_{ au}$	A, ~ 0.16	-7.84	0.0075
A^f_{FB}	$A_e A_b$	$\frac{3}{4}A_eA_b \sim 0.11$	-5.59	0.015
	$A_e A_c$	$rac{3}{4}A_eA_c\sim 0.08$	-4.39	0.001
	$A_e A_\ell$	$\frac{3}{4}A_eA_\ell \sim 0.02$	-1.92	0.017
\tilde{A}^{f}_{FB}	A_b	$\frac{3}{4} P_e A_b \sim 0.44$	-0.63	0.0018
	A_c	$\frac{3}{4} P_e A_c \sim 0.32$	-3.42	0.0008
	A	$\frac{3}{4} P_e A_\ell \sim 0.08$	-7.84	0.06

If we take into account the symmetric acceptance of a real particle detector which extends to some $\cos \theta_0$, this expression acquires a $\cos \theta$ dependence:

$$A_{FB}^{f} = \frac{\int_{0}^{\cos\theta_{0}} d\cos\theta \frac{d\sigma_{f}}{d\Omega} - \int_{-\cos\theta_{0}}^{\theta} d\cos\theta \frac{d\sigma_{f}}{d\Omega}}{\int_{-\cos\theta_{0}}^{\cos\theta_{0}} d\cos\theta \frac{d\sigma_{f}}{d\Omega}}$$
(2.69)

which yields, upon integration,

$$A_{FB}^{f} = \frac{3}{4} \frac{4\cos\theta_{0}}{3 + \cos\theta_{0}^{2}} A_{e} A_{f} . \qquad (2.70)$$

Since in the case of A_{FB}^{f} the product $A_{e}A_{f}$ is measured, one cannot disentangle the initial and final state couplings. The experiments at LEP have used these asymmetries as a method to derive information on $\sin^{2}\theta_{W}^{eff}$ by using the full $\sin^{2}\theta_{W}^{eff}$ dependence of both the initial and final state couplings. As will be seen below, the quantity A_{f}

is large for quark final states, which implies that the statistical sensitivity of this method to $\sin^2 \theta_W^{eff}$ can be large if one can efficiently select samples of quarks whose charge and production angle can be measured. In terms of raw sensitivity to $\sin^2 \theta_W^{eff}$ we have

$$\frac{dA_{FB}^{d}}{d\sin^{2}\theta_{W}} = -5.58$$

$$\frac{dA_{FB}^{u}}{d\sin^{2}\theta_{W}} = -4.39$$

$$\frac{dA_{FB}^{\ell}}{d\sin^{2}\theta_{W}} = -1.92$$

for d-type and u-type quarks, and leptons, respectively. The simplest decay channel in which to measure A_{FB}^{f} is $Z \rightarrow \ell \bar{\ell}$. The small branching fraction of the Z^{0} to leptons, coupled with the small sensitivity to $\sin^{2} \theta_{W}^{eff}$ and the small asymmetry size (~ 2%) limit the statistical sensitivity of this method. The asymmetry in the case of 6 quarks is large (~ 11%), and it also has a large sensitivity to $\sin^{2} \theta_{W}^{eff}$; the use of this channel is only limited by the ability to correctly determine the production direction of the 6 (as opposed to \bar{b}) quark. Techniques for accomplishing this determination will be discussed in Section 2.8, below.

By alternating the longitudinal polarization of the electron beam, we can form a double asymmetry, the so-called left-right forward-backward asymmetry $\tilde{A}_{FB}^{f}[49]$:

$$\tilde{A}_{FB}^{f} = \frac{\sigma_{F}^{L} + \sigma_{B}^{R} - \left(\sigma_{F}^{R} + \sigma_{B}^{L}\right)}{\sigma_{F}^{L} + \sigma_{B}^{R} + \sigma_{F}^{R} + \sigma_{D}^{L} + \sigma_{B}^{R}} = \frac{3}{4} \left|P_{e}\right| A_{f}$$

$$(2.71)$$

In the above expression, σ_F^L denotes the cross section for producing the fermion f in the forward hemisphere ($\cos \theta > 0$) with a left-handed electron beam. This shows the experimental advantage in using a longitudinally polarized beam: by forming the double asymmetry, the Z^0 production asymmetry caused by the parity violation in the initial state coupling has been removed, leaving experimentally accessible the final state coupling. In addition, for substantial net beam polarization, the raw asymmetry \tilde{A}_{FB}^{I} is quite large compared to the standard forward-backward asymmetry A_{FB}^{I} , since the multiplicative factor P_e can, in principle, be ~ 1, whereas $A_e \sim 0.16$. This makes

for a smaller fractional error on the measurement of \tilde{A}_{FB}^{f} compared with that for A_{FB}^{f} for a similar data set.

For completeness, we give here the angular dependence of the asymmetry as derived from Eq. 2.56:

$$\tilde{A}_{FB}^{f} = |P_e| A_f \frac{2\cos\theta}{1 + \cos^2\theta} .$$
(2.72)

This thesis presents the first measurement of A, using exclusively reconstructed charm mesons.

We turn now to a discussion of the radiative corrections that effect the measurement of A_{FB} .

2.6 Radiative Corrections to A_{FB}

We divide the description of the effects of radiative corrections on \tilde{A}_{FB} into two sections: the electroweak corrections, which are small, and the QCD corrections, which can be large and must be considered in detail.

2.6.1 Electroweak Corrections

The primary contribution to the electroweak corrections of \tilde{A}_{FB} (Hereafter denoted A_f , since that is what is actually being measured) arises from $\gamma - Z$ interference, which introduces some center-of-mass energy dependence to A_f . As argued in Refs. [50] and [35], contributions from the weak interact ion vertex corrections and the box diagrams are very small. The oblique correct ions to A_f can be related to those for $A_{LR} = A_e$ by

$$\begin{aligned} (\Delta A_{f})_{oblique} &= (\Delta A_{e})_{oblique} \\ (\Delta A_{u,c})_{oblique} &\simeq \frac{12}{25} (\mathbf{A} A_{e})_{oblique} \\ (\Delta A_{d,s,b})_{oblique} &\simeq \frac{4}{45} (\Delta A_{e})_{oblique} \,. \end{aligned}$$

Those corrections due to QED initial state radiation have the effect of shifting the effective center-of-mass energy. To demonstrate the size of these effects, we show in



Figure 2.9: Electroweak-corrected values of A_b and A, vs. \sqrt{s} near the Z^0 pole. QCD radiative corrections are not included. The tree-level values are $A_{,} = 0.67$ and $A_b = 0.93$ for comparison. These values have been calculated using the ZFITTER package[51].

Figure 2.9 the values of A, and A_b vs. \sqrt{s} near the Z^0 pole with the electroweak radiative corrections applied These values have been calculated using the ZFITTER package[51]. One can see that the variation of these asymmetry parameters with respect to the center-of-mass energy, as well as the magnitude of the corrections, is small, as expected.

2.6.2 QCD Corrections

$\mathcal{O}(\alpha_s)$ Corrections

First order QCD corrections to the heavy quark forward backward asymmetry have been calculated by a number of groups[52, 53]. The full differential cross section for massive quarks including $\mathcal{O}(\alpha_s)$ corrections is given by[52]

$$\frac{d^{2}\sigma_{f}}{d\Omega} = \frac{3}{4} \left(\frac{\alpha}{4\sin^{2}2\theta_{W}}\right)^{2} \frac{\beta}{\Gamma^{2}} \times \left\{ \left(v^{2} + a^{2} - 2avP_{-}\right) \cdot \left[\left(v_{f}^{2} + a_{f}^{2}\right)\left(1 + \frac{\bar{m}_{f}^{2}}{2}\right)\left(1 + \frac{4\alpha_{s}}{3\pi}F_{1}(\bar{m}_{f})\right) - \frac{3}{Z}a_{f}^{2}\bar{m}_{f}^{2}\left(1 + \frac{4\alpha_{s}}{3\pi}F_{4}(\bar{m}_{f})\right)\right](1 + \cos^{2}\theta) + \left(v^{2} + a^{2} - 2avP_{-}\right) \cdot \left[\left(v_{f}^{2} + a_{f}^{2}\right)\left(\frac{\bar{m}_{f}^{2}}{2} + \frac{4\alpha_{s}}{3\pi}F_{2}(\bar{m}_{f})\right) - a_{f}^{2}\frac{\bar{m}_{f}^{2}}{2}\left(1 + \frac{4\alpha_{s}}{3\pi}F_{5}(\bar{m}_{f})\right)\right](1 - 3\cos^{2}\theta) - 4\left(2va - \left(v^{2} + a^{2}\right)P_{-}\right)a_{f}v_{f}\beta\left(1 + \frac{4\alpha_{s}}{3\pi}F_{3}(\bar{m}_{f})\right)\cos\theta\right\}, \quad (2.73)$$

where P- is the longitudinal polarization of the electron beam, and

$$\frac{4\alpha_s}{3\pi}F_1(\bar{m}_f) = \frac{\alpha_s}{\pi}(1 + 3\bar{m}_f^2), \qquad (2.74)$$

$$\frac{4\alpha_s}{3\pi}F_2(\bar{m}_f) = \frac{\alpha_s}{\pi} \frac{2}{3} \left[-\frac{\pi^2}{6}\bar{m}_f - \bar{m}_f^2 \left(\frac{1}{3} + \ln\frac{\bar{m}_f}{2} + \frac{1}{3}\ln^2\frac{\bar{m}_f}{2} \right) \right], \qquad (2.75)$$

$$\frac{4\alpha_s}{3\pi}F_3(\bar{m}_f) = \frac{\alpha_s}{\pi} \left[\frac{8}{3} \,\bar{m}_f + \bar{m}_f^2 \left(\frac{7}{3} + \frac{\pi^2}{18} - \frac{2}{3} \ln \frac{\bar{m}_f}{2} + \frac{1}{3} \ln^2 \frac{\bar{m}_f}{2} \right) \right], (2.77)$$

$$\frac{4\alpha_s}{\pi}F_1(\bar{m}_s) = \frac{\alpha_s}{\pi} \left(3 + 4 \ln \frac{\bar{m}_f}{2} \right)$$
(2.77)

$$\frac{3\pi}{3\pi}F_4(m_f) = \frac{\pi}{\pi} \left(3 + 4 \ln \frac{\pi}{2} \right), \qquad (2.77)$$

$$\frac{4\alpha_s}{3\pi}F_5(\bar{m}_f) = \frac{\alpha_s}{\pi} \left(5+4 \ln \frac{m_f}{2}\right). \tag{2.78}$$

The quantity \bar{m}_f is the scaled heavy quark mass, with velocity $\beta^2 = 1 - \bar{m}_f^2$. The expressions in Equation 2.74 are correct to order \bar{m}_f^2 .

hing the asymmetry
$$\tilde{A}_{FB}$$
 results in

$$\tilde{A}_{FB}(0) = \frac{3}{4}\beta \frac{2va}{v^2 + a^2} 2v_f a_f \left(1 + \frac{4\alpha_s}{3\pi}F_3(\bar{m}_f)\right)$$

$$\times \left\{ \left(v_f^2 + a_f^2\right) \left(1 + \frac{\bar{m}_f^2}{2}\right) \left(1 + \frac{4\alpha_s}{3\pi}F_1(\bar{m}_f)\right)$$

$$\frac{2}{2} \frac{3\bar{m}_f^2}{2} a_{\bar{f}} \left(1 + \frac{4\alpha_s}{3\pi}F_4(\bar{m}_f)\right) \right\}^{-1}.$$
(2.79)

We write it in the form

Form

$$\tilde{A}_{FB}(0) = \tilde{A}_{FB}^{0}(0)(1 - A),$$
 (2.80)

with the exact lowest order asymmetry

$$A_{FB}^{0}(0) = \frac{3}{4}\beta \frac{2av}{v^{2} + a^{2}} \times \frac{2a_{f}v_{f}}{v_{f}^{2} \frac{3-\beta^{2}}{2} + \beta^{2}a_{f}^{2}},$$
(2.81)

and the radiative correction to first order in α_s and second order in \bar{m}_f , given by

$$\Delta = \frac{\alpha_s}{\pi} \left[1 - \frac{8}{3} \bar{m}_f - \frac{\bar{m}_f^2}{3} \left(7 + \frac{\pi^2}{6} - 2\ln \frac{m_f}{2} + \ln^2 \frac{\bar{m}_f}{2} \right) + 3\bar{m}_f^2 \frac{v_f^2 - 2a_f^2 \ln \frac{\bar{m}_f}{2}}{v_f^2 + a_f^2} \right].$$
(2.82)

to $\mathcal{O}(\alpha_s)$ and \bar{m}_f^2 . Note that the correction has the expected sign: gluon radiation must *dilute* the asymmetry, as it tends to smear out the initial quark directions, in most cases smearing them away from the region at high cost9 where A_{FB} is large. This is easily seen in the limit where the radiated gluon is very energetic. In this case, the observed quark directions are essentially unrelated to their directions before gluon radiation, and the quarks can end up anywhere in the detector.

These analytic formulae have been incorporated numerically into the differential cross section that is used in our measurement of $A_{,.}$ In Figure 2.10 we show the effects of the $\mathcal{O}(\alpha_s)$ corrections for massless and massive quarks.

$\mathcal{O}(\alpha_s^2)$ Corrections

Second order calculations for the QCD radiative corrections (with massive quarks) to A_f have been performed[54]. The total corrections have the form

$$A_{c}^{meas}\left(\mathcal{O}(\alpha_{s}^{2})\right) = A_{c}^{0}\left(1 + \left(\frac{\alpha_{s}}{\pi}\right)^{2}\left[-(4.4 \pm 0.4) - (26 \pm 6)f_{D_{B}}\right]\right), \qquad (2.83)$$



Figure 2.10: The $\mathcal{O}(\alpha_s)$ QCD corrections to A_b and A, for massless (solid line) and massive quarks (dotted line). The quantity Δ_{QCD} relates the uncorrected asymmetry A^0 to the measured asymmetry by $A^{"'"} = A^0(1 + \Delta_{QCD})$.

for c quarks, and

$$A_b^{meas}\left(\mathcal{O}(\alpha_s^2)\right) = A_b^0 \left(1 + \left(\frac{\alpha_s}{\pi}\right)^2 \left[-(1.9 \pm 0.4) - (3.5 \pm 0.7)f_{D_B}\right]\right),$$
(2.84)

for b quarks. These corrections are just the next order contributions: they must be applied in addition to the leading order calculation shown above. The second order correction depends in detail on the acceptance of analysis cuts for four-jet events wherein a radiated gluon emerges as a pair of heavy quarks. In principle, in the case where the two primary quarks are one of the light flavors, these events should be treated as QCD corrections to the light quark production cross section. Since the



Figure 2.11: Radiative Corrections to A_b from virtual top quark loops.

heavy quarks can cause the event to be tagged as part of the heavy flavor sample, however, they need to be considered, as they will modify the measured asymmetry by a factor proportional to their probability to pass the event selection criteria. This experimental acceptance is parametrized by the quantity f_{D_B} and must be determined by Monte Carlo simulation. The errors quoted arc due to the ambiguity of what value the quark masses should take in the calculation, where the ranges assumed are $4 \text{ GeV}/c^2 < m_b < 5 \text{ GeV}/c^2$ and $1.1 \text{ GeV}/c^2 < m_c < 1.7 \text{ GeV}/c^2$.

2.7 Possible Modifications to A_f from New Physics

The values of the fermion asymmetries can be modified by any new physics that can change the final state through corrections to the $Zf\bar{f}$ vertex. The value of A_b already receives a correction of this type from the presence of diagrams containing virtual top quarks, as shown in Figure 2.11, whose effects are proportional to $m_t^2/M_W^2|V_{tb}|^2$. The magnitude of these corrections is less than 0.1% for A_b [35], and is negligible in the case of $A_{,,}$ since the corresponding diagrams contain strange quarks instead of top quarks. It is through diagrams such as these, however, that the influence of other, heavy particles can enter^{||}.

The sensitivity of the parameter A_f to modifications of the electroweak coupling

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can be parametrized in terms of the difference of the observed AJ and the standard model prediction A_f^{SM} . If the variable η_f is introduced, where

$$\frac{A_f}{A_f^{SM}} \equiv 1 + \eta_f , \qquad (2.85)$$

variations in the asymmetry parameters can be expressed in terms of changes in the vector (v_f) and axial vector (a_f) couplings of the fermions. A short calculation yields

$$\eta_{f} = \frac{2(1-c_{f}^{2})}{c_{f}(1+c_{f}^{2})} [\delta v_{f} - c_{f} \delta a_{f}]$$

$$= \frac{2(1-c_{f})}{c_{f}} \frac{(1+c_{f})^{2}}{1+c_{f}^{2}} \left[\delta g_{R}^{f} + \left(\frac{1-c}{1+c}\right) \delta g_{L}^{f} \right], \qquad (2.86)$$

where $c_f \equiv v_f/a_f = 1 - 8/3 \sin^2 \theta_W$, and the couplings have been reexpressed as $v_f = g_L^f + g_R^f$ and $a_f = g_L^f - g_R^f$ so as to give expressions independent of the value for $\sin^2 \theta_W^{eff}$. Assuming $\sin^2 \theta_W = 0.2325$, we obtain

$$\eta_{c} \simeq 5.43[\delta_{R}^{c} + 0.45\delta g_{L}^{c}] \eta_{b} \simeq -1.74 [\delta_{R}^{b} + 0.18\delta g_{L}^{b}]$$
(2.87)

This shows that, in fact, A, is more sensitive to deviations of the couplings if the changes were to occur equally for b and c. Since charm quarks have the light strange quark in their weak isospin doublet, however, most effects that arise from new physics in the vertex couplings are markedly smaller for charm quarks when compared to those for b's. In the following sections, we briefly review how some of the potential extensions of the standard model could modify the values of the heavy quark asymmetries A_f , with particular attention to those processes which could result in large modifications of the Zcc coupling.

2.7.1 Two-Higgs Doublet Models and Supersymmetry

When the single higgs doublet of the minimal standard model as described above is extended to two doublets, the single "standard model" higgs particle is replaced by five higgs bosons, each of which can participate in vertex loops and hence, can effect

^{||}The considerations for this section are drawn from Reference [55].

the heavy quark asymmetries[56]. Typically, three neutral higgs h^0 , H^0 , and A^0 , and two charged higgs H^{\pm} , comprise the higgs spectrum of these theories. Theories of this type are parametrized by the ratio of the vacuum expectation values of the two doublets $v_2/v_1 \equiv \tan \beta$. Charged higgs and (in supersymmetry) chargino exchange in $Z \rightarrow b\bar{b}$ decays can have potentially large effects on the effective $Zb\bar{b}$ coupling[57]; these modifications are suppressed by $|V_{cb}|^2 < 2 \times 10^{-3}$ for charm. Corrections from neutral higgs exchange can be proportional to m_f , and thus are small as well. Other contributions from neutralino exchange and squark contributions in supersymmetric theory are also expected to be small for charm.

Alternate models with two higgs doublets can contain flavor-changing neutral currents[58], where the effects are largest in the t - c - higgs vertex. This is advantageous, as it allows a direct coupling of c to the "heavy" physics opened up by m_t . Depending on the strength of the Yukawa coupling in these theories, large (> 5% in $\delta A_c/A_c$) could be observed.

2.7.2 Extended Gauge Sectors

Of the various models for extending the gauge structure in which the Standard Model is embedded in order to achieve a unification of couplings, the E(6) models[59] are interesting to our present discussion because they can contain "Alternate Left-Right" models that effectively switch the quantum numbers of the ordinary fermions[60]. This leads to large vertex corrections for charm quark final states which are similar to those seen for *b* quarks in the Standard Model, which could potentially add large corrections to the Zcc vertex. The ratio of the higgs mass to the mass of the W_R boson that mediates the interaction is the determining factor in the size of the modification of *A*, due to a model of this type. For $m_{W_R} > 0.3$ TeV and higgs masses greater than ~ 1 TeV, 10% deviations from the expected Standard Model value is possible.

2.8 Measurements of the Heavy Quark Forward-Backward Asymmetries

As shown in Table 2.4, the quarks produced in Z^0 decays have large asymmetry parameters $A_{,,}$ which implies that the product $\frac{3}{4}A_eA_f$ is also relatively large, or at least large enough to be easily measurable. The large values of the parameters A_f also imply that the sensitivity of A_{FB} to variations in the value of A, is nearly maximal (see Table 2.4), which motivates the effort to measure each separately in order to search for differences in the initial and final state couplings. The statistical boon provided by the large hadronic width of the Z^0 provides extra incentive to attempt measurements of A_{FB} for quarks.

In order to be sufficiently sensitive to the quark forward-backward asymmetries, however, as pure sample as possible of the particular quark flavor is needed. With the techniques presently available in modern particle detectors, the only two primary quarks whose events can be isolated in a "clean" (read, "with small systematic errors") manner are c and b. These considerations have led to many measurements of A_{FB} for heavy quarks by the LEP experiments in order to gain as much information as possible about the value of $\sin^2 \theta_W^{eff}$. We now turn to a discussion of the various techniques used to obtain samples of heavy quarks for asymmetry measurements. As will be seen, all of these techniques exploit one or more of the features of heavy quark production and decay:

- Hard Fragmentation Functions',
- · Large Masses,
- Long Lifetimes.

Results of the various methods will be discussed in Chapter 1.

*See Chapter 3.



Figure 2.12: The p and p_{\perp} distributions for leptons in hadronic Z^0 decays.

2.8.1 High p_{\perp} Lepton Tags

The first inclusive method used to isolate heavy quark events[61] was to require the presence of high-energy, isolated leptons as a tag. This technique exploits the relatively large (- 20%) branching fraction of heavy hadrons to leptons plus a hadronic system $(Q \rightarrow \ell + X)$ the hard heavy quark fragmentation functions, and the large energy available in the heavy hadron decay. The latter two properties of heavy quark events result in leptons with large total momentum and large p_{\perp} relative to the initial heavy hadron direction, which is commonly approximated by the jet axis[†] or the event thrust[62] axis.

The possible sources of leptons in hadronic decays of the Z^0 are[‡]:

•
$$b \rightarrow \ell^-$$

• $c \rightarrow \ell^+$

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.
$$b \to c \to \ell^+$$

. $b \to \bar{c} \to \ell^-$, where the \bar{c} is produced from the branching of the W^-
. $b \to \tau^- \to e^-$

. punchthrough, decays of πs or Ks in flight

and the decays that proceed through mixing:

 $b \to \overline{b} \to \ell^+$ • $b \to \overline{b} \to c \to \ell^-$, etc.

. b

The sign of the outgoing lepton is noted here explicitly, as it is used to determine the sign of the parent heavy quark in order to determine whether the fermion was produced in the forward or backward hemisphere. From the list above, it is clear that there is the possibility of confusion, as a lepton can come from any of the sources. As shown in Figure 2.12, however, those leptons with high momentum and high transverse momentum relative to the jet axis actually do come predominantly from heavy quark decays. For b events in SLD, the purity of the remaining sample after a cut which excludes leptons that fall within the ellipse

$$\left(\frac{p}{16}\right)^2 + \left(\frac{p_{\perp}}{1.2}\right)^2 = 1 \tag{2.88}$$

is approximately 72% [166]. The efficiency for b's in the hadronic event sample to pass this cut is about 5%. As the momentum requirements are lowered, a significant portion of the leptons from charm decays are allowed into the sample. As long as the proper corrections for $B - \bar{B}^0$ mixing are included, one can perform a maximum likelihood fit to the lepton spectra vs. polar angle and determine A^b_{FB} or A^b_{FB} and A_{FB}^{c} simultaneously. This method has been employed by the LEP experiments[63] to measure A_{FB}^{b} and A_{FB}^{c} . It should be mentioned that the errors on A_{FB}^{c} derived from an analysis of this type tend to be significantly larger than those on A_{FB}^{b} due to the effective background subtraction that occurs when fitting the asymmetry of a charm sample with small signal-to-noise.

[†]See Chapter 8.

[‡]Throughout this thesis all statements are assumed to apply equally to charged-conjugate modes.

2.8.2 Exclusive Reconstruction of Charmed Mesons

It is possible to obtain a relatively pure sample of events containing primary charm quarks, but at the price of requiring the presence of specific final states in the charm meson decay. Since methods of this type will be the primary focus of this thesis, we will discuss them in some detail here. These analyses usually consist of kinematically reconstructing topologically simple decay modes of the D^0 meson, then requiring that the D^0 was produced in the decay cascade $D^{\star +} \rightarrow \pi^+_* D^0$. This second requirement reduces the tagging efficiency, but, enhances the tag purity through a clever kinematic trick[64]. Since the masses of the $D^{\star+}$ and D^0 are 2010 MeV/ c^2 and 1864 MeV/ c^2 , respectively, the mass difference between the $D^{\star+}$ and the $\pi^+_s D^0$ combination is essentially a b-function at zero§. The distribution is smeared out slightly due to the extremely small π^+_{\bullet} momentum (39 MeV/c)[¶] in the D^{*+} rest frame and from tracking resolution effects. The kinematics of this decay result in a strong correlation between the angle of the π_s^+ with respect to the D^0 direction and the energy of the π_s^+ . Since tracking systems tend to measure angles better than momenta, the width of the peak in the mass difference (Am) distribution is much narrower than that for the D^0 mass. Significantly more background is excluded by cutting on the reconstructed D^0 mass and the Am distribution simultaneously. The added advantage in using this technique in a forward-backward asymmetry measurement is that the sign of the π_{\bullet} gives the sign of the $D^{\star+}$ and hence the sign of the charm quark without any of the ambiguities associated with leptons.

In addition, one can take advantage of the hard charm fragmentation function to obtain a sample of even higher purity, as the background events and Ds produced in b decay cascades typically have lower energy. Requiring the fractional energy $x_D = 2E_D/E_{CM}$ of the charm system to be more than 0.4, for example, reduces the background contamination by another large factor.

To estimate the usefulness of this tagging technique, let us first calculate the fraction of $Z^0 \to c\bar{c}$ events which contain the decay chain $D^{\star +} \to \pi_s^+ D^0$... This last

Table 2.5: The largest branching fractions in D^0 decay. We show here the branching fractions into topological final states ignoring any resonant substructure.

Decay Mode	Branching	Fraction	(%)
Semi-leptonic Modes:			
$D^0 \rightarrow K^- e^+ \nu_e$	3.80	± 0.22	
$D^0 \rightarrow K^- \mu^+ \nu_\mu$	3.2	2 ± 0.4	
$D^0 \rightarrow K^0 - \pi^0 e^+ \nu_e$	1.	$.6^{+1.3}_{-0.5}$	
$D^0 \to \overline{K}^0 \pi^- e^+ \nu_e$	2.	$8^{+1.7}_{-0.9}$	
Hadronic Modes:			
$D^0 \rightarrow K^- \pi^+$	4.01	± 0.14	
$D^0 \to \overline{K}^0 \pi^0$	2.05	5 ± 0.26	
$D^0 \to \overline{K}^0 \pi^+ \pi^-$	5.3	± 0.6	
$D^0 \rightarrow K^- \pi^+ \pi^0$	13.8	8 ± 1.0	
$D^0 \to \overline{K}^0 \pi^0 \pi^0$	~1.	0 ± 0.2	
$D^0 \rightarrow K^- \pi^+ \pi^+ \pi^-$	8.1	± 0.5	
$D^0 \longrightarrow \overline{K}^0 \pi^+ \pi^- \pi^0$	9.8	±1.4	
$D^0 \rightarrow K^- \pi^+ \pi^0 \pi^0$	18	5 ± 5	
$D^0 \to K^- \pi^+ \pi^+ \pi^- \pi^0$	4.3	± 0.4	
$D^0 \rightarrow \overline{K}^0 \pi^+ \pi^- \pi^0 \pi^0 (\pi^0)$	10	$.6^{+7.3}_{-3.0}$	
$D^0 \rightarrow \pi^- \pi^+ \pi^0$	1.6	± 1.1	
$D^0 \to \pi^- \pi^- \pi^+ \pi^+ \pi^0$	1.9	± 0.4	

step tells us that even if we can reconstruct every last D^0 , only 7% of all charm events will fall into our event sample using this technique if we make no other cuts. Unfortunately, we need to impose more cuts, and only a small fraction of D^0 decays are simple enough to be reconstructed in a straightforward manner. In Table 2.5 we list the largest branching fractions of the D^0 for reference[66].

In general one would choose to concentrate effort on the D^0 decay modes with the largest branching fractions to maximize the statistical power of the result. As mentioned above, however, the cleanest signals are obtained when one can use both the reconstructed D^0 mass and the Am distribution, so modes with invisible neutrals are disfavored. Since the inefficiencies involved in π^0 reconstruction are often large, the

[§]The actual mass difference used in analyses is $Am = m(D^{*+}) - m(D^0)$, which results in **a** narrow peak about $m(x) + 6 \text{ MeV}/c^2$ instead.

⁴for this reason, this pion is often referred to as the "bachelor", "slow", or "spectator" pion[65].

modes containing more than one π^0 also are often excluded. The special kinematics of the decay $D^0 \rightarrow K^- \pi^+ \pi^0$ make it suitable for inclusion since the π^0 need not be reconstructed to obtain a narrow mass peak suitable for event selection. This comes about in the following manner[65]. The decay $D^0 \to K^- \pi^+ \pi^0$ can proceed through two intermediate resonance decays $D^0 \to K^- \rho^+$ and $D^0 \to K^{\star-} \pi^+$. Since the D^0 is a pseudoscalar, these decays are both from a state of J = 0 to states of J = 1 + J = 0. In this case, the vector meson must be in the $J_2 = 0$ state to conserve angular momentum. This results in the decay products being emitted preferentially along the vector meson direction $(m_{1,0} \propto \cos \theta)$. In each case, approximately half the time, the K^- and the π^+ will be emitted in opposite directions in the D^0 rest frame, which results in their combination having a high invariant mass. The other \sim 50% of the combinations yields a peak at lower mass. The peak is sufficiently narrow for one to select these decays by imposing a mass cut in a band below the true D^0 mass. Typically, then, the decay modes $D^0 \to K^- \pi^+, D^0 \to K^- \pi^+ \pi^0$, and $D^0 \rightarrow K^- \pi^+ \pi^+ \pi^-$ are included in analyses which attempt to isolate a sample enriched in primary charm events using exclusive decay modes.

The fraction of charm events remaining after each step in the decay of a charm quark for these decay modes can be summarized by

$$c \xrightarrow{80\%} D \text{ mesons } \xrightarrow{55\%} D^{\star +} \xrightarrow{68\%} \pi_s^+ D^0 \xrightarrow{4\%} \pi_s^+ K^- \pi^+ = 1.2\%$$

$$\xrightarrow{14\%} \pi_s^+ K^- \pi^+ \pi^0 = 4.2\%$$

$$\xrightarrow{8\%} \pi_s^+ K^- \pi^+ \pi^+ \pi^- = 2.4\%$$
(2.89)

where the branching fractions for $a \rightarrow b$ are given above the arrow for each step in the cascade. So, the net charm-selection efficiency of these analyses could be as high as 8% if multiple decay modes are included (assuming 100% event selection efficiency; typical event selection efficiencies are as high as 20%). Event samples derived with these selection techniques have also been used at LEP[67, 68, 69] to measure A_{FB}^{c} , assuming A_{FB}^{b} is given by its standard model value.

2.8.3 Lifetime Tag Combined with Jet Charge

With the advent of precision microvertex detectors in collider physics, it has become possible to reliably separate events containing heavy hadrons from those with only light hadrons by using the lifetime signatures of the decaying hadrons. Heavy, weaklydecaying hadrons such as B and D mesons have "long" lifetimes on the order of 1 ps, which, at Z^0 energies, implies that their average decay lengths will be one to several millimeters. Their decay products will tend to have large p_{\perp} relative to the meson flight direction, which makes the secondary decay vertices easily resolvable using a vertex detector whose typical impact parameter resolution is less than 100 μ m for these tracks. Heavy quark events can be tagged in an inclusive manner by requiring the presence of a decay vertex or vertices in an event or hemisphere that is well separated from the interaction point. Requiring some number of tracks to appear not to originate from the interaction point, either by simple counting or by a likelihood function based on this information, can also be used^{||}. A lifetime-significance likelihood function tag was used to measure $\Gamma(Z^0 \to b\bar{b})/\Gamma(Z^0 \to hadrons)$. in Ref. [70]. Up to this time, tags of this type have concentrated on isolating pure samples of b hadrons, as their typical lifetimes and decay multiplicities are largest. Efficiencies in excess of 60% for purities above 90% have been achieved for these sorts of tags when they are applied to select b events.

Once a pure sample of b events (in this case) has been selected, an inclusive way must be found to measure the charge of the outgoing quarks to discern whether the fermion was produced in the forward or backward hemisphere. One way to proceed is to employ the momentum-weighted jet charge Q[71]:

$$Q = \sum_{\text{tracks}} -q_i \cdot \text{sign} \left(\vec{p}_i \cdot \hat{T} \right) \left| \vec{p}_i \cdot \hat{T} \right|^{\kappa}$$
(2.90)

Here, q_i is the sign of the *i*th particle, which has S-momentum $\vec{p_i}$, \hat{T} is a unit vector along the thrust axis direction, and κ is an adjustable parameter which must be optimized for each situation. The sign of the vector \hat{T} is chosen such that overall sum Q is negative, making \hat{T} an estimate of the *b* quark direction. Q is defined such that

For details of a simple counting tag, see Chapter 7

high momentum particles close to the thrust axis contribute the most information to the charge sum in an attempt to maximize the weight of the *b* decay products in the sum. This is successful since the hard *b* quark fragmentation function results in the decay products of the *b* having large average momentum along the original *b* quark direction in addition to carrying the majority of the energy in each thrust hemisphere. The LEP experiments have measured A_{FB}^b using this technique[72].

2.9 Measurements of A_b and A, at SLD

The analysis that is presented in this thesis is a single panel in a triptych of measurements of the parity-violating couplings A_b and A, performed using the SLD detector at SLC and the techniques described above.

Using the 1993 data sample of Z^0 decays, measurements of A_b and A, have been performed using high- p_{\perp} muons and electrons as a tag of heavy quark events. As an illustration of the results, the asymmetry vs the polar angle of the outgoing fermion is shown in Figure 2.13 for leptons selected with the elliptical cut of Eq. 2.88. The



Figure 2.13: The left-right forward-backward asymmetry for high p, p_{\perp} leptons.



Figure 2.14: The outgoing b quark direction as determined by jet charge for the two beam helicities. Note the large raw asymmetry.

curve drawn through the points is the asymmetry function from Eq. 2.72. Note that the observed asymmetry is quite large, as would be expected with the large absolute beam polarization (63%) of this data set. From a maximum likelihood analysis of the lepton p and p_{\perp} spectra, they obtained the following results[73] for A_b and A_i :

$$A_b = 0.91 \pm 0.14(stat) \pm 0.07(syst)$$

$$A_b = 0.37 \pm 0.23(stat) \pm 0.21(syst)$$

Note that the comment made above in reference to the typical error size on A_{FB}^c from measurements using the lepton tag technique applies equally well to measurements of A_r .

A jet-charge technique has also been employed to measure A_b after a sample of b events was selected using an impact parameter tagging technique like that described in Chapter 7. To demonstrate the power of this technique, we show in Figure 2.14

the raw distributions of the polar angle of the outgoing fermion for the two different longitudinal polarization states of the beam. The raw asymmetry is quite large, even if the identification of the fermion direction is not quite perfect**. A binned fit to the expression in Eq. 2.72 after background subtraction yields[74]

$$A_b = 0.87 \pm 0.11(stat) \pm 0.09(syst)$$

All of the above results are consistent with the standard model values of $A_{h} = 0.93$ and $A_{\rm r} = 0.67$. Comparison of these values with those obtained by the LEP collaborations will be made in Chapter 1.

We now turn to a description of the methods used in this thesis to measure A,.

The Measurement of A, Using Exclusively 2.10 **Reconstructed** Mesons

In this thesis, we concentrate on a measurement of the parameter A, using an enriched sample of $Z^0 \rightarrow CC$ events. As discussed above, the only way such a sample can be obtained is through the reconstruction of exclusive decay channels of D mesons, and this will be our method. Since the charm-tagging efficiency of this method is so small and we begin with a data sample 50 times smaller than that of the LEP experiments, we need to design an analysis that yields as many charm events as possible in order to arrive at a competitive measurement for A.. To this end, we have created a new set of selection criteria for isolating D^0 decays that make extensive use of the SLD precision vertex detector by requiring that the those particles used to reconstruct the D^0 meson also form a secondary decay vertex well separated from the interaction point. The sample of D^0 decays used to measure A, is a combination of events selected by the standard kinematic analysis and those chosen by the new decay length selection criteria. In addition, to maximize the statistical power of the measurement we include a sample of D^+ decays. Note that sign of the D^+ also unambiguously determines the sign of the charm quark produced in the Z^0 decay. Since there is no corresponding Am trick to improve the signal to noise in the case of D^+ production*, we are only interested in those D^+ decay modes with a large probability for correct reconstruction+. We list the major branching fractions [194] of the D^+ in Table 2.6 for comparison to Table 2.5; we have chosen to consider only the mode $D^+ \rightarrow K^- \pi^+ \pi^+$ for this analysis as it offers a large branching fraction and contains no π^0 s. The D⁺ selection analysis also makes extensive use of the vertex detector to obtain a sample of events with very little background.

A detailed description of the D^0 and D^+ selection techniques will be presented in Chapter 13. Discussion of the likelihood fitting method used to extract A, follows in Chapter 1, with a discussion of systematic errors on the measurement (Chapter 1) and conclusions (Chapter 1) serving to close the electroweak part of this thesis.

^{**}It is estimated that the outgoing quark is tagged properly as a b or \overline{b} in 63% of b events

^{*}Using the decay $D^{\star 0} \to \pi^0 D^+$ necessitates reconstruction of the π^0 . *One advantage to using the D^+ modes is that the presence of a $D^{\star} \to D$ cascade is not required, so that the analysis can start with a larger pool of charm events. Slightly harder cuts still result in a large sample of events

Table 2.6: The largest, branching fractions in D+ decay. We show here the branching fractions into topological final states ignoring any resonant substructure.

Decay Mode	• Branching Fraction (%)
Semi-leptonic Mode	s:
$D^+ \rightarrow \overline{K}^0 e^+ \nu_e$	6.6 ± 0.9
$D^+ \rightarrow \overline{K}^0 \mu^+ \nu_\mu$	$7.0^{+3.0}_{-2.0}$
$D^+ \rightarrow K^- \pi^+ e^+ \nu_e$	$4.2_{-0.7}^{+0.9}$
$D+ \rightarrow K^- \pi^+ \mu^+ \nu_\mu$	3.2 ± 1.7
Hadronic Modes:	
$D+ \rightarrow \overline{K}^0 \pi^+$	2.74 ± 0.29
$D^+ \rightarrow K^- \pi^+ \pi^+$	9.1 ±0.6
$D^+ \rightarrow \overline{K}^0 \pi^+ \pi^0$	9.7 ± 3.0
$D+ \rightarrow K^-\pi^+\pi + \pi^0$	6.4 ± 1.1
$D+ \rightarrow \overline{K}^{0}\pi^{+}\pi^{+}\pi^{-}$	7.0 ± 1.0
$D^+ \rightarrow K^- \pi^+ \pi + \pi^0 \pi$	τ^0 2.2 ^{+5.0} _{-0.9}
$D^+ \rightarrow \overline{K}^0 \pi^+ \pi^+ \pi^- \pi^0$	$5.5^{+3.0}_{-1.4}$
$D^+ \to \overline{K}^0 \overline{K}^0 K^+$	3.1 ± 0.7
Chapter 3

The Strong Interaction: An Introduction

The quest for the correct theoretical formulation of the nuclear force has continued since the clarification of nuclear structure that resulted from Chadwick's discovery of the neutron in 1935 [75]. Yukawa's pion [76], although it provided a key concept in the development of field theory, proved to be the first of an enormous number of meson states instead of the particle solely responsible for carrying the strong interaction. The incredible simplification realized in the quark model of Gell-Mann and Zweig [77] and the introduction of "color" by Han and Nambu[78] represent the cornerstone of our "modern" understanding of strong interactions, as their supposition of flavor and color SU(3) as the correct group symmetries has been borne out by experiment. The early 1970's witnessed the transition of quarks from phantoms to physical entities [79] in the deep-inelastic electron scattering experiments performed at SLAC. The discovery[80] that the non-Abelian nature of Yang-Mills theories[10] allowed the ultraviolet freedom of quarks necessary to explain the Bjorken scaling seen in these experiments allowed the theory of the strong interactions that we now know as Quantum Chromodynamics (QCD) to reach its final form. The observations at PETRA [81] of jets of particles interpreted to be from the radiation of hard gluons served to cement QCD as the "correct" theory of the strong interactions. Since that time, all other tests of the assumptions of QCD have reinforced this conclusion. Since part of the topic of this

thesis is another of these tests of QCD, it is proper to review in more detail the theoretical formulation of OCD and the manner in which it can be tested.

3.1 Quantum Chromodynamics

This section presents a discussion of the theory of the strong interactions, proceeding from its basic group-theoretical symmetries through the question of renormalizability. The phenomenological successes of QCD will be pointed out at the appropriate points.

3.1.1 The QCD Lagrangian

We begin with the lagrangian density for Quantum Chromodynamics* of one quark field q:

$$\mathcal{L} = \tilde{q}(i\gamma^{\mu}\partial_{\mu} - m)q - g_{s}(\tilde{q}\gamma^{\mu}T_{a}q)B^{a}_{\mu} - \frac{1}{4}G^{a}_{\mu\nu}G^{\mu\nu}_{a}$$
(3.91)

gluon fields, and q is the coupling introduced in requiring local SU(3)-color gauge invariance. The generators T satisfy

$$[\mathbf{T}_{a}, \mathbf{T}_{b}] = i f_{abc} \mathbf{T}_{c}, \tag{3.92}$$

where the factors f_{abc} are the SU(3) structure constants. The kinetic energy term is given by

$$G^a_{\mu\nu} = \partial_\mu B^a_\nu - \partial_\nu B^a_\mu - g_s f_{abc} B^b_\mu B^c_\nu \tag{3.93}$$

This last term, required by local gauge-invariance, separates the behaviour of QCD from that of Quantum Electrodynamics (QED), as it implies that the gluons themselves carry the color charge. This addition arises due to the non-Abelian nature of the gauge group; similar terms occur in the Electroweak Lagrangian, Eq. 2.2. The interactions associated with each of the terms in the Lagrangian of Eq. 3.91 are shown schematically in Figure 3.1.



 a^2B^4 aB³ qā BB aBaā

Figure 3.1: A symbolic representation of the interactions contained in the QCD lagrangian. The diagrams labeled qq and BB are the quark and gluon propagators. The vertices to order g are shown next, and finally, the four-gluon vertex is shown.

It has become customary to refer to the quantity $\alpha_s = \frac{g_a^2}{4\pi}$ as the strong coupling. All further discussion of the strong coupling in this thesis will follow this convention.

3.1.2 Simple QCD Processes

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In this section, we will discuss the processes $e^+e^- \rightarrow q\bar{q}$ and $e^+e^- \rightarrow q\bar{q}g$ in order to exhibit some of the properties of the strong interaction formulated in QCD. Figure 3.2a shows the Feynman diagram for $e^+e^- \rightarrow q\bar{q}$ at the tree level, which, except for the required factor of N_c (the number of colors) in the cross section, requires no input from QCD[†]. To lowest order in the strong coupling α_s , the rate for $e^+e^- \rightarrow q\bar{q}$ is modified by virtual corrections due to internal gluon loops, as shown in Figure 3.2b. Each of these diagrams contains ultra-violet and/or infra-red divergences, but the sum of the squared amplitudes yields a finite result for the total cross section* when the $\mathcal{O}(\alpha_s)$ correction due to the emission of a real gluon (Figure 3.3a) is added. For completeness, Figure 3.2c contains the Feynman graphs of processes of order α_{\star}^2 included in a higher order calculation of $e^+e^- \rightarrow q\bar{q}$. The addition of these diagrams to those for all other $\mathcal{O}(\alpha_*^2)$ processes, including those in Figure 3.3b and Figure 3.4 also results in a finite total cross section for $e^+e^- \rightarrow$ hadrons after the coupling constant is renormalized (see the next section).

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^{&#}x27;the discussion of the field-theoretical properties of QCD presented in this and the following sections has been synthesized from a number of references. See Refs. [26, 25, 82) for details.

[†]The electroweak differential cross section for fermion production was presented in Eq. 2.56. [‡]See, e.g., Ref. [82, p. 50].



Figure 3.2: Feynman diagrams for $e^+e^- \rightarrow q\bar{q}$, showing (a) the $\mathcal{O}(\alpha_s^0)$ process, (b) $\mathcal{O}(\alpha_s)$ processes, and (c) $\mathcal{O}(\alpha_s^2)$ processes.



Figure 3.3: Feynman diagrams for $e^+e^- \rightarrow q\bar{q}g$, showing (a) the $\mathcal{O}(\alpha_s)$ process, and (b) $\mathcal{O}(\alpha_s^2)$ processes.

Let us consider explicitly the process $e^+e^- \rightarrow q\bar{q}g$ where, instead of being interested only in the total cross section for $e^+e^- \rightarrow$ hadrons, we would like to observe the radiated gluon. To lowest order in QCD, the differential cross section for radiating a



Figure 3.4: Feynman diagrams for $e^+e^- \rightarrow q\bar{q}q\bar{q}$ and $e^+e^- \rightarrow q\bar{q}gg$.

single gluon in events of the type $e^+e^- \rightarrow q\bar{q}g$ (see Figure 3.3a) is given by [83]

$$\frac{da}{dx_1 dx_2} = \sigma_0 \frac{2}{3} \frac{\alpha_s}{7r} \frac{x_1^2 + x_2^2}{(1 - x_1)(1 - x_2)},$$
(3.94)

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where σ_0 is the lowest order Born cross section,

$$\sigma_0 = \frac{4\pi\alpha^2}{3s} N_c \sum_f q_f^2 , \qquad (3.95)$$

and $x_i = 2E_i/E_{CM}$ for $i = 1 : \vec{q}, i = 2 : q, i = 3 : g$. Here, $N_c = 3$ is the number of colors, α is the electroweak coupling, and q_f is the charge of quark flavor f. Of course, the values of x that are kinematically allowed in this formulation are $0 < x \leq 1$.

The lowest-order cross section of Eq. 3.94 diverges as x_1 or x_2 approaches unity. The case $x_1 \rightarrow 1$, $x_2 \rightarrow 1$ corresponds to the emission of a gluon of vanishingly small momentum. This is an infrared divergence, in that the cross section diverges when the momentum of the emitted quantum approaches zero. The other case, where $x_1 \rightarrow 1$, $x_2 \neq 1$ corresponds to the situation where one of the quarks emits a soft, collinear gluon, and is known as a collinear divergence. In either case, the gluon is not resolvable from the q or \bar{q} . As stated above, however, the total cross section for $e^+e^- \rightarrow$ hadrons is finite, so the probability of radiating a "resolvable" final-state gluon must also be finite. We can define the term "resolvable" by any number of different criteria. For example, the energy of the outgoing gluon and its angles with respect to the two quarks must be above some threshold [84], or, equivalently, the scaled invariant mass squared

$$y = \frac{m_{ij}^2}{s} \tag{3.96}$$

must be greater than some resolution parameter y_{min} for all pairs of quarks and gluons[85]. Integration over the region of phase space that the gluon is resolvable yields a finite cross section for $e^+e^- \rightarrow q\bar{q}g$ events[§]. Note that the cross section for $e^+e^- \rightarrow q\bar{q}g$ (and that for $e^+e^- \rightarrow q\bar{q}g$, for that matter) must now be functions of the resolution criteria[89]:

$$\sigma^{qq}(y_{min}) = \sigma_0 \left[1 + \frac{4}{3} \frac{\alpha_s}{2\pi} \left(-2\log^2 y_{min} - 3\log y_{min} + 4y_{min}\log y_{min} - 1 + \frac{\pi^2}{3} \right) \right]$$

$$\sigma^{qqg}(y_{min}) = \sigma_0 \frac{4}{3} \frac{\alpha_s}{2\pi} \left(2\log^2 y_{min} + 3\log y_{min} - 4y_{min}\log y_{min} + \frac{5}{2} - \frac{\pi^2}{3} \right) \quad (3.97)$$

Note that the cross section for the three-parton final state is proportional to α_s ; we will use the measurement of this quantity to derive the magnitude of the strong coupling.

Similar calculations have been performed to obtain matrix elements for more complicated final states of quarks and gluons (henceforth referred to collectively as "partons"). In particular, results exist for the cross sections for four $(e^+e^- \rightarrow q\bar{q}gg)$ or $e^+e^- \rightarrow q\bar{q}q'\bar{q}'$ [87] and 5-jet $(e^+e^- \rightarrow q\bar{q}ggg)$ or $e^+e^- \rightarrow q\bar{q}q'\bar{q}'g$ [88] parton final states. In principle, these calculations can be compared with the experimental measurements of the same quantities as a check of the validity of higher-order QCD calculations.

As mentioned above, the individual Feynman amplitudes *for* the first order QCD corrections to any process contain divergences. Since the "bare" coupling of the theory appears to be infinite when the first order corrections are added, it is useful as a computational tool to invoke a renormalization procedure to redefine the coupling a, **as** a finite experimental observable and absorb the infinities of the individual graphs. This will be discussed in the next section.

[§]The infra-red divergences cancel those in the $\mathcal{O}(\alpha_s)$ calculation for $e^+e^- \to q\bar{q}$, so that the cross section for $e^+e^- \to q\bar{q}$ is also finite, as it must be to preserve unitarity at this order in perturbation theory[86].

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Figure 3.5: The leading order corrections to the quark-gluon vertex, i.e., to α_s , showing: (a) the bare qqg vertex; (b) the one loop vertex corrections, which multiply the bare vertex by the factor $(Z_1^{-1} - 1)$; (c) the quark self-energy corrections, which multiply the bare vertex by the factor $-2(Z_2^{-1} - 1)$; (d) the gluon self-energy corrections, which multiply the bare vertex by the factor $(Z_3 - 1)$. The solid lines correspond to quarks, the helical curves represent gluons, and the dashed lines are the "ghost" distribution necessary for removing non-physical gluon polarization states from the calculation.

3.1.3 The Renormalization of the Strong Coupling

Before moving on to a discussion of the various predictions of QCD, the behaviour of the theory at the one-loop level needs to be explored. It is the renormalized coupling which exhibits all of the salient features of QCD phenomenology, and we turn now to a discussion of its derivation.

The lowest order diagrams and virtual corrections to the quark and gluon propagators are shown in Figure 3.5. As is shown in the figure, the lowest order vertex factor $-ig_s^0 \gamma_{\mu} T_a$ is modified[¶] by the various one-loop corrections:

 $-ig_s^0 \gamma_\mu T_a (Z_1^{-1} - 1)$ by the vertex corrections,

- $(-ig_s^0\gamma_{\mu}\mathbf{T}_a)(-2)(Z_2^{-1}-1)$ by the quark self-energy corrections,
- $-ig_s^0 \gamma_\mu T_a(Z_3 1)$ by the corrections to the gluon propagator.

Each of the Z_i terms contain ultra-violet divergences resulting from the infinite momenta which are possible in the parton loops. These infinities must be absorbed into an effective coupling constant in order to render the theory calculationally useful. Including the proper normalization for external fermion lines and the gluon source, we can add the amplitudes for all of the diagrams to arrive at

$$-ig_{s}\gamma_{\mu}\mathbf{T}_{a} = \frac{-ig_{s}^{0}\gamma_{\mu}\mathbf{T}_{a}}{Z_{2}\sqrt{Z_{3}^{*}}} \Big[1 + (Z_{1}^{-1} - 1) - 2(Z_{2}^{-1} - 1) + (Z_{3} - 1)\Big] \quad (3.98)$$
$$= \frac{Z_{2}\sqrt{Z_{3}}}{Z_{1}} (-ig_{s}^{0}\gamma_{\mu}\mathbf{T}_{a})$$

or

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$$g_{s} = \frac{Z_2 \sqrt{Z_3}}{Z_1} g_{s}^{0} . \tag{3.99}$$

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This coefficient can be written as a power series in terms of α_s^0 , the bare coupling, where $\alpha_s^0 = g_s^0/2\pi$. The new, renormalized effective coupling α_s^{eff} is given by

$$\alpha_s^{eff}(Q^2) \equiv \alpha_s(Q^2) = \alpha_s^0 \left(1 + \alpha_s^0 B_{QCD}(Q^2) + \alpha_s^0 B_{QCD}(Q^2) \alpha_s^0 B_{QCD}(Q^2) + (3:) \right) 00$$

= $\frac{\alpha_s^0}{1 - \alpha_s^0 B_{QCD}(Q^2)}$

Here, B_{QCD} is the actual functional form of the result, and Q^2 is the gluon momentum squared and hence the energy available for the interaction. In general, the specific form of B_{QCD} depends on the method chosen to regulate the divergences that arise in calculating the Z_i factors while preserving gauge invariance. One way to do this, called dimensional regularization[90, 91] is to calculate the loop diagrams in N dimensions, where N is typically set to N = 4 + ϵ , and then take the limit of the results as N \rightarrow 4 ($\epsilon \rightarrow 0$).

Equation 3.101 implies that

$$\frac{1}{\alpha_s(Q^2)} = \frac{1}{\alpha_s^0} - B_{QCD}(Q^2) . \qquad (3.101)$$

[¶]Expressions for these corrections can be found in, e.g., Ref [82]. Note that Figure 3.5d contains an extra diagram due to the "ghost" terms which must be added to remove non-physical polarization states in the gauge chosen for this calculation.

The function $B_{QCD}(Q^2)$ diverges, however, and a standard renormalization procedure must be followed in order to obtain a finite, observable coupling in terms of a finite difference of infinite terms. Since $B_{QCD}(Q^2)$ is proportional to $\log Q^2$, we cannot define the coupling in the limit of zero momentum transfer. Instead, we must choose a finite reference scale $Q^2 = \mu^2$ and define the "experimental coupling" as

$$\alpha_s \equiv \alpha_s(\mu^2) \tag{3.102}$$

as a replacement for α_s^0 where is appears. The relationship between the experimental coupling and the bare coupling is given, then, by

$$\frac{1}{\alpha_s^0} = \frac{1}{\alpha_s(\mu^2)} + B_{QCD}(\mu^2), \qquad (3.103)$$

so that at other values of Q^2 , we find the coupling is

$$\frac{1}{\alpha_s(Q^2)} = \frac{1}{\alpha_{\rm ss}(\mu^2)} - \left(B_{_{QCD}}(Q^2) - B_{_{QCD}}(\mu^2)\right).$$
(3.104)

Since the Q^2 dependence of $B_{QCD}(Q^2)$ only appears in a logarithmic term,

$$B_{QCD}(Q^2) - B_{QCD}(\mu^2) \propto \log(Q^2/\mu^2)$$
(3.105)

To lowest order [82, p. 242], the exact expression is

$$B_{QCD}(Q^2) - B_{QCD}(\mu^2) = -\frac{\beta_0}{4\pi} \log(Q^2/\mu^2), \qquad (3.106)$$

where

$$\beta_0 = \frac{11}{3}C_A - \frac{4}{3}C_F = 11 - \frac{2}{3}n_f . \qquad (3.107)$$

Here, C_A and C_F are the Dynkin indices of the adjoint and fundamental representations of $SU(3)^{\parallel}$, and n_f is the number of quark flavors accessible at energy range Q^2 . The difference in the two infinite B_{QCD} terms is finite and doesn't depend on the constants introduced by dimensional regularization. Inserting this expression into Equation 3.104 yields

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$$\alpha_s(Q^2) = \frac{\alpha_s(\mu^2)}{1 + \alpha_s(\mu^2)\frac{\beta_0}{4\pi}\log(Q^2/\mu^2)}$$
(3.108)

It is common to rewrite the above equation in terms of a single parameter, often denoted A. If we make the replacement

$$\log(\Lambda^2) = -\frac{4\pi}{\beta_0 \alpha_s(\mu^2)} + \log(\mu^2) , \qquad (3.109)$$

we can rewrite Eq. 3.108 as

$$\alpha_s(Q^2) = \frac{4\pi}{\beta_0 \log(Q^2/\Lambda^2)} \tag{3.110}$$

This particular choice of which constants are absorbed into A is known as the Minimal Subtraction (MS) scheme[92]. Other schemes, such as the Modified Minimal Subtraction ($\overline{\text{MS}}$) scheme[93], also exist.

Note that there two effects of the renormalization of the coupling. First, the coupling α_s now is function of Q^2 , i.e. α_s is a running coupling. Second, the renormalization procedure has re-expressed the coupling of QCD as a function of a single parameter, A, which is now the fundamental parameter of QCD.

3.1.4 The Altarelli-Parisi Equations: The Leading Logarithm Approximation

It was shown by Altarelli and Parisi[94] that, to leading order in powers of the logarithms $\log Q^2/\Lambda^2$, the results of $\mathcal{O}(\alpha_s^2)$ QCD calculations for the behaviour of the structure functions in deep-inelastic scattering can be reproduced by a consideration of only the basic vertices of QCD coupled with parton probability densities in an infinite momentum frame $(p \gg m_q)$. Their description of hard scattering is valid to all orders in $\alpha_s(Q^2) \log Q^2/\Lambda^2$. These corrections to the tree level diagrams can be summed to infinite order, resulting in the so-called Leading Logarithm Approximation (LLA)[96] to QCD. ł

These factors arise in calculating the traces in obtaining the loop amplitudes. The quantity C_A comes from the trace $\text{Tr}(\mathbf{F}_a\mathbf{F}_b) = C_A\delta^{ab} = 3\delta^{ab}$, where \mathbf{F}_a are the 8×8 matrices defined by $(\mathbf{F}_a)_{bc} = -if_{abc}$. So, $C_A = 3$, the number of colors. In the fundamental representation, $\text{Tr}(\mathbf{T}_a\mathbf{T}_b) = C_F\delta^{ab} = \frac{1}{2}\delta^{ab}$ for each fermion in the fundamental representation. All told, the factor due to C_F is given by $\frac{1}{2}n_f$, where n_f is the number of accessible fermions.

The Altarelli-Parisi formulation of the LLA in terms of the more intuitive parton model concepts allows a much simpler basis for calculating those quantities where the leading logarithms add important contributions. In their description, the basic vertices of QCD are expressed in terms of probability densities for the branching of one parton into two partons, $P_{a \rightarrow bc}$. The evolution of a parton density function $f_b(z, Q^2)$ is given by

$$\frac{df_a(z, Q^2)}{dt} = \frac{\alpha_s(Q^2)}{2\pi} \int_z^1 \frac{dy}{y} P_{a \to bc}(y) f_b(\frac{z}{y}, Q^2) .$$
(3.111)

Here, t describes the evolution of the system, e.g. $t = \log(Q^2/\Lambda^2)$, and z gives the fraction of the initial parton energy retained by parton b^{**} . The $(Q^2$ -independent)^{††} functions $P_{a \rightarrow bc}$ are given by:

. 0

$$P_{q \to qg} = \frac{4}{3} \frac{1 + (1 + z)^2}{2}$$

$$P_{g \to qg} = 6 \left(\frac{1 - f}{z} + \frac{z}{1 - z} + z(1 - z) \right)$$

$$P_{g \to q\bar{q}} = -\frac{4}{3} \frac{1 + z^2}{1 - z}.$$
(3. .2)

The usefulness of this formulation in the generation of parton distributions within a MC simulation of hadron production will be discussed below in Section 3.2.3.

3.1.5 Higher Order Corrections:

The Renormalization Group Equation

It is possible to perform the calculation of the renormalized coupling to higher order in α_s by considering the increasingly complex Feynman graphs which occur. Some of these for $\mathcal{O}(\alpha_s^2)$ are shown in Figure 3.2~.

The results of these calculations can be conveniently expressed in terms of the β function $\beta(\alpha_s)$ contained within the so-called *Renormalization* Group Equation:

$$\mu \partial \alpha_s / \partial \mu \equiv \beta(\alpha_s) = -\frac{\beta_0}{2\pi} \alpha_s^2 - \frac{\beta_1}{4\pi^2} \alpha_s^3 - \frac{\beta_2}{64\pi^3} \alpha_s^4 + \cdots$$
(3.113)

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The coefficients are given by

$$\beta_1 = (306 - 38n_f)/6 \tag{3.114}$$

and

$$\beta_2 = 2857 - \frac{5033}{9}n_f + \frac{3225}{27}n_f^2 \tag{3.115}$$

in minimal subtraction (MS) schemes. Note that the definition of the coefficients β_n for $n \ge 2$ in the series depend on the specific renormalization procedure used to subtract the infinities, i.e. they are renormalization-scheme-dependent. The so-called renormalization scale μ is a measure of the mean gluon energy taken in the higherorder corrections that modify the coupling. It is the same parameter introduced in the calculation of Section 3.1.3 to define the scale of the leading order corrections.

With this definition of the β function, the definition of the scale A can be given as

$$\Lambda = \mu \exp\left[-\int_{s}^{\alpha} (Q^2)_A \frac{d\alpha_s}{\beta(\alpha_s)}\right].$$
(3.116)

The parameter A is arbitrary and serves to specify the boundary conditions on the choice of scale.

If we could perform the calculation to all orders, any dependence of the answer on μ must vanish, as it is an artifact of the truncated calculation. We can see from combining Eq. 3.113 with Eq. 3.116 that, at least to this order,

$$\partial \Lambda / \partial \mu = 0 \tag{3.117}$$

though A still depends on the renormalization scheme through the exact values of the coefficients in the β function series.

The integration of the renormalization group equation yields an expression for α_{\bullet} in terms of μ , or equivalently, from Eq. 3.116, in terms of A. One choice of a representation for this result is [98]:

$$\begin{aligned} \alpha_{s}(Q^{2}) &= \frac{4\pi}{\beta_{0} \log(Q^{2}/\Lambda^{2})} \left[1 - \frac{\beta_{1}}{\beta_{0}} \frac{\ln\left[\ln(Q^{2}/\Lambda^{2})\right]}{\ln(Q^{2}/\Lambda^{2})} \right. \\ &+ \frac{4\beta_{1}^{2}}{\beta_{0}^{4} \log^{2}(Q^{2}/\Lambda^{2})} \times \left(\left(\ln\left[\ln(Q^{2}/\Lambda^{2})\right] - \frac{1}{2} \right)^{2} + \frac{\beta_{2}\beta_{0}}{8\beta_{1}^{2}} - \frac{5}{4} \right) \right] (3.118) \end{aligned}$$

^{**}Parton c necessarily gets the fraction 1 - z.

^{††}These terms arc: Q^2 independent to all orders of α_s . A suitable rescaling of the energy scale A suffices to include higher-order (in a.) contributions[97].

Having proceeded through the discussion of first- and second-order calculations of the modification of the strong coupling, we should pause and verify that the coupling we have obtained is relatively small, so that our use of a power series in these calculations can be shown to be valid, as least a posterior?. We would expect that the strong interaction becomes "strong" when the energies of the quarks are on the order of the binding energy of a simple hadron, like the pion or kaon. Inserting a value of A of the order f_{π} (A = 100 MeV/c), we find that, for values of Q^2 larger than about (5 GeV/c)², $\alpha_s < 0.2$, so at least there is some hope that there is a perturbative regime in which calculations based on the power series will converge. We will now discuss several of the properties exhibited by the renormalized value of the strong coupling α_s .

3.1.6 Asymptotic Freedom and Quark Confinement

The renormalized coupling shown in Equation 3.108 or Equation 3.118 possesses a remarkable property: as the momentum transfer of the reaction, Q^2 , approaches infinity, the strength of the coupling drops to zero[80]. This is known as "Asymptotic Freedom", and was surmised to be a necessary part of the theory of QCD after the experimental data from deep-inelastic lepton scattering on hadrons at SLAC exhibited so-called "Bjorken scaling" [99]. The observed behaviour of the structure functions in deep-inelastic scattering implied that the quarks within the nucleons were quasi-free particles when struck by a hard photon, and thus any theoretical description of the strong force must include this property.

The form of Equation 3.113 gives us a convenient way to discuss the properties of the theory. Consider the lowest order term

$$\mu \partial \alpha / \partial \mu \equiv \beta(\alpha) = -\frac{\beta_0}{2\pi} \hat{\alpha} . \qquad (3.119)$$

Since

$$\beta_0 = 11 - \frac{2}{3}n_f \,, \tag{3.120}$$

for $n_f < 16$, the coupling α_s decreases as $\mu \to co$, *i.e.*, SU(3) QCD is asymptotically free. This was shown to be a general property of Yang-Mills theories; indeed, only non-Abeliau theories have a chance of being asymptotically free[100].

Let us compare QCD and QED. The analogous equation to Eq. 3.108 in QED is

$$\alpha_{EM}(Q^2) = \frac{\alpha_{EM}}{1 - (\alpha_{EM}/3\pi \log(Q^2/\mu^2))} .$$
(3.121)

The coefficient we would identify as β_0^{QED} is then

$$\beta_0^{QED} = -\frac{4}{3} , \qquad (3.122)$$

which implies from Eq. 3.119 that the coupling in QED increases as $Q^2 \to \infty$, the diametrically opposite behaviour. It is the contribution of the gluon loops in the gluon self-energy correction, which are present only because the three-gluon vertex exists, that add the extra, positive term in β_0^{QCD} that makes the theory asymptotically free.

A more "physical" picture can be obtained by considering a picture of the "charge screening" that occurs at small distances in the two theories. In OED, an electron is surrounded by a cloud of virtual electron-positron pairs. The pairs are polarized in such a manner that the positrons are closer to the pole electron, and thus a probe photon of relatively low energy observes a lower, shielded charge, as it cannot resolve the polarized cloud. As the energy of the photon increases, the distance scale of its interaction shrinks, allowing it to resolve an area containing fewer of the shielding particles, and hence observing a higher charge. Something completely different hap pens in QCD. We can again for simplicity consider a photon probe impinging on a quark of a given color, say blue^{t‡}. The same vacuum polarization effects that occur in QED will also be present, i.e., the quark will be surrounded by pairs of blue-anti-blue quarks from the vacuum, which will shield its charge. However, if the blue quark were to emit a virtual gluon of charge, say, blue-anti-red, it changes for a time into a red quark. Since this radiation of virtual gluons occurs continually, the color charge due to the quark is distributed and diffuse, so that the higher the resolution of the photon probe, the less effects of the color charge it can "see". In effect, the fact that 8 colored gluons can be radiated overcomes the charge screening behaviour.

If we now look at the opposite Q^2 extreme, we can see that, as Q^2 approaches Λ^2

^{‡‡}The "colors" of color SU(3) have come to be referred to as red, green, and blue, the primary colors of light.

in

$$\alpha_s(Q^2) = \frac{\Delta \pi}{\beta_0 \log(Q^2/\Lambda^2)}, \qquad (3.123)$$

the coupling becomes infinite. Of course, the actual scale at which this happens is somewhat arbitrary, as it depends on the exact value of A chosen, but it is a general feature of QCD that perturbative calculations are no longer valid when the energy of the quarks and gluons approaches zero $(Q^{"} \rightarrow A)$. This behaviour implies that free, individual, bare quarks are not found in nature. They must be confined [101] within color-neutral hadronic matter, since it would require an infinite amount of energy to liberate them; this is in stark contrast to the behaviour of electrons in the more familiar QED. At the present time, no rigorous proof exists that SU(3)QCD contains confinement. It can be shown[78], however, that the presence of the gluon self coupling results in the lowest-energy states being those which are color singlets. For a discussion of dynamic symmetry breaking in non-abelian theories and its relationship to confinement, see Ref [102].

The confinement of quarks and gluons within hadrons renders the comparison of theoretical calculations with experimental results difficult. Any experiment designed to test the tenets of QCD must begin and end with observable particles. Regardless of the energy of the probe used as a measure of the interaction, the parton distributions produced by a process calculable in perturbative QCD necessarily pass through the incalculable non-perturbative regime as the asymptotically free partons are "dressed" up to become hadrons which can be observed in a particle detector. In practice, experimentalists rely on several different Monte Carlo-based models of this "hadronization" process in order to relate the observed distributions of hadrons back to the underlying parton processes responsible for them. The observed distributions and the simulation techniques used to reproduce them are the subject of the next section.

3.2 Production of Hadrons in e^+e^- Collisions

So far in this discussion of QCD, we have concentrated in a somewhat abstract manner on the field-theoretical properties of a hadronic theory based on color SU(3)

and some general phenomenological implications. Here, we present an overview of hadron production in e^+e^- collisions including the QCD processes that result in the experimental observables.

3.2.1 The Fragmentation Process

As the previous section implied, when a quark-anti-quark pair is produced in the process $e^+e^- \rightarrow q\bar{q}$, the two bare quarks are not observed as two individual charged tracks emerging from the production point. Instead, a process similar to that shown in Figure 3.6 occurs. The two quarks, slowed by the increasing magnitude of the strong force, radiate gluons, each of which can split into two quarks which can themselves reradiate. This "showering" continues until the available energy is degraded towards Λ^2 . As mentioned above, at this time the coupling becomes very large, and the partons are no longer able to exist as free entities. It is here where the uncalculable hadronization process occurs, and we are left with many different hadronic states. Many of these states can be unstable resonances, which quickly decay into more stable hadrons. It is this final distribution of stable hadrons which is observed in particle detectors.

3.2.2 Jets

As will be shown explicitly below, the shower process for a single quark is analogous to electromagnetic bremsstrahlung. Many of the radiated gluons are thus emitted traveling close to the initial quark direction. If the primary quark has sufficient relativistic boost, the radiated partons, and hence their hadronic decay products, will form a collimated cone around the initial quark momentum direction. These cones of particles are referred to as "jets", and are ubiquitous in hadronic interactions at high energy colliders *. Evidence for jets was first observed by the MARK I experiment[103] at SPEAR, using an analysis that showed that the hadronic events became less spherical (and hence more "jetty") as the energy of collisions was increased from 3 to 7.4 GeV. At higher center-of-mass energies, jets are an obvious feature of hadronic events. In 1979, experiments at PETRA[81]obscrvctl planar events with 3 well-separated jets,

^{&#}x27;A precise definition of what is meant by a "jet" will be given below, in Section 3.3.1.



Figure 3.6: An overview of an $e^+e^- \rightarrow q\bar{q}$ event. Four stages of the event evolution are indicated. Stage (i) is the perturbative phase, where the Q^2 of the interactions are sufficiently large that perturbation theory is valid. Stage (ii) is fragmentation/hadronization, where the initial partons are combined by some prescription, and color-singlet hadrons are formed. In stage (iii), the short-lived hadronic resonances decay; it is their decay products and the other stable hadrons that reach stage (iv), experimental observation.

the first "direct" observation of hard gluon radiation. As an example of the observed jet structure, Figure 3.7 shows events in SLD with 2 and 3 jets. As will be explained in the following sections, a number of properties of jets can be related to the underlying parton distributions and hence can be used to test the predictions of perturbative QCD.

3.2.3 Simulations of Fragmentation

The cascade of quarks and gluon produced in a parton shower is far too complicated to lend itself to an analytical calculation of the exact final state with present-day 68



Figure 3.7: Views of 2- and 3-jet events in SLD. The curved lines originating from the center of the detector are charged tracks; the irregular rectangular boxes which point back to the origin are energy deposits in the calorimeter system. See Chapter 4 for a description of the various detector elements.

techniques+. Instead, various approaches based on Monte Carlo simulation have been developed in an effort to relate simple processes containing small numbers of partons to the properties of the hadronic final state[104]. They differ in the details of the showering process, but all models contain a picture of the hadronic interactions that looks something like Fig. 3.6.

To begin the showering process, the initial state consisting of the four-momenta of the original partons must be specified. Typically, this is done by either providing the calculated transition matrix elements of the original distributions for some number of partons, known simply as the Matrix Element approach, or beginning with two partons and allowing them to shower based on splitting probabilities, known as the Parton Shower approach*

Matrix Elements for Parton Production

As discussed in Section 3.1.2 the differential cross section for radiating a single gluon in events of the type $e^+e^- \rightarrow q\bar{q}g$ (see Figure 3.3) is given by

$$\frac{d\sigma}{dx_1 dx_2} = \sigma_0 \frac{2}{3} \frac{\alpha_s}{\pi} \frac{x_1^2 + x_2^2}{(1 - x_1)(1 - x_2)},$$
(3.124)

where σ_0 is the lowest order Born cross section, and $x_i = 2E_i/E_{CM}$ for i = 1: \bar{q} , i = 2: q, i = 3: q.

Operationally, the writer of a Monte Carlo simulation can avoid the infrared singularities of the matrix element by imposing some sort of cut-off in the minimum invariant mass between any two of the partons. This corresponds to modifying the the definition of what constitutes a 2- vs. a 3-parton event depending on the energy of the radiated gluon.

As discussed above, some of the second- and third-order QCD matrix elements have also been calculated, yielding the cross sections for 4-jet $(e^+e^- \rightarrow q\bar{q}gg)$ or $e^+e^- \rightarrow q\bar{q}q'\bar{q}'$ [87] and 5-jet $(e^+e^- \rightarrow q\bar{q}ggg)$ or $e^+e^- \rightarrow q\bar{q}q'\bar{q}'g$ [88] events. These can also be used to generate the initial-state parton distributions for comparison with experiment.

The Parton Shower Model

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Rather than attempt to calculate the matrix elements for all of the possible final states to as high an order as possible, one can use the Altarelli-Parisi (LLA) formulation* and consider the possible branchings contained within QCD: $q \rightarrow qg$, $g \rightarrow q\bar{q}$, and $g \rightarrow gg$. Beginning with the initial quark pair, the system can be evolved given the probabilities that gluons are radiated, etc., as a functions of the parton energies and the energies of the radiated particles. This process would look something like that shown in Figure 3.8. Typically, α_s is allowed to "run" during the shower as the energy of the partons decreases.

The shower process is initiated with a value of the energy of parton a and a suitably chosen value of t, where t describes the energy evolution of the system. The virtuality t of the shower is allowed to run down until a branching occurs. Partons b and c are then evolved following the same prescription, and so on. Branchings are allowed until the available Q^2 reaches some cut-off value Q_0^2 , typically of $\mathcal{O}(\Lambda)$. Note that the cutoff Λ_{PS} depends on the actual definition of Q^2 chosen in implementing the parton shower and Λ_{PS} is generally not equal to the A introduced in renormalization.

Different implementations of this model are distinguished by the definitions oft, z, and Q^2 . The implementation developed by the Lund group[105], for example, treats the shower process in the center-of-mass frame, with z being the energy fraction in that frame. The evolution parameter t is given by $t = \log(m^2/\Lambda^2)$, where m is the mass of the parton. To account properly for the interference among the many soft gluons that are emitted[106], an "angular ordering" can be imposed on the shower such that successive gluon emissions occur at smaller relative angles.

[†]Of course, striking advances in field theory could change this situation.

[‡]Note that wc have been using the term "parton shower" to refer to the cascade of partons formed as the initial quarks fly apart from their production point. The Parton Shower model of this cascade refers to a specific prescription for gluon radiation and splitting into quark pairs during the showering process.

^{&#}x27;See Section 3.1.4.



Figure 3.8: Parton Shower evolution, showing two primary quarks and a possible set of their "branching8 .

Models of Hadronization/Fragmentation

Once the initial parton distributions are generated from either the Matrix Element calculations or the evolution of a Parton Shower they must be "converted" into the color-singlet hadrons that can be observed by a particle detector. The process of this conversion is called fragmentation, [†] and must be implemented by some model motivated by appropriate theoretical assumptions, as the processes that form hadrons take place in the non-perturbative regime and are therefore incalculable. There exist three major classes of models for fragmentation:



Figure 3.9: A schematic representation of Independent Fragmentation, showing the quark-antiquark pairs pulled from the vacuum and the energy fractions carried by each successive hadron.

- The Independent Fragmentation model
- The String Fragmentation model
- The Cluster Fragmentation model.

Each of these will be described in some detail in the following paragraphs.

Independent Fragmentation: In this case, each energetic parton in the event corresponding to a well-separated "jet" is fragmented in isolation. Small adjustments are made to the fragmented systems to conserve energy and momentum and to guarantee that quark flavor is conserved. The fragmentation process occurs in the following iterative manner (see Figure 3.9): an initial quark jet q with energy W becomes paired into a hadron $q\bar{q}_1$ carrying energy fraction $W \cdot z_1$, leaving behind a remainder jet q_1 with energy $(1 - z_1)W$. This jet in turn is split into another hadron, $q_1\bar{q}_2$, carrying energy $z_2(1 - z_1)W$, leaving again a remnant jet q_2 . The process continues until the energy is too low to form the lightest hadron. The energy sharing between daughter products is given by a *fragmentation* function, often denoted J(r), which is assumed to be the same at each step in the fragmentation process. The Field-Feynman model[107] was one of the earlier versions of this fragmentation scheme to be implemented in Monte Carlo calculations.

[†]We use the term "fragmentation" here to refer to the specification of both the underlying **parton** distribution and the means by which the observable hadrons are formed, as these different models can be distinguished by their description of both of these processes.

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This simple approach was made increasingly complicated by the introduction of resonance production and decay, strangeness and baryon production, and mechanisms to generate an appropriate amount of transverse momentum relative to the initial quark direction. Various other independent fragmentation schemes of increasing complexity have followed this initial formulation[108, 109]. The most commonly used of these today is the ISAJET Monte Carlo[110].

String Fragmentation: The idea of using a string stretched between the outgoing partons to represent the linear confinement of QCD was first introduced in a Monte Carlo model by Artru and Mennessier[111]. Since then, the Lund group has been responsible for developing the model of string fragmentation as we know it today. Their Monte Carlo simulation, JETSET 6.3[113]**is** used as the basis for the results of this thesis.

In the string fragmentation model, the axis of the color flux tube[‡] which is stretched between two partons moving away from each other is represented by a semiclassical massless relativistic string with no transverse degrees of freedom[114]. The string is assumed to have a uniform energy density per unit length. This gives both a linearly rising confinement potential and a constant probability per unit length to split into new $q\bar{q}$ pairs that allow hadron formation. The splitting is done in such a way that energy, momentum, and all internal quantum numbers are conserved. A schematic view of the string splitting process for massless quarks is shown in Figure 3.10. Massive quarks must actually be produced some distance apart so that the field energy between them can be transformed into mass and transverse momentum. This can be accomplished by creating them at a point and allowing them to quantum-mechanically "tunnel" out to the allowed region. This tunneling process occurs with a probability proportional to [115]

$$e^{-\kappa E_T^2} = e^{-\kappa p_T^2} e^{-\kappa m^2}$$
(3.125)

where κ is inversely proportional to the string constant. This formulation serves to generate the transverse momentum spectrum of hadrons relative to the original quark



Figure 3.10: A schematic representation of String Frngmentation. In (a), the two primary quarks are shown as they move away from their common production point. The shaded region is the physical extent of the non-vanishing color field. The first three breakings of the string are shown in (b), with the quark-anti-quark pairs labeled in the order that they appeared. The further evolution of the shower is shown in (c). As the quarks which form the observable mesons are produced with some transverse momentum relative to each other, they oscillate about their center of mass, which produces the sequential rectangles seen as the mesons move towards the top of the diagram. The actual meson trajectory is the line drawn through the rectangle vertices.

[‡]Since the gluons exchanged between the two quarks are able to interact with each other, one can visualize the color "field lines" as pulled into a narrow tube rather than spreading out to infinity as would be the case in QED.



Figure 3.11: A $q\bar{q}g$ event in the string fragmentation scheme. The heavy line is the string stretched between the gluon and the two quarks.

direction, and specifies by virtue of the relative quark masses the flavor composition of the $q\bar{q}$ pairs created when the string is broken. Charm quarks, for example, are suppressed by a factor of ~ 10⁻¹¹ relative to the light *u* and *d* quarks and are essentially never produced in the string breakup.

These simple assumptions of the string picture lead to a number of interesting features. Since the hadron production vertices are causally disconnected, the time-ordering of the production is Lorentz frame dependent. This implies that all of the vertices must be treated identically, since no one branching is more primary than another, and leads to a an explicit form for the fragmentation function f(z)[§]. The kinematics of string-breakup also lead to an ordering (OR average) of hadrons in rapidity with respect to the primary quarks, as the slower hadrons are produced earliest.

The treatment of gluons is worth noting. The gluon is treated as a "kink" in the string that is stretched between the two quarks (see Figure 3.11). This makes intuitive sense, since the gluon carries two color indices[¶]. One consequence of this formulation is that the multiplicity of hadrons will tend to be larger between the gluon jet and the two quark jets compared with the region between the two quark



Figure 3.12: A diagram of Cluster Fragmentation.

jets. The first observation of this so-called "string effect" was reported by the JADE collaboration in 1980[116], and, after more years of study, is now an established part of the fragmentation phenomenology[117].

Cluster Fragmentation: In the Cluster Fragmentation model, the Parton Shower evolution is given the primary role of producing the underlying structure of the final state hadrons by being allowed to evolve into many gluons and quarks. After forced splitting of the gluons into quarks, the quarks are combined into colorless "clusters", which are then decayed isotropically in their rest frames, typically using some simple phase-space model. The assumption of the Cluster Fragmentation model, then, is that all of the properties of the hadronic final state should be determined by the perturbative QCD, rather than some ad hoc model of non-perturbative fragmentation. A diagram of Cluster Fragmentation is shown in Figure 3.12.

Fragmentation functions in general will be discussed below.

⁴The ratio of the strengths of the color fields between quarks and gluons should be $2/(1-1/N_C^2) = 9/4$, where N_C is the number of colors. The value of 2, then corresponds to an infinite number of colors.

The first widely successful model of this type was developed by Marchesini and Webber[118], which included a "coherent" parton shower approach[119] containing previously ignored gluon interference effects[106]. The present version of this program is known as HERWIG[120].

Fragmentation Functions

We introduced above the function f(z) which governs the sharing of energy between the daughter hadron and the remnant jet that remains as the fragmentation process progresses. Here, we present without derivation the two fragmentation functions used to generate the Monte Carlo events analyzed for this thesis. The Lund symmetric fragmentation function[121] can be derived by requiring the left-right symmetries involved in breaking a string starting at the quark or the antiquark side, and is given by

$$f(r) \propto z^{-1} (1 - \frac{bE_z}{z}) c^{-T}$$
 (3.126)

Here, a and 6 are parameters which must be determined from experimental data. A fit, to the observed hadron spectra in $e+e^{-}$ data gives typical values of a = 0.5, $6 = 0.9 \,\text{GeV}^{-2}[104]$.

For events containing heavy quarks, experimental results indicated that a "harder" fragmentation function was necessary. The so-called Peterson form[122] has become the standard parametrization:

$$f(z) \propto \frac{1}{z \left(1 - \frac{1}{z} - \frac{\epsilon}{1 - z}\right)}$$
(3.127)

Typical values for ϵ are $\epsilon_c = 0.06$ and $\epsilon_b = 0.006[123]$. These values and the Peterson fragmentation function were used to generate the heavy quark Monte Carlo events for this thesis.

The remainder of this chapter is divided somewhat arbitrarily into sections on techniques for measuring the strong coupling α_s and methods for testing the properties of QCD outlined above. The division is arbitrary in the sense that measuring α_s with different techniques and comparing the results actually provides a test of the calculational methods used to derive the theoretical predictions used to make the measurement[124].

3.3 Measurements of α_s

To begin the discussion of techniques for measuring the strong coupling, we remind the reader of the matrix clement for $e^+e^- \rightarrow q\bar{q}g$ events shown in Eq. 3.94

$$\frac{da}{dx_1 dx_2} = \sigma_0 \frac{2\alpha_s}{37r(1-x_1)(1-x_2)}.$$
(3.128)

As might be expected, the rate for the radiation of a gluon is indeed proportional to the coupling α_s . So, a measurement of the rate of single gluon emission should, at this simplistic level, allow the derivation of a value for the strong coupling. From the above discussion on QCD phenomenology, it should be obvious that the single gluon is not observable, and that instead the final state will contain some number of hadronic jets. So, measuring the rate of single gluon emission corresponds experimentally to measuring the rate of events that appear to have three jets. A discussion of what this actually means will follow. We present this first as it is not only the most intuitive method, it is the one used in this thesis.

3.3.1 Jet Rates

As was mentioned in the discussion following Eq. 3.94, the differential cross section diverges as x_1 or x_2 , the fraction of the beam energy carried by the outgoing quarks, approaches unity. This corresponds to the instance when soft gluons are emitted along the quark direction but are still counted as independent partons. When combined with the lowest order propagator and vertex corrections, the result actually is finite, but this behaviour points to a common resolution to the problem of infrared divergences. To sidestep the collinear divergences, a mass cut-off can be added to the theory, such that any two partons whose combined invariant mass is less than some parameter y_{min} are considered as an unresolvable single parton. This makes particular **sense** when one is dealing with a final hadronic state containing jets rather than individual partons, as only jets coming from partons with high relative transverse momentum are distinguishable experimentally. The first. measurement of α_s using the rate of "resolvable" 3-jet events was published by the PLUTO collaboration in 1980[125].

All of the jet algorithms currently in use are based on a procedure of iterative cluster-finding. First, the user specifies the manner in which the invariant "mass" y_{ij} should be calculated. An example, from the algorithm developed by the JADE collaboration[126], is

$$Y_{ij} = \frac{2E_i E_j (1 - \cos \theta_{ij})}{s}, \qquad (3.129)$$

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where E_i is the energy of the i^{th} particle, and θ_{ij} is the angle between particles *i* and *j*. The algorithm then considers all of the particles in the event and combines the two with the lowest y value into a single pseudo-particle following some prescription, like $p^{\mu} = p_i^{\mu} + p_j^{\mu}$. The process is repeated using the remaining particles until all of the (pseudo)- particles have an invariant mass greater than y_{min} (also referred to as y_{cut}) when combinations with all other clusters are tried. The number of clusters remaining is then the number of jets in the event.

Initially, this manner of measuring α_s was limited by the omission of large nextto-leading terms from the theoretical calculations. Results up to second order in α_s have now been included, and attempts have been made to reduce the dependence on the renormalization parameter μ as well on the uncalculated higher-order terms[127]. The general form of the predictions for the 3-jet rate σ_{3jet}/σ_0 is[128, 129]

$$\frac{\sigma_{3jet}(y_{cut})}{\sigma_0} = \frac{\alpha_s(\mu)}{2\pi} A(y_{cut}) + \left(\frac{\alpha_s(\mu)}{2\pi}\right)^2 \left[B(y_{cut}) + A(y_{cut})2\pi b_0 \log f\right].$$
(3.130)

Here, $f = \mu^2/s$, and $b_0 = (33 - 2n_f)/(12\pi)$, where n_f is the number of active flavors $(n_f = 5 \text{ at } \sqrt{s} = M_Z)$. The terms proportional to *B* arise in part from 4-jet events where two of the jets were unresolvable. The 4-jet rate σ_{4jet}/σ_0 has only been calculated to lowest order[87], and has the expected form:

$$\frac{\sigma_{4jet}(y_{cut})}{\sigma_0} = \left(\frac{\alpha_s(\mu)}{2\pi}\right)^z C(y_{cut}) \quad . \tag{3.131}$$

In general, the coefficients $A(y_{cut})$ and $C(y_{cut})$ are the same for all algorithms, since there should be no dependence of the lowest order result on the choice of some specific jet-finding algorithm. It should be noted that these results are a comparison of the second-order jet production rates with the usual Born cross section σ_0 , where

$$\sigma_0 = \frac{4\pi\alpha_s^2}{3s} N_C \sum_j Q_j^2 . \qquad (3.132)$$

The tree-level cross section is also modified by QCD corrections, which have actually been calculated to third order in α_s for studies of the Z^0 hadronic width[128]. The second-order result in the $\overline{\text{MS}}$ scheme (A $\equiv \Lambda_{\overline{\text{MS}}}$) is[130]

CHAPTER 3. THE STRONG INTERACTION: AN INTRODUCTION

$$\frac{1}{\sigma_0}\sigma_{tot} = 1 + \frac{\alpha_s(Q^2)}{\pi} + \left(\frac{\alpha_s(Q^2)}{\pi}\right)^2 (1.986 - 0.115n_f), \qquad (3.133)$$

where the corresponding second-order value of a, was given above in Eq. 3.118. An experiment measuring jet production rates necessarily measures $\sigma_{njet}/\sigma_{tot}$, so these corrections should be taken into account when measurements of α_s arc derived from these quantities.

3.3.2 Event Shape Variables and Particle Correlations

Various other properties of hadronic events in e^+e^- annihilation events can also be calculated in perturbative QCD. Theoretical predictions similar in form to Eq. 3.130 exist for quantities related to various distributions of particle momenta about the jet (or some other) axis and the flow of hadrons in the event; measurements of α_s can then be derived from the observed distributions. Rather than list and describe all of the other techniques for measuring α_s here, we refer the reader to Ref. [128], which presents an overview of hadronic observables in e^+e^- annihilation at the Z^0 pole. A recent paper from SLD [131] contains a comprehensive set of measurements of α_s using these techniques.

3.4 Tests of QCD

The emergence of Quantum Chromodynamics as a viable theory of hadronic interactions (marked by the discovery of asymptotic freedom) coincided almost exactly in time with the final piece of evidence that the quark model of hadronic structure was correct (the discovery of the charmonium states in 1974[132]). The predictions of QCD were now subject to experimental scrutiny, as new experimental results were compared to all of the different possible models of the strong force in order to determine the "correct" theory of the strong force*.

In addition to the crucial success in describing the phenomenology of deep inelastic scattering, the first observation of jets in e^+e^- annihilation[103] provided another early triumph of the theory, as jets are essentially a result of quark confinement. As mentioned above, the observation of 3-jet events at PETRA[81] gave the first direct evidence for the gluon, though neutral, chargeless hadronic constituents had been required earlier to make the quark-parton model agree with deep inelastic scattering data[134]. It was also shown quickly from the angular distribution of the gluon jets that the radiated particle was consistent with having spin 1[135]. Another important result from PETRA was the first evidence that α_s "runs" (is Q^2 dependent) from studies of jet-production rates at different center-of-mass energies[136]. Direct evidence for gluon self-coupling was found in the orientation of the jets in 4-jet events found at TRISTAN[137]. Even as these basic tenets of QCD have been verified, efforts continue in testing the theory on many fronts.

3.4.1 Flavor Independence of α_s

One active area of investigation is tests of the flavor independence of the strong coupling. In principle, the QCD Lagrangian, instead of looking like Eq. 3.91, could have a sum over flavors:

$$\mathcal{L} = \sum_{j}^{n_{f}} \left[\bar{q}_{j} (i\gamma^{\mu}\partial_{\mu} - m_{j})q_{j} - g_{s}^{j} (\bar{q}_{j}\gamma^{\mu}T_{a}q_{j})G_{\mu}^{a} \right] - \frac{1}{4} n_{f}G_{\mu\nu}^{a}G_{a}^{\mu\nu} .$$
(3.134)

However, this formulation violates the Ward-Takahashi identity[138], which, by virtue of color conservation, requires that the coupling in the three-gluon vertex be the same as that in the $gq\bar{q}$ vertex. A brief discussion of the implications of the Ward-Takahashi identity and its importance is in order.

The Ward-Takahashi Identity

To introduce the identity, we return to the more familiar environs of QED. Gauge and Lorentz invariance of the electromagnetic field A, requires that free photons have only transverse polarization states, i.e., that the condition

$$k_{\mu}\epsilon^{\mu}(k) = 0 \tag{3.135}$$

can be reduced to

$$\vec{k} \cdot \vec{\epsilon}(k) = 0 , \qquad (3.136)$$

where k is the photon momentum, and $\epsilon(k)$ is the photon polarization vector. This can be seen in the following manner. The free photon field has the form $A_{,} = \epsilon(k)_{\mu}e^{-ik\cdot x}$, and $A_{,}$, must be invariant under the transformation

$$A_{,,} \to A'_{\mu} = A_{,,} + \partial_{\mu}\chi \tag{3.137}$$

where χ satisfies $\partial^2 \chi = 0$ to insure that $\partial^2 A_{\mu} = 0$. We are free, then, to choose a gauge parameter

$$\chi = iae^{-ik \cdot x} \tag{3.138}$$

where a is some constant. Under the conditions of Eq. 3.137, the physics must be unchanged if the polarization vector ϵ_{μ} is replaced by

$$\epsilon_{\mu} \to \epsilon'_{\mu} = \epsilon_{\mu} + ak_{\mu} \tag{3.139}$$

This implies that two polarization vectors which differ by a multiple of k_{μ} describe the same photon. This degree of freedom implies that there are only two independent polarization vectors to describe the photon; we can require the time component ϵ^0 to vanish identically, $\epsilon^0 \equiv 0$. Then, requiring the field *A*, to satisfy Maxwell's equations for a free field gives us the further condition that the polarization must be perpendicular to *k*, which results in the condition of Eq. 3.136.

We can now examine some of the consequences of the relationships derived above from gauge invariance. Given any amplitude \mathcal{M}_1 including an external photon of

^{&#}x27;An early indication that color (or some other new quantum number allowing three independent degrees of freedom) was necessary in hadronic interactions was provided by the rate of the reaction $\pi^0 \rightarrow \gamma\gamma$. A calculation of this rate proved to be incompatible with experimental observations unless each "quark" that contributed to the final amplitude was counted three times[133]. However, this was realized before the full applicability of QCD became apparent.

polarization $\epsilon_{\mu}(k)$, the following relation must hold[†]:

$$k_{\mu} \cdot \mathcal{M}_{\mathbf{1}}^{\mu} = 0 \tag{3.142}$$

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We have used the subscript 1 here to denote the amplitude including the external photon. Although we have not proven it here, this result is true to all orders of perturbation theory. If we define the amplitude \mathcal{M}_0 as the same set of diagrams *without* the photon, then we obtain \mathcal{M}_1 by summing the diagrams \mathcal{M}_0 with the additional photon inserted at, every possible point. For an amplitude \mathcal{M}_1 with n incoming and *n* outgoing fermion lines (and an arbitrary number of external photons), the general form of the Ward-Takahashi identity is given by[‡]

$$k_{\mu}\mathcal{M}_{1}^{\mu} = q_{f} \sum_{i} \left[\mathcal{M}_{0} \left(p_{1} \cdot p_{n}; q_{1} \dots (q_{i} + k) \cdot \cdot \right) - _{\mathsf{MO}} \left(p_{1} \cdot \cdot (p_{i} + k) \cdot \cdot ; q_{1} \cdot \cdot \cdot q_{n} \right) \right]$$
(3.143)

Here, q_f is the fermion charge, the p_l (l = 1, n) are the momenta of the incoming fermions, and the q_l are the momenta of the outgoing fermions. In the case that \mathcal{M}_1 is an amplitude for some physical process, the two amplitudes \mathcal{M}_0 do not contribute, and the original form of the identity (Eq. 3.142) is recovered.

The simplest example is shown in Figure 3.13a, where we have a free quark propagating as the fundamental process Mo. If we add an external photon, this gives the diagram on the right hand side, which we denote \mathcal{M}_1 . In this case, there is only one diagram in the sum over insertion points, since there is only one place we can attach the photon to the quark line. From Equation 3.143, we find the situation shown in Fig. 3.13b. The components of these graphs are simple, as they are given

$$\mathcal{M}_{\mu}(p) \propto \int dx e^{-ip \cdot x} \langle f|j_{\mu}(x)|i\rangle$$
 (3.140)

Since $\partial_{\mu} j^{\mu} = 0$ by current conservation,

$$k_{\mu}\mathcal{M}^{\mu} \propto \int dx e^{-ip \cdot x} \langle f | \partial_{\mu} j^{\mu}(x) | i \rangle = 0 . \qquad (3.141)$$

[‡]For a straightforward derivation, see Ref. [140], Section 7.4



Figure 3.13: An illustration of the amplitudes involved in a discussion of the Ward-Takahashi identity. In each case, the diagrams labeled \mathcal{M}_0 represent the amplitude for some process without the addition of an external photon or gluon. The diagrams labeled \mathcal{M}_1 make up the amplitude after the external photon or gluon has been added. In (b), a diagramatic representation of the identity is given for the simplest case. See the text for a description.

by fermion propagators and the vertex function. Let us discuss the ramifications of the Ward-Takahashi identity on the renormalized mass and charge in QED by using the renormalized fermion propagator and vertex functions. Then, we haves

$$S(p) = \frac{i}{p - m - \Sigma(p)}$$
(3.144)

for the fermion propagator. Here, C(p) is the sum of all (infinite) one-particle irreducible virtual corrections to the fermion line, and m is the fermion mass. The multiplicative renormalization coefficient Z_2^{\P} is defined by

$$Z_2^{-1} = 1 - \frac{d\Sigma}{d\psi} \qquad (3.145)$$

[†]An alternate way to view the same statement is to consider the actual form of the amplitude \mathcal{M}_1 . It must be proportional to a matrix element containing the electromagnetic current j,,(s), so that we have

[§]These expressions can be found in any advanced field theory text; see, *e.g.*, Ref [140], Chapter ⁷. [¶]See Section 3.1.3 for the lowest order diagrams.

and is the residue of the propagator pole

$$S(p) \sim \frac{iZ_2}{\not p - m} \tag{3.146}$$

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at the physical mass m. The vertex function is denoted by the symbol $\Gamma^{\mu}(p + k, p)$, which represents the sum of all of the corrections to the bare vertex. The renormalization constant Z_1 is given by

$$\Gamma^{\mu}(p=0) = Z_1^{-1} \gamma^{\mu} \tag{3.147}$$

If we set p near the mass shell, an expansion of Eq. 3.143 to lowest order in k gives

or

$$Z_1 = Z_2$$
 (3.149)

This implies that the renormalization of the vertex function and the fermion propagator arc not, independent. If wc recall from our discussion of the renormalization of the strong coupling, Eq. 3.98, that the renormalization of the charge is given by

$$e = \frac{Z_2 \sqrt{Z_3}}{Z_1} e_0 \tag{3.150}$$

we can easily see that the renormalization of the charge is given only by the contribution from the vacuum polarization of the photon, represented by the factor Z_3 . That is, the electric charge renormalization must be universal, a result which is valid to all orders in perturbation theory. Note that different fermions are allowed to have different electric charges; the ratio of the charges can be expressed in terms of a universal constant e.

The next most complicated case is shown in Fig. 3.13b, in which we add an extra photon to a $qq\gamma$ vertex to obtain quark Compton scattering. In this case, the same cancellation must occur. This can be compared to the diagrams of Fig. 3.13c, which show the same process in QCD, with photons replaced by gluons. Although the derivation of the Ward-Takahashi identity is more difficult in Yang-Mills theories[139], the qualitative result is the same: $k_{\mu}^{q}\mathcal{M}_{1}^{\mu} = O_{r}$ where k_{μ}^{g} is the momentum of the

external gluon. In QCD, the presence of the three-gluon vertex here demands that the coupling between gluons must be the same as that between gluons and quarks, or else QCD ceases to be a unitary, renormalizable theory. Thus, the strong "charge" or the magnitude of the strong coupling must be the same for all quark flavors. Again, this must be true to all orders in perturbation theory. It is through this mechanism that modifications to the lagrangian like that of Eq. 3.134 are forbidden.

Other Means of Introducing Flavor Dependence

Undaunted by field-theoretical considerations, we can still entertain the possibilities of different strong coupling strengths for each of the quark flavors, since, if the experimentally measured couplings were to be unequal, new physics must be present. Vastly different couplings among the light quarks can be dismissed by the consideration of the spectrum of hadrons seen in e^+e^- interactions, as a large discrepancy between α_s^u and α_s^d would violate the apparent SU(3) flavor and isospin symmetries observed in nature.

More possible, perhaps, is a difference in the couplings of the light and heavy quarks, through some mechanism involving powers of the quark masses which would suppress the effect for the light quarks. If, for example, a quark had some sort of chrome-magnetic moment, its effective coupling to the strong interaction would be different than quarks lacking this extra vertex correction. This sort of effect has been proposed in the context of studies of the top quark[141], but it is equally applicable here. In this sort of theory, the second term of the QCD lagrangian would be replaced by

$$\mathcal{L} = g_s \bar{q}_j T_a \left(\gamma^\mu + i \frac{F_2(k^2)}{2m_j} \sigma_{\mu\nu} k^\nu \right) q_j G^a_\mu . \tag{3.151}$$

Here, k is the outgoing gluon momentum from the $q_j \bar{q}_j g$ vertex, and $F_2(k^2)$ is a form factor describing the strength of the anomalous coupling. This term does not violate the Ward-Takahashi identity in QCD, and thus is permissible. This is usually parametrized by a constant such as κ , where $F_2(k^2 = 0) \equiv \kappa$. Performing the

calculation of the 3-jet rate with these terms in the lagrangian yields

$$\frac{\sigma_{q\bar{q}g}}{\sigma_{q\bar{q}}} \simeq \frac{\alpha \ \kappa^2 s \ v^2 + 1.25a^2}{\pi \ 18m_q^2} \tag{3.152}$$

where a (v) is the axial (vector) coupling of the quarks in the electroweak interaction. If κ is non-zero for some quark, the measured value of α_s for that quark will be larger than those with $\kappa = 0$. New couplings of this type do not arise in the "usual" extensions of the standard model, so the observation of an anomalous strong coupling for a specific quark flavor would be a sign of truly "new" physics.

Tests of the flavor independence of QCD have been made possible by the accumulation of relatively large data sets containing five quark flavors at PETRA, TRISTAN, and at LEP/SLC. To perform a test of this sort, an event sample containing predominantly the quark flavor of interest must be obtained and an analysis performed on this sample to yield a value of the strong coupling for that quark flavor. Methods of flavor tagging have been developed which allow the isolation of events containing primary quarks of a particular flavor. The separation of heavy (b and c) and light (uds) quark events is the most straightforward due to the long lifetimes, hard fragmentation functions, and large masses of the hadrons that contain heavy quarks. A chronological discussion of the previous tests of flavor independence will serve both to introduce the various flavor-tagging techniques and to summarize the existing results.

The TASSO Collaboration carried out the first tests of the flavor independence of α_s at PETRA. Events containing charm quarks were identified by exclusive reconstruction of $D^* \rightarrow \pi D^0$ meson decays11 and the EEC distribution was used for the measurements of the strong coupling. They obtained[142] $\alpha_s^c / \alpha_s^{all} = 0.91$ f0.38f0.15. It is worth noting that this particular technique of selecting charm decays necessarily requires the charm meson to be carrying a large fraction of the energy available to it, which limits the acceptance for events with very hard gluons. Also, charm events containing the appropriate $D^* \rightarrow \pi D^0$ cascades only make up a small fraction (~ 7%) of the total and the exclusive decay modes used to reconstruct the D^0 have small branching fractions. These factors lead to an undesirably small efficiency for tagging charm events and hence to an undesirably large statistical error. Events containing *b* quarks were tagged by searching for the presence of decay vertices well- separated from the beam interaction point. They obtained [143] $\alpha_s^b/\alpha_s^{all} = 1.17 \pm 0.50 \pm 0.28$ from the EEC distribution. Using lifetime information to tag heavy quarks is made easier due to the long typical flight distances of the heavy mesons and the large available energy in the meson decay. These factors result in sufficient numbers of charged tracks emanating from the meson decay point with relatively large transverse momentum relative to the meson flight direction to make searching for decay vertices, for example, an efficient process. This method of tagging is also much less dependent on the kinematic properties of the event, as the vertex flight. distance resolution is typically much smaller than the mean meson decay length".

The era of experimentation at the Z^0 pole with e^+e^- colliders has allowed tests of the flavor independence of QCD with unprecedented precision. All of the early results from the LEP experiments were produced without the present set of precision microvertex detectors that allow the type of b tagging done with TASSO. Instead, kinematic tags were used to separate events of different flavors. The most common is to use identified leptons with large momentum and large transverse momentum relative to the nearest jet axis as a tag of b or c quarks^{††} Depending on the cuts placed on the lepton momenta, this method could also bias the event sample away from those events containing hard gluon radiation. The lepton tag approach was followed by the L3 and DELPHI collaborations, who found $\alpha_{*}^{b}/\alpha_{*}^{udsc} = 1.00 \pm 0.05 \pm 0.06$ and $\alpha_s^b/\alpha_s^{udsc} = 1.00 \pm 0.04 \pm 0.03$, respectively, from the measured 3-jet rate. The OPAL collaboration[145] performed the first truly comprehensive study of the flavor independence of α_s . They used high p_{\perp} leptons for a b tag, exclusively reconstructed D mesons for the charm tag, fast K_s^0 for a strange tag ($x_K = 2E_K/E_{CM} > 0.4$), and fast pions, protons, and kaons $(0.7 < x_{\pi Kp} < 1.07)^{\ddagger}$ as a tag of light (uds) flavors. The S-jet rate is measured, and a grand unfolding is done to obtain the ratios of

See the other half of this thesis for a discussion of these techniques.

^{**}Although, this was not the case for the TASS0 vertex detector, which explains their large statistical error and relatively small tagging efficiency.

^{††}See Section 2.8.1.

^{‡‡}The upper limit of 1.07 is one standard deviation on the measurement of the momentum of a particle with x = 1.

couplings:

$$\begin{aligned} \alpha_s^b / \alpha_s^{all} &= 1.021 \pm 0.013 \pm 0.023, \\ \alpha_s^c / \alpha_s^{all} &= 0.912 \pm 0.067 \pm 0.061, \\ \alpha_s^s / \alpha_s^{all} &= 1.141 \pm 0.043 \pm 0.142, \\ \alpha_s^d / \alpha_s^{all} &= 0.933 \pm 0.087 \pm 0.175, \\ \alpha_s^d / \alpha_s^{all} &= 0.951 \pm 0.103 \pm 0.182. \end{aligned}$$
(3.153)

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The large systematic errors on the light quark couplings are due to the uncertaintics in basing a tag on the identity of the fastest particle in an event, as this is uncharted territory for many of the Monte Carlo models of hadron production. The large statistical errors on these quantities is a consequence of the inefficiency of the light flavor tags. A recent paper by OPAL has extended the repertoire of techniques used in the measurement of α_s for the flavor tagged samples [146]. They use a detached vertex *b* tag similar to the one developed at TASS0 to obtain a pure sample of *b* events. A large number of measurements of α_s are then performed on this *b* sample, and the results are averaged to state a single value: $\alpha_s^b/\alpha_s^{all} = 0.994 \pm 0.005^{+0.010}_{-0.012}$. This is by far the most stringent test of the flavor independence of α_s for *b* quarks.

One feature of all of these analyses is that a ratio of couplings like $\alpha_s^j/\alpha_s^{all}$ is measured. This has the advantage that it reduces the effects of a number of errors that can plague the determination of α_s . For example, if corrections due to detector resolution or acceptance are essentially the same for each flavor, the uncertainties on the ratio due to these corrections are smaller. Uncertainties due to the choice of renormalization scale μ should also mostly cancel, since the evolution of the parton showers are identical in all events up to the effects of quark masses.

In all but the OPAL analysis involving five flavors, assumptions need to be made about the relative strengths of the strong coupling for the other quark flavors, as no sample of events is 100% pure, and the background being subtracted has some 3-jet rate. The most, common approach is to assume that $\alpha_s^u = \alpha_s^d = \alpha_s^s = \alpha_s^b$ or c and then proceed with the background subtraction.

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Methods Chosen for This Analysis

In an attempt to obtain the best possible statistical precision for all flavors, we have chosen to pursue an "inclusive" analysis to test the flavor independence of the strong interaction. The method is inclusive in that all hadronic events which pass the selection cuts (Chapter 6) are used. We use the precision microvertex detector to separate the flavors based on the charged multiplicity of secondary decays*: there are no decays of heavy secondary particles in uds events, and the average number of secondary tracks in c and b quark events differs by almost 3 tracks. The tag is based on track counting; reconstructed decay vertices are not required, as this results in tags of lower efficiency. Using the number of tracks that are not consistent with originating from the precisely determined position of the interaction point, we can obtain highly enriched samples of events containing primary **uds** and b quarks, with a somewhat enriched sample of c quark events in the remainder. This tagging process is relatively insensitive to the underlying event kinematics, and results in event tagging efficiencies that are quite high. We note in passing that it is the tiny, stable beam spot of the SLC that allows us to tag the *uds* sample. We do not need to obtain the position of the interaction point for every event, a procedure which could introduce large systematic effects for a tag of this sort.

We also perform a jet rates analysis on each of the tagged samples, and the results are unfolded to arrive at values for $\alpha_s^b/\alpha_s^{all}$, $\alpha_s^b/\alpha_s^{all}$, and $\alpha_s^{uds}/\alpha_s^{all}$. Note that by using all events and unfolding to arrive at values for $\alpha_s^j/\alpha_s^{all}$ we only make the weak assumption that $\alpha_s^u = \alpha_s^d = \alpha_s^s$, which, as discussed above, is consistent with the approximate isospin and SU(3) flavor symmetry observed in hadronic structure.

To determine a, ratios from the jet rates measurements, we use six of the commonly used jet-finding algorithms in order to compare the effects of the uncalculated higher-order terms in the the 3- and 4-jet cross sections (Chapter 8). Our results will be discussed and compared with those mentioned above in Chapter 12.

^{*}The method of flavor tag will be explained in great detail in Chapter 7.

Chapter 4

Experimental Apparatus: the SLC and SLD

This chapter presents a description of the unique facility that exists at SLAC for the study of Z^0 boson physics. Added detail will be included in the sections that are relevant to the analyses presented here.

4.1 The SLAC Linear Collider

The SLAC Linear Collider (SLC) is the world's first linear collider, a an accelerator concept which gained prominence [147] in the 1970's as perhaps the only economically feasible way of pushing e^+e^- colliders to the high energy frontier. The prototypical linear collider consists of two linear accelerators (linacs) producing high-intensity beams which impinge on a common interaction point (IP) which is surrounded by the particle detector that will record the interactions produced in the collisions. In their inimitable SLAC fashion, the designers of the SLC, when faced by budgetary and spatial constraints, produced a linear collider folded back on itself, as shown in Figure 4.1. Both the e^+ and e^- beams are accelerated by the linac and then the beams are split, bent 90°, and forced to collide head-on. One of the salient features of a linear collider is that the beam bunches cannot be immediately re-used



Figure 4.1: The Polarized SLC, showing the overall layout of the accelerator complex. The orientation of the electron spins is given by the arrows along the path of the beam. SLD sits at the IP.

for collisions as they are in a storage ring collider due to the violence of the beambeam interaction. Some schemes exist for recirculation of the beams[148], but here the spent bunches are absorbed in a beam dump after each collision. Another "feature" is the low repetition rate of the accelerator, which is limited by power constraints as well as the maximum operational frequency of the pulsed-magnet "kickers" which perform nearly instantaneous steering of the beams along the desired orbits. In order to obtain a sufficiently high interaction rate in the beam collisions, it is necessary to focus the particle beams to incredibly tiny cross-sectional areas, which is accomplished by a complex optical system known as the "Final Focus" immediately preceding the interaction point. Some of these features will be discussed in more detail in the following sections.

4.1.1 The Polarized Electron Source

The SLC is unique in its ability to accelerate, transport, and collide a longitudinally polarized beam of electrons. This is made possible by the introduction of Gallium-Arsenide (GaAs) as the active material[149] in a photocathode- based electron gun[150]. A circularly polarized laser is used to selectively excite transitions into longitudinally-polarized states in the conduction band. An energy state transition diagram is shown in Figure 4.2. In 1992, a bulk GAs cathode was used which had a theoretical maximum polarization of 50%. The average polarization measured[151] was 22%. In the 1993 run, a new strained-lattice cathode[152] using GaAs grown on a GaAsP (Gallium-Arsenide-Phosphide) substrate was employed, yielding an average polarization of ~ 65% at, the source (see Figure 4.2). A new cathode with a thinner GaAs layer is currently yielding average polarizations of ~ 80%.

4.1.2 Beam Transport

A single accelerator cycle begins with the production at the source of two bunches of electrons each containing approximately 5×10^{10} particles. These are captured by the accelerator, and accelerated up to an energy of 1.19 GeV at which time they are stored in the North damping ring. The damping ring reduces the size of the

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Figure 4.2: The energy state diagram for bulk GaAs (top) and the changes it undergoes when the lattice is strained (bottom). The relative sizes of the matrix elements governing the interstate transitions are given in the circles. For excitation of electrons to the $m_j = \pm 1/2$ state, the theoretical maximum polarization is 50% for bulk GaAs (top). For the strained lattice, the degeneracy between the $|m_j| = 3/2$ and the $|m_j| = 1/2$ valence states is broken, allowing the potential for 100% polarization.

beam phase-space through the emission of synchrotron radiation and the application of radiofrequency power. After a large number of turns, they are extracted from the ring and follow a positron bunch from the South damping ring down the linac. Approximately two-thirds of the way down the linac, the second electron bunch is diverted onto a target to produce positrons, which are then captured and sent back to the South damping ring to await the next accelerator cycle. After reaching their maximum energy of 46.7 GeV^{*}, the production electron and positron bunches are split apart in the Beam Switchyard by a bending magnet and follow the curved, terrain-following arcs around to the IP, where they collide. Each outgoing beam is then steered onto a beam dump, and the cycle begins again. The repetition rate is 120 Hz. More description of the accelerator and the accelerator-detector interface and interaction will be provided in Appendix **C**.

4.1.3 Spin Transport

To preserve the longitudinal polarization of the electron beam, a series of spin rotation solenoids were installed. The one immediately upstream of the North damping ring is crucial, as it rotates the electron spins into the vertical plane so that they will not precess during their time in the ring. Any precession of this type would result in a large decrease in the average polarization of the beam, since particles with different energies have different spin-precession rates. Since the time spent in the damping ring is relatively long, any non-zero precession rate would have a long time to depolarize the beam.

The other two spin-rotator solenoids were used in 1992 to provide an arbitrary orientation of the electron spin at the IP. However, in 1993 the operational mode changed to so-called "flat beams", whose vertical size is much smaller than the horizontal[209]. Since the spin rotator solenoids downstream of the damping ring would mix the horizontal and vertical beam motions, potentially destroying the small vertical spot, they were turned off. Instead, the fact that the SLC arcs have a betatron oscillation frequency very close to a spin-precession resonance frequency is used to manipulate the





Figure 4.3: A diagram of the WiSRD energy spectrometer, showing the beam components and the WiSRD detector.

spin orientation[153]. This has been extremely successful.

4.1.4 Beam Energy Measurement

The beam energy is measured on every pulse by two spectrometers that are placed just prior to the outgoing beam dumps where the spent beams arrive after the collision. The actual energy measurement is performed by deflecting each beam horizontally, then vertically by a precisely-calibrated bend magnet, then horizontally again. The horizontal bends produce two synchrotron radiation swaths whose positions are measured by the Wire Imaging Synchrotron Radiation Detector (WISRD)[154]. The vertical distance between the two stripes is inversely proportional to the beam energy, which can be extracted given the integrated field of the precision bend and the distance to the detector. A schematic view of the WISRD spectrometer is shown in Figure 4.3. The center of mass energy for the 1993 run was $91.26 \pm 0.02 \text{ GeV}[47]^{\dagger}$.

^{*}energy loss due to synchrotron radiation in the arcs reduces the beam energy to the desired value of $M_Z/2$.

[†]Since this is not exactly the Z^0 mass, we will have to account for the difference in energy as part of the electroweak radiative corrections to the charm asymmetry measurement.

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4.1.5 Beam Polarization Measurement

All of the electroweak asymmetries measured at the SLC depend on a precise knowledge of the longitudinal polarization of the electron beam. The polarization is measured immediately downstream of the SLC IP by a Compton Polarimeter, which uses the difference in the Compton scattering cross sections[155] for the $J_z = \frac{3}{2}$ and $J_z = \frac{1}{2}$ combinations of circularly polarized light and the longitudinally polarized electron beam to extract the electron beam polarization. The measurement is based on the dependence of the scattering cross section (σ_C) on helicity and on the electronphoton collision angle in the center-of-mass frame' When the electron-photon collision is boosted back into the laboratory frame, the differences in scattering probability as a function of scattering angle become observable as an asymmetry in the scattered electron energy spectrum. The shape of the asymmetry function does not depend on the polarizations of the electron or photon beams, as it is only a function of the energy in the center-of-mass system. The unknown electron beam polarization P_e can be extracted from

$$A_{meas} = \frac{\sigma_C(J_z = \frac{3}{2}) - \sigma_C(J_z = \frac{1}{2})}{\sigma_C(J_z = \frac{3}{2}) + \sigma_C(J_z = \frac{1}{2})} = a_d P_e P_\gamma A_C(E)$$
(4.154)

where P_{γ} is the (measured) laser polarization, E is the (measured) scattered electron energy, A_C is the (calculated) Compton scattering asymmetry, $A_{,,,,}$ is the observed asymmetry, and a_d is the (calculated) analyzing power of the detector which actually measures the scattered electron energy and determines the asymmetry. The errors on the beam polarization measurement come primarily from uncertainties in the net circular polarization of the laser and the analyzing power of the tletcetor.

The polarimeter system[156] is shown schematically in Figure 4.4. The photon beam is provided by a frequency-doubled YAG laser which produces 2.33 eV photons which are transported down into the SLC to collide with the 45.6 GeV electron beam at the Compton IP, which is 33 meters downstream from the SLC IP. The electron energy is measured by a spectrometer consisting of an analyzing bend magnet and a multi-channel Čerenkov detector. A tremendous amount of effort, has been



Figure 4.4: A schematic diagram of the Compton Polarimeter system, showing the positions of the components relative to the SLD detect or.

expended in understanding the measurement's provided by the polarimeter system to a precision of better than 1%. The results of t his work are shown in a somewhat symbolic form in Figure 4.5, where it can be seen that the values of the Compton asymmetry measured in each channel of the Čerenkov detector agree absolutely with the theoretical calculation and a Monte Carlo simulation of the detector response to within 0.5%, channel-by-channel.

The measured polarization for the 1993 run is $63.0 \pm 1.1\%$.

[†]The $J_z = \frac{3}{2}$ state has a larger probability for complete backscattering collisions.



Figure 4.5: The measured Compton asymmetry, compared with the theoretical calculation as modified by the EGS[213] simulation of the Čerenkov detector response. The lower plot shows the absolute channel-by-channel residual between the measurements and the calculated values.

4.1.6 SLC Performance History

Construction of the SLC began in 1983 and concluded in 1987. After a somewhat difficult commissioning period, the first Z^0 boson ever produced in e^+e^- collisions was observed by the Mark II detector in April of 1989. Over the course of two runs spanning 1989 and 1990 the Mark II recorded 826 Z^0 decays, with peak luminosities of the SLC reaching 4 $Z_{\rm S}/{\rm hr}^{\rm s}$. The Mark II era ended on November 21, 1990.



Figure 4.6: The history of the SLC Luminosity performance, along with the achieved polarizations to date. The arrow labeled "Design" is the proposed SLC luminosity, with the accelerator running at 180 Hz colliding bunches of 7 x 10^{10} particles.

The SLC Large Detector (SLD) program began with an engineering run in 1991, during which luminosities of 6 Zs/hr were achieved during 60 Hz operation. The improvement in luminosity is mostly due in this instance to the increased magnetic field provided by the new superconducting final quadrupole triplets which perform the final focussing of the beams just prior to the IP. In 1992, SLD had its first physics running, recording approximately 1000 Z^0 decays with unpolarized beams, and 10,000 Z^0 events with an average polarization of 22%. Peak luminosities were around 30 Zs/hr. During a machine physics program at the conclusion of the 1992 run which was aimed at studying final focus systems for the next generation of linear colliders[158], it was realized that the linac was capable of transporting the naturally flat[¶] beam from the damping rings without much growth in size due to added random

[§]So as to not be depressed by the actual luminosity achieved by the collider, SLC luminosities are typically quoted in "Zs/hr" instead of particles/cm²/s. Assuming a peak Z⁰ production cross section of 30 nb, a Z/hr is roughly equivalent to 9.2×10^{27} cm⁻²s⁻¹.

[¶]Due to the horizontal bends of the damping ring, the beam size in the vertical plane σ_y is typically much smaller than σ_x unless the x and y motions are mixed.

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Figure 4.7: The history of the SLC beam size/area, showing the horizontal and vertical sizes of the beams, and the combined beam area. The design area is also shown.

beam motion. This reduced the area of the beam cross section by approximately a factor of two with essentially no changes to the accelerator, and was chosen as the default operating mode starting with the 1993 run. The SLC produced approximately 50,000 Z^0 bosons that were recorded by the SLD during the 1993 run, and peak luminosities approached GO Z/hr. A history of the luminosity provided by the SLC is shown in Figure 4.G. The exponential growth of the luminosity continues, as, at the time of this writing, peak luminosities near 90 Zs/hr have been seen^{||}. This next jump has been made possible by the installation of new final focus elements that cancel some of the chromatic aberrations that limited the vertical spot size. The historical progression of spot sizes in the horizontal and vertical planes is shown in Figure 4.7. The small, stable luminous region of the SLC will play a large part in the analysis of this thesis. The history of the integrated luminosity for the 1991 through

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Figure 4.8: The Integrated Luminosity for SLD running from 1991 to 1993, displayed as luminosity per week (left axis) and total luminosity (right axis) as delivered by SLC and recorded by SLD.

1993 runs of the SLD experiment is shown in Figure 4.8.

4.2 The SLC Large Detector

The SLC Large Detector (SLD) was first proposed[159] in 1984 and was completed in 1990. It was rolled onto the SLC beamline in time for the summer 1991 engineering run, and has been recording data ever since. The SLD possesses the cylinder-endcap geometry typical of most collider detectors. An isometric view of the SLD is shown in Figure 4.9, where the separate elements which make up the detector as a whole are labeled. The individual detector elements, beginning at the center and working

It is heartening to see that the SLC is now capable of producing in a day more Z^0 events than the Mark II recorded in its grueling two years of running.



Figure 4.9: An isometric view of the SLD detector, showing the layout of the detector components.

outwards, are:

- The Vertex Detector (VXD)
- . The Luminosity Monitors (LUM)
- . The Central Drift Chamber (CDC)
- The Cerenkov Ring Imaging Detector (CRID)
- The Liquid Argon Calorimeter (LAC)
- The Solenoidal Magnet



Figure 4.10: A quadrant view of the SLD detector, giving the overall dimensions of the detector components.

. The Warm Iron Calorimeter (WIC)

The endcaps which complete the hermetic coverage of the detector are organized in a similar fashion, except that the VXD (and the solenoid!) has no endcap components. Each of these detector elements will be described briefly in the following sections with

emphasis on the barrel systems, as the endcaps (excepting the LAC) were not used in these analyses. Special attention will be focused on the tracking systems, as their performance lies at the heart of the analyses presented in this thesis. A quadrant' view of the detector including the physical dimensions is shown in Figure 4.10.

The SLD coordinate system takes z to be along the beam axis, with z = 0 defined as the center of the CDC. The xy plane is perpendicular to the beam axis, with the xaxis parallel to the ground. The angle θ will be used often in this thesis, and is defined as the angle between a vector or line drawn from the IP to the object in question and the beam axis. We will define positive $\cos\theta$ to mean that the angle in question is smaller than 90" when measured with respect to the outgoing electron beam, although the official coordinate system of SLD signs this in the opposite manner.

4.2.1 The VXD and Beampipe

The beampipe section which lies at the very center of the SLD is fashioned from a thin beryllium cylinder. This section of pipe was made 25 cm long so that the full tracking volume of the SLD would be preceded by a minimum of material. The beampipe's outer diameter is 25 mm. Including the VXD cooling jacket, the total thickness in radiation lengths before the first layer of the VXD is 0.71%, which is smaller than the thickness of each VXD layer. This material is kept as small as possible to minimize the amount of scattering each track undergoes before the first position measurement can bc performed.

The VXD uses Charge-Coupled Devices (CCDs) as the medium for detecting the deposition of ionization from through-going charged particles. Since a single CCD is composed of a large number of tiny pixels, a detector based on this technology is a source of three dimensional space points along the track trajectory.

The VXD is constructed [160] from sixty 9.2 cm-long ladders arranged in four concentric cylinders which are held in place by a beryllium shell. An end-on view of the detector is shown in Figure 4.11. Eight CCDs are mounted on each ladder, with four on each side to maintain full coverage in $\cos \theta$. One can see from Figure 4.11 that most of the tracks passing through the VXD will acquire two VXD hits. Some



Figure 4.11: An end view of the VXD, showing the four layers of CCD-carrying ladders. The gaps between ladders result in what is effectively a two-layer tracking device. The innermost layer is 28.5 mm from the center of the detector.

tracks can hit more than two CCDs, and the average number of VXD hits per track is 2.3. Two hits are possible on any track for $|\cos\theta| < 0.74$. The inner layer of CCDs is 29.5 mm from the IP, and the outer layer is 41.5 mm away. Each layer comprises 1.1% of a radiation length (X_0) in material.

Each CCD is approximately 1 cm square, and contains 375x578 pixels, each 22 μ m-square. Each pixel has a depletion depth of 20 μ m, which allows excellent position resolution even for tracks passing through the detector at large dip angles. This fine granularity provides robustness against tracking inefficiency due to large backgrounds in the detector. This is a useful feature in the SLC environment, since the slow readout speed of the CCDs necessitates a readout time of approximately 160 ms (19 beam crossings), during which time any charge deposition in the detector is recorded. Even in this high-noise environment, occupancies above 0.1% are extremely rare. Upon installation, it was discovered that two of the ladders and an assorted CCD or two were completely inoperative due to inaccessible bad connections. This resulted in

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approximately 4% of channels remaining inoperative for the lifetime of the detector.

An internal alignment of the positions of all of the CCDs relative to their ladders and the relative positions of the ladders themselves was performed using charged tracks from Z^0 hadronic decays. To check this procedure, tracks with three VXD hits were used to obtain the intrinsic position resolution (including alignment errors) of each hit. The measured single-hit resolutions are 5.5 μ m in the xy plane, and 5.5 μ m up to 9 μ m for central tracks and those with $|\cos \theta| > 0.55$, respectively[160]. These values are consistent with those found by using the two-track miss distance in $Z^0 \rightarrow \mu^+\mu^-$ and $Z^0 \rightarrow e^+e^-$ events fitted with the VXD hits only. The slight degradation in the resolution at larger $\cos \theta$ is most likely due to radial alignment errors and bowing of the CCDs.



Figure 4.12: A beam's eye view of the Luminosity Monitor (LUM), showing the pad structure and the readout cards. The inner diameter is indicated in the center.

4.2.2 The LUM

The luminosity monitors for the SLD are silicon-tungsten calorimeters arranged in a pad configuration surrounding the beampipe[161]. Their primary purpose is to measure the absolute integrated luminosity by recording all small angle Bhabha scattering events. The inner edge of the fiducial acceptance is defined by a tungsten mask that covers the region of $|\cos \theta| < 25$ mr. The pads are arranged into a projective tower geometry. A beam's eye view of the LUM is shown in Figure 4.12, where the tower segmentation can be seen. Each tower is divided longitudinally into an EM1 section of 5.5 radiation lengths in thickness and an EM2 section of 15.6 radiation lengths. The energy resolution has been measured[162] to essentially agree with the design of 3% at 50 GeV.

4.2.3 The CDC

The CDC is the primary tracking detector for the SLD[163]. It is constructed in the form of a cylindrical annulus of inner radius 20 cm, outer radius 1 m, and length 2 m. The active elements are 5120 sense wires interspersed throughout the chamber volurne. The shell structure which bears the tension of the wires was designed to be as thin as possible. It consists of "dished" 5 mm-thick aluminum endplates and inner and outer cylinders made from an aluminum sheet-Hexcel fiberboard laminate. The inner cylinder has 1.8% of a radiation length of material.

The sense wires are arranged in 80 layers which arc organized into 10 superlayers of 8 wires each. Six of the superlayers have a 40 mrad stereo angle with respect to the beam axis to allow a measurement of the z position of the track hits. A subsection of the CDC endplate showing the layer geometry is shown in Figure 4.13. The superlayers are broken into jet cells, where each set of sense wires is radially oriented". The wire layout of a single cell is shown in Figure 4.14. The field-shaping and guard wires consist, of 150 μ m diameter gold-coated aluminum wires, and the sense wires are 25 μ m diameter gold-coated tungsten. Each set of sense and guard

^{**}This structure was chosen for mechanical stability, not ease of tracking, as it guarantees that there ix an ambiguity as to whether a hit originated to the left or right of the sense wire plane (this is commonly called the "left-right" ambiguity).



Figure 4.13: A portion of the CDC endplate, showing the layout of axial (A) and stereo (U, V) layers.

wires is assembled into individual Lexan blocks which can be precisely positioned and tensioned in *situ* in slots in the aluminum endplate.

To read out the signals from the chamber[164], the charge from each sense wire is clocked on each beam crossing into a switched capacitor array, which functions as an analog memory unit (HAMU). If the chamber is to be read out, the charge is digitized and sent up to a FASTBUS waveform sampling module (WSM), which performs zero suppression, corrects each HAMU channel for linearity, and subjects



Figure 4.14: A diagram of an individual CDC cell, showing the sense, guard and field wire layout. The two sense wires at the end of each row are dummy wires and are not read out.

the data to a waveform- finding algorithm which extracts, for each wire end, the time, charge, height, and width of each pulse found in the data. Both ends of the chamber are instrumented, so that the z coordinate of each hit can be measured to approximately 2% of the wire length using the charge division of the pulse.

The gas mixture was chosen to provide maximum precision on the drift distance

measurement. To this end, CO_2 was chosen as the primary gas in the mixture, as it has a low drift velocity and a low diffusion constant. The low drift velocity allows higher resolution on the shape of the arriving pulse for a fixed sampling time, and the low diffusion limits the smearing of the pulse over long distances. The other components of the gas are: isobutane, added as a quencher; argon, added to increase avalanche gain; and trace amounts of water to suppress the effects of wire aging[165]. The gas composition in percent is CO₂:Ar:Isobutane:H₂O::75:21:4:0.2. The drift velocity of the final gas mixture is 7.9 μ m/ns at the mean drift field of 0.9 kV/cm. It is worth noting that the drift velocity of the gas depends strongly on gas density, gas composition, and the electric field in the drift region. All of these must be monitored in order to guarantee stability and precision of the drift distance measurement. The local resolution, which is calculated by the difference in resitluals of adjacent hits on a track within individual drift cells, and the global resolution, which is given by the residuals of all of the hits on a track, are shown in Figure 4.15 as a function of drift distance from the sense wire plane. The difference in local and global resolutions can be attributed to small inter-cell alignment errors^{††}. Also plotted is the resolution curve one would expect for diffusion, given a resolution of 68 μ m at 1 cm. The mean resolution of the hits between 0.5 and 2.5 cm is 82 μ m, making this the most precise large drift chamber ever constructed.

Two other measures of the chamber performance on the local level are the hitfinding efficiency and the two-hit resolution. The hit-finding efficiency is defined as the fraction of time a hit in one of the 80 sense wire layers is found on a track^{‡‡} which passed through that layer. This is shown in Figure 4.16, along with the MC expectation for the same quantity. High hit-finding efficiency and good agreement with the MC have been achieved. The two-hit resolution is given by the minimum resolvable separation between hits on adjacent tracks. This is shown in Figure 4.17, where an efficiency of 50% is obtained at a hit separation of 1 mm. This resolution is determined by the length of each pulse after diffusion, *etc.*, and the specific portion of the waveform algorithm that searches for second hits on the falling tails of pulses.



Figure 4.15: The local and global CDC drift distance resolution measured with tracks in hadronic events from the 1994 data sample. Also shown is the resolution curve expected due to diffusion effects for the specific CDC gas composition. The diffusion curve has been normalized to give the correct minimum resolution.

4.2.4 The CRID

To distinguish among species of charged particles, it is necessary to determine, either directly or indirectly, the masses of each particle under scrutiny. With particles moving at relativistic speeds, it becomes a much simpler matter to measure the momentum of a particle and its velocity, thereby determining the mass indirectly*. This has been accomplished in past collider experiments by precise measurements of the travel times of particles between two fixed points ("time-of-flight"), and measurements of the ionization per unit path length through *a gas* ("*dE/dx*"), which is related to the Lorentz $\beta\gamma$ of the particle). Čerenkov detectors, which measure the Lorentz β of

^{††}The 30-40 μ m alignment errors implied by Figure 4.15 are consistent with the resolution of the

cell-to-cell alignment procedure, described in Ref. [166].

^{‡‡}The tracking reconstruction algorithm will be discussed in the next chapter.

^{&#}x27;Of course, particle species can also be distinguished by their unique interactions with matter as they travel through a detector. SLD's muon system and calorimetry both use this other technique for particle identification.



Figure 4.16: The efficiency for finding a single hit on a track which passes through a given layer in the CDC. Shown are the efficiencies in the data (points) and the MC simulation (solid histogram). The detailed geometry of the cell locations determines the layer-to-layer variations.

a particle, are heavily used in fixed-target experiments for particle identification but have only recently been adapted for the cylindrical geometry of a collider detector. The SLD CRID is one of only two such detectors ever constructed+ These devices allow measurements of the Cerenkov angle for all particles in an event whose velocity is large enough to allow them to emit Ccrcnkov radiation. The active elements of the CRID detector are long[‡] time projection chambers (TPCs) which drift and cap ture the photoelectrons that have been liberated from the TPC gas volume by the Čerenkov photons. The CRID has two radiators to generate Cerenkov photons, a liquid (C_6F_14) and a gas (C_5F_12) . Photons from the liquid radiator impinge directly on the drift boxes, while a system of 400 mirrors focuses the photons from the gas radiator back onto the photo-sensitive medium. Given the direction and momentum of a particle as determined by the CDC, the radius of the "ring" of photoelectrons created by the cone of Čerenkov photons determines the Čerenkov angle and hence

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Figure 4.17: The number of hits found in a single CDC layer as a function of the distance after the first hit. The horizontal line is drawn at full hit-finding efficiency. The 50% efficiency point corresponds to a hit separation of 1 mm.

the particle velocity. A schematic view of the CRID detector is shown in Figure 4.18. At this time, the CRID is becoming well-enough understood to be useful in physics analyses. The observed number of photoelectrons and the local resolutions on the measured Čerenkov angles have been measured to be close to the design values.

4.2.5 The LAC

Calorimetric energy measurements are provided by the LAC, a sampling calorimeter whose basic module consists of lead plates immersed in liquid argon[169]. The argon is the active medium, as it is ionized by charged particles passing through the calorimeter. The lead serves to induce particle showers and as the collection medium for the charge liberated in ionizing the argon. The basic structure of the LAC is shown in Figure 4.19. The layers of lead are broken up into alternating grounded

[†]The other is the DELPHI Ring Imaging Čerenkov (RICH) detector[167].

[‡] We will only discuss the barrel CRID here. The endcap CRIDs operate in a similar manner.

plates and tiles held at high voltage; it is here where the charge collection occurs. As the lead is stacked, the tiles are arranged into projective towers whose longitudinal depth depends on how many layers have been ganged together for readout.

The LAC is segmented into two layers, the electromagnetic (EM) and hadronic (HAD) sections. The EM section is formed of 2 mm thick lead plates separated by 2.75 mm of liquid argon, and is subdivided into two layers, EM1 (6 X_0 , 0.24 interaction lengths (λ_0)) and EM2 (15 X_0 , 0.60 λ_0). The HAD section is formed of 6 mm thick lead plates, but the 2.75 mm argon gap is maintained. The HAD section is also subdivided into HAD1 (13.9 X_0 , 1 λ_0) and IIAD2 (13.9 X_0 , 1 λ_0). The EM towers subtend one quarter of the solid angle of the HAD towers, so that each set of four EM towers is backed by a single HAD tower. The EM sections contain approximately 99% of the energy from a 45 GeV electron, while the LAC as a whole contains 85-90% of the total energy in a hadronic Z^0 decay. The energy resolutions for the LAC have been measured[170] using $e^+e^- \rightarrow e^+e^-$ and $Z^0 \rightarrow q\bar{q}$ events to be $12\%/\sqrt{E}$ and $65\%/\sqrt{E}$ for the EM and HAD sections, respectively.



Figure 4.18: A schematic diagram of the principle of CRID operation. A charged particle entering the liquid radiator emits Čerenkov photons which impinge on the drift boxes containing a photo-sensitive medium. Čerenkov photons from the gas radiator arc focussed back onto the photo-sensitive medium for detection.



Figure 4.19: The internal structure of the LAC, showing the alternate layers of plates and tiles which form the readout towers.

4.2.6 The Solenoidal Coil

Surrounding the LAC is the SLD solenoid, an aluminum core magnet. It provides a field of 0.6 Tesla along the beam axis. The field has been mapped[171] to be uniform to 3% over the tracking volume of the CDC; the lowest order polynomial expansion to the non-axial field in a finite solenoid is used during track reconstruction to parametrize the field non-uniformity. The polynomial agrees with the the measured field to better than 0.05% over the CDC volume.

4.2.7 The WIC

The WIC serves four functions within the SLD: flux return for the solenoid, a backing calorimeter to measure the residual hadronic energy which has leaked out of the LAC, a muon-identification system, and the structural support for the rest of the detector components. The WIC is made of 18 layers of Iarroci (limited-streamer) tubes sandwiched between 5 mm thick steel plates[172]. The tubes are instrumented with square pad readout for calorimetric purposes and long strips for reading out the individual tubes in order to use the WIC as a muon tracker. The WIC geometry is



Figure 4.20: The WIC structure, showing the single layers containing longitudinal strips for muon tracking as well as pad tower readout. Also shown are double layers with crossed strips for tracking in the other plane.

shown in Figure 4.20. The WIC strips which provide the muon-tracking information are arranged in two separate arrays 90° from each other to enable the trajectory of a muon in two dimensions. The endcap chambers in particular have half of the strips oriented vertically and half horizontally. This information will be useful in the context of the backgrounds discussion in Appendix C. The barrel WIC has proved to function as intended and has been used to identify muons from heavy quark semi-leptonic decays[73].

4.2.8 The SLD Detector Simulation

The interactions of all particles with the materials of the SLD are simulated by the GEANT 3.15[173] package of routines. The composition of each of the detector volumes is specified, and the simulation deposits the appropriate amount of energy in each volume, scatters through-going particles by angles consistent with the material thickness, and allows electromagnetic and hadronic showers to occur. Hadronic interactions are simulated by the GHEISHA package; electromagnetic showers are deposited in the detector using a parametrized shower shape rather than running a

detailed simulation.

To simulate the accelerator-induced backgrounds that arise in the various detector elements, all of the signals from random-trigger events close in time to the Z^0 decays are overlayed on top of the generated Monte Carlo events. This is an attempt to insure that the proper luminosity-weighted background levels are properly simulated.

In addition, the observed complement of dead channels is reproduced in the simulated detector response, so that the time-dependent configuration of the simulated detector matches that of the real SLD over the course of the run.

A detailed description of the SLD modifications to the heavy flavor decay package in JETSET6.3 can be found in Ref. [166].

Chapter 5

Tracking System Performance

As will be seen in the next chapters, a thorough understanding of the minute details of the tracking system is crucial to the analyses presented here. This chapter gives an overview of the excellent performance of the SLD tracking system and our understanding of its workings, which is evidenced by measurements of data events as well as our ability to obtain agreement between data and the Monte Carlo simulation. In the previous chapter describing the detector, various details of our understanding of the intrinsic local resolutions of the CDC and VXD were given. Below, we will discuss the global performance of these detectors. As charged tracks are reconstructed using both the VXD and CDC, this global performance directly impacts the physics analyses that rely on the precision of the tracking system.

5.1 Track Reconstruction

Charged track reconstruction begins in the individual drift cells of the CDC[•]. A "Vectored-Hit Finder" searches for straight or slightly-curved strings of hits on the eight sense wires in each cell. Shorter combinations not contained within longer ones are also sought. Each of these miniature track segments is reduced to a vector giving its direction and position. These serve as the input to the patter recognition software,

which attempts to find actual tracks in the CDC. This stage begins by looking first at patterns of hits which form circles in the four axial layers, then proceeds by adding hits from the stereo layers to form long tracks. Shorter tracks not contained within longer ones are also found[174], down to tracks spanning 4 CDC superlayers. After the tracks are found, the hits which they contain are recorded, and a preliminary track fit is performed. This fit produces initial track parameters which can be input into the actual iterative track fitting algorithm. The output of the track fitter is a track defined by two sets of track parameters, one set evaluated at its innermost radius, one at its outermost. Extrapolation inward to the VXD or outward to the CRID, LAC, or WIC begins with these sets of track parameters.

The CDC performance can be studied by using the pairs of tracks from $Z^0 \rightarrow \mu^+ \mu^$ and Z" $\rightarrow e^+e^-$ events, as well as those from cosmic rays[†]. The measured resolution for tracks found in the CDC only for the quantities relevant to physics analyses are shown in Table 5.1. For tracks to be used in analyses requiring precise knowledge of the track trajectory near the IP, the resolutions of $doca_{xy}^{IP}$, the distance of closest approach to the IP in the plane transverse to the beam axis, $doca_{rz}^{IP}$, the distance of closest approach to the IP along the beam axis, and λ , the polar angle of the track with respect to the vertical, are critical determinants of how well tracks can be extrapolated to the VXD in an attempt to find associated hits. The other parameters shown in the table are the resolution on the track direction ϕ in the z-y plane, and the infinite momentum (a) and momentum-dependent (b) terms in the momentum resolution, where $\sigma(p_{\perp})/p_t^2 = \sqrt{a^2 + (b/p_{\perp})^2}$. Plots of the x-y and r-z two track miss distances are shown in Figure 5.2. Except for the momentum, the resolutions for a single track are deconvolved from the distributions of the differences between the two tracks for each of the measured quantities. The larger resolutions for those quantities measured in the T-Z plane are typical of chambers that use stereo wire layers to measure the track positions along the cylindrical axis of the detector. As the VXD does not have a large enough lever arm to measure track momentum itself, it must rely on the CDC measurement as a starting point. A plot of q/p for $Z^0 \rightarrow \mu^+ \mu^-$

^{&#}x27;As the VXD only provides two hits per track, it is not possible to do stand-alone tracking within the VXD itself. Tracks are first found in the CDC and projected inward to candidate hits in the VXD.

[†]In principle, one could use $2\gamma \rightarrow \overline{Z}$ -prong events to study the low-momentum behaviour of the measured quantities, but this suggestion has not been followed up to this point.


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Figure 5.1: The signed inverse momentum of muons in $Z^0 \rightarrow \mu^+ \mu^-$ events as determined by the CDC alone. The solid curves are gaussian fits to this distribution. The muons used for this measurement have sufficiently small polar angles that the distribution is well-approximated by a gaussian.

and $Z^0 \rightarrow e^+e^-$ events is given in Fig 5.1 shows the expected gaussian shape. The momentum resolution of the CDC is $\sigma(p_\perp)/p_t^2 = \overline{0.0050^2 + (0.010/p_\perp)^2}$.

To match the tracks reconstructed in the CDC to hits they possibly left behind in the VXD, the tracks are extrapolated inward to the outermost layer of the VXD. The extrapolated track errors give an indication of an appropriate search region in which to look for VXD hits that could have belonged to that track. A Billoir fit[175] is performed[‡] for each possible combinations to select those VXD hits (if any) which



Figure 5.2: The two-track miss distances of muons in $Z^0 \rightarrow \mu^+ \mu^-$ events as determined by the CDC alone. The solid points represent the measured distribution in the z-y plane, the open squares represent that in the *r*-*z* plane.

provide the best match to the track. Hits are not allowed to be shared by tracks. A final pass is made using a weak vertex constraint to attempt to link tracks to those regions of the VXD where only one hit is available due to the dead CCDs.

The efficiency for linking "quality" CDC tracks[§] to hits in the VXD is shown in Figures 5.3 and 5.4a-b plotted against track momentum, track polar angle θ , and track azimuthal **angle** ϕ . Also shown in these plots are the **same** quantities and the fraction of the links that are incorrect as predicted by the MC. One can see from these plots the quality of the links, as well as the quality of the MC simulation, which agrees with the data to better than 1% for quality tracks. Figure 5.5 shows the fit quality $\sqrt{2\chi^2} - \sqrt{2n_{d.o.f} - 1}$ of the combined CDC+VXD Billoir fit for data and MC. The basis.

^tThe Billoir technique for track fitting uses a simple matrix multiplication to handle the extra extrapolation error introduced by passing through material instead of an iterative "swimming" method, which steps the tracks along their helical path and considers material on a much more local

[§]The cuts to define a "quality" track will be discussed in the next chapter.

Table 5.1: The resolutions measured for various track parameters with the CDC and CDC+VXD tracking systems. The variables are defined in the text. Resolutions quoted are those from single tracks, as deconvolved from the two-track difference distributions.

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Parameter	CDC Only Resolution	CDC+VXD Resolution
$doca_{xy}$	$150 \ \mu m$	11 μ m
$doca_{rz}$	$1000 \ \mu m$	$38 \ \mu m$
ϕ	0.45 mrad	0.32 mratl
λ	3.7 mrad	2.2 mrad
a	0.005	.0026
b	0.01	0.0095



Figure 5.3: The efficiency for linking "good" CDC tracks (see Chapter 7 for a definition) to the VXD as a function of track momentum. The points are the measured efficiencies in the data, the solid histogram is the same quantity as determined in the MC simulation.

good agreement here demonstrates that our understanding of the various resolutions and alignment problems is approximately correct.



Figure 5.4: The efficiency for linking "good" CDC tracks (see Chapter 7 for a definition) to the VXD as a function of (a) track polar angle and (b) track azimuthal angle. The points are the measured efficiencies in the data, the solid histogram is the same quantity as determined in the MC simulation. Note the effect of the 0.4 cm shift of the VXD along the beam axis, which leaves the efficiency in $\cos \theta$ somewhat asymmetric.



Figure 5.5: A comparison of the $\sqrt{2\chi^2} - \sqrt{2n_{d.o.f.}} - 1$ distribution for the CDC+VXD fit for data (points) and MC (histogram) tracks.

As would be expected, adding VXD hits to the CDC track dramatically improves the resolution on most of the track parameters over those measured by the CDC alone. The CDC+VXD results for these resolutions are also shown in Table 5.1. The momentum resolutions certainly benefits from the addition of the precise VXD points close to the IP; the combined momentum resolution is $\sigma(p_{\perp})/p_t^2 = 0.00262 + (0.0095/p_{\perp})^2$. Due to the small lever arm of the VXD, the combined measurements are still affected by the precision of the CDC. In particular, the single track impact parameter resolutions in the *r-z* plane are more than three times worse than those in the z-y plane due to the lower precision of the CDC measurement of the track dip angle.

An estimate of the impact parameter resolutions for lower momentum tracks can be made from the widths of the distributions of $doca_{zy}^{IP}$ and $doca_{rz}^{IP}$ by using the MC to unfold the effects of heavy hadron decays and IP motion. The results of this study are shown in Figure 5.6 for tracks at $\theta = 0$ " and tracks at $|\theta| = 70$ ". The impact parameter resolutions obtained for tracks with momentum 1 GeV/c are $\delta(doca_{xy}^{IP}) = 76\mu m$ and $\delta(doca_{rz}^{IP}) = 80\mu m$ at $\theta = 0^{\circ}$.

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Figure 5.6: The impact parameter resolutions as a function of track momentum and polar angle, for the z-y and r-z planes. The solid lines are the resolutions obtained from the MC; the points are those from the data obtained by correcting the observed distributions for the presence of long-lived mesons. Two sets of resolutions are shown: those for tracks with normal incidence on the VXD ladders (0") and those with a relatively large polar angle (70"). Note the worsening of the impact parameter resolution due to the extra multiple scattering in the effectively thicker detector at large angles.

5.2 Determination of the

Primary Vertex (IP) Position

The inclusion of this section breaks the flow of the discussion of tracking performance, but is is necessary to define what is meant by the event primary vertex or IP position with respect to which all track impact parameters are measured. Up to this point, we have used the term IP to mean loosely "where the Z^0 decay occurred". Here, we will make the distinction between the SLC II', which is defined as the center of the luminous region of the overlapping beams, and the event Primary Vertex (PV), which is where the tracks in the event appear to originate, given the information provided by the SLD tracking system.

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5.2.1 Transverse Position

In principle, one could find the PV for each event from fitting all of the tracks which appear to come from a single, central point to a common vertex. This results in an error on the spatial position of the PV in the *x-y* plane that looks like an ellipse due to the inevitable presence of collimated jets of particles. This is shown schematically in Figure 5.7. Typically, the size of the errors are 100 μ m along the ellipse major axis, and 15 μ m along the ellipse minor axis. As these values are large compared to the resolutions discussed above for the SLD tracking system, they would introduce unacceptably large errors on the track impact parameters if this were the only way to find the PV for each event.

Fortunately, the SLC luminous region is stable and extremely small. As the beams must actually collide to produce useful luminosity, the average position of the collision point is stable over times of at least one hour during smooth running. Steering correctors in the SLC final focus driven by feedback loops make small corrections to the beam positions over this time to maximize the luminosity, but the amount of induced motion is small. So, rather than attempt to find the PV event-by-event, wc can substitute the average position of the SLC IP. In addition to drastically reducing the magnitude of the position error, averaging over events mitigates several other

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problems: the error ellipse on the mean IP will be roughly circular, removing the asymmetry induced by the jet axis; the bias on the PV position due to single tracks will be greatly reduced; and possible correlations between the PV position and those of decay vertices in events with many secondary decays will be essentially eliminated.

To find values for the average IP (IP) position over the entire run, the SLD data sample is divided[¶] into sets of 30 sequential hadronic Z^0 decays (events selected with loose criteria), except when a run ended more than 20 events into a set, at which point the set is ended^{||}. A trial (IP) is determined for each set by fitting for the PV, then averaging the result. The $\langle IP \rangle$ for each set is then derived by fitting all tracks which have VXD hits and which come within 30 of the trial (IP) to a common vertex. The fit (IP) is then used as a new trial (IP), and the process is iterated until it converges. Typically, 330 tracks are used in a fit, and the fit converges after 5 iterations. The χ^2/dof for the fit and the fraction of tracks within 3σ of the fit (IP) arc monitored for each set and compared with those from previous and subsequent sets to identify sets which might span a major shift in IP position. Information from the SLC correctors is used to help determine exactly when a major shift occurs. When a major shift is found within a set, the boundaries of the set are changed to coincide with where the IP shift occurs while still maintaining ~ 30 events per set whenever possible, and the fitting procedure is repeated. For a set to be used in the analysis, the χ^2/dof is required to be < 1.3, and the number of tracks used in the fit is required to be more than eight times the number of events in the set. The fit (IP) position for the set in which an event resides is then used as the best estimate of the PV position in the z-y plane for that event.

The uncertainty in the (II') (σ_{IP}) is the combination of the statistical error from the fit $(\sim 3\mu m)^{**}$, the extent of the SLC luminous region $(\sim 1\mu m)$, and the motion of the IP within a set $(\sim 6\mu m)$. This totals $\sim 7\mu m$ when added in quadrature. There

⁽This work was done and the description graciously provided by Steve Wagner. A similar section appears in Ref. [182].

¹Major jumps in the SLC IP position almost always occur due to the reestablishment of collisions in a different place than before after a period of inactivity or outage, for whatever reason. As SLD data-taking **runs** are **usually** ended after 20 minutes of SLC inactivity, these major changes in IP position tend to happen between SLD runs.

^{**}Compare this to 15 µm for the single event PV determination, above.

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Figure 5.7: A diagram showing the definition of the variables used in the discussion of the $\langle IP \rangle$ fit quality. Shown arc the possible results of fitting the IP position for a given event, with the error ellipse of this fit (which has major (minor) axis length σ_x (sigma,)). Also shown is the mean IP position $\langle IP \rangle$ for this set of runs, and the two coordinates specifying its location relative to the fit event IP. The $\langle IP \rangle$ position is x_T away from the fit event IP along the major axis of the error ellipse, and y_T along the minor axis.

are several ways to estimate σ_{IP} using the data. The distribution of track impact parameters with respect to (IP) in $Z^0 \rightarrow \mu^+ \mu^-$ events is shown in Figure 5.8. As $Z^0 \rightarrow \mu^+ \mu^-$ events are not used in any way for the determination of (IP), they can provide an independent check on the precision and accuracy of the beam position measurement. The σ of the distribution is 12.7 μ m; when the track extrapolation error is subtracted in quadrature, this gives $\sigma_{IP} = 6.7 \mu$ m, which confirms that our estimates of the uncertainties in extracting the (IP) are approximately correct. In addition, in hadronic events where most, tracks fit to a common vertex, the distance (y_T) between the (IP) and the fit vertex, projected onto the minor axis of the fit vertex error ellipse, also contains information on σ_{IP} .

The quantity y_T is calculated for each event using the fitted primary vertex for that event. The primary vertex for each event is found by beginning with the four tracks whose impact parameters with respect to the (II') are smallest. Tracks which make the χ^2 vertex probability the largest are added until the overall χ^2 probability drops below 1%. The position of the event *PV* and the error ellipse parameters are



Figure 5.8: The impact parameters of muons from $Z^0 \rightarrow \mu^+ \mu^-$ events to the (IP). As they are not used in the determination of the $\langle IP \rangle$ position, these tracks provide an independent measure of the resolution on the (IP) position.

then calculated. Figure 5.9 shows the distribution of y_T for the events where

- \geq 7 tracks are included in the primary vertex
- . \geq 70% of the tracks in the event are used in the primary vertex fit
- The error on y_T is less than 15 μ m
- . the overall χ^2 vertex probability is greater than 1%.

The data points arc overlayed on a MC simulation in which 7 μ m of IP motion has been added. The non-gaussian tails are similar in both MC and data. This distribution for all events, while having high statistics, includes some contamination from events containing heavy quark decays, which could bias the fit PV position. We can produce another sample by requiring



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Figure 5.9: The y_T distribution for all events satisfying moderate vertexing criteria (see text). The points are the data, the solid histogram is the MC with an input smearing of 7 μ m in the (IP) position. Small tails can be seen on both distributions.

- 100% of tracks with VXD hits are used in the primary vertex
- the overall χ^2 vertex probability is greater than 10%.

These events contain a very pure sample of light quark events, but the resulting y_T distribution is lower in statistics, as shown in Figure 5.10. Nonetheless, the y_T distributions for both these samples were considered; all distributions agree with $\sigma_{IP} = 7\mu \text{m} \pm 2\mu \text{m}$ for the 1993 data.

Even though great care has been taken to minimize the potential for large beam motion within a set of events, it is still possible that some sets contain events which occur a large distance away from the (IP). This would lead to large numbers of tracks in each of these events that miss the (IP) by a significant amount, which would lead to a much higher probability for these events to be tagged as b quark events. This would effect any analysis that attempts to use the precision tracking as a flavor tag.



Figure 5.10: The y_T distribution for events with all tracks in the primary vertex (see text). The points are the data, the solid histogram is the MC with an input smearing of 7 μ m in the (IP) position.

In an attempt to set an upper limit on how often this could happen, the same distributions which are used to estimate σ_{IP} are searched for evidence of non-gaussian tails. The $\mu^+\mu^-$ impact parameter distribution shows no evidence for non-gaussian tails, but the statistics of this sample are not large enough to rule out non-gaussian tails in the much larger sample of hadronic events. The highest statistics check, the y_T distribution in all hadronic events, shows similar small non-gaussian tails in both the MC and data. The MC indicates that this is caused by the occasional inclusion of B or D decay tracks in the vertex fit, which then pulls the PV vertex position slightly. All other distributions show smaller non-gaussian tails than this one. We take the y_T distribution in all hadronic events as a conservative limit on the size of these tails, which leads us to include a second IP extent (σ'_{IP}) of 100 μ m in < 0.25% of the data to simulate this effect. The full shift in the results obtained with and without this additional tail will be taken as the systematic error due to this effect.

5.2.2 Longitudinal Position

Since the SLC bunches are approximately 1 mm long, the luminous region along the z axis has a much larger physical extent than in the x-y plane^{\dagger}. There is nothing to be gained by averaging many events together in this plane to obtain the mean IP position, since the tracking resolution is much finer than the possible distribution of interactions. The best estimate of the PV z position for each event comes from a technique using only the median z of the tracks in the event itself. Each track with associated VXD hits is extrapolated to the point of closest approach to the (IP) in the z-y plane, and the z coordinate of the track at this point is denoted as z_{doca} . A selection of tracks is then made to require the track x-y impact parameter to be less than $500\mu m$ and the track to pass within 3σ of the (IP) based on the estimated track impact parameter error and $\langle IP \rangle$ error. The event IP z location is simply defined as the median of the z_{doca} values from the selected tracks. For the small fraction of events with no tracks passing this selection, all tracks with VXD hits are used. The choice of the median z method instead of the more common approach involving vertexing is based on the result of a MC study showing that the median z is more robust against the biases due to the inclusion of tracks not originating from PV. The typical resolution for locating the PV z as derived from MC are $(32,36,52)\mu m$ for (uds, c, b) events respectively. The tails of the PV z residual distributions can be characterized by the fraction of events with residual > $100 \mu m$. The fraction of such events are (0.8%, 1.6%, 7.5%) for (uds, c, b) events according to the MC simulation.

5.3 The Impact Parameter Distribution

The final element in describing the performance of the SLD tracking system is also the most relevant. Shown in Figure 5.11 is the distribution of the normalized impact parameters $doca_{xy}^{IP}/\delta(doca_{xy}^{IP})$. The agreement between data and MC is quite good, even out to 20σ on the tails. The agreement is equally good in the *r*-*z* plane. The



Figure 5.11: A comparison of the distributions of normalized impact parameters in the x-y plane for all hadronic events in the data (points) and MC (histrogram). The tracks included here are "good" tracks (see Chapter 7). The under- and overflow bins are also shown.

only "smearing" which has been added to the MC simulation is the residual alignment errors of the VXD and the CDC internal and global alignments. The local alignment smearing was done on a CCD-by-CCD, ladder-by-ladder, and cell-by-cell basis in order to get the χ^2 distributions to match between data and MC. An average amount of smearing was added to each element, and at no time was there any tuning done on the tails of the distributions in Fig. 5.11 to attempt to make the tails agree. In this version of the MC, the original guess of the CCD alignment position precision was underestimated, and there was in addition an error in the simulation, both of which conspired to make the core of these distributions wider in MC than in data. Given the "true" track parameters from MC, the distributions were actually *unsmeared* to

^{††}We have actually used the distribution of z PV positions to measure the bunch length, effectively turning the SLD into the world's most expensive bunch length monitor. We measure a value consistent with the expectation of 1 mm.



Figure 5.12: The ratio of the distributions of the normalized impact parameter for the unsmeared MC and data. (see text)

remove this effect. The ratio of unsmeared MC to the data is shown in Figure 5.12; the unsmearing correction was only applied in the "core" region, and its magnitude was calculated solely from the MC. The agreement between MC and data in the core region reflects the true quality of the MC simulation, not a conspiracy to force agreement between the simulation and the data. The full change in the results of the analysis by applying this correction will be taken as a systematic error.

In the following sections, we discuss aspects of the tracking system performance that are relative to the reconstruction and isolation of events containing D mesons for the charm asymmetry measurement. Two methods are used to find Ds: one requires reconstruction of high-momentum D mesons from 2 or 3 tracks, the other requires the isolation of the D decay vertex. Thus, the proper simulation of the mass resolution of the detector and its vertexing capabilities are necessary in order for us to have confidence in the MC results. It is worth noting that we are concerned here about the tracking performance as applied to multi-track combinations as opposed to the single variable distributions we have been discussing.

5.4 Mass Resolution

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To study the mass resolution of the SLD tracking system under conditions similar to those encounter4 in reconstructing D mesons, we can examine high-momentum K_{\bullet}^{0} and A decays (collectively called V^{0} decays), and compare the widths of the mass peaks in data and MC. If the tracking performance is properly simulated at high momentum, we should obtain reasonable agreement between the data and MC distributions. Candidates for K^0_* and A decays are found by pairing all combinations of oppositely-signed tracks and assigning either both tracks the π mass (K_{\bullet}^{a}) or one the proton mass and one the π mass. At high momentum, the faster of the two particles in Λ decay is almost always the proton; using this information for the mass assignments reduces the combinatoric background. Several cuts are applied to each V^0 candidate to obtain a clean sample [176]. The pair of tracks forming the candidate V^0 must pass within 2.5 mm of each other, and their actual distance of closest approach must be less than 3σ of the uncertainty on their common point of origin in the plane perpendicular to the sum of their momenta. The difference between the V^0 candidate momentum direction and its flight direction must be less than 1" in the x-y plane, and less than 2" overall. Figures 5.13a and b show the invariant masses of the K_s^0 and A candidates for $x_{V0} > 0.2$, which is the lower of the two momentum cuts that are used in the D meson reconstruction analysis. There is reasonable agreement between the data and MC simulations, which gives us some confidence that we will be able to reconstruct accurately the masses of D^0 and D^+ mesons for the charm asymmetry analysis.

5.5 Vertex Finding

There are two concerns about tracking performance that are relevant to an analysis that depends on the isolation of a decay vertex involving two or more tracks. The first is whether or not such a vertex can be found, including whether or not the tracks are properly reconstructed, which is essentially a tracking (vertexing) efficiency question. The second is whether or not, given that a vertex is found, the vertex can



Figure 5.13: The invariant mass distributions for (a) K_s^0 and (b) A candidates with a total momentum greater than 9 GeV/c ($x_{V^0} > 0.2$). The points (histogram) represent the data (MC) distribution. A reasonable aggreement between the data and the simulation can be seen.



Figure 5.14: The distribution of the number of two-track vertices found in hadronic events, comparing the data (points) and MC simulation (histogram).

be isolated from the other vertices around it which arise from random combinations of unassociated tracks, which is a vertex resolution question. We will discuss each of these concerns in turn.

Rather than give an explicit exposition of the vertexing efficiency as a function of vertex momentum and vertex decay length, we present here evidence that both the tracking efficiency and the track resolution is properly simulated in the MC'. To do this, we can compare the data and MC distributions of the number of two track vertices found in all hadronic events. We must make some cut to consider only those vertices which are some distance from the If'; this we choose to be 3σ . This distribution then contains the efficiencies for track finding and correct track reconstruction, as well as a check that the errors on the vertex fit are calculated correctly. The tracks used in making this plot, arc "good" tracks as defined in Chapter

[•]For an asymmetry measurement, the quality of the simulation is also more important than the actual numerical value of the vertexing efficiency.



Figure 5.15: The flight distance resolution of the SLD tracking system as simulated in MC for $D^+ \rightarrow K^- \pi^+ \pi^+$ decays. This fit curve is a single gaussian of width 200 μ m.

7. The vertex fit is performed using the SLD vertexing routines ZXTWO followed by ZXFIT, which perform a 1-C fit for the vertex position and error matrix. The momenta of the tracks are reevaluated at the vertex position. Figure 5.14 shows the number of vertices with flight distances greater than 3σ for all (selected) hadronic events in the 1993 data sample and the corresponding MC simulation. Good agreement is seen between the two distributions, which implies that the track- and vertex-finding efficiencies are well simulated in the MC.

The resolution of the tracking system in measuring decay lengths can be estimated from the MC simulation, as we have no way of knowing the exact decay lengths of secondary vertices in the data[†] We have used the MC to estimate the decay length resolution for $D^0 \rightarrow K^-\pi^+$ and $D^+ \rightarrow K^-\pi^+\pi^+$ decays, as these two and three-track modes are those used in the analysis. The MC distribution of the measured minus the true decay length for the D^+ decay is shown in Figure 5.15. The width of the single

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gaussian fit is 200 μ m. Since the mean D^0 (D^+) decay length is approximately 1.5 mm (2.5 mm), requiring found vertices to be separated from the IP by 3σ or so will still result in an efficient vertex tag while suppressing the large number of random vertices that occur from tracks passing close to the IP.

[†]In the process of measuring the lifetime of the τ lepton using the vertex decay length, the decay length resolution can be estimated by the extent of the tail at negative vertex Right distance[177]. Results from this estimate are consistent with the MC simulation.

Chapter 6

Hadronic Event Selection

One of the advantages of doing physics in the environment of an e^+e^- collider is that one expects the background processes that could mimic an event of physics interest to occur infrequently enough that the trigger algorithm that causes events to be recorded can be relatively simple and indiscriminant, and yet not be overwhelmed by a flood of data. This approach also leads to a very efficient trigger algorithm, where the efficiency for actually recording the events of interest is for all practical purposes 100%. In order to be this efficient, however, the trigger algorithm must be made sufficiently indiscriminant, which results in essentially all events that register any reasonable amount of energy in the detector being written to tape. Nonetheless, this approach has worked well, even at the SLC, where accelerator-induced backgrounds can be more of a problem that at a typical storage ring collider. Since magnetic tape is a relatively cheap storage medium we can afford to write to tape a data stream which contains less than 1% Z^0 events and sift through later to select out the decays that we want for physics analysis purposes. This chapter describes in some detail the selection criteria applied to the raw events to pass the trigger, the initial Z^0 selection filter, and our definitions of a hadronic Z^0 decay.

6.1 The SLD Trigger

The trigger criteria are described in detail in Ref [178], and will be summarized briefly here for completeness. During the 1993 SLD run, there were six main classes of triggers: a total energy trigger, a charged track trigger, a "hadron" (HAD) trigger, a wide-angle bhabha (WAB) trigger, a muon pair trigger, and a "bhabha" trigger*. The energy trigger required a minimum total energy of 8 GeV^{\dagger}, where the sum is taken only over those EM (HAD) calorimeter towers containing more than 60 ADC counts (120 ADC counts) of deposited energy, which corresponds to an energy threshold of 246 MeV (1.296 GeV) per tower. Only the calorimeter systems were read out when this trigger fired. The charged track trigger was based on a pattern map of the cells that might be hit as a charged track of momentum greater than 250 MeV/c passed through the CDC[179]. A hit cell was defined as a cell where pulses consistent with hits from tracks fired a discriminator on 6 of the possible 8 wires. Those events containing two tracks passing through at least 9 superlayers of the CDC and lying roughly 120" apart were passed by this trigger, and all of SLD was read out if it fired. The HAD trigger is a combination of the previous two triggers, as it required one charged track of 9 or more superlayers and a large energy deposition in the LAC. As it was expected that most Z^0 events would satisfy this trigger, the entirety of SLD was read out when this trigger was satisfied. The WAB trigger was designed to insure that all wide-angle e^+e^- pairs were recorded, even those at angles where the track stubs in the CDC were not long enough to satisfy the track trigger requirements. This trigger also initiated readout of all of SLD. The muon trigger required a combination of a charged track in the CDC and hits in opposite WIC octants, as this would be

^{*}Here, what we refer to as trigger classes are technically known as the "slots", which are combinations of the individual "evaluators" that represent the trigger decisions from the individual trigger processors in the data acquisition system.

[†]the energies quoted in reference to the trigger are in units of the energy that would be left in the calorimeter by a minimum-ionizing muon, and are thus uncorrected for the different interactions that could occur in hadronic vs. electromagnetic electromagnetic showers and the corresponding differences in calorimeter response. Within SLD, this way of quoting energy is commonly referred to as the *min-I* scale. A minimum-ionizing particle passing through the LAC leaves an average of 360 MeV in the EM sections, and OGO MeV in the HAD sections. The conversion factors from ADC counts to energy are 4.1 MeV/ADC count (10.8 MeV/ADC count) for the EM (HAD) layers of the calorimeter.

the signature for a $Z^0 \rightarrow \mu^+ \mu^-$ event. Once again, all of SLD was read out if this trigger was satisfied. The bhabha trigger required a total energy (EM scale) in the EM2 section of the LUM to be above 12.5 GeV in both the north and south detectors, where the sum is made over towers with more than 1.25 GeV of energy. This energy deposition is what one might expect as a signature for an e^+e^- pair above some splash of background. This only triggered the readout of the LUM/MASiC system, and thus will not concern us further.

6.1.1 The Tracking Trigger Veto

A number of vetos are added to the trigger system in order to prevent excessive detector dead-time caused by processing huge events where the detector is sprayed with accelerator-induced backgrounds. Accelerator hiccups could happen in many different ways (see Appendix C), but the effect on SLD is more or less independent of the cause, as the detector will end up completely blasted with noise. One of these vetos, the one employed in the track trigger, has a large impact on this analysis. The tracking trigger veto causes an event with more than 275 cells having 6 out of 8 wires hit to not be read out. In the nominal configuration, this has a relatively small effect on the number of Z^0 decays that lack tracking information. Unfortunately, during approximately two-fifths of the 1993 run, a hardware problem effectively shifted the number of hit cells up from a minimal value of 50 up to a minimum value of 100^{\ddagger} . The effect of this is that in approximately 12% of hadronic events the CDC information is missing when it would ordinarily be present. For the a, analysis, this could have severe consequences, as it is easy to see that the veto preferentially removes 3-jet events, as they tend to have higher charged multiplicity and would hit more cells. In particular, since b events have a higher average multiplicity, J-jet b events are vetoed more often than 3-jet events from other quark flavors. All of these effects could potentially bias the flavor-independence test if they were to he ignored[§]. A simulation of the track trigger veto does in fact exist, however, and can be used on the MC to provide the

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efficiencies for trigger and event selection for the given quark flavors and event jet multiplicities that should match the data. This has been discussed in great detail in Ref. [166]. Variations of the rate at which events are vetoed will be considered below as a systematic error.

6.2 The Hadronic Event Filter

The initial event selection is done using calorimetry information only, as it is much faster to process compared with the tracking data. The filter algorithm used at this stage, called "EIT pass-1"[178], requires that the events are relatively spherical and have good forward-backward momentum balance. It is based on three LAC quantities:

- *NEMHI*, the number of LAC EM towers with signals above 60 ADC counts (~ 250 MeV min-I)
- *EHI*, the sum of the energy deposited in all EM (HAD) towers with signals greater than high thresholds of 60 (120) ADC counts. (This is equivalent to 250 MeV (I.3 GeV) min-I.)
- *ELO*, the sum of the energy deposited in all EM (HAD) towers with signals greater than low thresholds of 8 (12) ADC counts, (This is equivalent to 33 MeV (130 GeV) min-I.)

The filter requires that each event satisfy:

- 1. NEMHI ≥ 10
- 2. EHI > 15 GeV min-I
- 3. ELO < 140 GeV min-I
- $4.2 \times EHI > 3 \times (ELO-70)$
- 5. The north and south hemispheres of the detector must each have NEMHI > 0

[†]The minimum value is a constant that is adjusted depending on the number of cells that are inoperative at a given time.

[§]Note that, since the veto is forward-backward symmetric, it has no effect on the A_c analysis.

Cuts 3 and 5 tend to remove beam-wall events and other beam burp events; cut 4 removes those events with large numbers of SLC muons (see Appendix C) passing through the LAC. The combined efficiency for a hadronic Z^0 decay to pass the trigger and the EIT pass-1 filter is approximately 92%[180].

6.3 Hadronic Event Selection Cuts

After applying the EIT pass-1 filter, 63553 events remain in the data sample from the 1993 run. In order to reduce this to a sample of hadronic events useful for physics analysis, we apply cuts designed to select events contained within the fiducial region of the SLD detector. Since we use charged tracks as a basis for much of the analysis, this restricts us to the barrel region of the SLD, as the CDC starts to lose tracking efficiency outside of $|\cos \theta| > 0.8$. We also wish the energy flow of the event to be well contained within this fiducial region so as to guarantee that we are not missing large portions of the event that have ended up in the detector endcaps. In addition, we should note that, since it possesses the simplest geometry, the barrel region of the detector is also the easiest to model in the MC, and can be considered "well understood" in this context relative to the rest of the SLD. Since a number of the cuts are based on charged tracks, we first set forth a definition of a "charged track", which is any track satisfying:

- . transverse momentum relative to the beam axis $p_{\perp} > 200 \text{ MeV}/c^{\P}$
- $|\cos\theta| < 0.8$
- . distance of closest, approach in the xy plane $(doca_{xy}^{IP}) < 5$ cm
- . distance of closest approach in the rz plane $(doca_{rz}^{IP}) < 10$ cm

Since slightly different event selection criteria were used for the two analyses presented here, we present two descriptions of the event selection cuts.

6.3.1 Event Selection for the α_s Analysis

The selected events for the flavor independence analysis are required to satisfy:

- \geq 7 charged tracks
- $|\cos \theta_{Thrust}| < 0.71$
- $E_{vis} > 18$ GeV, where E_{vis} is calculated from the charged tracks, assuming all tracks to be pions.

One can see that these cuts meet the requirements mentioned above for defining a set of fiducial hadronic events with minimal background from other processes.

The one exception to the exclusive use of charged-track information in this analysis is the determination of the thrust axis direction, where we use the thrust axis calculated from calorimeter clusters. The reason for preferring the calorimeter determination of the thrust direction is shown in Figure 6.1, where the thrust direction calculated for each event from calorimeter and track information is plotted. One can clearly see the cases at high calorimeter cos θ_{Thrust} where the tracking determination fails because of lost information. Since one expects the global event properties of these events to be modeled less well than those where many tracks are present in the fiducial volume, it is best to exclude them from the analysis sample.

Figures 6.2 show the data sample before cuts, with the hadronic event MC overlayed. After applying these cuts, 28036 events remain. Table 6.1 shows the efficiency for the different event flavors to pass these selection cuts for the period prior to the trigger veto problem (henceforth called the "Pre-Veto" period), the period of running during which the track trigger veto was incorrectly set (the "Veto" period), and for the later run time after it had been fixed (the "Non-Veto" period). Any small flavor bias in event selection will be explicitly corrected in the analysis, but, as can be seen, any event selection bias is small. The veto period does show some bias for rejecting b events. This is expected, as they tend to have a higher charged multiplicity, and thus should trigger the veto more often. The surviving backgrounds are small and are neglected in the analysis.

[¶]The A_c analysis used a minimum transverse momentum of $p_{\perp} > 200 \text{ MeV}/c$.

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It is interesting to break this down further into the efficiencies for 2- and 3-jet events to pass the selection cuts (and the trigger veto). Table 6.2 contains more detailed cut information. The effect of the tracking trigger veto is readily seen, as the efficiencies for 3-jet events to pass the veto and cuts are from 7% to 12.5% lower during the Veto period compared with the rest of the run.



Figure 6.1: A comparison of the direction of the thrust axis as determined by the CDC (tracks) and the LAC (calorimeter clusters). The effects of the limited solid angle coverage of the CDC are readily apparent, as there are many events where the majority of the energy is deposited in the LAC endcaps (and hence at large $\cos \theta$) in which the thrust axis direction is miscalculated by using the CDC-only information.

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Figure 6.2: Distributions of the event selection variables. To make each plot, all of the other cuts have been applied except for the one being plotted. The arrows indicate the values of the cut on each of the variables.

Table 6.1: The efficiencies for hadronic events of different primary quark flavors to pass the event selection cuts.

Run Period	uds efficiency	c efficiency	b efficiency
Pre-Veto Period Veto Period Non-Veto Period	$\begin{array}{c} 0.5749 \pm 0.0020 \\ 0.5457 \pm 0.0017 \\ 0.5934 \pm 0.0029 \end{array}$	$\begin{array}{c} 0.5738 \pm 0.0037 \\ 0.5392 \pm 0.0033 \\ 0.5969 \pm 0.0055 \end{array}$	$\begin{array}{l} 0.5694 \ \pm \ 0.0033 \\ 0.5115 \ \pm \ 0.0029 \\ 0.5926 \ \pm \ 0.0049 \end{array}$

Table 6.2: The efficiencies for hadronic events of different primary quark flavors and different numbers of jets to pass the event selection cuts. The E algorithm with a y_{cut} value of 0.8 was used to obtain these values.

Run Period	uds efficiency	c efficiency	b efficiency	
Pre-Veto Period: 2-jet events 3-jet events	0.6145 ± 0.0023 0.4619 ± 0.0039	0.6177 ± 0.0043 0.4534 ± 0.0073	0.6113 ± 0.0038 0.4512 ± 0.0065	
Veto Period: 2-jet events S-jet events	0.6013 ± 0.0020 0.3868 ± 0.0033	0.5979 ± 0.0037 0.3776 ± 0.0061	$\begin{array}{c} 0.5751 \pm 0.0033 \\ 0.3299 \pm 0.0053 \end{array}$	
Non-Veto Period: 2-jet events 3-jet events	0.6283 ± 0.0033 0.4945 ± 0.0058	0.6329 ± 0.0063 0.4964 ± 0.0109	0.6354 ± 0.0056 0.4753 ± 0.0097	

6.3.2 Event Selection for the A, Analysis

Since, as will be seen, the A, analysis is driven to obtain the largest possible sample of D decays, we can loosen the event selection cuts substantially. We are not concerned about the energy flow of the event, as the reconstructed D meson direction is used instead of a jet or thrust axis to approximate the initial quark direction. Since the asymmetry is also largest at large $|\cos \theta|$, we would also like to be able to reconstruct as many D'S at large angles as is possible given the acceptance of the vertex detector. The requirement that three tracks form a D^* or D^+ also tends to suppress backgrounds from mismeasured tracks, so that any degradation in the tracking performance at large

 $|\cos \theta|$ amounts only to an overall efficiency loss. To this end, we define the hadronic sample for the A_{ℓ} analysis to be those events which have

- 5 or more charged tracks
- • $\cos \theta_{Thrust} | < 0.8$
- $E_{vis} > 18 \text{ GeV}$.

We use charged tracks in all cases, even to find the thrust axis direction^{\parallel}. These looser cuts result in an event sample of 33,524 hadronic decays.

[#]We choose to keep those events at large $|\cos \theta_{Thrust}|$ which would have failed the containment cut had the thrust axis direction been calculated using calorimeter information.

Chapter 7

Flavor Tagging

As mentioned in the introduction, we have chosen to use the differences in lifetime and decay multiplicity among the different hadronic species as a tool to separate events containing different flavors of primary quarks. In principle, the tag should be relatively insensitive to the kinematics of the tagged event (whether it had 2 or 3 jets, in particular), should be efficient, and should provide maximal purity for all of the tagged samples. Simplicity can also be considered a virtue in this particular case. These criteria rule out the more complicated tags mentioned in Section 3.4.1, as well as attempts to select heavy flavor events by the exclusive reconstruction of charm or bottom mesons. As will be described below, a reasonable compromise is to look only at the number of tracks in each event that miss the IP by a significant distance, as this efficiently separates high-purity samples of b and **uds** events, **as** well as a c sample of enriched purity. This chapter presents a detailed explanation of this tagging method.

7.1 Selection of Quality Tracks

As will be described in Chapter 9, this analysis depends in a critical manner upon our ability to understand the performance of the SLD tracking system in as much detail as possible. In particular, the entire process of flavor separation requires the MC to tell us the efficiencies for tagging an event of a given quark flavor for each

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of the tags. As the tags are based on track multiplicity, it is imperative that the ensemble of tracks used in the tags are well understood and well measured. To this end, we need to apply further cuts on the charged tracks used in the analysis so as to guarantee that they are well-described by the MC simulation and do not have some parameter or other lying way out on a mysterious tail of some distribution. These cuts have the other beneficial effect of removing some tracks from light quark events that carry lifetime information, such as those from K^0 decays, which can contribute to the light quark contamination of the heavy flavor samples. Tracks considered as "Tagging Tracks" must have:

- at least one VXD hit
- . at least 40 CDC hits, with the first hit at a radius less than 39 cm
- . a combined CDC+VXD fit quality $\sqrt{2\chi^2} \sqrt{2n_{d.o.f} 1} < 8.0$
- total momentum greater than 0.5 GeV/c
- $doca_{xy}^{IP} < 0.3$ cm and $doca_{rz}^{IP} < 1.5$ cm.
- an error on $doca_{xy}^{IP}$ of < 250 μ m

The cuts are placed on the track parameters calculated from the combined CDC+VXD fit. In addition, tracks from candidate K^0 and Λ decays and -y-conversions found by kinematic reconstruction of two-track vertices are removed from this analysis[181] to eliminate tracks from *uds* events that seem to come from particles with long lifetimes. All of these cuts tend to discard poorly measured or mis-reconstructed tracks, since they tend to have larger extrapolation errors, fewer hits, and larger IP miss distances than properly reconstructed, well-measured tracks. Figure 7.1 shows the distributions to which the cuts are applied, as well **as** the cut values, for the three classes of tracks containing different sources of lifetime information. Immediately following, Figure 7.2 shows the distributions of these same variables for MC and data tracks, with all cuts applied except the one being examined. Within the cut values, the shapes of these distributions are well-modeled by the MC.



Figure 7.1: Distributions of six of the variables upon which cuts are placed to select quality tracks. The curves shown are solid for those tracks from *uds* events and/or the fragmentation tracks in c and b events (uds + frag), dotted for those from K^0 and Λ decays and T-conversions (Vees), and dashed for those from c and b hadron decays (c + b). Some salient features are the harder momentum spectrum for those tracks from heavy quark decays and the long tails on the impact parameters for strange decays and other Vees.



Figure 7.2: Distributions of six of the variables upon which cuts are placed to select quality tracks, comparing data (points) and MC (histogram).

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Figure 7.3: The distributions of the number of quality tracks for those events passing the hadronic selection cuts in the three running periods. The discrepancy between the charged multiplicities between data and MC can easily be seen.

A potentially serious problem is masked by the normalization of these plots, however, **as** data and MC have been normalized so that the total number of tracks is the same. Figures 7.3a-c display the problem in a more telling manner, showing the number of quality tracks n_{good} in the data and MC for those *events* passing the hadronic event selection cuts. For this generation of MC, the tracking efficiency in simulated events is noticeably higher than that in the data. This is due to correlated hit loss at the edges of the CDC drift cells, where an entire superlayer's worth of hits can be lost if the track passes too close to the field wires. The details of the charge deposition near the edges of the CDC cells were oversimplified in the MC, resulting in a difference



Figure 7.4: Diagram of a B Meson Decay Cascade

between data and MC in the number of hits available for track finding and fitting*. Studies have shown [182] that the track loss occurs in an essentially random manner and specifically does not depend on track p_{\perp} , ϕ , θ , or the angle of the track with respect to the jet axis. This is almost certainly due to the randomly uniform location of cell edges within the CDC. So, to insure agreement between the MC and data in terms of the number of quality tracks available for tagging, we apply a track-by-track efficiency correction to all MC events such that the mean values of n_{aood} agree for data and MC events. The error on this correction could be arbitrarily small, since the large number of tracks in each sample results in a tiny error on $\langle n_{aood} \rangle$. To be conservative, however, we compare our MC multiplicity to the world average, since we don't know a priori that the MC has the correct multiplicity. The mean total charged multiplicity at the Z^0 has been best determined by by ALEPH[183], who measured 20.85 \pm 0.24 and by OPAL[184], who measured 21.4 \pm 0.43. One can see that these two measurements differ by 0.6 tracks, yet are consistent within errors. The average total charged multiplicity in the SLD MC, 21.1 tracks, is consistent with the these other results, but we will assign a f0.3 tracks error to the efficiency correction to cover any possible error.

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7.2 The Normalized Impact Parameter

From this point on, we will refer to the xy distance of closest approach to the IP for a charged track $doca_{xy}$ as the impact parameter d; the associated error on this quantity is defined as σ_d . The basis of the tag method, originally suggested by Hayes [185], is that tracks from the decays of mesons containing c and especially bquarks tend not to extrapolate back to the IP. This is due to two factors, the long meson lifetimes and the significant p_{\perp} acquired relative to the meson flight direction, a result of the large available energy in the heavy quark decay. Figure 7.4 illustrates this point. In principle, a vertex detector with sufficient resolution should be able to measure the impact parameters of these tracks well enough to allow a sample of heavy quark events to be selected based on the number of tracks that miss the IP. As an improvement to this technique, the Mark II collaboration [186] introduced a signed impact parameter, where now d is positive if the track crosses the meson flight direction downstream of the IP, and negative if it crosses the meson flight direction upstream or behind the IP[†]. This creates an asymmetric distribution, since the tracks from heavy quarks tend to populate the region of large positive values of d. Figure 7.5 shows the signing procedure in cartoon form. Since we wish to be insensitive to track mis-reconstruction, the absolute distance scale, and our imprecise knowledge of the meson lifetimes, we will use instead the significance of the distance by which each track rnisses the IP, which we refer to as the normalized impact parameter $d/\sigma_d \equiv d_{norm}$. The distributions of d and d_{norm} were shown in Section 5.3 during the earlier discussion of tracking performance. Here, we show in Figure 7.6 the d_{norm} distribution for all tracks from uds, c, and b events in the MC as an illustration of the different regions that are populated by tracks from events with long-lived heavy mesons.

In this analysis, we have completely decoupled the tagging from any of the jetfinding studies by choosing a fixed definition of the procedure used to sign the impact parameters of the tracks. For the event tags, we employ the JADE jet-finding algorithm[126] with calorimeter clusters as input. The jet.resolution parameter y_{cut}

^{*}This has been corrected in the latest round of MC event generation.

[†]In practice the jet axis is used as an approximation to the meson flight direction.



Figure 7.5: Definition of the signed impact parameter. Tracks which cross the jet axis "downstream" of the IP are given positive impact parameters.

of 0.02 is used to define these jets. Several studies by the SLD B-physics group[182] have shown this to provide the maximum fraction of correctly-signed tracks in the sample of hadronic decays.

7.3 Definition of the Flavor Tags

From the distributions shown in Figure 7.6, one can see that b and c events do tend to have a larger fraction of tracks missing the IP by a few sigma. The impact parameter tagging method attempts to capitalize on this difference to separate events with primary light quarks from those with primary heavy quarks. The two variables that define a tag of this type are the impact parameter significance (i.e., how far away from the IP does the track have to be to be have missed it significantly?) and the number of tracks of that significance in each event (how many tracks significantly miss the IP?), which we will refer to as n_{sig} . To choose the tag parameters, we undertook a study of the efficiencies ε^{\ddagger} , purities $\Pi^{\$}$, and tag biases[¶]. Unfortunately, the search of a 27-dimensional space doesn't lend itself to graphical representation. Since tags of this manner are essentially guaranteed to be of high efficiency, we searched for tags

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Figure 7.6: The impact parameter distribution for all tracks in hadronic events, and all tracks in uds, c, and b events. The distributions shown are normalized to the total number of tracks in each sample. The long tails at high impact parameters are readily visible in the c and b events.

[†]the fraction of the total number of events of a chosen flavor that end up in the tagged sample which is supposed to select that flavor

[§]the fraction of events in the tagged sample that are of the desired flavor

^{&#}x27;the difference in efficiency for tagging a 2-jet event and a 3-jet event of a given flavor, to be tliscussed in Chapter 8 $\,$

that gave the highest possible purities with the smallest tag bias. High purity samples are desired because large background 'subtraction" that occurs when unfolding the properties of the pure flavor sample from the mix of flavors in a given tag is reduced, as' are the correlations between the results for each flavor once the unfolding is completed. It is interesting to note that, as the significance cut of the tag is made harder, the *b* sample becomes more pure, but the *uds* purity is diluted. The opposite happens when the significance cut is lowered unless a large n_{sig} is required. An intermediate cut of a significance of 3.0 balances the purity in the *uds* and *b* samples. Figure 7.7 shows the number of tracks with $d_{norm} > 3.0$ in each event. The leftmost bin, the one with no significant tracks, is almost 90% *uds* events, while the bins with 4 or more significant tracks are extremely pure in *b* events. The middle bins contain an enriched sample of *c* events. Once again, the MC provides a good description of the data for this distribution.

The event sample is divided accordingly into three mutually exclusive tagged subsamples, each of which will be referred to as the i^{th} tagged sample:

- . those events with $n_{sig} = 0$ were defined to be the uds-tagged sample (i = 1)
- . those with $1 \le n_{sig} \le 3$ were defined to be the c-tagged sample (i = 2)
- . those with $n_{sig} \ge 4$ were defined to be the btagged sample (i = 3).

For illustration, we list here a sample set of tag efficiency matrices, for 2- and J-jet events 11 which have passed the selection cuts. The matrices ε_n^{ij} are the efficiencies for tagging an event with n = 2, 3 jets of true flavor j (j = 1 : uds, 2 : c, 3 : b) with tag *i*. Hence, the columns contain the probabilities that an event of a given flavor will be tagged by each of the tags; each column should thus sum to unity.

$$\varepsilon_{2}^{ij} = \begin{pmatrix} 0.771 \ 0.374 \ 0.044 \\ 0.228 \ 0.593 \ 0.439 \\ 0.001 \ 0.033 \ 0.517 \end{pmatrix} \varepsilon_{3}^{ij} = \begin{pmatrix} 0.732 \ 0.411 \ 0.072 \\ 0.266 \ 0.560 \ 0.512 \end{pmatrix}$$
(7.1)

^{IJ}Jets in this example have been found with the 'E' jet-finding algorithm with $y_{cut} = 0.08$, to be defined in the following chapter.

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For completeness, we show also the purity matrices for 2- and J-jet events. Here, the rows give the purities of each tag:

$$\Pi_{2}^{ij} = \begin{pmatrix} 0.865 & 0.117 & 0.018\\ 0.413 & 0.300 & 0.287\\ 0.005 & 0.047 & 0.948 \end{pmatrix} \qquad \Pi_{3}^{ij} = \begin{pmatrix} 0.839 & 0.131 & 0.030\\ 0.438 & 0.257 & 0.304\\ 0.013 & 0.049 & 0.939 \end{pmatrix}$$
(7.2)

For such a simple tag, we actually do quite well in terms of absolute performance: the b tag is about 51% efficient, with a high purity of about 96%. The excellent resolution of the tracking system makes this quality of tag easy to achieve.

We move on now to a description of the jet-finding that is performed on these tagged samples to make possible the extraction of the strong coupling.



Figure 7.7: The number of tracks n_{sig} per event which miss the origin by more than 3σ in the z-y plane. The solid histogram represents the total MC sample, while the points represent the data distribution. The flavor composition of the distribution is also shown. (See text.)

Chapter 8

Jet Finding

In order to extract the strong coupling from the final state hadronic structure of Z^0 decays, we must **first** set forth the definition(s) of the jets into which we will organize them.

8.1 Jet Algorithm Definitions

As mentioned in Section 3.3.1, jets are defined operationally by iterative clustering algorithms. When these algorithms are applied, a measure y_{kl} , such as the invariant mass, of all pairs of particles k and 1 in a given event is calculated, and the pair with the smallest y_{kl} is combined into a single (pseudo)-particle. This process is repeated until all pairs have y_{kl} exceeding a value y_{cut} . At this point, each pseudo-particle that remains is called a "jet", and the jet multiplicity of the event is just the number of jets.

The different algorithms are distinguished by their definition of the measure y_{kl} and the manner in which the two particles k and 1 are combined (known as the "recombination scheme"). In the analysis presented here, we use six of these jet algorithms: the 'E', 'EO', 'P', and 'PO' variations of the JADE algorithm[126], as well as the Durham ('D') and Geneva ('G') schemes[187]. The definitions of y_{kl} and the recombination schemes for each algorithm are presented below*

^{&#}x27;In the text to follow, fax-momenta will lx represented by the symbol p_m , where m refers to

8.1.1 The E Scheme

In the E scheme, y_{kl} is defined as the square of the invariant mass of the particles k. and 1 normalized by the visible energy in the event, E_{vis} ,

$$y_{kl} = \frac{(p_k + p_l)^2}{E_{vis}^2} \tag{8.1}$$

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The two particles k and 1 are combined into a new (pseudo)-particle m according to

$$p_m = p_k + p_l av{8.2}$$

where pion masses are assumed in calculating the energies. Total energy and momentum are explicitly conserved in this scheme.

8.1.2 The EO Scheme

In the EO scheme, y_{kl} is defined by Equation 8.1, while the recombination scheme is given by

$$E_{m} = E_{k} + E_{l}$$

$$\vec{p}_{m} = \frac{E_{m}}{|\vec{p}_{k} + \vec{p}_{l}|} (\vec{p}_{k} + \vec{p}_{l})$$
(8.3)

Here, \vec{p}_m is rescaled so that particle *m* has zero invariant mass. This scheme does not conserve the total momentum in an event.

8.1.3 The P Scheme

In the P scheme, y_{kl} is given by Equation 8.1 and the recombination is defined by

$$p_m = \vec{p}_k + \vec{p}_l$$

$$E_m = |\vec{p}_m| . \qquad (8.4)$$

This scheme does not conserve the total energy of an event, but does conserve the total momentum sum.

8.1.4 The PO Scheme

The PO scheme identical to the P scheme, except that the visible energy E_{vis} in Equation 8.1 is recalculated after each iteration according to

$$E_{vis} = \sum_{m} E_m . \tag{8.5}$$

8.1.5 The D Scheme

In the D scheme, the parameter y_{kl} is defined as

$$y_{kl} = \frac{2\min(E_k^2, E_l^2)(1 - \cos\theta_{kl})}{E_{vis}^2}, \qquad (8.6)$$

where θ_{kl} is the angle between particles k and l. The recombination scheme is given by Equation 8.2. From an inspection of Equation 8.6, one can see that the D algorithm will tend to cluster soft particles with other nearby soft particles, a property not shared with the JADE-derived algorithms, which often associate soft particles with a jet in the opposite thrust hemisphere.

8.1.6 The G Scheme

The definition of y_{kl} for the G scheme is

$$y_{kl} = \frac{8E_k E_l (1 - \cos \theta_{kl})}{9(E_k + E_l)^2} .$$
 (8.7)

The recombination scheme is also given in this case by Equation 8.2. Note that the G scheme has the same properties *vis-à-vis* soft particles as the D scheme, but in this case y_{kl} only depends on the energies of the two particles under consideration, not on the total E_{vis} .

8.2 Definition of the n-Jet Rate and D_2

In determining a, from QCD predictions, it **is** of course the *rate* of multi-jet events that matters, not the number of jets in each event. So, we define the n-jet rate

the m^{th} particle. Three-momenta will be distinguished by the appearance of a vector symbol over the $p, e.g., \vec{p}_m$. The energy of particle m will be represented by the symbol E_m .

 $R_n(y_{cut})$ as the fraction of events that have n jets, where the jets have been found using a jet algorithm with mass cut-off y_{cut} . For purposes of illustration, we show in Figure 8.1 $R_n(y_{cut})$ for n = 2, 3, 4, 5 for all of the algorithms, where the jets have been found using charged tracks in the detector.

Since each of the points at each value of y_{cut} in Figure 8.1 is determined using the entire data set, all the points show are highly correlated. One can remove this correlation by defining the differential 2-jet rate, D_2 [188]

$$D_2(y_{cut}) = \frac{R_2(m) - R_2(y_{cut} - \Delta y_{cut})}{\Delta y_{cut}}$$
(8.8)

An event will only contribute to the D_2 distribution if it is redefined as a 2-jet event instead of a S-jet, event as the value of y_{cut} is changed from y_{cut} to $y_{cut} - \Delta y_{cut}$. Sample D_2 distributions will be shown in Section 9.3 and at the end of the present chapter.

8.3 QCD calculations for Jet Rates and D_2

As mentioned in the introduction, perturbative QCD calculations which relate observable quantities to the strong coupling α_s have been performed for all of the jet-finding algorithms and the D_2 distribution defined above. In the next two sections, we discuss these calculations and their implications for the prospects of determining α_s .

The general form of the predictions for the 3-jet rate $R_3[128, 129)$ was presented in Section 3.3.1 and is reproduced here for reference:

$$R_{3}(y_{cut}) = \frac{\alpha_{s}(\mu)}{2\pi} A(y_{cut}) + \left(\frac{\alpha_{s}(\mu)}{2\pi}\right)^{2} \left[B(y_{cut}) + A(y_{cut})2\pi b_{0} \log f\right]$$
(8.9)

Here μ is the renormalization scale defined in Section 3.1.3, and $f = \mu^2/s$. Here, $b_0 = (33 - 2n_f)/(12\pi)$, where n_f is the number of active flavors ($n_f = 5$ at $\sqrt{s} = M_Z$). The terms proportional to α_s^2 include contributions (B) from four-jet events where two of the jets were sufficiently close to be unresolvable at the chosen value of y_{cut} . The calculation relating D_2 to α_s is somewhat messy, as it requires differentiating the expression in Equation 8.9 with respect to y_{cut} and thus has no analytic closed form. The 4-jet rate R_4 has the form has only been calculated to lowest order[87], and has



Figure 8.1: Plots showing the fractional jet rates for each of the 6 jet algorithms. The various points represent the n-jet rates, and the solid curves are the same quantities derived from the MC simulation of hadronic Z^0 decays.

the expected form:

$$R_4(y_{cut}) = \left(\frac{\alpha_s(\mu)}{2\pi}\right)^2 C(y_{cut}) . \tag{8.10}$$

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In general, the coefficients $A(y_{cut})$ and $C(y_{cut})$ are the same for all algorithms, since there should be no dependence of the lowest order result on the choice of some specific jet-finding algorithm. The three coefficients **A**, **B**, and C were evaluated using the EVENT program developed by Kunszt and Nason[128], and tables of the values obtained are presented in Appendix B. Typical values of **A**, **B**, and C, for example for the **E** algorithm at y_{cut} of 0.08, are Note again that the renormalization scale μ appears in the formula for R_3 . Its effect on the analysis will be discussed in Chapters 9 and 11.

We find in examining the results of the calculations that the 4-jet rate predicted by Equation 8.10 is not that observed in the data for all but the G algorithm. We need to increase the coefficient $C(y_{cut})$ by \sim a factor of two for all of the other jet algorithms to obtain good agreement with the data. The full range of this variation will be taken as the systematic error due to the uncertainties in this calculation.

8.4 The Need for Corrections

It should be stressed again that the perturbative QCD calculations described in the previous section refer to quantities calculated at parton-level in the theory. Ideally, the measurement of some global property of hadronic events using the final state hadrons should not be effected by the hadronization process itself. Unfortunately, losses due to such things as detector acceptance, non-perfect tracking and calorimetry, and the need to select fiducial events from a sample containing background all imply that the global properties measured with the actual detector apparatus will not be the same as those exhibited by the underlying partons. The typical remedy[131] to this situation is to apply corrections, either bin-by-bin over the distribution in question or using the full bin-to-bin correlation matrix! to "correct" the measured data distribution for all of its mistreatment, resulting in a sample that should resemble the true parton distribution. This requires, of course, a believable detector MC, as well as some model

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of the hadronization process. Examples of corrected D_2 distributions, which can then be compared directly with the theoretical predictions, are shown in Section 9.3. In this analysis, we will also need to make corrections of this kind, and thus require good agreement between MC and data for the distributions of interest. In order to demonstrate the quality of the MC description of the detector, we show in Figures 8.2-8.4 the D_2 distributions for the **tagged** samples in the three run periods. The MC and data agree quite well over the many decades of these distributions. We will not actually use the D_2 distribution to extract α_s for the different quark flavors, but it serves here as a good cross check and gives us confidence that we understand the effects of the detector on the jet rates well enough to attempt the needed corrections. The methods used for the corrections will be described in the next chapter. 171



Figure 8.2: The D_2 distribution for the data (points) and MC (line) uds, c, and b-tagged samples. The data and MC are both taken from the Pre-veto run period.



Figure 8.3: The D_2 distribution for the data (points) and MC (line) uds, c, and b-tagged samples. The data and MC are both taken from the Veto run period.



Figure 8.4: The D_2 distribution for the data (points) and MC (line) uds, c, and b-tagged samples. The data and MC are both taken from the Non-veto run period.

Chapter 9

Analysis Method

Now that the tools for flavor tagging and jet-finding have been presented, we can move on to a discussion of the method by which we can extract a. for each quark flavor. This chapter presents a detailed discussion of the choice of method and the mathematical formulation of the unfolding and parton-level-correction procedure, and ends with the results for the jet rates determined for each quark flavor.

9.1 Considerations

From the inception of this project, it has been clear that the most precise test of the flavor-independence of strong interactions must come from measuring a **ratio** of the strong couplings of one quark flavor to another or to the entire sample. The ratio must be formed in a way such that the uncertainties due to detector effects, hadronization models, and renormalization scale uncertainty largely cancel. If this is not the case, then the test of flavor-independence will be plagued by the same large sources of theoretical uncertainty that effect an absolute determination of $\alpha_{\mathfrak{o}}$. These problems were discussed in some detail in Chapter 3. Until this work, the standard practice' has been to isolate an enriched sample of the quark flavor of interest and then to apply some combination of the techniques discussed in Chapter 3 to measure $\alpha_{\mathfrak{o}}$ for that particular quark flavor, which is then compared to a, measured from

See Section 3.4.1

the whole sample of hadronic decays using the same techniques. This requires one to make the parton-to-hadron and hadron-to-detector corrections separately for the tagged sample and the whole set of hadronic decays, and then hope that they (and hence there uncertainties) cancel to a large extent when the ratio is formed. This may not be the case for events containing heavy quarks, for example, which then have to be given special treatment. In addition, some assumption must be made about the value of α_s for the background in the tag sample, since no tag produces samples of 100% purity. One can attempt, as have the OPAL collaboration[145], to loosen these assumptions by combining the results from an exclusive, low efficiency tag for each flavor and unfold the ratios of couplings for pure quark samples. This works, but has some large errors from the unfolding procedure associated with it.

Instead of pursuing this method, WC have chosen to pursue an inclusive tagging analysis, where we use all hadronic events in our fiducial sample. We have seen in Chapter 7 that the multiplicity of significant-impact-parameter tracks provides a clean tag of uds and b events and some separation for c events when the entire event sample is broken into the three tagged samples defined there. This tag definition is based on an "inclusive" variable like the number of tracks, rather than something "exclusive", like requiring fast protons. Hence, one can be more certain that the efficiency matrices used to unfold the tagged distributions into results for pure flavor samples are insensitive to the minute details of the MC.

9.2 The Choice of R_3

as the Analysis Variable

Once the flavor separation has been performed, one is left with the task of actually determining a, or α_s ratios from the tagged samples. In principle, fitting the D_2 distribution would be the obvious choice, since it maximizes the statistical weight of each event. However, to avoid uncertainties due to the large corrections necessary to obtain the parton level distributions for each tagged sample and the global event sample separately, we really would like to fit, the *ratio* of the corrected values as a

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function of y_{cut} . This proves to be a somewhat tangled undertaking, as both distribution shape and normalization information become relevant to the fit. In addition, we would then be fitting the ratio of two second derivatives, which doesn't exactly lead to a clear, intuitive grasp of the problem. If we add to this the difficulty that the fitting apparatus becomes extremely complicated and hard to understand in a straightforward manner, we are driven to search for a more elegant means of extracting a, from the data.

This is provided by returning to the 3-jet rate R_3 as the basis for the analysis. Calculation of R_3 for each of the tagged samples is a simple task, and if jets are found at a single value of y_{cut} for each algorithm, there is no double counting of events. We have in Equation 8.9 a perturbative QCD calculation for R_3 relating the parton-level Jet rate to α_s , so we can actually perform a measurement with this observable, and a ratio of the three-jet rate for a tagged sample to that of the entire event sample can easily be formed. This also provides us with a more intuitive variable space in which to perform the analysis. In the case of jet rates, calculations of the effects of heavy quark masses on the phase space for hard gluon emission have been performed, and we can use them to correct our observed distributions for the kinematic differences of heavy quark events. As we shall see, this choice of method allows us to proceed in series of well-defined steps from the tagged distribution to the final ratios of α_s values.

9.3 Choosing the Values of y_{cut}

In order to perform the jet-rates analysis, we need to define what we mean by a "jet" for each jet-finding algorithms discussed in Chapter 8. This involves the choice of the jet resolution parameter y_{cut} , which gives the maximum allowed invariant mass of a pair of (pseudo)-particles such that they will be combined into a single cluster. Originally, we worried about the uncalculated higher orders in the QCD predictions for $R_3(y_{cut})$, thinking they would introduce large theoretical uncertainties. To ameliorate this problem, we chose at that time y_{cut} values such that the 4-jet rate $R_4(y_{cut})$ was less than 1%. After more reflection, we redefined the 3-jet rate $R_3(y_{cut})$ so that

 $R_3(y_{cut}) \equiv R_3(y_{cut}) + R_4(y_{cut}) + R_5(y_{cut}) + ...^{\dagger}$, as it turned out that the additional uncertainty in doing this is small[‡]. Now, we are free to try and minimize the statistical error on R_3 , by choosing a y_{cut} value that gives the largest number of 3-jet events. We' must be wary, though, of selecting y_{cut} in a region of parameter space where the data are not well described by the $\mathcal{O}(\alpha_s^2)$ calculations that we wish to use to extract α_s , or where the decay products of heavy quarks are split into different jets. A solution to the problem is provided by the parton-level-corrected D_2 distributions calculated by Ohnishi [131]. Figures 9.1 through 9.6 show the parton-corrected data plotted on top of the full $\mathcal{O}(\alpha_{\bullet}^2)$ calculation. The deviations of the data and the theory at low y_{cut} occur when the contributions to the 2-jet rate from the emission of many soft gluons becomes important. One can see that this happens at different values of y_{cut} for each algorithm, as can be expected. Since it is this $\mathcal{O}(\alpha_s^2)$ calculation that we would like to use to extract α_s for this analysis, we are free to choose the lowest value of y_{cut} where the data is well-described by the theory. The chosen y_{cut} values are denoted by the label "this analysis". These values also appear in Table 9.3. These figures also display the hadronization (C_H) and detector acceptance and resolution corrections (C_{D}) that were applied to the data to obtain the parton-level distributions. The fit ranges that were used for extracting α_s from fitting the D_2 distributions are also shown. For some algorithms, these are smaller than the full range over which the theory appears to agree with the data due primarily to the size and uncertainty in the applied corrections C_H and C_D . We are free to explore a larger range of y_{cut} as we expect the uncertainty due to these corrections to be relatively small in this analysis.

9.4 Fitting the Data for R_3

Here, we present the mathematical formulation of the fit which allows us to extract R_3 at the parton-level for each quark flavor from R_3 in each of the tagged samples. The parton-level 3-jet rate R_3^j for each of the j quark types (j = 1 : uds, j = 2 : c,



Figure 9.1: The D_2 Distribution and $\mathcal{O}(\alpha_s^2)$ Calculation for the E algorithm. The grey bands show the range of uncertainty on the corrections which would be applied to the data to obtain the proper "hadron'-level (C_D) and parton-level $(C_H) D_2$ distributions if one wanted to extract a value of α_s from this distribution. The solid curve in the top plot is the $\mathcal{O}(\alpha_s^2)$ QCD calculation of D_2 . The arrow shows the value of y_{cut} selected for this analysis (see text).

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[†]Note that this implies that $R_2 = 1 - R_3$.

[‡]This will be discussed in more detail in Chapters 10 and 11, but the effect is $\sim 0.3\%$



Figure 9.2: The D_2 Distribution and $\mathcal{O}(\alpha_s^2)$ Calculation for the EO algorithm. See the caption of Figure 9.1 for details.



Figure 9.3: The D_2 Distribution and $\mathcal{O}(\alpha_s^2)$ Calculation for the P algorithm. See the caption of Figure 9.1 for details.



Figure 9.4: The D_2 Distribution and $\mathcal{O}(\alpha_s^2)$ Calculation for the PO algorithm. See the caption of Figure 9.1 for details.



Figure 9.5: The D_2 Distribution and $\mathcal{O}(\alpha_s^2)$ Calculation for the D algorithm. See the caption of Figure 9.1 for details.



Figure 9.6: The D_2 Distribution and $\mathcal{O}(\alpha_s^2)$ Calculation for the G algorithm. See the caption of Figure 9.1 for details.

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and j = 3: b) and for each jet algorithm was extracted from a maximum likelihood fit to the following expressions for n_2^i and n_3^i , the number of 2-jet and 3-jet events, respectively, in the *i*th tagged sample:

$$n_{2}^{i} = \sum_{j=1}^{3} (\varepsilon_{(2 \to 2)}^{ij} (1 - R_{3}^{j}) + \varepsilon_{(3 \to 2)}^{ij} R_{3}^{j}) f^{j} N$$

$$n_{3}^{i} = \sum_{j=1}^{3} (\varepsilon_{(3 \to 3)}^{ij} R_{3}^{j} + \varepsilon_{(2 \to 3)}^{ij} (1 - R_{3}^{j})) f^{j} N \qquad (9.1)$$

Here N is the total number of selected events corrected for the efficiency for a hadronic Z^0 decay to pass all event selection cuts, which effectively makes N the total number of parton-level events that would correspond to observing $n_2^i + n_3^i$ data events after cuts and detector acceptance. The symbol f^j is the Standard Model fractional hadronic width for Z^0 decays to the j^{th} quark type, i.e.,

$$f^{j} = \frac{BR(Z^{0} \to q_{j}\bar{q}_{j})}{BR(Z^{0} \to \text{hadrons})} .$$
(9.2)

The matrices $\varepsilon_{(2\to2)}^{ij}$ and $\varepsilon_{(3\to3)}^{ij}$ are the efficiencies for an event of type *j* containing 2- or 3-jets at the parton level to pass all cuts and be tagged as a 2- or 3-jet event, respectively, of type *i*. The matrices $\varepsilon_{(2\to3)}^{ij}$ and $\varepsilon_{(3\to2)}^{ij}$ are the efficiencies for an event of type *j* containing 2- or 3-jets at the parton level to pass all cuts and be tagged as a 3- or 2-jet event, respectively, of type *i*. We remind the reader that, by definition, $R_2 = 1 - R_3$. The efficiency matrices ε were calculated from the MC by performing the jet-finding at the parton level, then subjecting the events to the flavor tag and jet-finding at the detector level. One can think of the matrices ε for 2- and 3-jet events as being the multiplicative product of three efficiency matrices

$$\varepsilon_{n \to m}^{ij} = \varepsilon_{cuts} \cdot \varepsilon_{n \to m}^{jets} \cdot \varepsilon_{tag}^{ij}$$
(9.3)

where the first two matrices are diagonal, since only the flavor tagging mixes the different quark species. The first matrix, ε_{cuts} , contains the efficiencies for the events of the different quark flavors to pass the trigger and event selection cuts. The second matrix keeps track of the migration of the parton-level jets into the observable final states in the detector, and the third matrix is a measure of the efficiency for tagging

the appropriate final state. In general, all three of these matrices are different for 2- and 3-jet events. Values for ε_{cuts} were given at the end of Chapter 6, and sample tagging efficiency matrices ε_{ij}^{tag} were given in Chapter 7. For completeness, the four efficiency matrices for each jet algorithm and the mathematics for performing the fit are contained in Appendix A. The determination of these matrices are the heart of this analysis; virtually all systematic errors can be expressed as a variation in some or all of their elements.

One can see that this formalism explicitly accounts for modifications of the partonlevel 3-jet rate due to hadronization, detector effects, and tagging bias. Since the jet rates for all three flavors are fit simultaneously, there are also no assumptions made about the values of the strong couplings for each quark type. In fact, the total 3-jet rate R_3^{all} is just given by

$$R_3^{all} = f_{uds} \cdot R_3^{uds} + f_c \cdot R_3^c + f_b \cdot R_3^b , \qquad (9.4)$$

so it, too, comes from the fit.

Note that, by choosing to formulate the problem in terms of efficiency matrices, we do not depend on the MC to properly reproduce the rate of 3-jet production and thus this analysis is only weakly dependent on the choice of α_s used in generating the MC sample.

9.5 Tag Bias

We pause briefly in the description of the analysis to discuss the bias of these tags towards selecting Z-jet events preferentially over 3-jet events. This can happen for a number of reasons depending on the tag in question. For example, requiring a high-p, high- p_{\perp} lepton to tag **b** events tends to tag those events where the **b** has acquired a large fraction of the available energy from the Z^0 decay. Since the **b** hadrons in events with hard gluon radiation are less energetic, they are tagged less frequently by this tag, and the tag is therefore biased toward selecting 2-jet events. To quantify this, we define the tag bias B^i as the difference from unity of the ratio of the *i*th diagonal

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elements of the 2- and 3-jet efficiency matrices:

$$B^{i} \equiv 1 - (\varepsilon_{(2 \rightarrow 2)}^{ii} + \varepsilon_{(2 \rightarrow 3)}^{ii}) / (\varepsilon_{(3 \rightarrow 3)}^{ii} + \varepsilon_{(3 \rightarrow 2)}^{ii})$$

$$(9.5)$$

Averaged over the six jet algorithms, the biases are:

$$B^{uds} = 0.057 \pm 0.005, \quad B^c = 0.083 \pm 0.016, \quad B^b = 0.303 \pm 0.037$$
, (9.6)

where the errors represent the r.m.s variations. In contrast to the previously employed tags[§], the tag biases for *uds* and c events are relatively small. In general, lifetime tags of the type we use here are sensitive to the average impact parameter resolution of the ensemble of tracks in the desired event. This provides the key to understanding the tag biases mentioned above. As the gluon in a 3-jet **b** event becomes more and more energetic, the recoiling **b** hadrons carry less and less energy. Since the impact parameter resolution rapidly becomes large for lower momentum particles, the tracks from the decay of these softer *bs* will be less likely to have a significant impact parameter because the measurement of their trajectories will be less precise. The smaller boost of the inefficiency for tagging 3-jet events. The differences in the case of **uds** events can be mainly attributed to event containment effects, as there is some loss of efficiency when one of the jets exits the detector along the beam pipe.

As mentioned above, these biases are explicitly corrected by the matrix unfolding procedure. It is interesting to note the variation in bias among the different algorithms and among the Pre-Veto, Veto, and Non-Veto run periods, Table 9.1 lists all of these values. One can see that the inherent bias in this method, given by the Pre-Veto and Non-Veto samples, is much smaller than those quoted in Equation 9.6, even for the \boldsymbol{b} tag. The track trigger veto has drastically increased the bias of each of the tags by preferentially vetoing 3-jet events. Since we believe that the effects of the veto are properly modeled in the MC, the corrections we apply in the unfolding fit should still yield the correct parton-level distributions.

See Section 3.4.1.

Table 9.1: The tag biases, as defined in Eq. 9.5 for each algorithm and each running period.

Jet Algorithm	B^{uds}	B^c	B^b
Pre-Veto:			
Е	1.012	1.041	1.244
EO	1.019	1.017	1.203
Р	1.024	1.041	1.172
РО	1.027	1.039	1.172
D	1.028	1.049	1.266
G	1.039	1.057	1.214
Veto:	Ť		
Е	1.175	1.259	1.621
EO	1.159	1.204	1.511
Р	1.139	1.202	1.474
PO	1.143	1.215	1.477
D	1.182	1.257	1.670
G	1.135	1.213	1.456
Non-Veto:			
Е	0.969	0.975	1.151
EO	0.967	0.961	1.138
Р	0.979	0.949	1.137
РО	0.975	0.951	1.137
D	p.972	0.957	1.170
G	0.978	0.952	1.124

9.6 Results for R_3^j

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Table 9.2 fists the number of events for each run period that make up the inputs to the unfolding procedure defined in Equation 9.1. The MC sample used for this analysis consists of 286,764 events, divided into 102,360 Pre-Veto events, 137,773 Veto events, and 46,631 Non-Veto events. We solve Equation 9.1 for each of the six jet algorithms for the Pre-Veto, Veto, and Non-Veto periods separately'. The results are then combined by taking the proper weighted average. The extracted 3-jet rate ratios R_3^j/R_3^{all} are shown in Table 9.3, where the statistical errors on each R_3 value have been obtained using the full covariance matrix. Note that, since the same events are used to obtain each of these results, the values are highly correlated. Since we are unfolding these values from a single n_{sig} distribution containing all of the events, correlations are induced between the values obtained for each of the quark flavors. Averaged over all six jet algorithms, the correlation coefficients from the fit are: udsc: -0.76 ± 0.02 , uds-b : 0.30 ± 0.02 , c-b : -0.55 ± 0.01 , where the errors are the r.m.s. variations. They follow the pattern one might expect: since the **uds** and **b** samples are relatively pure, the correlation between the values obtained for these two samples should be small. The correlation between the c result and those for the other samples is large due to the much lower purity of the charm sample, which means that the result for the 3-jet rate for c quarks is extremely sensitive to the tails of the b and **uds** n_{sig} distributions in the c tag region.

The next chapter will bring us from these raw extracted jet rate ratios to the final result.

This is necessary because of the different multiplicity corrections that need to be applied.
n_2^{uds}	n_2^c	n_2^b	n_3^{uds}	n_3^c	n_3^b
3485	2082	804	963	734	220
3447	2102	815	1001	714	209
3157	1919	746	1291	897	278
3063	1856	711	1385	960	313
3428	2074	806	1020	742	218
2493	1494	636	1955	1322	388
5366	2981	1084	1169	722	209
5320	2971	1086	1215	732	207
4918	2760	1025	1617	943	268
4810	2685	994	1725	1018	299
5297	2960	1098	1238	743	195
3999	2274	874	2536	1429	419
3237	2014	794	1042	714	182
3217	2009	800	1062	719	176
20/0	1836	733	1339	892	243
2740	1000				
2823	1778	718	1456	950	258
2823 3177	1778 1980	718 789	1456 1102	950 748	258 187
	n2 ^{uds} 3485 3447 3157 3063 3428 2493 5366 5320 4918 4810 5297 3999 3237 3217	$\begin{array}{c ccccccccccccccccccccccccccccccccccc$	$\begin{array}{c ccccccccccccccccccccccccccccccccccc$	$\begin{array}{c ccccccccccccccccccccccccccccccccccc$	$\begin{array}{c c c c c c c c c c c c c c c c c c c $

Table 9.2: The number of 2- and 3-jet events tagged by each tag for each algorithm and run period. These are the n_2^i and n_3^i that are input to the fit in Eq. 9.1.

Table 9.3: The results for the ratios R_3^j/R_3^{all} obtained in inverting Equation 9.1 for each of the jet-finding algorithms used with the given y_{cut} value.

Algorithm	y_c	R_3^{uds}/R_3^{all}	R_3^c/R_3^{all}	R_3^b/R_3^{all}
E	0.080	0.941 ± 0.042	1.212 ± 0.173	0.980 ± 0.063
E0	0.050	0.975 ± 0.036	1.113 ± 0.145	0.981 ± 0.053
Р	0.030	1.001 ± 0.027	0.985 ± 0.109	1.007 ± 0.041
P0	0.030	1.014 ± 0.026	0.899 ± 0.102	1.037 ± 0.039
D	0.015	0.989 ± 0.035	1.096 ± 0.145	0.947 ± 0.049
G	0.030	1.032 ± 0.020	0.942 ± 0.079	0.952 ± 0.030

Chapter 10

From Jet Rates to α_s

This chapter brings to a close the description of the methods used in this analysis. Here, we will discuss the corrections that must be applied to the measured 3-jet rates to account for the suppression of hard gluon radiation due to the heavy quark masses. Then, we will present the final step in the extraction of the α_s values for each of the quark flavors, the conversion of the ratios R_3^j/R_3^{all} into the ratios $\alpha_s^j/\alpha_s^{all}$.

10.1 Phase-Space Suppression of Gluon Radiation in Heavy Quark Events

The 3-jet rate in heavy quark (b, c) events is expected to be reduced relative to that in light quark events by the diminished phase-space for gluon emission due to the heavy quark masses. This was first pointed out by Ioffe[189], and further explored most recently by Ballestrero, Maina, and Moretti[191]. This is exactly analogous to the mass effects on[192] the rates for the processes $e^+e^- \rightarrow \mu^+\mu^-\gamma$ and $e^+e^- \rightarrow \tau^+\tau^-\gamma$ compared to that for $e^+e^- \rightarrow e^+e^-\gamma$ in QED. In the QED calculation, the cross section for the radiative decays is a function of the "observability" of the photon in a detector, i.e., its energy and angle from the e, μ or τ , whereas in the QCD calculation the jet resolution y_{cut} is used. The implication of this effect is that one would tend to measure a lower 3-jet rate (and hence a lower α_s value) in heavy quark samples than

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for those from light (read massless) quarks.

We have used computer code provided by E. Maina that allows us to calculate the magnitude of the phase suppression effects for each algorithm at each y_{cut} for c and b quarks[193]. Figures 10.1 and 10.2 show the expected change in the ratios of R_3^c/R_3^u and R_3^b/R_3^d for each of the jet algorithms as a function of y_{cut} . The quark mass enters into the jet-rate definition only in terms of the definition for y_{kl} , The algorithms E, EO, P, and PO, then, are effected in an identical manner, since their definitions of y_{kl} are identical*. An arrow on each plot indicates the chosen value(s) of y_{cut} . We have assumed the c (b) quark mass is 1.5 (4.75) GeV/c². The dotted lines on each plot are the results for R_3^c/R_3^u and R_3^b/R_3^d when the quark masses are varied by f0.25 GeV/c². The suppression factors are also listed in Table 10.1.

In order to obtain corrected jet rates whose ratios can be related directly to the $\mathcal{O}(\alpha_s^2)$ calculation the 3-jet rate we have included these corrections in the unfolding procedure, such that the proper set of expressions for unfolding the R_3^j s is given by

$$n_{2}^{i} = \sum_{j=1}^{3} (\varepsilon_{(2 \to 2)}^{ij} (1 - P^{j} R_{3}^{j}) + \varepsilon_{(3 \to 2)}^{ij} P^{j} R_{3}^{j}) f^{j} N$$
(10.1)

$$n_3^i = \sum_{j=1}^3 (\varepsilon_{(3\to3)}^{ij} P^j R_3^j + \varepsilon_{(2\to3)}^{ij} (1-P^j R_3^j)) f^j N , \qquad (10.2)$$

where P^j is the phase-space suppression for R_3^j/R_3^{all} . The values of P^j are obtained from those in Table 10.1 with a small amount of algebra, the values of f^j , and the assumption that $R_3^u = R_3^d = R_3^g$. The values of R_3^j/R_3^{all} obtained with the phasespace-corrected unfolding procedure are given in Table 10.1.

10.2 Obtaining the Values for $\alpha_s^j/\alpha_s^{all}$

Now that we have removed any bias due to heavy quark masses, we can continue with the analysis and translate the ratios R_3^j/R_3^{all} into the desired ratios $\alpha_s^j/\alpha_s^{all}$. The needed relation is given by combining Equations 8.9 and 8.10; since we have defined $R_3 \equiv 1 - R_2$, we include contributions from the calculated rates of 3- and

^{*}See Section 8.1.

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Figure 10.1: The ratio R_3^c/R_3^u (solid curve) as calculated using the work of E. Maina. The arrows show the y_{cut} values used for the analysis. The dashed line above the solid curve is the same ratio, but for $m_c = 1.25 \text{ GeV}/c^2$. The dashed curve below the solid curve is the same ratio, but for $m_c = 1.75 \text{ GeV}/c^2$.



Figure 10.2: The ratio R_3^b/R_3^d (solid curve) as calculated using the work of E. Maina. The arrows show the y_{cut} values used for the analysis. The dashed line above the solid curve is the same ratio, but for $m_b = 4.5 \text{ GeV}/c^2$. The dashed curve below the solid curve is the same ratio, but for $m_b = 5.0 \text{ GeV}/c^2$.

Table 10.1: The values for R_3^j/R_3^{all} derived from applying the phase-space suppression correction P^j as shown in Eq 10.1. The calculated suppression factors $\mathbf{R}_3^{cu} = R_3^c/R_3^u$ and $\mathbf{R}_3^{bd} = R_3^b/R_3^c$ are also shown.

Algorithm	R_3^{uds}/R_3^{all}	R_3^c/R_3^{all}	R_3^b/R_3^{all}	\mathbf{R}_3^{cu}	\mathbf{R}_3^{bd}
Е	0.934 ± 0.041	1.200 ± 0.171	1.006 ± 0.064	0.995	0.958
EO	0.967 ± 0.036	1.098 ± 0.143	1.016 ± 0.054	0.994	0.945
Р	0.989 ± 0.027	0.968 ± 0.107	1.053 ± 0.043	0.992	0.929
РО	1.001 ± 0.026	0.883 ± 0.100	1.084 ± 0.041	0.992	0.929
D	0.977 ± 0.035	1.075 ± 0.142	0.997 ± 0.052	0.991	0.921
G	1.017 ± 0.020	0.924 ± 0.078	1.005 ± 0.032	0.989	0.915

4-jet production This yields the ratio

$$\frac{R_3^{i}(y_{cut})}{R_3^{all}(y_{cut})} = \frac{\frac{\alpha_s^2(\mu)}{2\pi}A(y_{cut}) + (\frac{\alpha_s^2(\mu)}{2\pi})^2[B(y_{cut}) + C(y_{cut}) + A(y_{cut})2\pi b_0 \log f]}{(\frac{\alpha_s^{all}(\mu)}{2\pi}A(y_{cut}) + (\frac{\alpha_s^{all}(\mu)}{2\pi})^2[B(y_{cut}) + C(y_{cut}) + A(y_{cut})2\pi b_0 \log f]}$$
(10.3)

The values of the ratios R_3^j/R_3^{all} from Table 10.1 are input into this equation and it is inverted, yielding the values for $\alpha_s^j/\alpha_s^{all}$ for each jet-finding algorithm. It is worth noting that there is a residual dependence on the value of α_s^{all} as it appears outside of the ratio when this expression is inverted. (This is easily seen by multiplying numerator and denominator by $1/(\alpha_s^{all})^2$.) To handle this correctly, we have used the value of α_s^{all} and the corresponding value of the renormalization scale f determined using the D_2 fit in Ref [131] for each algorithm. These are listed in Table 10.2.

The results for $\alpha_s^j/\alpha_s^{all}$ are given in Table 10.3 along with their statistical errors. The statistical errors here are the result of the propagation of the statistical errors from the R_3 ratios obtained in the previous chapter. The results for each algorithm for each quark flavor do not differ significantly from unity.

The extent that the ratios R_3^j/R_3^{all} and $\alpha_s^j/\alpha_s^{all}$ differ varies from algorithm to algorithm due to the differences in the coefficients A, B, and C. The difference in the value of $\alpha_s^j/\alpha_s^{all}$ from unity that comes from of inverting Equation 10.3 versus the input difference $R_3^j/R_3^{all} - 1$ is shown in Figure 10.3.

The next chapter will consider the systematic errors on these results.



Figure 10.3: Conversion Factors for R_3^j/R_3^{all} to $\alpha_s^j/\alpha_s^{all}$ for each of the 6 jet algorithms. In each plot, the horizontal axis is the input value of R_3^j/R_3^{all} and the vertical axis is the resultant value of $\alpha_s^j/\alpha_s^{all}$.

Algorithm	α^{all}_{s}	f
Е	0.118 ± 0.012	0.4
EO	0.128 ± 0.021	0.25
Р	0.116 ± 0.008	0.5
РО	0.114 ± 0.007	0.6
D	0.125 ± 0.010	0.5
G	0.108 ± 0.005	0.8

Table 10.2: The values for a_{i}^{all} , their errors, and the values of f at which they were derived using the D_2 distribution in Ref [131].

Table 10.3: The results of solving Eq. 10.3 for $\alpha_s^j/\alpha_s^{all}$ using the R_3^j/R_3^{all} values from Table 10.1. The errors shown are statistical only, but the same events are used by each of the jet-finding algorithms, so the results are highly correlated.

Algorithm	$\alpha_s^{uds}/\alpha_s^{all}$	$\alpha_s^c/\alpha_s^{all}$	$\alpha_s^b/\alpha_s^{all}$
Е	0.954 ± 0.031	1.154 ± 0.114	1.015 ± 0.047
EO	0.974 ± 0.029	1.081 ± 0.110	1.019 ± 0.043
Р	0.991 ± 0.023	0.974 ± 0.091	1.047 ± 0.036
PO	1.002 ± 0.021	0.901 ± 0.086	1.073 ± 0.033
D	0.982 ± 0.028	1.060 ± 0.111	1.002 ± 0.041
G	1.016 ± 0.018	0.931 ± 0.072	1.006 ± 0.029

Chapter 11

Systematic Errors and Consistency Checks

In this chapter, we consider the potential systematic effects that could change the results for $\alpha_{\bullet}^{i}/\alpha_{\bullet}^{all}$. In general, these effects break down into two classes: experimental errors and theoretical uncertainties. Experimental errors result from possible inaccuracies in detector modeling, extra tails on important detector resolution parameters that cannot be constrained by the data, and errors on the experimental measurements that function as the input parameters to the modeling of the underlying physics processes. The last category includes such things as errors on the measured values of the total charged multiplicity in B meson decays, the average b hadron lifetime, etc.. The MC simulations are tuned to reproduce the measured distributions of these quantities, so the possible variation in the experimental results must be considered. Theoretical uncertainties include such things as the possible range of heavy quark masses, the renormalization scale μ , and effects due to different models of hadronization. Each of these and their effect on the final results will be discussed in turn in the sections below. We will also consider a number of cross checks that we can perform with the data to guarantee that the results are as insensitive as possible to the specific choice of analysis method.

11.1 Experimental Errors

The experimental systematic errors can be separated into three categories: those due to uncertainties in the modeling of *b* hadron decays, those due to uncertainties in modeling c hadron decays, and those due to imperfect detector simulation. In general, all of these effects could change the elements of the tagging efficiency matrices ε . The influence of the various uncertainties on the final results is evaluated in each CaSe by varying the appropriate parameter in the MC, recalculating the matrices ε , performing a new fit to Equation 9.1, and redcriving $\alpha_s^j/\alpha_s^{all}$ based on the same data sample.

11.1.1 Event Re-weighting

Generating a new set of MC events to evaluate the effects of the variation in each of the parameters governing the production and decay of c and b hadrons would be impossible. Instead, we use an event re-weighting scheme to produce the correct distributions in the MC. This works as follows: given a normalized* distribution F(x) from the default MC and the desired normalized distribution G(x), one can assign a weight to each event of G(x)/F(x), where the weight is either calculated analytically using the two distributions (as is the case in the re-weighting of the lifetimes) or is taken from a look-up table for a given bin of x. The totals weights of the events are then used to calculate the correct ε matrices for this set of parameters. All of the re-weighted distributions were checked at the parton-level to insure that they had mean weight of unity and that the normalizations are correct.

11.1.2 Errors due to b-Hadron Modeling

As the systematic errors due to the variation of each of the parameters are approximately the same size for each of the jet-finding algorithms, we have chosen to merely average the results and present a single value for each error in the tables. The errors on the ratios $\alpha_s^j/\alpha_s^{all}$ due to the effects discussed below appear in Table 11.1. The are listed here in decreasing order of the effect on $\alpha_s^b/\alpha_s^{all}$.

- b Hadron Decay Multiplicity: The mean charged multiplicity in each bhadron decay is varied by ±0.20 tracks, assuming the b hadron decay in each event hemisphere is uncorrelated. The value for the average B meson decay multiplicity comes from the Y(4s) measurements of Argus[195], who obtained (n_ch) = 5.39f0.15 tracks. Our variation corresponds to a slightly larger than 1σ change, where we have enlarged the error to allow for slightly different multiplicities in B_s and B baryon decays, which are as yet unmeasured. This should be the largest systematic error due to uncertainties in modeling b physics, as we are relying on the entire spectrum of the number of significant tracks to be properly described by the MC. Obviously, any change in the number of tracks in B meson decay must change the calculated tagging efficiency.
- . b Fragmentation: The mean energy fraction $\langle x_E \rangle$ received by weakly-decaying b hadrons is varied by $\langle x_E \rangle = 0.695 \pm 0.011$, a value consistent with measurements from LEP[196]. This is done by changing the ϵ parameter in the Petersen fragmentation function. The tag should also be fairly sensitive to the distribution of quark energies, as we have already shown that the **bias** for tagging P-jet bevents is larger than for the other flavors. This is easily understood, as softer bs produce decay tracks with smaller total momentum and hence larger extrapolation errors due to multiple scattering, which makes the track less likely to miss the IP by a significant amount. Since the b has less boost, it also doesn't travel quite as far before decaying, which also makes that hemisphere harder to tag. Some care must be taken in the re-weighting procedure that is applied in order to calculate the effects of changing the fragmentation. The energies of the bhadrons in events with hard gluon radiation are typically lower than the mean energy of the distribution and the energies are reduced in a correlated manner. So, any re-weighting scheme that attempts to weigh each event hemisphere separately arrives at an improper weight for these events. The actual effects of hemisphere weighting in this case are even more serious, as 2-jet events and 3-jet events are shifted in opposite directions when the fragmentation function is changed. This is clearly not the desired result of the re-weighting scheme! To

^{*} $\int F(x)dx = 1.$

remedy the situation, we created a routine to calculate an event weight based on the values of the two most energetic weakly-decaying *b* hadrons in the event. Large MC samples of *b* events were generated with different Petersen ϵ values,' and the ratios of the 2-dimensional distributions of the two hadron energies were formed to calculate the weights. Figure 11.1 shows the set of weights calculated in this manner for shifting $\langle x_E \rangle$ up or down, as well as the default distribution of hadron energies.

- *b Hadron Lifetimes:* As the case of tagging *b* events is made possible by the long lifetimes of the *b* hadron species, any change in the lifetimes will have an effect on the tagging efficiency. The values of the *B* meson (baryon) lifetimes in the MC sample are 1.55 ps (1.10 ps). The *B* meson lifetime was varied by fO.1 ps, and the *B* baryon lifetime by f0.3 ps. Given the small errors on the average *b* hadron lifetime[194], these errors may seem a bit generous. However, given the past history of this quantity[199], it is best to be conservative in this particular case.
- b baryon Production Rate: The rate of b baryon production in b events was changed from its central value[197] of 9% by $\pm 3\%$. A change in the b baryon production rate could effect the b-tagging efficiency due to the low lifetime measured for the baryons.
- $R_b = BR(Z^0 \rightarrow b\bar{b})/BR(Z^0 \rightarrow hadrons)$: Without the standard model hadronic branching fractions, this analysis could not have begun. The fraction of b events in Z^0 decays gives the normalization for the unfolding procedure, so any change in that value has a direct effect on the values for the other quark flavors, due to the correlations in the unfolding. The high purity of the b-tag sample largely reduces the effects of these correlations, however. We varied R_b by ± 0.004 about the standard model prediction of $R_b = 0.220$. This range of error is consistent with the current world average[194].
- $B \to D^+$ Fraction: Changes in the fraction of $B \to D^+$ decays could have a large effect on the tagging efficiency for b events, as the long decay path of the



Figure 11.1: (a) The distribution of x_1 vs. x_2 for weakly decaying *b* hadrons, showing the pronounced peaking at higher **x** values. Figure (b) shows the computed event weights used to simulate a *b* fragmentation function of Peterson $\epsilon = .0.003$. The relative size of the boxes in each cell represents the relative magnitude of the event weights across the weight distribution. Note the weights have the expected distribution, emphasizing those events where both quarks have large x values. Figure (c) shows the weights to simulate a *b* fragmentation function of $\epsilon = .0.011$.

Table 11.1: Systematic Errors Due to b Hadron Modeling

Source of Error:	Central Value	Variation	$\Delta\left(\frac{\alpha_s^{uds}}{\alpha_s^{all}}\right)$	$\Delta\left(\frac{\alpha_{e}^{c}}{\alpha_{e}^{all}}\right)$	$\Delta\left(\frac{\alpha_{\bullet}^{b}}{\alpha_{\bullet}^{all}}\right)$
B decay multiplicity	$\langle n_{ch} \rangle = 5.39$	f0.20 tks.	0.005	0.046	0.023
<i>B</i> fragmentation	$\langle x_b \rangle = 0.700$	± 0.011	0.003	0.010	0.016
B meson lifetime	$\tau_B = 1.55 \text{ ps}$	± 0.1 ps	0.005	0.030	0.012
B baryon lifetime	$\tau_B = 1.10 \text{ ps}$	±0.3 ps	0.002	0.012	0.004
B baryon prod. rate	$f_{\Lambda_b} = 9\%$	$\pm 3\%$	0.009	0.013	0.007
R_b (bottom fraction)	0.220	f0.004	0.002	0.003	0.006
$B \rightarrow D^+ + X$ fraction	0.17	f0.07	0.002	0.005	0.001

 D^+ downstream of the *B* meson decay point enhances the efficiency for tagging that event. The uncertainty of this branching fraction is[194] ±4% currently; we assign a ±6% error to account for the uncertainties in the production rates of all of the charmed hadrons in *B* decays.

11.1.3 Errors due to c-Hadron Modeling

The errors on the ratios $\alpha_s^j/\alpha_s^{all}$ due to the effects discussed below appear in Table 11.2. The are listed here in decreasing order of the effect on $\alpha_s^c/\alpha_s^{all}$.

. $R_c = BR(Z^0 \rightarrow c\bar{c})/BR(Z^0 \rightarrow hadrons)$: As the c-tag is of relatively low purity and contaminated by both *uds* and *b* events, the fraction of c events assumed for that sample has a large effect on the outcome for $\alpha_s^c/\alpha_s^{all}$. This is due to the large background subtraction that occurs to derive the J-jet rate for charm; the effect is shown graphically in Figure 11.2, where the value for R_c input to the analysis is varied, and the results for $\alpha_s^j/\alpha_s^{all}$ for all three flavors are plotted. As the amount of charm in the sample is assumed to be less and less, the 3-jet rates for *uds* and *b* adjust to maintain the total S-jet rate in the charm tag sample, though the χ^2 of the fit (not shown) does get worse. We varied R_c by ±0.016 about the standard model value of 0.170. The variation is sightly



Figure 11.2: The results for $\alpha_s^j / \alpha_s^{all}$ when the input value of R_c is varied. The excess branching fraction is assumed in this case to come from *uds* events.

- larger than the error on the world-average[194] for $R_{,.}$ The central value of the world average is approximately 1.50 low at this point, at $R_c = 0.156$. Since the measurements are still consistent with the standard model, we choose to use that central value.
- . c fragmentation The mean energy fraction $\langle x_E \rangle$ received by weakly-decaying c hadrons was varied (in the same manner as the *b* fragmentation case, above) by $\langle x_E \rangle = 0.501 \pm 0.012$, a value consistent with measurements from LEP[198].
- $c\bar{c} \rightarrow D^+$ Fraction: Once again, due to the long D^+ lifetime, we would expect that the fraction of charm decays containing significant tracks depends on the rate of D^+ production in $c\bar{c}$ events. We assign a $\pm 5\%$ error to account for the uncertainties in the production rates of all of the charmed hadrons.

Table 11.2: Systematic Errors Due to c Hadron Modeling

Source of Error:	Central Value	Variation	$\Delta\left(\frac{\alpha_s^{uds}}{\alpha_s^{all}}\right)$	$\Delta\left(\frac{\alpha_{s}^{c}}{\alpha_{s}^{all}}\right)$	$\Delta\left(\frac{\alpha_s^b}{\alpha_s^{all}}\right)$
$\begin{array}{l} R_c \mbox{ (charm fraction)} \\ c \mbox{ fragmentation} \\ c\overline{c} \rightarrow D^+ + X \mbox{ fraction} \\ c \mbox{ decay multiplicity} \end{array}$	0.170 $\langle x_c \rangle = 0.494$ 0.20 $\langle n_{ch} \rangle \sim 2.34$	$ \begin{array}{c} f0.017 \\ f0.012 \\ f0.04 \\ \pm 0.20 \text{ tks.} \end{array} $	0.016 0.004 0.003 0.003	-0.055 0.016 0.01 0.008	0.008 0.003 0.002 0.006

• The Decay Multiplicity: The mean decay multiplicities of the D^0 , D^+ , D_s , and Λ_c are varied by ± 0.06 , ± 0.10 , ± 0.31 , ± 0.40 , respectively. The measurements of these decay multiplicities are from MARK-III[200]. The multiplicities for D^0 and D_s in our MC agree exactly with the measured values, so the variations are just the experimental errors. For the D^+ , our MC multiplicity distribution is deficient in 1-prong decay modes when compared with MARK-III's measurement. So, the variation for this decay mode is the difference in the mean decay multiplicity between our MC value and the measurement. The variation for A, is estimated from varying the parameters that can effect charm baryon production in the MC.

11.1.4 Other Experimental Errors

The errors on the ratios $\alpha_s^j/\alpha_s^{all}$ due to the effects discussed below appear in Table 11.4.

• Limited MC Statistics: Since the ε efficiency matrices are evaluated using the MC, each matrix element has a statistical error due to the finite number of MC events available. We estimated the effects of this uncertainty in the tags by generating the multinomial distribution for the elements of each column of each ε matrix based on the MC statistics used to create the matrix. This was repeated 1000 times for each algorithm, with results for $\alpha_s^j/\alpha_s^{all}$ derived each time using the numbers of tagged events in the data. The standard deviations

of the $\alpha_{ss}^{j}/\alpha_{ss}^{all}$ distributions were taken as the error due to finite statistics. This procedure results in errors that are approximately what one expects based on the total number of MC and data events.

- *Tracking Efficiency:* Any variation of the number of tracks available for use in tagging will have some effect on the tag efficiency. As discussed in detail in Section 7.1, the average multiplicity for tagging tracks differs between MC and data. We have applied a correction to the MC to remove this discrepancy, and have assigned a conservative error of f0.3 tracks on the correction. The variation in the final results due to changing this correction are relatively small.
- Tails in the IP Position Distribution: If the position of the IP is incorrectly determined, then it is possible that a large number of tracks in any given event could miss the IP by a significant amount and cause the event to be included in the *b* tag sample. Since we cannot rule out the presence of a 100 μ m tail at the level of 0.25%, we add this amount of tail to the MC events to simulate this possible effect.
- *Tracking Resolution* We discussed in Chapter 5 the "unsmearing" correction which was applied to tracks in the MC to remove errors in simulating the CDC positions. We take here the full correction as a systematic error and perform the analysis without the unsmearing applied. This turns out to be a small effect. The effects of changing the effective tracking resolution by requiring two VXD hits on each track were also minimal, leading us to believe that the tracking is well modeled by the MC.
- "Vee"-Finding We can place an upper limit on the possible effects of a variation in the frequency of $s\bar{s}$ production in hadronic showers by allowing K^0 and Λ decays into our sample of charged tracks. Even though the efficiency for reconstructing their decays becomes small in the region close to the IP where they could be more easily mistaken for a heavy quark decay, the fraction of tracks from these sources that remains in the analysis when they are not excluded changes by far more than the 10% variation one can estimate for the

uncertainty on ss production. The overall effect is also small.

Variation of Event and Track Selection Cuts: Changing the event and track selection cuts checks the validity of the MC simulation of the detector response. and, to a lesser extent, the physics included in the c and b hadron decay models. As we have already considered the systematic errors due to changing the physics parameters, we will try not to double-count sources of error, choosing for consideration cuts that should be relatively independent of those parameters[†]. The effect of changing these cuts is to modify the total efficiency for an event to pass the event selection cuts; since the efficiencies for events of each quark flavor to pass the cuts are virtually identical, we would expect these effects to be small. Due to the small size of these effects relative to the statistical errors (and fluctuations) on the measured values, we varied each of the cuts in steps away from the central value and used the slope derived from these points to calculate the size of the error. The derivatives of the ratios $\alpha_s^j/\alpha_s^{all}$ are given in Table 11.3. Due to the small number of events that are affected by the analysis changes, all of these slopes differ from zero by less than 0. 1σ ; the effects of these variations on the analysis will be ignored.

Thrust Axis Containment: the cut on $|\cos \theta_{Thrust}|$ was varied over $0.67 \leq Cut(|\cos \theta_{Thrust}|) \leq 0.75$, a range consistent with our estimates of a 2" smearing in the thrust axis direction due to shower fluctuations in the calorimeter response.

Visible Energy: the minimum visible energy required for an event to be selected was varied by 16 GeV $\leq Cut(E_{vis}) \leq$ 20 GeV. The effect of this variation is to subtly modify the probabilities for a parton-level event to switch assignments between a 2-jet event and a 3-jet event and to change the probability of a parton-level 3-jet event to pass the selection cuts.

Table 11.3: The derivatives of $r^j = \alpha_s^j / \alpha_s^{all}$ with respect to the varied analysis parameter 2.

$ \begin{array}{ c c c c c c c c c c c c c c c c c c c$	Source:	Central Value	Variation	$\partial r^{uds}/\partial x$	$\partial r^c/\partial x$	$\partial r^b / \partial x$
	$E_{vis} \\ \cos \theta_T \\ n_{trk} \\ \cos \theta_{trk} \\ p_{tag} \\ p_\perp$	$\begin{array}{l} E_{vis} > 18 \; \mathrm{GeV} \\ \cos \theta_T < 0.71 \\ n_{trk} \geq 7 \\ \cos \theta_{trk} < 0.8 \\ p_{tag} \geq 0.5 \; \mathrm{MeV} \\ p_{\perp} \geq 0.2 \; \mathrm{MeV} \end{array}$	$\pm 2 \text{ GeV} \\ \pm 0.04 \\ \pm 1 \\ \pm 0.04 \\ \pm 0.1 \\ \text{fO.1}$	<.01/GeV 0.805 0.003 -0.310 -0.152/GeV 0.074/GeV	<.01/GeV -3.112 0.031 0.248 0.083/GeV -0.388/GeV	<.01/GeV 0.987 0.017 1.185 0.265/GeV 0.065/GeV

Charged Track Multiplicity: the number of charged tracks required for an event to **pass** the selection cuts was varied between 6 and 8 tracks to check the effects of allowing some variation in the amount of background in the sample.

Charged Track Definition: The definition of the allowed track parameters for a track to be included in the total charged multiplicity of an event was varied within reasonable bounds to study the effects of additional soft tracks on jet definitions, total visible energy, and the backgrounds in the sample. The allowed p_{\perp} was varied between 100 MeV/c and 300 MeV/c. The allowed track $|\cos \theta|$ was varied between 0.76 and 0.84. Variation of the allowed track $doca_{zy}$ and $doca_{rz}$ produced negligible changes.

• Tagging *Track* Momentum *Cut:* The possible variations in the ε matrices due to the uncertainties in modeling the momentum spectra of the *band* c hadron decay products was not explicitly considered above in the physics modeling errors. We can estimate the effect on this analysis by varying the allowed momentum of the tagging tracks used in the analysis. Since the minimum momentum cut is already set to the relatively high value of 0.5 GeV/c, we should be insensitive to the low-momentum region where most of the uncertainty lies. A variation of $\pm 100 \text{ MeV}/c$ in this cut gives the expected small effects on the final result.

[†]Variation of cuts or tag definitions that are directly affected by the physics model can be used to cross-check our estimates of the inherent modeling errors. These checks will be discussed in Section 11.3.

Table	11.4:	Systematic	Errors	Due	to	Detector	Modeling

Source of Error:	Variation	$\Delta\left(\frac{\alpha_s^{uds}}{\alpha_s^{all}}\right)$	$\Delta\left(\frac{\alpha_s^c}{\alpha_s^{all}}\right)$	$\Delta\left(\frac{\alpha_s^b}{\alpha_s^{all}}\right)$
M.C. Statistics tracking efficiency IP Position Tails tracking resolution Vee finding	± 0.3 tracks 0.25% 100 μ m no unsmear none	$\begin{array}{c} 0.011 \\ 0.003 \\ \thicksim 0.001 \\ < 0.001 \\ < 0.001 \end{array}$	$\begin{array}{c} 0.048 \\ 0.031 \\ \thicksim 0.005 \\ 0.017 \\ 0.013 \end{array}$	$\begin{array}{c} 0.014 \\ 0.017 \\ \thicksim 0.005 \\ 0.014 \\ 0.007 \end{array}$

11.2 Theoretical Uncertainties

Uncertainties in the understanding of the exact physical processes underlying the production and decay of hadrons and the emission of hard gluon radiation can lead to varying predictions for the theoretical inputs to our models of hadronic interactions. Some physical parameters, such as the quark masses and the rate for gluon splitting to heavy quark pairs, have only weak experimental constraints. The sheer difficulty of calculation in perturbative QCD and the size of α_s result in uncertainties in the current calculations due to the effect of as-yet-uncalculated higher-order processes. The MC simulations that reproduce the observed data have tunable parameters that are not directly observable experimentally and are thus allowed to vary over some range. All of these elements contribute to the theoretical uncertainties associated with this analysis. Again, we present in Table 11.5 for each of the quark flavors an average of the systematic errors for all of the jet-finding algorithms.

• Hadronization Modeling We checked the variation of the parton-level 3-jet rates as functions of the various parameters within JETSET 6.3 that control the string fragmentation process. Their effects on the final results were then evaluated by changing the 3-jet rates in the MC by the variation seen in changing the parameters. Since the ratio $\alpha_s^j/\alpha_s^{all}$ is not sensitive to the absolute 3-jet rate in the MC, the effects of these variations are negligible. The parameters studied included

 Q_0 (PARE(22)), which controls the parton shower cut-off ($Q_0 = 1.0^{+1.0}_{-0.5} \text{ GeV}$)

 σ_q (PAR(12)), which varies the width of the gaussian distribution of transverse momentum generated in the parton shower relative to the original quark direction ($\sigma_q = 0.39 \pm 0.04$ GeV)

a and b (PAR(31) and PAR(32)), the coefficients in the symmetric Lund fragmentation function (these were checked at the hadron-level)

- The rate for $g \rightarrow QQ$: A large rate for gluons to split into heavy quarks could change the frequency with which *uds* events are tagged by the heavy flavor tags. As there are very few experimental constraints on this rate, we have estimated the effect by varying the rate for $g \rightarrow QQ$ by $\pm 50\%$ from the JETSET default value. This has a relatively small effect on the results.
- The Uncertainty OR α_s^{all} : As mentioned in Section 10.2, the results for the α_s ratios depend on the value of α_s^{all} used when converting from jet rates to α_s using the $\mathcal{O}(\alpha_s^2)$ QCD calculation. As the experimental values for α_s^{all} that were used also have errors associated with them, their variation must also be considered as a source of uncertainty. This is included **as** a theoretical uncertainty because the dominant error on the values for α_s^{all} is due to the inability to fix the value of the renormalization scale $f = \mu^2/s$, which arises from a cut-off put into the theoretical calculations to avoid the need to calculated higher-order processes. Thus, to estimate the effects of this error correctly, we need to vary the renormalization scale and α_s^{all} simultaneously in the expression for R_3 (from Section 8.3):

$$R_{3}(y_{cut}) \equiv \frac{\alpha_{s}(\mu)}{2\pi} A(y_{cut}) \equiv (\frac{\alpha_{s}(\mu)}{2\pi})^{2} [B(y_{cut}) \equiv A(y_{cut}) 2\pi b_{0} \log f] = (11.1)$$

This generates up to a 1% error on $\alpha_s^j/\alpha_s^{all}$, which is far smaller than the 10% error this effect has on the absolute determination of α_s^{all} from the D_2 distribution[131]. The variation of each of the values of α_s^{all} for the individual jet algorithms (see Table 10.2) was used to estimate the error.

• Uncertainties in the R_4 Calculation: As discussed in Section 8.3, the calculated 4-jet rate does not match that observed in the data by approximately a factor

of two for all algorithms except for G. This is presumably due to uncalculated coefficients of $\mathcal{O}(\alpha_s^3)$ and higher. We take the full 100% variation in the value of the coefficient $C(y_{cut})$ as the systematic error due to this source of uncertainty. It has an almost negligibly small effect on the final results.

- Variation of the Heavy Quark Masses: The correction applied in the previous chapter to account for the smaller phase-space for hard gluon emission in heavy quark events relies on knowing the values of the *b* and c quark masses. As this correction changes the results for R_3^b/R_3^{all} by as much as 5%, it is important to consider possible variations in its magnitude. The quark masses used in calculating the central values of the corrections were 4.5 GeV/ c^2 (1.5 GeV/ c^2) for the *b* (c) quark mass, and we considered variations of $\pm 0.25 \text{GeV}/c^2$ about these values to estimate the error on the parton-level quark masses used in the calculation.
- Uncalculated Higher-Orders in the Phase-Space Suppression Calculation: The calculation upon which our phase-space correction is based[193] has only been performed to leading order in α_s . We have as yet made no estimate of the possible size of the next-to-leading terms; they should be of order α_s/π times the first order correction, which would make them small.
- Jet-Algorithmic Dependence of Results: All of the results quoted in this chapter have been the averages of the shifts in the values for $\alpha_s^j/\alpha_s^{all}$ due to the variations of the different parameters which can affect the results. We have chosen to take a simple mean of these results, as there is no reason a *priori* to prefer one algorithm over any other. However, there is significant scatter in the final values, which is possibly not all statistical. Since all of the jet algorithms are being used to perform the same measurements on the same data set, and the results should be independent of jet-algorithm to first order, we attribute the scatter to uncalculated higher-order terms that contribute to the 3-jet rate. To estimate the size of this effect, we calculate the r.m.s. of the results for each flavor relative to the mean value, and quote this as the theoretical uncertainty associated with

Source of Error:	Central Value	Variation	A (\$	$) \Delta \left(\frac{d}{d}\right)$	$\frac{\alpha_{q}^{c}}{\alpha_{s}^{a} l l} \Delta \left(\frac{\alpha_{q}^{b}}{\alpha_{s}^{a} l l} \right)$
Hadronization	Q_0 , etc.		0.002	0.007	< 0.001
$g \rightarrow QQ$ rate		$\pm 50\%$	0.001	0.002	0.002
α_s^{all}			0.003	0.009	0.010
$\Delta C(y_c)$			0.001	0.003	0.003
Δm_b	$4.5 \; { m GeV}/c^2$	$\pm .25 GeV/c^2$	0.001	0.002	0.004
Δm_c	$1.5 { m GeV}/c^2$	$\pm .25 GeV/c^2$	< 0.001	0.002	0.001
Jet Algorithm		_	0.021	0.096	0.027

Table 11.5: Theoretical Uncertainties

Table 11.6: Summary Table of Systematic Errors

Source of Error:	$\Delta\left(\frac{\alpha_s^{uds}}{\alpha_s^{all}}\right)$	$\Delta\left(\frac{\alpha_s^c}{\alpha_s^{all}}\right)$	$\Delta\left(\frac{\alpha_s^b}{\alpha_s^{all}}\right)$
Experimental Errors: B Physics C Physics Detector Modeling MC Statistics	0.0084 0.017 0.003 0.011	0.060 0.060 0.032 0.048	0.033 0.011 0.017 0.014
QCD Theoretical Uncertainty Jet Algorithm	0.003 0.021	0.011 0.097	0.012 0.028

the QCD calculations describing the jet-finding algorithms. This is by far the largest theory error.

All of the systematic errors are summarized in Table 11.6, where the total contribution of each category to the error is listed.

11.3 Checks on the Validity of the Results

In this section we present two cross-checks that give us confidence that the analysis method we have used and our estimates of the associated uncertainties are robust.

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11.3.1 Variation of the Chosen y_{cut}

To insure that the dependence of our results on the choice of y_{cut} was minimal, we repeated the entire analysis at different y_{cut} values for all of the jet algorithms. The average values are shown in Figures 11.3 and 11.4. Note that all of the points are highly correlated. No significant deviation from our results is observed.

11.3.2 Variation of the Tag Definition

We defined as a significant track one that misses the IP by more than 3σ . We can change this requirement and observe the variation in our results. This serves to verify that our estimates of the physics and tracking modeling errors are correct, since differences between the true parameters and our MC simulation could manifest themselves in a dependence of the result on what is chosen for the significance cut. Figure 11.5 shows the results for $\alpha_s^j / \alpha_s^{all}$ as the significant cut is varied. There seems to be a systematic slope in all of the results. It is interesting to consider this result for the Veto period compared with the Pre- and Non-Veto periods, which is shown in Figure 11.6. It can be seen that virtually all of the variation with respect to the significance cut occurs in the Veto period. As the only substantial difference in the MC simulation between the Veto period and the others is in the effects of the track trigger, we can only surmise that improper simulation of this effect is the root cause of this dependence. To check this, we analyzed the Pre-Veto period data with the Veto period MC without applying the trigger veto to the MC, and compared these results with the same analysis performed with the correct MC and data match. The results are consistent, which implies that nothing is wrong with the MC simulation excluding the simulation of the trigger veto. We can check the validity of the veto simulation in two ways. First, we can measure the branching fraction for $Z^0 \rightarrow b\bar{b}$ in the Preveto and Veto periods separately, and see if they compare. Using the MC veto simulation to correct for the fact that b events during the Veto period are less likely to pass the trigger[‡], we obtain $R_b^{PV} = 0.233 \pm 0.007$ and $R_b^V = 0.226 \pm 0.006$, where the errors are statistical only. The first value is in excellent agreement with

[‡]See Table 6.2.



Figure 11.3: The Results $vs. y_{cut}$ Choice for E, EO, and P



Figure 11.4: The Results $vs. y_{cut}$ Choice for P0, D, and G



Figure 11.5: The results of the analysis obtained for different cuts on the track significance used to form the n_{sig} distribution used as the flavor tag.

that of the official SLD R_b measurement[182] of $R_b = 0.231 \pm 0.004$ using the same technique and the Preveto and Nonveto data samples. Since the values for the two run periods are within 1σ , we cannot measure any difference introduced by the veto simulation using this technique. Another method is to vary the rate of the veto itself. We considered the Veto period with a tag requirement of $b_{norm} > 4.0$, as this is where the maximum deviation occurs in Figure 11.6. No shift in the values of the three $\alpha_s^2/\alpha_s^a ll$ ratios is seen when the veto rate is changed by $\pm 10\%$. We conclude that the small slope seen on this plot is due to the statistics of the impact parameter tag, and not to some systematic problem with the veto simulation.

11.3.3 Sensitivity to Anomalous Flavor-Dependent Strong Couplings

To investigate the sensitivity of this analysis method, we used a subsample of the MC as a mock data sample and removed different fractions of 3-jet b or uds and c events so as to vary R_3^b . This mock data sample was then subjected to the standard

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Figure 11.6: The results of the analysis obtained for different cuts on the track significance used to form the n_{sig} distribution used as the flavor tag. The three different run periods are shown for comparison.



Figure 11.7: Sensitivity to Input Anomolous Strong Coupling

analysis. The results are shown in Figure 11.7. One can see that if there were large flavor-dependent differences in the strong coupling we would have observed them in this analysis.

We now move on to the presentation of the final results, and some comments and conclusions about this analysis.

Chapter 12

The Flavor Independence of α_s : Conclusions

12.1 Final Results and Discussion

The final results for all of the jet algorithms are shown in Figure 12.1. For the average values shown on this figure, the r.m.s. of the results over all algorithms is added in quadrature to the other theoretical uncertainties, as was discussed in Section 11.2. The average values, including this additional error, are:

$$\frac{\alpha_s^{uds}}{\alpha_s^{all}} = 0.987 \pm 0.027(stat) \pm 0.022(syst) \pm 0.022(theory)$$

$$\frac{\alpha_s^c}{\alpha_s^{all}} = 1.012 \pm 0.104(stat) \pm 0.102(syst) \pm 0.096(theory) \quad (12.1)$$

$$\frac{\alpha_s^b}{\alpha_s^{all}} = 1.026 \pm 0.041(stat) \pm 0.041(syst) \pm 0.030(theory)$$

As a first point of discussion, it is clear from the above results that α_s is independent of quark flavor within our present experimental sensitivity. Our novel means of separating the different quark flavors inclusively by differences in the lifetime content of the charged tracks has proven to be free of large systematic biases while allowing us to make use of the full statistics of our data sample. In particular, the measurement



Figure 12.1: Values of $\alpha_s^j/\alpha_s^{all}$ derived for each of the jet algorithms used in the analysis for each of the quark flavors (see text). The error bars on the average values include the statistical and systematic errors and the total theoretical uncertainty.

of ratios has minimized the effects of the renormalization scale uncertainty and the other theoretical effects that plague an absolute determination of α_{g} . We are still sensitive to effects relating to the different jet algorithms. The high purity of the

uds and *b* tags results in smaller total errors for these two results due to reduced sensitivity to fluctuations in the backgrounds which are being subtracted in the unfolding procedure. The result for charm, on the other hand, suffers from having a signal-to-noise ratio of 0.5; any small change in the background composition leads to large excursions of the extracted value for $\alpha_s^c/\alpha_s^{all}$. This is the source of the large systematic and theoretical errors on the charm result. Possible improvements on this method will be discussed below.

It is interesting to note the large theoretical uncertainty, which is dominated by the different values obtained for the α_s ratios when using the six jet algorithms. As mentioned in the previous chapter, this is most likely due to uncalculated higher-order terms in the theoretical predictions for the jet rates. This *large* variation of the results was expected from previous work[131] on measurements of α_s , but this is the first time a systematic study of this effect has been done checking the behaviour of the jet algorithms for quarks of flavors (masses). The possible variation of different observables in measuring the ratios $\alpha_s^j/\alpha_s^{all}$ and the associated theoretical uncertainties have been largely ignored in all of the other experimental results on flavor independence.

To compare our limits on flavor independence with those obtained by other experimental methods, we show in Figure 12.2 the above results along with those from the LEP experiments * discussed in Chapter 3. As the only other analysis which has attempted to measure the strong coupling to each quark flavor independently is the one from OPAL[145], it will be the focus of much of the comparative discussion. It should be mentioned that the errors on the measured ratios are necessarily enlarged by loosening the assumption that all of the. quark flavors have the same strong coupling except the one under study; both our results and the ones from the OPAL study share this consequence of a more general analysis.

The advantages of our method over those used by OPAL for the light quark tags is immediately obvious, as our results have smaller errors even though the analysis is based on a factor of 20 less data. This is mostly due to the use of the precision tracking and our knowledge of the stable IP position as a method of excluding heavy quark events from the sample of selected events, rather than selecting light quark events by



Figure 12.2: Values of $\alpha_s^j / \alpha_s^{all}$ measured in this thesis compared with other experimental results. References for other values are provided in Chapter 3. For a discussion of results, see text.

^{*}The TASSO results have been superseded in precision and will not be considered here.

some kinematic method from a sample with large, relatively unknown backgrounds. Even with the somewhat larger errors on the result for charm, our errors are more or less comparable to those obtained by OPAL due to the inefficiency of their charm tag (reconstructing D mesons) and its larger systematic bias. For b quarks, larger statistics and the smaller systematic errors associated with a high- p_{\perp} lepton tag produce a more precise result from OPAL.

In contrast, the advantages of obtaining a high statistics high purity sample and using that as a testing ground arc exemplified in the newest result from OPAL[146] for $\alpha_s^b/\alpha_s^{all}$. Although they fail to consider properly the additional theoretical error associated with combining correlated measurements of the same quantities, their result is still by far the most precise test of flavor independence in the strong interaction.

12.2 Possible Improvements of the Method

The primary weaknesses of the method used in this analysis are the low purity of the charm sample and the reliance on the Monte Carlo simulation of heavy quark decays to represent the true decay properties. The latter introduces the largest systematic errors, while the former serves to magnify all of the errors associated with $\alpha_s^c / \alpha_s^{all}$. Systematic errors due to uncertainties in the heavy quark decay model can be drastically reduced by performing a "double tag" using both hemispheres in the event to actually measure the tagging efficiency for a given tag[†]. Given the stability of the SLC beam spot, a technique of this sort could also be used in tagging the light flavors instead of just *b* events. A more complicated version of the efficiency matrices used in the unfolding could be created that accounts for the flavor composition of doubly-tagged vs. mixed-tag events. Since the tag efficiency could then be measured for *uds* and *b* events, this should result in smaller errors on those results. The charm result would also benefit, since the allowed fluctuations in the charm sample backgrounds would be smaller.

The only way to effectively decouple the charm result from those for uds and b is

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to obtain a charm sample of high purity. The established methods of reconstructing D mesons could be employed, but, as mentioned above, these also have large associated systematic errors. The ideal solution would be If a more efficient way of tagging charm could be realized, and it has smaller kinematic bias than the present D reconstruction analyses. Work is in progress on such a technique, but it has not matured at this time. The future of tests of flavor independence could well be in the isolation of high-purity samples of different flavors rather than the inclusive approach taken in this analysis.

12.3 Conclusions

We have presented a test of the flavor independence of the strong interaction by isolating event samples enriched in quarks of specific flavors (uds, c, and b) and measuring the strong coupling in each sample using a jet rates analysis. This is the first test of flavor independence to use a precision microvertex detector to select event samples of all quark flavors based on the lifetime content of charged tracks. The strong coupling for each quark flavor was unfolded from the composition of quark flavors in each sample, resulting in

$$\frac{\alpha_s^{uas}}{\alpha_s^{all}} = 0.987 \pm 0.027(stat) \pm 0.022(syst) \pm 0.022(theory)$$

$$\frac{\alpha_s^c}{\alpha_s^{all}} = 1.012 \pm 0.104(stat) \pm 0.102(syst) \pm 0.096(theory) \qquad (12.2)$$

$$\frac{\alpha_s^b}{\alpha_s^{all}} = 1.026 \pm 0.041(stat) \pm 0.041(syst) \pm 0.030(theory) .$$

The effects of quark masses on the jet rates and the full $\mathcal{O}(\alpha_s^2)$ calculations for the jet rates have been employed in this measurement. The dominant experimental systematic errors arise from the uncertainties in the physics of heavy quark decay. The largest theoretical uncertainties are due to the variation in the results obtained with the six different jet algorithms used in the analysis, which are presumably caused by uncalculated higher-order terms in the predictions for jet production rates. These results are the most precise test of the *flavor* independence of the strong interaction for light quark events; the other measurements compare favorably with similar results obtained at LEP.

[†]This has been aptly demonstrated, *e.g.* by t ho ALEPII Collaboration in their measurement of $\Gamma(Z^0 \to b\bar{b})/\Gamma(Z^0 \to hadrons)$ [70].

Part II

A Measurement of the Parity Violation in the Z^0cc Coupling

Chapter 13

Analysis Method

This chapter contains a description of the reconstruction and selection techniques used to isolate hadronic events containing charm mesons' These techniques are based upon the reconstruction of D mesons, which are identified by the combined invariant mass of their decay products. At this point, we introduce the symbol *RCBG*, which will be used throughout this thesis to denote Random *Combinatoric BackGround* events or other quantities, such as asymmetries, related to or derived from these events[†].

Before we begin the discussion of the analysis techniques, we mention several points that are common to all of the charm selection analyses to avoid repeating them later:

. We are allowed to be less restrictive in the quality of tracks required in this analysis than in the inclusive tagging analysis. The invariant masses obtained when a **mismeasured** track is combined with other tracks is essentially random and thus does not contribute preferentially to background under any of our signal mass peaks. So, we use in this analysis all charged tracks defined by the criteria listed in Chapter 6^{\ddagger} namely, those tracks which have

^{*}It should be stated that this analysis has been a collaborative effort between myself and Steve Wagner, to whom the credit should go for many of the clever insights into charm physics that appear here.

[†]This was suggested by my wrists and hands as one way to stave off carpal tunnel syndrome until after this thesis is typed.

[†]The requirement on transverse momentum has been changed here to 150 MeV/c, rather than the 200 MeV/c used in the QCD analysis.

- transverse momentum relative to the beam axis $p_{\perp} > 150 \text{ MeV}/c$
- $-|\cos\theta| < 0.8$
- distance of closest approach in the xy plane $(doca_{xy}) < 5$ cm
- distance of closest approach in the rz plane $(doca_{rz}) < 10$ cm

with the one additional constraint that they be well-measured enough to have at least one VXD hit.

- The set of tracks combined to form the charm meson candidates is taken from one thrust hemisphere at a time. This is done to reduce *RCBG* from combinations of tracks from opposite sides of the event that certainly do not come from the same charm meson decay. It also speeds up the analysis.
- As will be discussed extensively in Section 14.2.1, the mixing-corrected asymmetry in events where a D^{*+} or D^+ is produced in b decays is almost identical to the standard model expectation for $A_{,.}$ Since we wish to measure only $A_{,.}$ we have to assume A_b is its standard model value, which, for a large contamination of b events in our signal, could result in large systematic errors due to the possible variation of A_b . So, to minimize the dependence of our results on the exact value of A_b , we will attempt to specifically remove those D^{*+} and D^+ mesons which result from the decays of b hadrons.

13.1 The Selection of D^0 Candidates

As mentioned in Chapter 2, the sample of charm events containing D^0 decays is selected using two independent procedures, a kinematic reconstruction analysis and a decay length analysis. The two selected samples are then combined to form the final sample used for the asymmetry measurement. Also discussed in the opening chapter was the physics of the decay $D^0 \rightarrow K^-\pi^+\pi^0$, specifically the peaking of the $K^-\pi^+$ mass due to the angular momentum constraints of the decay. Since the reconstructed $K^-\pi^+$ mass in both the $D^0 \rightarrow K^-\pi^+\pi^0$ and $D^0 \rightarrow K^-\pi^+\pi^0$ decays has a relatively narrow peak, the two selection analyses for these decay modes are identical, except for the central value of the mass region allowed for the D^0 candidates in each case. The two selection techniques are presented below.

13.1.1 The Kinematic Technique

This is the "standard" $D^{\star +} \to \pi_s^+ D^0$ reconstruction technique as mentioned in Chap ter 2. It relies solely on the kinematic properties of $D^{\star +}$ and D^0 decays for background rejection.

To begin, all pairs of oppositely charged tracks are combined to form a D^0 candidate by assigning the K^- mass to one of the particles and the π^+ mass to the other.



Figure 13.1: The invariant mass of $K\pi$ pairs after the cuts described in the text. The MC shapes from true $D^0 \rightarrow K\pi$ and $D^0 \rightarrow K\pi\pi^0$ decays are also shown for reference. The relative normalizations of signal and MC are arbitrary. Also shown are the definitions of the signal regions for the two decay modes.

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Figure 13.2: Definition of the helicity angle θ^* , which is the angle between the K direction and the D" flight direction in the D^0 center-of-mass frame.

We shall refer to the combined invariant mass as $m(K\pi)$. This provides the basic sample of candidate decays on which to base the rest of the selection procedure. The mass peak formed by a sample of signal events from the MC is shown in Figure 13.1, along with the data after the following cuts for the two D'' decay modes. The shape and width of the peaks suggest that the symmetric region of acceptance for candidate events should be 1.765 GeV/ $c^2 < m(K\pi) < 1.965$ GeV/ c^2 for $D^0 \rightarrow K^-\pi^+$ and 1.500 GeV/ $c^2 < m(K\pi) < 1.700$ GeV/ c^2 for $D^0 \rightarrow K^-\pi^+\pi^0$ decays.

Next, we impose a helicity cut on the decaying D^0 system which rejects backgrounds inconsistent with the decay of a spin-zero meson[201]. This proceeds as follows:

- . the lorentz boost of the D^0 candidate is calculated based on the total momentum of the K and π , and the $K\pi$ system is boosted back into the D^0 rest frame
- the angle θ^* , the opening angle between the K momentum and the D^0 direction in the D^0 rest frame, is calculated.

This situation is shown pictorially in Figure 13.2. For the isotropic decay of a spinzero particle into two spin-zero particles, one would expect the distribution of θ^* to be flat. Tracks from RCBG that happen to form a D^0 candidate with the appropriate mass tend to have their directions parallel to the boost direction in the D^0 rest frame, as they are unlikely to possess the large transverse momentum of the true D^0

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Figure 13.3: Helicity angle distribution for $D^0 \to K\pi$ decays (dashed line) and background (solid line). The arrows show the cut $|\cos \theta^*| < 0.9$ placed on this variable.

decay products. Thus, some fraction of the background events can be removed with little loss of signal by eliminating events where the decay products lie along the D^0 direction. We have chosen to place our cut at $|\cos \theta^*| < 0.9$. The distribution of θ^* for MC signal and background events is shown in Figure 13.3.

The D^0 candidate is then combined with a π_s^+ candidate track having momentum p > 1 GeV/c and charge opposite to the K^- candidate. Requiring a momentum greater than 1 GeV/c removes some of the contamination due to slow pions from b decay, as they have a softer momentum spectrum. A comparison of the π_s^+ momentum spectra for D^{*+} decays in b and c events is shown in Figure 13.4. The $\pi_s^+D^0$ combination formed here will be used to create the Am (Am $\equiv m(D^{*+}) - m(D^0)$) distribution which will be used as the "signal" for the D^0 analysis.

To further suppress RCBG and D^{*+} mesons from b events, the energy E_{D^*} of the D^{*+} is calculated, and any D^{*+} candidate with $x_{D^*} = 2E_{D^*}/E_{CM} < 0.4$ is rejected. This cut takes advantage of the harder charm fragmentation function for



Figure 13.4: Momentum of the π_s^+ from $D^{\star+}$ decays in **b** and c events compared with the momentum spectrum of all tracks.

Table 13.1: The efficiencies for D^0 decays in c and b events to pass the kinematic selection cuts described in the text.

	$D \rightarrow K\pi$		$D \to K \pi \pi^0$		
cut	$c { m eff}$	b eff	c eff	b eff	
D^0 mass	0.842	0.912	0.462	0.577	
θ^*	0.813	0.810	0.744	0.737	
$p_{\pi} > 1 \text{ GeV}/c$	0.769	0.458	0.667	0.338	
$x_D > 0.4$	0.796	0.492	0.559	0.259	
$\Delta(m)$	0.938	0.943	0.746	0.827	
Total	0.393	0.157	0.096	0.031	

 D^{*+} formation. A comparison of the x_{D^*} distributions for D^{*+} mesons produced in b and c decays is shown in Figure 13.5. The efficiencies for each of these cuts for selecting candidate events for the two D^0 decay modes are given in Table 13.1.

The Am distribution for this set of selection criteria is shown in Figure 13.6. for



Figure 13.5: Distributions of x_{D^*} for b and c decays.

the two decay modes with the expected signal shape overlayed on the data[§]. We will define the "signal" region of the Am distribution to be Am <0.150 GeV/ c^2 . For this selection, the signal region contains 65 events for the $D^0 \rightarrow K^-\pi^+$ mode and 100 events for the $D^0 \rightarrow K^-\pi^+\pi^0$ mode which fall within the D^0 mass acceptance windows for the two decays. We will also need to measure properties of the background from the data to reduce our dependence on the MC. We define a higher mass "sideband" region of 0.160 GeV/ c^2 < Am < 0.200 GeV/ c^2 for this purpose'.

13.1.2 The Decay Length Technique

In the complementary decay length analysis, we rely on the fact that $D^0 \pm c\bar{c}$ events have a long decay length ((L) ~ lmm) and are produced at the Z^0 Primary Vertex (PV) as opposed to being created in a *b* decay cascade. Since the decay length resolution is $\langle \sigma_L \rangle \sim 200 \mu$ m, a clean separation of events containing D^0 decays from

[§]The D^0 mass cut specified above has been applied to make this plot.

[¶]Technically, this should probably be called an "upper band" since there is no similar region on the other side of the peak in the Am distribution. For the ease of reference, we will continue to call this a generic "sideband".



Figure 13.6: The Am distribution for the kinematic selection analysis, with the MC pure signal shape overlayed on the data. The normalization of the MC shape is arbitrary and is meant to indicate the expected signal distribution. The signal and sideband regions, as well as the number of events they contain, are also shown.

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Figure 13.7: The χ^2 vertex probability for fake vertices formed only from tracks originating at the primary vertex. Note the peak near a probability of zero.

RCBG is possible, even at low x_D .

The analysis begins with the same pairing of all possible combinations of oppositely charged tracks to form D^0 candidates as in the kinematic analysis, above. The same restrictions are applied on the reconstructed invariant mass.

Next, a constrained vertex fit is performed^{||}; only the track combinations with χ^2 vertex probabilities greater than 1% are retained as D^0 candidates. This corresponds to a 1% loss of "true" vertices, which seems like a rather loose cut, except that the false combinations are peaked sharply at a probability of zero. An example of the vertex probability for combinatoric background is shown in Figure 13.7.

A decay length significance cut of $L/\sigma_L > 2.5$ is then applied. The values of σ_L are obtained from the fit for each of the vertices under consideration. This cut drastically reduces the number of background vertices that occur close to the PV due to the large apparent track overlap. This is shown pictorially in Figure 13.8, where the distributions of L/σ_L are shown for true vertices and those from combinatoric background.

Another cut using vertex information can be applied to remove a large fraction of

^{II}The fit is performed using the SLD fitting routine ZXFIT, which fits any number of tracks to a common vertex with the constraint that they pass through their mutual point of closest approach.



Figure 13.8: The normalized decay length for true secondary vertices and those fake vertices formed from fragmentation tracks due to the confusion near the IP.

the D^0 mesons produced in b decays. Since the D^0 's in charm events are produced at the primary event vertex, their flight path should intersect the PV. Since D's produced in the b-decay cascade acquire significant p_{\perp} relative to the b-hadron flight direction, they may not appear to originate from the PV. The different possible situations are shown in cartoon form in Figure 13.9. Given the precision with which the position of the PV is known within SLD^{**}, this is a powerful tool for selecting charm decays. The impact parameters in the s-y plane for the D^0 candidates from c, b, and RCBG events are overlayed with those measured in the data in Figures 13.10a and 13.10b. Notice that this cut even works for the only-partially-reconstructed $D^0 \rightarrow K^-\pi^+\pi^0$ decays; the hard charm fragmentation function insures that the charged tracks from the D^0 decay have sufficient energy to provide an accurate determination of the D^0 direction. A relatively tight cut on the 2-D impact parameter of the D" momentum vector to the PV of $d_{PV}^{xy} < 20 \ \mu m$ is then appiied to further reject RCBG and contamination from b events.

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Figure 13.9: Illustration of several of the vertexing situations that can occur when all tracks are used to form two-prong vertices. Many vertices are formed close to the IP as in (a) from random track overlap. The vertex momentum tends not to point directly back at the IP in this case. For a D decay in charm events, as in (b), the D is produced at the origin, and the vertex momentum will point back to the IP within measurement errors. The Ds produced in a b decay cascade, (c), acquire some transverse momentum from the b hadron decay, and thus do not necessarily point back at the IP.

Finally, we require only $x_{D^*} > 0.2$, and make no cuts on $\cos \theta^*$ or minimum π_s^+ momentum, since the charm purity of the sample remaining after the previous cuts is sufficiently high. Table 13.2 lists the efficiencies for charm events to pass each of the preceding cuts. Figure 13.11 shows the Am distributions for the decay length selection analysis. After applying the D^0 mass requirements, 49 events for the $D^0 \to K^-\pi^+$ mode and 73 events for the $D^0 \to K^-\pi^+\pi^0$ mode appear in the signal region.

Comparison of Figures 13.6 and 13.11 shows that the kinematic and decay length selection techniques yield D^0 samples with essentially the same signal to noise in both of the reconstructed decay modes. To maximize the available data, we choose to combine the results of these two analyses by taking the union of the two selected

^{**}See Chapter 5.



Figure 13.10: The impact parameters of D^0 mesons to the IP for the modes (a) $D \to K\pi$ and (b) $D \to K\pi\pi^0$. Shown are the data points, and the contributions for vertices from RCBG, *c*, and *b* decays. The arrows indicate the cut on this variable, placed at 20 μ m.

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Figure 13.11: The Am distributions for the decay length analysis, data points only. The signal and sideband regions, **as** well **as** the number of events they contain, are also shown.

	$D \rightarrow K \pi$		$D \rightarrow K \pi \pi^0$		
cut	c eff	beff	$c\mathrm{eff}$	b eff	
D^0 mass	0.842	0.912	0.462	0.577	
vertex χ^2	0.996	0.999	0.883	0.788	
L/σ_L	0.524	0.731	0.402	0.643	
vertex doca	0.758	0.204	0.581	0.163	
$x_D > 0.2$	0.975	0.859	0.938	0.640	
$\Delta(m)$	0.989	0.973	0.798	0.650	
Total	0.321	0.113	0.071	0.020	

Table 13.2: The efficiencies for D^0 decays in c and b events to pass the decay length selection cuts described in the text,

samples and using this as the set of candidate D^0 decays for the asymmetry measurement and background studies. Any errors introduced by not treating potential slight differences between the two samples is expected to be negligible compared to the other measurement errors. The overlap of the two samples is $(28.4 \pm 5.7)\%$ for the $K\pi$ mode and $(28.2 \pm 4.6)\%$ for the $K\pi\pi^0$ mode^{††}. The total number of events that appear in the signal region for the combined samples are given in Table 13.4.

13.2 The Selection of D^+ Candidates

In order to select a clean enough sample of events containing D^+ mesons, the selection criteria comprise several of the features of the previous two methods.

Candidates for $D^+ \rightarrow K^- \pi^+ \pi^+$ decays are formed by combining two tracks of the same sign with one track of the opposite sign, where all three tracks are required to have p > 1 GeV/c. The minimum momentum requirement serves to insure that the track parameters are well-measured, and results in little loss of efficiency for D^+s from charm decay due to the hard fragmentation function^{‡‡}. The two like-sign tracks

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Figure 13.12: The $\cos \theta^*$ distribution for D^+ decays (dashed line) and background (solid line). The arrow shows the cut $\cos \theta^* > -0.8$ placed on this variable.

are assigned π^+ masses and the opposite sign track is given the K^- mass. A series of cuts is then applied to reject RCBG and D^+ decays from b events.

First, we require $\cos \theta^* > -0.8$, as the helicity angle cut also works on D^+ events. The distribution of $\cos \theta^*$ in D^+ events is shown in Figure 13.12. To reject D^{*+} decays, the differences between $m(K^-\pi^+\pi^+)$ and $m(K\pi)$ are formed for each of the two pions and are required to be greater than 0.160 GeV/ c^2 , as this is where the mass difference Am must peak for combinations of a slow pion with $D^+ \rightarrow K^-\pi^+$ and other partially reconstructed D^0 decay modes. To remove RCBG, we require $L/\sigma_L > 3.0$ for the D^+ decay length, which has the same effect as the cut applied in the D^0 selection criteria, above. The cut is 3.0 σ in this case because the backgrounds in this decay mode are more severe; the longer lifetime of the D^+ allows us to tighten this cut without much loss of signal.

 $^{^{\}dagger\dagger}MC$ studies show that the exact value of the overlap fraction is very sensitive to the background contamination in the sample. This result will be used in our systematic studies, which appear in Chapter 16.

^{‡‡}We will apply a cut on x_{D^+} later, anyway.



Figure 13.13: The angle between the D^+ momentum and its flight direction. The contributions from c decays, and the combination of RCBG and b decays is shown with the **data** points.

Then, a "collinearity" cut is applied, requiring the vertex momentum direction to be collinear with its flight direction. This is essentially the same restriction as the impact parameter cut used in the D^0 decay length analysis described above, and is made to reject D^+s from b events and other RCBG vertices. the angle between the D^+ momentum vector and the vertex flight direction is required to be less than 5 mrad in the z-y plane, and less than 20 mrad in the r-z plane. In contrast to the fixed value of the impact parameter cut, the collinearity cut corresponds to an impact parameter cut whose allowed value increases with the decay length; this type of cut was found to yield better overall signal-to-noise than a standard impact parameter **cut** in this particular case. The value of 5 mr, though, is more restrictive than the simple impact parameter requirement of 20 μ m, as this requires an impact parameter

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Figure 13.14: The distribution of $K\pi\pi$ masses after cuts, with a pure $D^+ \rightarrow K\pi\pi$ signal shape overlayed on the data points. The MC shape has an arbitrary normalization and is shown to provide an indication of the location and shape of the D^+ signal peak. Also shown are the signal and sideband regions.

of the vertex momentum relative to the IP of 5 μ m at a decay length of 1 mm. Most of the reconstructed D^+ mesons in charm events fall within this cut, however. Plots of the collinearity angle in the *x*-*y* and r-z planes for MC signal events are shown in Figure 13.13.

	$D^+ \rightarrow K \pi \pi$		
cut	$c \mathrm{eff} b \mathrm{eff}$		
p > 1 GeV/c	0.655 0.494		
- θ* ·	0.850 0.850		
L/σ_L	0.828 0.976		
xy angle	0.452 0.832		
<i>rz</i> anglr	0.369 0.551		
A(m) > 0.16	0.892 0.968		
$x_D > 0.4$	0.436 0.193		
D^+ mass	0.869 0.800		
Total	$0.026 \ 0.028$		

Table 13.3: The efficiencies for each of the cuts in the D^+ analysis for c and b events.

Finally, x_{D^+} is required to be > 0.4; even with the above cuts, this hard x cut is necessary to achieve a high-purity sample of D^+ events. Table 13.3 gives the efficiency of each of the cuts. It is unclear at this time why the efficiencies for retaining D^+ decays from c and b events are equal. This may just be a computational error.

The mass peak for pure signal events from MC is shown in Figure 13.14 along with the events from the data that pass the above cuts. After all selection criteria, $D^+ \rightarrow K^- \pi^+ \pi^+$ candidates fail in the range 1.800 GeV/ $c^2 < m(K^- \pi^+ \pi^+) < 1.940$ GeV/ c^2 , while the sideband regions are defined as 1.640 GeV $< m(K^- \pi^+ \pi^+) < 1.740$ GeV/ c^2 and 2.000 GeV/ $c^2 < m(K^- \pi^+ \pi^+) < 2.100$ GeV/ c^2 . The total number of events that appear in the signal region is given in Table 13.4.

13.3 Determination of the Background and Signal Fractions

The relative fractions of signal and RCBG remaining after all analysis cuts are estimated using the MC to provide the correct shapes, then normalizing to the observed data sample. The statistics of a study using wrong-sign combinations to estimate

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the background are poor; we will use the wrong-sign distribution (among other techniques) to estimate the systematic error on this procedure.

To obtain the "asymmetry-carrying" charmed hadron signal shape in both the signal and sideband regions, all MC events with three tracks from a "correctly-signed" charm decay are taken as signal. For the D^0 analysis, this is a combination, then, of properly reconstructed $D^{\star+}$ decays, partially reconstructed $D^{\star+}$ decays where the π_{\bullet}^{+} (and hence the sign of the charm quark) is properly identified, and a small fraction of correctly reconstructed three-prong D^+ decays. If the π^+_s is misidentified in partially reconstructed $D^{\star+}$ decays where the D^0 decays into two charged tracks, there is essentially a 50% probability that the sign is chosen correctly from the other two decay products, so that the "incorrectly-signed" partial D^{*+} decays contribute no asymmetry. The contributions from four-prong D^0 decays are small, since the probability for mass misassignments to boost the reconstructed mass near the signal region is small. The five-prong branching fractions for charged D decays are significantly smaller, so they contribute even less to the signal fraction. We estimate the relative fractions of the different contributions in the signal sample to be: $D^{\star+}$:Partial $D^{\star+}:D^{\pm}:95\%:4\%:1\%$. Virtually all of the signal in the sideband region comes from charm events other than the chosen modes, however. For the D^+ analysis, all charm decays where the three tracks are correctly associated with a decaying charm meson are counted as signal, whether they originate from $D^{\star+}$, D^+ , or from other charmed hadrons, including D_{+}^{+} . The three tracks from modes other than the D^{+} decays must properly determine the sign of the charm quark. Again, contributions from five-prong charged D decays are expected to be small due to the small branching fractions for these modes. The probability for selecting three tracks from the decay $D^{\star +} \rightarrow \pi_{\star}^{+} D^{0}$ where the D^0 decays via $D^0 \to K^- \pi^+ \pi^+ \pi^-$ and misassigning the masses is small. The contributions of the different signal types are: $D^+:D^{*+}:"D^+"::98\%:2\%:<1\%$. The signal contributions to the sidebands do, however, come exclusively from these other asymmetry-carrying modes, like $D_+ \rightarrow K^* \ell \nu$ or $D^+ \rightarrow K^- \pi^+ \pi^+ \pi^0$.

The background shape is then taken from all entries in the signal and sideband mass regions that do not come from a 3-prong charmed hadron decay.

To obtain the fractions of signal and background in the event sample, the relative



Figure 13.15: The Am and D^+ mass distributions for data (left column) and MC (right column) for the three decay modes. Figures (a) and (d) are the distributions for $D^0 \to K\pi$, (b) and (e) are those for $D^0 \to K\pi\pi^0$, and (c) and (f) are those for $D^+ \to K\kappa\kappa$.

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normalizations of the MC signal and MC background shapes are selected such that the overall distribution matches that observed in the data in the signal and sideband regions only. The "dead" regions between the signal and sidebands will be used for systematic studies. The reason for taking this approach can be seen in the six panels of Figure 13.15, where we show the distributions in Am and $m(K\pi\pi)$ for data and MC. It is obvious that the MC predictions for the signal shape does not match the data for the $D^+ \rightarrow K\pi\pi$ mode. The distributions for the $D^0 \rightarrow K\pi\pi^0$ mode are also different. This is due to several incorrect production cross sections and branching fractions in the MC, which will be discussed below. Adjusting the relative normalizations of the signal and background shapes frees us from needing the MC to predict the absolute level of the background. The errors on the relative fractions will be included as a systematic error in Chapter 16.

The results of this background calculation procedure are shown in Figures 13.16, 13.17, and 13.18, where the shapes of the D^0 Am and D^+ mass distributions and the estimated background and signal shapes are plotted along with the experimental data. We see that good agreement between the MC and data is now obtained. To make this agreement possible, 18% more signal was necessary in the $D^0 \rightarrow K\pi$ mode, 48% more signal was necessary in the $D^0 \to K\pi\pi^0$ mode, and the branching fraction for $c \rightarrow D^+$ had to be increased by 35% in the $D^+ \rightarrow K\pi\pi$ mode. Some explanation for the variation between MC and data should be put forth here. When the predictions of the MC disagree with the data, the MC parameters immediately become suspect. We have checked in detail the values of production rates and branching fractions for charm mesons in the SLD MC, and some of our findings serve as explanations for the normalization differences. One major flaw of the MC simulation is that the ratio of the rate for vector meson to that for vector+pseudoscalar meson production in charm decays $(R_v/(R_v+R_p))$ is set to 3/4 in the MC instead of the 0.55 measured in b decays'. This has the effect of suppressing the rate of D^+ production in charm decays, which explains the differences between data and MC in the $K^-\pi^+\pi^+$ mass distribution. Once a compensating 35% more $c \rightarrow D^+$ decays have been added, the agreement between MC prediction and the observed data distribution is good. The incorrect

^{*}See Chapter 16 for references.

value of $R_v/(R_v + R_i)$ should impact the D^0 distributions in the opposite direction, however, as it would imply that more $D^{\star+}$'s are being produced in the MC than in nature. We need more signal in the MC to match the observed distributions, though, which does not match our expectations. We have checked the $D^{\star+} \rightarrow D^0$ branching fraction, and the MC value of 67.9% matches well the PDG value of (68.1±1.3)%[194]. The $D^0 \to K^- \pi^+$ branching fraction is the MC is too small, 3.57% compared with the PDG (4.01 ± 0.14) %, which would explain the 18% difference in the signal fractions between data and MC if we believed that $R_{\nu}/(R_{\nu} + R_{\mu})$ is correct in the MC. We should be able to distinguish whether the relative fractions of b and c decays in the MC matches what we observe in the data by comparing the signal to background ratios in the decay length vs. the kinematic analysis while cuts are varied, especially the x_D and pointing cuts, as this should change the relative fractions of b and c events differently in the two samples depending on the ratio of 6 to c content of the D^0 sample. These studies are still under consideration, and no clear explanation for the difference between MC and data for the D^0 modes exists. The fraction of double-tagged events does lend credence to the results of our background calculations, however; see Chapter 16.

These plots also show that the mass resolution of the tracking system is more or less correctly simulated in the MC[†]. The fractions of events in the signal regions that come from charm $(f_{c\rightarrow D})$, b events $(f_{b\rightarrow D})$, and RCBG (f_{RCBG}) are shown in Table 13.4 for the signal region. The same quantities for the sideband regions are given in Table 13.5.

13.4 The Raw Asymmetry

At this point, we can display the left-right forward-backward asymmetry present in the selected data. Figure 13.19 shows the distributions of the outgoing fermion directions as determined by the sign of the D meson for the two different beam helicity

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Figure 13.16: The Am distribution for data, MC signal and RCBG for $D^0 \rightarrow K\pi$. The relative normalizations of the MC signal and RCBG shapes has been obtained by matching the number of data events in the signal and sideband regions.

[†]We have already chosen our mass acceptance windows to be large enough to make the analysis independent of small differences in the mass resolution between data and MC, so it is gratifying to see that they agree well.



Figure 13.17: The Δm distribution for data, MC signal and RCBG for $D^0 \rightarrow K\pi\pi^0$. The relative normalizations of the MC signal and RCBG shapes has been obtained by matching the number of data events in the signal and sideband regions.



Figure 13.18: The Δm distribution for data, MC signal and RCBG for $D^+ \rightarrow K\pi\pi$. The relative normalizations of the MC signal and RCBG shapes has been obtained by matching the number of data events in the signal and sideband regions, with the addition of 35% more $c \rightarrow D^+$ events. (See text.)

Decay Mode	# of Signal Even	nts f_{c-}	$\bullet_D f_{b \to L}$	$o f_{RCBG}$
$D^{*+} \rightarrow \pi^{+}_{s}(K^{-}\pi^{+})$	88	0. 52	0. 22	0. 26
$D^* + \rightarrow \pi^+_{\mathbf{A}}(K^-\pi^+\pi^0)$	131	0.50	0. 22	0. 28
$D^+ \rightarrow K^- \pi^+ \pi^+$	98	0. 70	0.14	0. 16

Table 13.5: Total number of sideband events and the estimated fractions of c and b signal and RCBG in the sidebands for the three charm decay modes

Decay Mode	# of Signal Events	$f_{c \rightarrow D}$	$f_{b \to D}$	f_{RCBG}
$D^{\star +} \rightarrow \pi^+_e(K^-\pi^+)$	230	0. 047	0. 036	0. 917
$D^{*+} \to \pi^{+}_{s}(K^{-}\pi^{+}\pi^{0})$	398	0. 073	0.044	0.883
$D^+ \rightarrow K^- \pi^+ \pi^+$	45	0. 162	0. 020	0. 660



CHAPTER 13. ANALYSIS METHOD

states. Even with the low number of events, a large forward backward asymmetry can be seen whose sign changes when the beam polarization is reversed. A more revealing manner of displaying the asymmetry is to reverse the plot for the right-handed beam and overlay this on top of that for the left-handed beam. This corresponds to forming the first set of sums in the formula for \tilde{A}_{FB} from Eq. 2.72

$$\tilde{A}_{FB} = \frac{\sigma_F^L + \sigma_B^R - (\sigma_F^R + \sigma_B^L)}{\sigma_F^L + \sigma_B^R + \sigma_F^R + \sigma_B^R + \sigma_B^R} = \frac{3}{4} |P_e| A_f .$$
(13.1)

This distribution is shown in Figure 13.20. The background asymmetry shown on this plot is taken from the sidebands, and has had the asymmetry from the signal contamination of the sidebands subtracted. Note the large raw asymmetry. This shows in a very visual manner the advantages of measuring the asymmetry using a longitudinally polarized beam.

We turn now to a discussion of the maximum likelihood analysis used to extract A, from the charm sample.



Figure 13.19: The direction of the outgoing $D^{(*)+}$ for Z^0 events produced with electron beams of left-handed (a) and right-handed (b) helicitics. The angle shown is the direction of the outgoing $D^{(*)+}$ meson containing the c (as opposed to the \bar{c}) quark.



Figure 13.20: The forward backward asymmetry, combining the two polarization states by adding the distribution of $\cos(-\theta)$ for those events produced with a right-handed electron beam (see Fig. 13.19) with that of $\cos \theta$ for those events produced with a left-handed electron beam. The asymmetry shown from the background is from the sidebands with the appropriate amount of MC signal asymmetry subtracted.

Chapter 14

The Asymmetry Likelihood Function

14.1 Formulation of the Likelihood Function

To extract the charm asymmetry parameter $A_{,,}$ we use an unbinned maximum likelihood fit[202] based on the Born-level cross section for fermion production in Z^{0} -boson decay from Eq. 2.56:

$$\frac{d\sigma}{d\cos\theta} \propto (1 - A_e P_e) (1 + \cos^2\theta) + 2A_f (A_e - P_e) \cos 19, \qquad (14.1)$$

The fitting procedure used here is similar to that used by the OPAL collaboration in Ref. [203]. In this section, we give an overview of the likelihood fitting method. A detailed discussion of the fit parameters will be presented below.

The likelihood function should have the general form

$$\log \mathcal{L} = \sum_{i=1}^{n} \log \left(P \cdot \Theta(\mathbf{y}_i) \cdot (1 + \mathbf{y}_i^2 + A_f \cdot \mathbf{y}_i) \right), \qquad (14.2)$$

where $y = q \cdot \cos \theta$, q is the sign of the D^{*+} or D^+ charge, and θ is taken here to be the angle between the D^{*+} or D^+ meson momentum and the electron beam. Here, n is the total number of events. The function P is a normalization factor, and Θ parametrizes any non-uniform acceptance as a function of y. Note that any geometric acceptance

asymmetry should be independent of the *D* charge, and is thus an even function of y_i . Terms depending on Θ and the overall normalization merely add a constant term in the above sum and thus cannot affect the asymmetry parameter given by minimizing \mathcal{L} . We will take up the question of non-uniform acceptance again in the discussion of the systematic errors.

The actual form of the likelihood function used in this analysis is

$$\log \mathcal{L} = \sum_{i=1}^{n} \log \left\{ P_{b}^{j}(x_{Di}) \cdot \left[(1 - P_{e}A_{c})(1 + y_{i}^{2}) + 2(A_{e} - P_{e})y_{i} A_{b}^{D} \cdot (1 - \Delta_{QCD}^{b}(y_{i})) \right] \right. \\ \left. + P_{c}^{j}(x_{Di}) \left[(1 - P_{e}A_{e})(1 + y_{i}^{2}) + 2(A_{e} - P_{e})y_{i} \cdot A_{c}^{D} \left(1 - \Delta_{QCD}^{c}(y_{i}) \right) \right] \right. \\ \left. + P_{RCBG}^{j}(x_{Di}) \left[(1 + y_{i}^{2}) + 2A_{RCBG}y_{i} \right] \right\}$$
(14.3)

where A_b^D and A_c^D are the asymmetries from D^* and D^+ decays in tagged $b\bar{b}$ and $c\bar{c}$ events, respectively. The index j is used to indicate that each of the three charm decay modes are considered separately, so that the total likelihood is actually the sum of three parts:

$$\log \mathcal{L} = \log \mathcal{L}_{D^0 \to K\pi} + \log \mathcal{L}_{D^0 \to K\pi\pi^0} + \log \mathcal{L}_{D^+ \to K\pi\pi}$$
(14.4)

The quantity $\Delta_{QCD}^{j}(y)$ is the $\mathcal{O}(\alpha_{s})$ QCD radiative correction to the asymmetry, including quark mass effects [52] (See Section 2.6.2). The terms P_{b}^{j} , P_{c}^{j} , and P_{RCBG}^{j} represent the probabilities that an event from one of the j = 1, 2, 3 decay modes is either signal from $b\bar{b}$ or $c\bar{c}$ events or random combinatoric background, depending on the D candidate energy. An event is uniquely assigned to a decay mode, and is thus only counted once; less than 0.5% of events are double-tagged, so this has a negligible effect on the analysis.

In the next session, the detailed assumptions involved in the specification of the functions P^{j} , and the other values which must be input to the likelihood function are discussed.

14.2 Inputs to the Likelihood Analysis

14.2.1 Determination of Asymmetries

Since the number of D events available for this measurement are limited, we have chosen to input the values of the asymmetries A_b^D and $A_{,,,,}$ to the fit, and then fit only for $A_c^D \equiv A_{,.}$ Another reason for adopting this approach is the value of A_b^D , which, after correcting for $B^0 \bar{B}^0$ mixing, is almost identical to the standard model value for $A_{,,}$ 0.67. Since the x distributions or the Ds from b events from Ds from c events overlap to a large extent and the difference in D energy is the only variable we use to separate b and c, this coincidence of values would lead to a far worse error on our value of A, than if we fixed A_b^D to its standard model value and considered suitable errors about this point. Future versions of this analyses may be sensitive enough to the differences between b and c events to allow A_b^D to float in the fit'.

Determination of A_b^D

In order to arrive at the proper value of A_b^D to be included in the fit, we must take into account the dilution of the *b* quark asymmetry due to *b* mixing and possible effects from charm mesons produced in *W* splitting. Our estimate A_b^D proceeds in a manner similar to that of Ref [203] and Ref [67], where the corrections to A_b to obtain the observed asymmetry A_b^D are parametrized by

$$A_{,} = A_{b} \cdot (1 - 2\chi_{mix})(1 - 2\chi_{W \to cs}) .$$
(14.5)

We begin with the Standard Model prediction of $A_b = 0.935$ [51], and allow this to vary by ± 0.105 to accommodate the measurements from LEP and SLD [204]. The fraction of $D^{\star+}$ or D^+ , denoted as $D^{(\star)+}$ collectively, coming from the spectator part of B_d decays (Figure 14.1a) is taken half way between its two extrema, which are that either all $D^{(\star)+}$ mesons in B decays come only from B_d decays, or that there is no preference for $D^{(\star)+}$ coming from any particular type of B decay (B_u , B_d , B_s , Λ_b)

^{*}See the discussion in Chapter 17.
(Figure 14.1b). In the first case, the mixing parameters $\chi_d = 0.16 \pm 0.025$ measured at CLEO[205] should be used to calculate the value of A_b after mixing. For the second case, the average mixing parameter $\bar{\chi} = 0.12 \pm 0.01$ measured at LEP[123] should be used. We choose to take the average of these two extremes, which results in a mixing correction factor $1 - 2\chi_{mix} = 0.72 \pm 0.09$. The error on χ_{mix} is obtained by taking the measured χ_d value, adding 1 standard deviation, and subtracting $\bar{\chi}$:

$$\delta\left(\chi_{mix}\right) = \chi_d + \sigma\left(\chi_d\right) - \bar{\chi} \tag{14.6}$$

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The correction for wrong-sign $D^{(\star)+}$ from the W^- in $b \to cW^-$, $W \to \bar{c}s$ (Figure 14.1c) is expected to be small, since the estimated branching fraction for this splitting to D mesons is estimated to be at most 1%[203]. We conservatively take $\chi_{W\to cs} = 0.025 \pm 0.025$ to encompass a large possible range for this effect. These two corrections combine to give a value

$$A_b^D = 0.64 \pm 0.11 , \qquad (14.7)$$

where the error is taken as an experimental systematic error. This value is consistent with those obtained by ALEPH[67] and OPAL[203].



Figure 14.1: Possible sources for D mesons in B decays. Ds can come specifically from B_d decays (a), nonspecific b hadron decays (b), or from $W \rightarrow$ cs splitting (c).

Determination of A,,,,

With proper subtraction of the fraction of signal in the sideband regions as given in Table 13.5, we measure an average sideband asymmetry of $A_{,,,,} = 0.05 \pm 0.10$; all of the sideband data from the three decay modes give consistent results+. Since the measured asymmetry in the sideband regions is consistent with zero, we assume $A_{RCBG} = 0$ for the central value and consider possible deviations from zero as a systematic error.

14.2.2 Formulation of the Functions P(x)

The determination of $P_{f}^{j}(x_{D})$ is based on the relative fractions and the x_{D} distributions for the three decay modes. The functions P can be factorized in the following manner:

$$P_{c}^{j}(x_{D}) = \frac{N_{\text{signal}}(x_{D})}{N_{\text{total}}(x_{D})} \frac{\omega_{c}d_{c}(x_{D})}{\omega_{c}d_{c}(x_{D}) + \omega_{b}d_{b}(x_{D})}$$

$$P_{b}^{j}(x_{D}) = \frac{N_{\text{signal}}(x_{D})}{N_{\text{total}}(x_{D})} \frac{\omega_{b}d_{b}(x_{D})}{\omega_{c}d_{c}(x_{D}) + \omega_{b}d_{b}(x_{D})}$$
(14.8)

$$P_{_{RCBG}}^j(x_D) = 1 - \frac{N_{\text{signal}}(x_D)}{N_{\text{total}}(x_D)}$$

Here, we have introduced the symbols ω_c and ω_b to represent the fraction of the signal that is due to $D^{(\star)+}$ from c and b events, respectively; the values are derived from Table 13.4.

The ratios $N_{\text{signal}}/N_{\text{total}}$ are calculated by

1. The distributions of x_D from the sidebands and from a MC sample containing only $D^{(\star)+}$ from b and c events are normalized so that the integral over the distribution yields unity.

[†]For a comparison, without removing the asymmetry from the signal contamination we obtain $A_{RCPG} = 0.12 \pm 0.09$.

2. The fractions of the x_D distribution from the MC signal is subtracted from that of the sidebands, with the relative proportion of signal and background given by Table 13.5. The sideband distribution is renormalized back to unity integral.

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- 3. The ratios $N_{\text{background}}(x_D)/N_{\text{total}}(x_D)$ are formed. The shape of $N_{\text{total}}(x_D)$ is taken from the data. The fractions of background used to form the ratio are taken from Table 13.4. The ratio is calculated in bins of width 0.05 due to limited statistics.
- 4. The ratios $N_{\text{signal}}(x_D)/N_{\text{total}}(x_D)$ are just given by $1-N_{\text{background}}(x_D)/N_{\text{total}}(x_D)$

Note that the MC is only used in this procedure to estimate the fraction of signal in the sidebands, and for the x_D distribution of the pure signal. Figures 14.2a through c show the contributions to the x_D distributions estimated from this procedure overlayed on the data. Note that the poor statistics in the sideband regions results in ratios $N_{\text{signal}}(x_D)/N_{\text{total}}(x_D)$ that have significant bin-to-bin changes. This will be considered later in the discussion of systematic errors.

The functions $d_c(x_D)$ and $d_b(x_D)$ are parametrizations of the energy distributions of the decaying $D^{(\star)+}$ mesons, where the shapes of the two energy distributions are taken from MC. Thus, for a given value of x_D , the ratio $\omega_c d_c(x_D)/(\omega_c d_c(x_D) + \omega_b d_b(x_D))$ gives the probability that a $D^{(\star)+}$ candidate is from a primary $c\bar{c}$ or $b\bar{b}$ event. The energy spectra are generated separately for the vector and pseudoscalar Ds. A fit to a parametrizing function is then performed for each of the spectra, and the fs are normalized to unit area by an overall constant. The charm sample is fit to the Peterson fragmentation function[122]:

$$d_c(x_D) \propto \frac{1}{x_D \cdot \left(1 - \frac{1}{x_D} - \frac{\epsilon_D}{1 - x_D}\right)^2}$$
 (14.9)

The fitting function for $D^{(\star)+}$ mesons from **b** decays was introduced by the OPAL collaboration[206]:

$$d_b(x_D) = A \cdot \exp\left(\frac{-(x_d - B)^2}{C}\right) \cdot \left(1 + \frac{D}{x_D} + E \cdot x_D^2 + F \cdot x_D^3\right) .$$
(14.10)



Figure 14.2: The x_D distributions for data (points), background (filled histogram), and MC signal (open histogram), for the decay modes (a) $D^0 \to K\pi$, (b) $D^0 \to K\pi\pi^0$, and (c) $D^+ \to K\pi\pi$. The structure at low x_D results from the lower x_D cut used in the decay length selection of D^0 decay modes.



Figure 14.3: The MC distributions of $x_{c\to D}$ and the fitted Peterson parametrizations for (a) $c \to D^*$ and (b) $c \to D^+$. Note that the fragmentation function for $c \to D^+$ is slightly softer, as one might expect. See Equation 14.9 for the functional form.

Variations of the shape of these functions will be considered in the discussion of the systematic errors. Figures 14.3 and 14.4 show the fitted *d* functions for charm and *b* samples, respectively, for $D^{\star+}$ and D^{+} fragmentation.

Note that, given the definitions above, the total probability for each candidate sums to unity, as it should.

14.2.3 Other Inputs

The value of the polarization is taken to be $63.0 \pm 1.1\%$, as discussed in Chapter 4.

We have taken A, to be equal to the value from the Particle Data Group[194], $A_e = 0.1617f0.012$, where the error is chosen to cover both the SLD and LEP central



Figure 14.4: The MC distributions of $x_{b\to D}$ and the fitted parametrizations for (a) $b \to D^*$ and (b) $b \to D^+$. See Equation 14.10 for the functional form.

values.

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14.3 Results of the Fit

Performing the maximum likelihood fit to the data sample, we obtain the value

$$A_{r} = 0.71 \pm 0.20 (\text{stat})$$
.

The minimization of the negative log-likelihood is shown in Figure 14.5.



Figure 14.5: The minimization of the likelihood function. The parabolic lit parameters are given by $y = P1 + P2 * (x - P3)^2$.

14.4 Checks of Fit Value

To cross-check the value of A, obtained from the fit, we formed the bin-by-bin forward backward asymmetry using the data as presented in Figure 13.20 from the previous chapter. We can fit this asymmetry to the cos o-dependent form of \tilde{A}_{FB} , which was given in Eq. 2.72:

$$\tilde{A}_{FB}^{c}(z) = \frac{\left(\sigma_{L}^{c}(z) - \sigma_{R}^{c}(z)\right) - \left(\sigma_{L}^{c}(z) + \sigma_{R}^{c}(z)\right)}{\sigma_{L}^{c}(z) + \sigma_{R}^{c}(z) + \sigma_{L}^{c}(z) + \sigma_{R}^{c}(z)} = |P_{e}| A_{c} \frac{2z}{1 + z^{2}} , \qquad (14.11)$$

where $z = \cos \theta$, and L (R) denotes that the Z^0 was produced with an electron beam of left-handed (right-handed) helicity. We performed this fit with four and eight bins in z to check the effects of binning statistics. The results are shown in Figures 14.6a and b. The result for the fits when the appropriate fraction of asymmetry from b events is subtracted are:

$$A_{,} = 0.68 \pm 0.23 \quad (4bin) \qquad (14.12)$$
$$A_{,} = 0.71 \pm 0.24 \quad (8bin)$$

which are in perfectly good agreement with the likelihood fit value.

We now turn briefly to a discussion of the radiative corrections that must be applied to this value to obtain the pole value of the asymmetry, A_c^0 , and then to a

treatment of the possible systematic errors that effect this measurement.



Figure 14.6: Results of fitting for A, with a binned asymmetry. The functional form is given by $\tilde{A}_{FB} = |z|/(1 + z^2) \cdot 3/4P_eA_c$, where $z = \cos\theta$, and we have taken A, as the free parameter. The fit result must be corrected for background fraction and b contamination to obtain the values for A, given in the text.

Chapter 15

Radiative Corrections: Obtaining A_c^0

Electroweak and strong interaction radiative corrections need to be applied to the preceding chapter's result to obtain the value of the fermion asymmetry parameter at the Z^0 pole, A_c^0 . The nature of these corrections was discussed in Chapter 2, so we will only remind the reader briefly here what these corrections entail.

15.1 Electroweak Corrections

Electroweak corrections are included from initial and final state radiation, vertex corrections, γ exchange, and $\gamma - Z$ interference^{*}. These were calculated using the program ZFITTER[51], with the inputs $m_{top} = 175 \text{ GeV}/c^2$ and $m_{Higgs} = 300 \text{ GeV}/c^2$. The sum of these corrections is an upward correction of 0.8% to the uncorrected result. The small size of this correction illustrates the robustness of this measurement to other electroweak uncertainties, as most of the large electroweak corrections occur in the Z^0 and γ self-energy terms, which are absorbed into $\sin^2 \theta_W^{eff}$ and thus do not affect the measurements of the final state couplings as much as they do other measurements; for the unpolarized A_{FB}^c , measured at LEP, this radiative correction is so

Table 15.1: The Fraction of QQQQ or qqQQ events, collectively denoted as 4Q events, which pass our event selection cuts, tabulated for 4Q events from each of the three types of primary quarks.

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Energy fraction	uds 4Q events	c 4Q events	b 4Q events
$x_Q > 0.2$	0.791	0.778	0.789
$x_{O} > 0.4$	0.402	0.397	0.374

small, we will ignore possible variations of the correction due to uncertainties in the Higgs and top quark masses, as these are expected to be negligible.

15.2 QCD Corrections

The final state QCD corrections for massive quarks have been calculated to second order [54]. The second order correction depends in detail on the acceptance of analysis cuts for four-jet events wherein a radiated gluon emerges **as** a pair of heavy quarks, and is given by

$$A_{c}^{meas}\left(\mathcal{O}(\alpha_{s}^{2})\right) = A_{c}^{0}\left(1 + \left(\frac{\alpha_{s}}{\pi}\right)^{2}\left[-(4.4 \pm 0.4) - (26 \pm 6)f_{D_{B}}\right]\right),$$
(15.1)

where f_{D_B} is the fraction of these events which **pass** our analysis cuts. The other uncertainties are due to the unknown values of the quark masses. We estimated f_{D_B} using the MC to study the energy fractions x_Q carried by the heavy quarks in these events. The results are shown in Table 15.1. In our data sample, there are 55 events with $x_D < 0.4$, so the proper value of f_{D_B} to take is

$$\langle f_{D_B} \rangle = \frac{55}{317} \cdot (0.8) + \frac{262}{317} \cdot (0.4) = 0.47,$$
 (15.2)

where we have taken the average value of 0.8 and 0.4 from the table. We assign an error of f0.25 on the quantity f_{D_B} to cover our relative ignorance of the rate of 4Q events in the MC. When combined with a value[194][†] of $\alpha_s(M_z^2) = 0.119 \pm 0.02$ and

^{*}See Section 2.6.1 for a detailed discussion

[†]We have allowed a generous error on α_s .

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the other theoretical uncertainties in the above equation, this value of f_{D_B} gives a value of $-(2.3 \pm 1.0)\%$ for the correction, meaning that $A_c^0 = A_c^{meas}/(1 - 0.023)$. Including all radiative corrections, the final result becomes

$$A_c^0 = 0.73 \pm 0.22$$
(stat) (15.3)

Chapter 16

Systematic Errors

In this chapter, we discuss the possible systematic errors on the result for A_c^0 . As we shall see from the size of the errors derived below, the dominant systematic errors arise from our inability to place strong constraints on the amount of background contamination to our $D^{(\star)+}$ signal due to the small number of events in the sidebands. Let us turn now to the presentation of the systematic errors, which will be summarized in Table 16.5.

16.1 Systematic Errors Related to the Background

This section collects all of the systematic errors associated with our imprecise knowledge of the properties of the background sample.

16.1.1 Determination of the Background Fraction

The fractions of signal and RCBG (f_{RCBG}) in the signal regions were calculated in Section 13.3 by adjusting the background and signal shapes in the signal and sideband regions until the two combined shapes fit the observed data. This calculation suffers from two problems: low background statistics and incorrect MC values for several of the branching fractions and production cross sections of the *D* mesons used in the analysis. To minimize the effects of these uncertainties, we used the only the shapes



Figure 16.1: The MC background distribution for the D^+ analysis. The sideband regions are those between the dashed lines. The sloping line shown is drawn between the mean values of the number of events in each sideband, with the assumed starting and ending points being the centers of each of the sideband regions. The grey region indicates the amount by which this method of determining the background shape overestimates the number of background events in the signal region.

of the signal and background samples and allowed their relative normalizations to fluctuate. Even this was not sufficient in the case of the the D^+ sample, as the cross section for the production of D^+ in charm events was too low in the MC simulation, so that 35% more pure D^+ signal needed to be added to the distribution to obtain reasonable agreement.

We can estimate the error on these calculations in a number of independent ways. First, there is the simple statistical error on the number of events in the sidebands for both MC and data. This amounts to 13.1% of f_{RCBG} . Second, we can use the measured asymmetry in the sidebands with no signal subtracted, 12%, and assume that this comes entirely from signal, so that the fraction of RCBG in the sidebands,

 f_{BCBG}^{SB} , is given by

$$0.12 = (1 - f_{RCBG}^{SB}) \ge 0.67 .$$
 (16.1)

This implies that f_{RCBG}^{SB} is $(6.4 \pm 13.5)\%$ less than what we calculate with the other method. Thirdly, we can compare the background shapes from the MC to the observed data distributions. For the D^{*+} modes, we can use the MC background shape to project the number of data events in the "dead" region between the signal and sideband regions of the Am plot. Any excess can then be attributed to an error in f_{RCBG} . We find a difference of $(4.5 \pm 21.0)\%$ in these two decay modes. We can also estimate the amount of background in the D^{*+} modes by looking at wrong-sign three track combinations, where the D^0 is formed using two same-sign tracks and is then combined with an oppositely-charged track as the slow **pion** candidate. Unfortunately, statistics on these combinations are limited in the data; this technique results in an estimate of $\delta f_{RCBG} = 41 \pm 23\%$. For the D^+ mode, there are a number of options:

- the procedure described in Section 13.3 gives $f_{RCBG} = 0.16$
- . we can take from the MC the difference in a linear extrapolation of the backgroundonly sideband populations underneath the D^+ mass peak and the number of events that actually are under the peak. This situation is shown in Figure 16.1. The linear extrapolation overestimates the number of background events in the D^+ sample by $(23 \pm 8)\%$ if the MC can be trusted at this level.
- We have also compared the predictions of a fit to the shapes of the background and signal with a single gaussian plus a third-order polynomial. The fit for the data is shown in Figure 16.2a and for the MC in Figure 16.2b^{*}. This fit to the MC gives an estimate of 101 signal events in the peak, which actually contains 145 D^+ decays. The fit is apparently pulled high by fluctuations just outside the mass peak gaussian, and the amount of signal is underestimated by 43%. Using the same (possibly flawed) technique for the data, we find that, when corrected for the amount of signal in the sidebands, we obtain an estimate that the mass

[•]The incorrect (too small) branching fraction for $c \to D^+$ is readily apparent in the figure from the MC.

Table 16.1: The number of double-tagged D^0 decays in data and MC

Decay Mode	f_{DT} (DATA)	f_{DT} (MC)
$D^{\star +} \to \pi^+_s (K^- \pi^+)$ $D^{\star +} \to \pi^+_s (K^- \pi^+ \pi^0)$	0.284 ± 0.057 0.282 ± 0.046	0.272 0.213

peak is 70% signal, or $f_{RCBG} = 0.3$ (which can be compared with the calculated value of 0.16). Since we know this technique errs towards underestimating the signal fraction by a large amount, this can be considered as an relatively hard upper limit on f_{RCBG} for the D^+ decay mode. This method serves then as more of a cross check than an estimation of the size of the variation in the background fraction.

Using the statistical errors on these measurements, we take a weighted average over all of the different techniques to obtain the error on the background fraction of the entire charm sample of $\delta f_{RCBG} = \pm 18.8\%$, or $f_{RCBG} = 0.239 \pm 0.05$.

An independent way to check the assumed background and signal fractions in the case of the D^0 decay modes is to examine the fraction of overlap between those events selected in the kinematic *vs.* the decay length analysis. The methods are sufficiently different in their demands on the signal events that we can hypothesize that only real signal events will actually be "double-tagged" by being selected by both sets of cuts. Let us first consider the double tag fraction f_{DT} found in the data compared with that from the MC with no modifications, These numbers are presented in Table 16.1.

Note that the MC events have a significantly lower fraction of double tags. We can use the MC simulation to check the hypothesis that only signal events are double-tagged. From the MC, we find that no background events are double-tagged, and the fraction of signal events that double tag (f_{DT}^{signal}) are 0.383 for the $K\pi$ mode, and 0.350 for the $K\pi\pi^0$ mode. Assuming that only signal events double tag, we have

$$f_{DT} = f_{DT}^{sig} (1 - f_{RCBG})$$
(16.2)

The double tag fraction is then directly proportional to the fraction of background in



Figure 16.2: A third order polynomial background fit to the D^+ signal and background for (a) the data, and (b) the MC. The first three fit parameters are those for the gaussian. Note that the width of the D^+ peak in the data is approximately 40% larger than that in the MC; the full extent of the peak still lies well within our signal region of $K\pi\pi$ mass.

Decay Mode	f_{DT} (DATA)	f_{DT} (MC)
$D^{\star +} \to \pi^+_s(K^-\pi^+)$ $D^{\star +} \to \pi^+_s(K^-\pi^+\pi^0)$	0.284	$0.284 \\ 0.251$

Table 16.2: The number of double-tagged D^0 decays in data and MC after correction

the signal region, making this a very sensitive test of our measure of the background fraction. If we take the values of f_{DT}^{signal} given above and use the fractions f_{ncng} derived from our procedure of normalization adjustment, we obtain the values of Table 16.2. These values support our conclusions on the relative amounts of signal and background in the two D^0 samples, as the double-tag fractions are now well within 1 σ errors of the data.

Assuming the allowed variation in the background fractions for the three decay modes in the likelihood fit leads to an overall error due to this source of 0.044 on A_c^0 . Changing of the background fractions modifies the amount of observed asymmetry that is assumed to arise from charm and *b* decays; an 18% change in the background fraction should have a large effect on the charm asymmetry.

16.1.2 The Background Asymmetry

When the asymmetry due to the presence of a small number of signal events in the sideband regions is subtracted, we obtain a net sideband asymmetry of $+0.05 \pm 0.10$. For the central value, we have assumed that the background asymmetry is zero, and have taken the la statistical deviation, 0.15, as an upper limit on the background asymmetry. This yields a change in A_c^0 of 0.044 when included in the likelihood fit. Since the background fraction in any of the samples is of the order 20%, we would expect that allowing a larger background asymmetry should have a non-negligible effect on the charm asymmetry.

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Figure 16.3: Comparison of signal and sideband acceptances, obtained by dividing the normalized distributions of the number of events vs. $|\cos \theta|$ for the signal and sideband data samples. The fit parameters are also shown.

16.1.3 Background Acceptance

From Figure 13.20, it can be seen that the acceptance for reconstructed Ds begins to fall off around $|\cos \theta| \sim 0.65$. In an attempt to check whether the acceptance for background events has the same behaviour, we plot in Figure 16.3 the ratio of the signal to background acceptance as a function of $|\cos \theta|$. As the ratio is symmetric about $\cos \theta = 0$, we have folded the distribution over to obtain smaller statistical errors. A simple linear fit yields

$$R_{\text{accept}} = -.41 \cdot |\cos \theta| + 1.16 , \qquad (16.3)$$

where R_{accept} is the ratio of signal to background acceptance. We assumed for the central value of A_c^0 that the signal and background acceptances were the same. If the background fraction in the fit is allowed to depend on $|\cos \theta|$ as specified by the above equation, the value of A_c^0 changes by 0.045. This error is expected to be one of the largest systematic errors, as it is difficult to establish that the signal and background acceptances are identical. Since the asymmetry is much larger at large $|\cos \theta|$, variations in the acceptance can lead to significant fluctuations if the background fraction is allowed to change where the event weights are largest.

16.1.4 Background *x* **Distribution**

Our sensitivity to statistical fluctuations in the energy spectrum of the random cornbinatoric background is checked by performing the analysis with the probability function P_{RCBG} derived from the Monte Carlo background instead of the data from the sideband regions. To do this, the x distribution from the MC sidebands is used instead of that from the data. The same procedure presented in Section 14.2.2 is then used to recalculate the ratios $N_{signal}(x)/N_{background}(x)$ for each of the three decay modes. This results in a change of 0.039 in A_c^0 , which is certainly consistent with the statistical errors on the number of events in the sidebands.

We can also check the size of this variation by using the MC instead of the data for all of the shapes including N_{tot} , the total shape of signal and background. Following the same procedure, we obtain a change in A_c^0 of 0.064. We know that this is an overestimate of the error, since the MC event sample has a smaller ratio of signal to background than is observed in the data.

As another check, we can ignore the presence of any amount of signal in the sidebands and just use the data distribution to calculate $N_{signal}(x)/N_{background}(x)$. This results in a change of 0.019 in A_c^0 . We would expect that this would be smaller than using the MC sideband shape, as the amount of signal in the sideband regions is small, so the net result of the signal subtraction is a minor change in the distribution.

16.2 Systematic Errors Related to the Signal

We collect in this section those systematic errors that result from the various assumptions made about signal production rates, composition, fragmenation functions, and asymmetries.

16.2.1 Determination of the Signal Composition

One of the potential *sources* of error *in* this analysis is the determination of the fraction of events from *b* or c decays that populate the signal region. One difficulty in deriving these quantities is that the branching fractions for $c \rightarrow D^+$ and the ratio

Table 16.3: Existing measurements of D^{*+} and D^{+} production in Z^{0} decays.

Source	$F_b \cdot P_{b \to D^*} / F_c \cdot P_{c \to D^*}$
DELPH[68]	1.119 ± 0.304
ALEPH[67]	1.119 ± 0.180
OPAL[69]	1.000 ± 0.201
OPAL[209]	1.367 ± 0.168
SLD MC	0.946
Source	$F_b \cdot P_{b \to D^+} / F_c \cdot P_{c \to D^+}$
DELPHI[68]	1.119 ± 0.437
SLD MC	1.539
OPAL[209] SLD MC Source DELPHI[68] SLD MC	$\begin{array}{c} 1.367 \pm 0.168 \\ 0.946 \\ \hline F_b \cdot P_{b \to D^+} / F_c \cdot P_{c \to D^+} \\ 1.119 \pm 0.437 \\ 1.539 \end{array}$

of the production rates of vector mesons to that for pseudoscalar mesons in charm decays used in generating our MC sample do not correspond to the measured values for these quantities.

If we make the simple argument that the physical processes involved in producing charmed mesons in the breakup of b hadrons should be similar to that which occurs in the random pairing of quarks during fragmentation, we would expect that the relative fractions of D mesons from these sources in the entire hadronic event sample to be given by $\Gamma(Z^0 \to b\bar{b})/\Gamma(Z^0 \to c\bar{c}) = 0.22/0.17 = 1.29$. Several measurements of D^{*+} and D^+ production have been performed by the LEP experiments. They typically measure the combined product $F_b \cdot P_{b\to D}/F_c \cdot P_{c\to D}$, where $F_q = \Gamma(Z^0 \to q\bar{q})$, and $P_{q\to D}$ is the branching fraction for quark type q to the D meson of choice. Their results are summarized in Table 16.3, along with the values used by our MC simulation. The values in the MC were adjusted in the tuning of the b physics modeling described in Chapter 4 and were the result of a conscious effort to make the modeling of b decays **as** accurate as possible. Since changing the MC is not possible at this stage, we can only estimate the effects of the variation of these production cross-sections. To aid this discussion, we introduce here the parameter $R_D \equiv F_b \cdot P_{b\to D}/F_c \cdot P_{c\to D}$.

The relevant quantity to the analysis is not actually the production cross section, but the fraction of the events from each quark flavor that pass the cuts, $f_{q\to D}$ =

Table 16.4: The variation of ω_b for different values of F_b . $P_{b\to D}/F_c \cdot P_{c\to D}$.

Assumption	ω_b	Change (%)
Default	0.254	_
$R_{D^*} = R_{D^+} = 1.12$	0.267	\$5.3
$R_{D^*} = R_{D^+} = 1.29$	0.294	+15.9
$R_{D^*} = 1.54$	0.329	+29.8
$R_{D^+} = 0.946$	0.237	-6.5

 $F_q \cdot P_{q \to D}$. ε_q^{cuts} , where ε_q^{cuts} is the probability that D mesons produced by the quark q pass the selection cuts. The quantity ω_q , the fraction of the signal from quark type q (q = c, b), was introduced in Chapter 13 and can now be defined as, e.g.,

$$\omega_b = \frac{f_{b \to D}}{f_{c \to D} + f_{-b} - b} \quad . \tag{16.4}$$

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The uncertainty in the relative production cross sections can then be reexpressed as an uncertainty on ω_b . If we calculate ω_b from our default MC simulation, we obtain $\omega_b = 0.254$ as the average for the entire charm sample. Table 16.4 gives the variation in ω_b for various assumptions on the relative production cross sections R_D .

The first assumption is that both values of R_D are equal to the world average value. The second corresponds to our naive guess of the relative cross sections. The third and fourth variations simply assume that $R_{D^*} = R_{D^+}$ and take the two different MC values as the possible extrema. We have chosen to take the largest possible variation as our systematic error, so we use $\delta \omega_b / \omega_b = 0.30$.

Using this variation of ω_b in the likelihood fit yields an overall error on $A_c^0 \ 0.005$ from this uncertainty. The insensitivity to the fraction of b events in the signal is due to the near identical values of A_b and A, for Ds after the dilution due to b mixing.

An independent cross check can be performed on the determination of the relative fractions of signal and background and the relative fractions of b and c in the signal for the case of the D^+ decay mode. Figure 13.18 shows a broad hump of signal events to the left of the main D^+ mass peak. If our estimate of the signal fraction in the sidebands is correct (it depends on both the addition of 35% of $c \rightarrow D^+$ production

and the sideband signal normalization procedure), the fitted asymmetry in this region should be non-zero. From the ratio of signal to background in the region from 1.0 $\text{GeV}/c^2 < m(K\pi\pi) < 1.64 \text{ GeV}/c^2$ which lies directly outside the sideband region, we calculate that the measured asymmetry should be 0.26. A four-bin fit yields a value of 0.26 \pm 0.18, which is certainly consistent. This implies that our treatment of the signal and backgrounds in the D^+ case is at least self-consistent.

16.2.2 Variation of the Value for $A_{b\to D}$

In Section 14.2.1, we discussed the procedure by which the value of $A_{b\to D}$ used in the likelihood fit, $A_{b\to D} = 0.64 \pm 0.11$, was determined. Allowing this value to fluctuate by the allowed error results in a 0.020 change in A_c^0 . As $A_{b\to D}$ moves away from being identically equal to $A_{,,}$ some fluctuation should result. The size of this error is about what is expected, given the relatively small fraction of b events in the signal sample.

16.2.3 Fragmentation Function Variation

A number of checks can be made on our formulation of the fragmentation functions which was presented in Section 14.2.2. First, we can consider variations in the "hardness" of the fragmenation functions for b and c decays based on the uncertainties on the measurements of the fragmentation functions themselves from the LEP collaborations[123]. The central values used in our MC sample are $\epsilon_c = 0.06$ and $\epsilon_b = 0.006$, where ϵ_f is the parameter in the Peterson fragmentation function that controls the shape of the quark energy spectrum. To account for the uncertainties in the fragmentation functions, we generated samples of charm and b decays to D^{*+} and D^+ mesons at values of $\epsilon_c = 0.035$, 0.095 and $\epsilon_b = 0.0035$, 0.011, which are the range of uncertainties recommended for use in studies of the electroweak physics of heavy quarks by the LEP working group on this topic. As an example of the different shapes that are obtained, Figure 16.4 shows the nominal, hard, and soft fragmentation shapes for $c \rightarrow D^+$ decays. For each change of fragmentation parameter, the functions $d_{c\rightarrow D}$ and $d_{b\rightarrow D}$ are refit, and the new parametrizations are input into the likelihood fit. Because of the difference in energy spectrum, the values of $N_{signal}(x)/N_{background}(x)$



Figure 16.4: The MC distributions of $x_{b\to D}$ for nominal ($\epsilon = 0.06$), hard ($\epsilon = 0.035$), and soft ($\epsilon = 0.095$) charm fragmentation.

are also adjusted accordingly in order to perform this check. This modification to the likelihood function results in maximal changes of 0.025 and 0.007 to A_c^0 for the changes in ϵ_c and ϵ_b , respectively. We expect a relatively large dependence on the assumptions of fragmentation functions for the charm decays, since the x_D distribution is the only tool we use to separate charm decays from c and b events.' Since c events contribute by far the largest fraction to the signal, the size of this variation is commensurate with their weight in the fit. The change in $d_{b\to D}$ for different values of ϵ is relatively small; since the Ds in b decay are produced in the decay cascade, changes in the fragmentation function produce only second-order effects on the D energy distribution. Modifying the fraction of the charm events that pass the cuts (ω_c) based on the changes of the fragmentation functions when the mean energies are

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shifted yields smaller changes to A_c^0 than the above values.

We have chosen to take the maximal errors as the uncertainty due to our imprecise knowledge of the fragmentation functions.

To check our assumptions on the shapes of the fragmentation functions, we can observe the variations on the value of A_c^0 when we change the fit inputs. For the **c** events, we can replace the parametrized fit of the Peterson function+ with a normalized bin-valued function based on the generated distribution. This will not treat events that lie on the rapidly changing parts of the spectrum properly, but it gives us some estimate of our sensitivity to the exact shape of the distribution. For the *b* events, we can replace the OPAL fitting function with a simple, exponential function times a gaussian. This yields acceptable fits in the regions above our cuts on x_D . The variations on A_c^0 due to these modifications are 0.023 and 0.005 for the c and *b* changes. We add these in quadrature to obtain an error due to our assumptions on the fitting shape.

16.3 Uncertainties from Outside Parameters

Finally, we have the uncertainties arising from the other inputs to the fitting function.

16.3.1 Uncertainty on A,

As mentioned in Section 14.2, we have chosen to use $A_{c} = 0.1617 \pm 0.012$. The resulting change in A_{c}^{0} due to allowing A, to vary over this range is 0.003. This bears out our statement that the results of the likelihood fit should yield a value of A_{c}^{0} that is statistically insensitive to A_{c} , as this method should be equivalent to forming the double asymmetry of Eq. 2.72 which is, by definition, independent of the initial-state coupling.

[†]The fit quality is not particularly good.

16.32 Polarization Uncertainty

Changing the value of the measured polarization by its error yields a variation of 0.013 on A_c^0 . This is a 1.9% error, which, though it is larger than the 1.1% error on the polarization value, is comparable to our expectations, since the error on A_c^0 is directly proportional to the polarization error.

16.3.3 Uncertainty in the Value of $\alpha_s(M_Z^2)$

For the QCD radiative corrections, a central value of $\alpha_s(M_Z^2)$ must be assumed. We have taken the Particle Data Group value of $\alpha_s(M_Z^2) = 0.119$ as averaged from measurements using jet rates and assigned a generous error of 0.02 in *our* calculation of the radiative corrections. This is justified, as estimates of the theoretical errors on the calculations of α_s range up to this size of error[131]. This variation in $\alpha_s(M_Z^2)$ results in a 0.005 change in A_c^0 .

16.3.4 Uncertainty in the $\mathcal{O}(\alpha_s^2)$ QCD Correction

The errors reported in Chapter 15 for the $\mathcal{O}(\alpha_s^2)$ QCD corrections are dominated by uncertainties in the quark masses as input to the QCD calculations for the size of the corrections. We also include the error on $\alpha_s(M_Z^2)$ and on the relative rates of the production of heavy quark events from gluon splitting ($f_{D_B} = 0.47 \pm 0.25$), to obtain a total relative error of 1% on this correction. The corresponding variation of A_c^0 is 0.007.

All of the systematic errors are summarized Table 16.5.

We now turn to a discussion of these results and a comparison with other existing measurements.

Table 16.5: Contributions to the estimated systematic error on A_c^0

SOURCE	RANGE	ERROR
$\begin{array}{c} f_{RCBG} & \text{total} \\ A_{RCBG} \\ RCBG & x & \text{distribution} \\ RCBG & \cos\theta & \text{acceptance} \end{array}$	0.239 ± 0.045 +0.05 ± 0.10 sideband vs MC sideband vs signal	0.038 0.044 0.039 0.045
$\begin{array}{l} f_{b \rightarrow D} / (f_{\sim_b \frown D^+} f_{c \rightarrow D}) \\ < x_c > \\ < x_b > \\ d(x) \text{ shape} \\ A_e \\ Polarization \\ a, (M_Z) \text{ value} \\ \mathcal{O}(\alpha_s^2) \text{ QCD} \end{array}$	$\begin{array}{c} 0.25 \pm 0.6 \not \pm \pm 0.0\% \ 0.11 \\ 0.035 \le \epsilon_{e} \le -0.095 \\ 0.003 \le \epsilon_{b} \le 0.011 \\ \text{bins vs. fit} \\ 0.1617 \pm 0.012 \\ 63.0 \pm 1.1\% \\ 0.118 \pm 0.002 \end{array}$	0.005 0.020 0.025 0.007 0.023 0.003 0.013 0.005 0.097
Total		0.095

Lept

SLD

Chapter 17

Conclusions

17.1 Final Result

Including the systematic errors discussed in the previous section, we arrive at a final result for A_c^0 of

$$A_c^0 = 0.73 \pm 0.22(\text{stat}) \pm 0.10(\text{syst}).$$
 (17.1)

17.2 Comparison with Existing Measurements

We present in Figure 17.1 a comparison of the results from measurements of A_c^0 made at SLD with those from the LEP collaborations; their measurements of A_{FB}^0 have been divided by their value for A, taken from their measurements in the lepton sector: $A_r = 0.1453 \pm 0.0057$. Using the world-average value of A_c only shifts the scale slightly. We can see immediately the added statistical power afforded by the polarized asymmetry. The presence of a large raw asymmetry in the polarized data allows our measurements to compete on equal footing with those from LEP using data samples containing forty times more Z^0 events. The total error on A, from the D analysis is actually smaller than that for all of the other measurements except the lepton result from OPAL and the ALEPH D^* analysis.



Figure 17.1: A comparison of the world's measurements of A_c^{o} . The results of this thesis and the other SLD result obtained by fitting the high p, p_{\perp} lepton spectrum have been averaged to quote an "SLD Average" value. (See text.) The vertical line is drawn at the standard model value of $A_c^0 = 0.67$, assuming $\sin^2 \theta_W^{eff} = 0.234$.

 $0.37 \pm 0.23 \pm 0.21$

The vertical line drawn on the figure is the standard model expectation for A, assuming $\sin^2 \theta_W^{eff} = 0.234$. Our result is certainly consistent with the standard model within our errors.

17.3 Determination of $\sin^2 \theta_W^{eff}$ from A,

As mentioned in Chapter 2, the value *A*, does have a strong dependence on the value of $\sin^2 \theta_{\mathbf{W}}^{eff}$:

$$\frac{dA_c}{d\sin^2\theta_W} = -3.42 \tag{17.2}$$

Given, this, we can invert the relationship

$$A_{c} = \frac{2v_{c}a_{c}}{v_{c}^{2} + a_{c}^{2}} = \frac{2\left[1 - \frac{8}{3}\sin^{2}\theta_{W}^{eff}(M_{Z}^{2})\right]}{1 + \left[1 - \frac{8}{3}\sin^{2}\theta_{W}^{eff}(M_{Z}^{2})\right]^{2}},$$
(17.3)

solving for $\sin^2 \theta_W^{eff}$. This yields the result

$$\sin^2 \theta_W^{eff} = 0.212^{+0.065}_{-0.133} \tag{17.4}$$

The reason for the asymmetric error can be seen in Figure 17.2, where the value of $\sin^2 \theta_W^{eff}$ is plotted against values of $A_{,.}$ Because of the functional dependence of the couplings, the lower bound is weaker than the upper bound. We can compare this result to that derived from $A_{LR}[47]$

$$\sin^2 \theta_W^{eff} = 0.2294 \pm 0.0010 . \tag{17.5}$$

It is easy to see that this measurement will certainly not ever compete with A_{LR} for a measurement of $\sin^2 \theta_W^{eff}$. Is is, however, a unique measurement of the Zcc coupling.

17.4 Possible Improvements

It is clear from the discussion of this measurement that event statistics are the limiting factor for this analysis. More data brings with it more analyzing power for the asymmetry as well as more statistics in the sideband regions whose uncertaintics dominate



Figure 17.2: The dependence of $\sin^2 \theta_W^{eff}$ on *A*,, showing the results of this measurement with combined statistical and systematic errors.

the systematic error on our measurement. To this end, we are at work on an inclusive method of tagging D^{*+} decays using a partial reconstruction of the D^0 vertex. Since the vertex flight direction can be measured extremely well by the precision tracking of SLD, the D^{*+} direction can be estimated with small errors. The π_s^+ from the D^{*+} decay will have extremely low p_{\perp} with respect to the vertex flight direction; tracks of this type will form a peak at $p_{\perp}^2 = 0$ in the distribution of transverse momentum of all tracks with respect to the vertex flight distances. With suitable kinematic selection criteria on the three-track combinations, a clean sample of $D^{*+} \rightarrow D^0$ decays can be selected without reconstructing the D^0 . We estimate that the efficiency for finding the two-prong vertex in these decays to be about 40%, so this tagging method should be much more efficient than the standard exclusive reconstruction.

In addition, we can also use the information provided by the precision vertex to help better distinguish between b and c decays in the sample. Once more statistics

are accumulated with the SLD detector, the errors due to *b* contamination may well be comparable to those introduced by the backgrounds. If we add to the likelihood function the probability distributions for vertex pointing, the proper lifetime of the decay, and any other variable we can derive that allows better separation of *b* and c events, we should be able to reduce the systematic errors even further. The new SLD vertex detector with hundreds of thousands of Z^0 decays should enable the world's best measurement of *A*, using these techniques.

If we wish to reach a precision of $\sim 1 - 2\%$ on the measurement of $A_{,,}$ which is the level at which deviations due to "new physics" could arise, this more efficient tag will be crucial. A simple extrapolation of the present result to $1 \times 10^6 Z^0$ decays only reduces the statistical error to 5%. If we wish to approach a 1% statistical error, we must somehow increase the selection efficiency for charm events by a factor of 25 without substantially worsening the sample purity. This will require the addition of more exclusive modes through the use of π^0 reconstruction, the use of the SLD CRID to tag D^0 decays without requiring the slow pion, and as many other techniques as we can create.

Appendices

Appendix A

Efficiency Matrices for R₃ **Unfolding**

Here, we list the four 3 x 3 matrices used to unfold the parton-level 3-jet rates from the tagged samples. We give the matrices for each of the three run periods separately. A sample matrix of statistical errors for the 2-jet to 2-jet tags looks like

$$\delta\left(\varepsilon_{2 \to 2}^{ij}\right) = \begin{pmatrix} 0.0023 \ 0.0035 \ 0.0011 \\ 0.0015 \ 0.0041 \ 0.0033 \\ 0.0001 \ 0.0012 \ 0.0035 \end{pmatrix} .$$
(A.1)

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E:

-

$$\begin{aligned} \varepsilon_{2 \to 2}^{ij} &= \begin{pmatrix} 0.4077 \ 0.1938 \ 0.0203 \\ 0.1173 \ 0.3169 \ 0.2279 \\ 0.0006 \ 0.0177 \ 0.2797 \end{pmatrix} & \varepsilon_{3 \to 2}^{ij} &= \begin{pmatrix} 0.1881 \ 0.1024 \ 0.0179 \\ 0.0650 \ 0.1355 \ 0.1140 \\ 0.0004 \ 0.0070 \ 0.1032 \end{pmatrix} \\ \varepsilon_{3 \to 3}^{ij} &= \begin{pmatrix} 0.2466 \ 0.1385 \ 0.0260 \\ 0.0909 \ 0.1911 \ 0.1787 \\ 0.0009 \ 0.0096 \ 0.1364 \end{pmatrix} & \varepsilon_{2 \to 3}^{ij} &= \begin{pmatrix} 0.1881 \ 0.1024 \ 0.0179 \\ 0.0650 \ 0.1355 \ 0.1140 \\ 0.0004 \ 0.0070 \ 0.1032 \end{pmatrix} \\ (A.2) \end{aligned}$$

EO:

P۰

$$\varepsilon_{2 \to 2}^{ij} = \begin{vmatrix} 0.4178 & 0.2025 & 0.0206 \\ 0.1201 & 0.3198 & 0.2327 \\ 0.0006 & 0.0182 & 0.2890 \\ 0.2454 & 0.1408 & 0.0238 \\ 0.0912 & 0.1928 & 0.1688 \\ 0.0010 & 0.0098 & 0.1407 \\ \end{vmatrix} \epsilon_{3 \to 3}^{ij} = \begin{pmatrix} 0.1874 & 0.0925 & 0.0174 \\ 0.0634 & 0.1396 & 0.1219 \\ 0.0003 & 0.0067 & 0.1074 \\ 0.0232 & 0.0105 & 0.0015 \\ 0.0072 & 0.0183 & 0.0112 \\ 0.0000 & 0.0008 & 0.0095 \\ \end{vmatrix}$$
(-4.3)

$$\varepsilon_{2\to2}^{ij} = \begin{bmatrix} 0.4181\ 0.1983\ 0.0201\\ 0.1174\ 0.3224\ 0.2304\\ 0.0005\ 0.0192\ 0.2890 \end{bmatrix} \varepsilon_{3\to2}^{ij} = \begin{pmatrix} 0.1701\ 0.0885\ 0.0139\\ 0.0563\ 0.1249\ 0.1087\\ 0.0002\ 0.0046\ 0.1103 \end{pmatrix}$$

$$\varepsilon_{3\to3}^{ij} = \begin{pmatrix} 0.2620\ 0.1450\ 0.0239\\ 0.0954\ 0.2033\ 0.1724\\ 0.0010\ 0.0111\ 0.1462 \end{pmatrix} \varepsilon_{2\to3}^{ij} = \begin{pmatrix} 0.0244\ 0.0116\ 0.0014\\ 0.0079\ 0.0193\ 0.0127\\ 0.0001\ 0.0007\ 0.0115 \end{pmatrix}$$
(A.4)

$$\varepsilon_{2 \to 2}^{ij} = \begin{pmatrix} 0.4166 & 0.1963 & 0.0194 \\ 0.1162 & 0.3208 & 0.2275 \\ 0.0005 & 0.0193 & 0.2887 \end{pmatrix} \varepsilon_{3 \to 2}^{ij} = \begin{pmatrix} 0.1619 & 0.0836 & 0.0128 \\ 0.0534 & 0.1179 & 0.1020 \\ 0.0002 & 0.0046 & 0.1060 \end{pmatrix} \\ \varepsilon_{3 \to 3}^{ij} = \begin{pmatrix} 0.2697 & 0.1486 & 0.0248 \\ 0.0986 & 0.2109 & 0.1805 \\ 0.0011 & 0.0108 & 0.1509 \end{pmatrix} \varepsilon_{2 \to 3}^{ij} = \begin{pmatrix} 0.0266 & 0.0138 & 0.0019 \\ 0.0083 & 0.0209 & 0.0139 \\ 0.0001 & 0.0009 & 0.0125 \end{pmatrix}$$
(A.5)

$$\varepsilon_{2 \to 2}^{ij} = \begin{pmatrix} 0.4213 \ 0.2016 \ 0.0190 \\ 0.1191 \ 0.3258 \ 0.2295 \\ 0.0005 \ 0.0185 \ 0.2919 \end{pmatrix} \varepsilon_{3to2}^{ij} = \begin{pmatrix} 0.1713 \ 0.0898 \ 0.0180 \\ 0.0578 \ 0.1293 \ 0.1058 \\ 0.0003 \ 0.0067 \ 0.0910 \end{pmatrix}$$

$$\varepsilon_{3 \to 3}^{ij} = \begin{pmatrix} 0.2589 \ 0.1449 \ 0.0276 \\ 0.0991 \ 0.1962 \ 0.1894 \\ 0.0012 \ 0.0104 \ 0.1470 \end{pmatrix} \varepsilon_{2 \to 3}^{ij} = \begin{pmatrix} 0.0210 \ 0.0101 \ 0.0019 \\ 0.0067 \ 0.0157 \ 0.0139 \\ 0.0000 \ 0.0003 \ 0.0093 \end{pmatrix}$$
(A.6)

G:

$$\varepsilon_{2 \to 2}^{ij} = \begin{pmatrix}
0.4030 \ 0.1868 \ 0.0153 \\
0.1074 \ 0.3142 \ 0.2088 \\
0.0004 \ 0.0185 \ 0.2891
\end{pmatrix} \varepsilon_{3 \to 2}^{ij} = \begin{pmatrix}
0.1348 \ 0.0715 \ 0.0099 \\
0.0440 \ 0.0980 \ 0.0770 \\
0.0002 \ 0.0063 \ 0.0884
\end{pmatrix}$$
(A.7)

$$\varepsilon_{3 \to 3}^{ij} = \begin{pmatrix}
0.2959 \ 0.1578 \ 0.0277 \\
0.1062 \ 0.2297 \ 0.2059 \\
0.0010 \ 0.0108 \ 0.1671
\end{pmatrix} \varepsilon_{2to3}^{ij} = \begin{pmatrix}
0.1348 \ 0.0715 \ 0.0099 \\
0.0440 \ 0.0980 \ 0.0770 \\
0.0002 \ 0.0063 \ 0.0884
\end{pmatrix}$$
(A.7)

E:

$$\varepsilon_{2 \to 2}^{ij} = \begin{pmatrix} 0.4151 \ 0.1886 \ 0.0191 \\ 0.1079 \ 0.3129 \ 0.2234 \\ 0.0004 \ 0.0170 \ 0.2660 \end{pmatrix} \varepsilon_{3 \to 2}^{ij} = \begin{pmatrix} 0.1735 \ 0.0920 \ 0.0114 \\ 0.0490 \ 0.1201 \ 0.0966 \\ 0.0002 \ 0.0063 \ 0.0824 \end{pmatrix}$$
(A.8)
$$\varepsilon_{3 \to 3}^{ij} = \begin{pmatrix} 0.2053 \ 0.1086 \ 0.0165 \\ 0.0668 \ 0.1468 \ 0.1238 \\ 0.0004 \ 0.0082 \ 0.0911 \end{pmatrix} \varepsilon_{2 \to 3}^{ij} = \begin{pmatrix} 0.1735 \ 0.0920 \ 0.0114 \\ 0.0490 \ 0.1201 \ 0.0966 \\ 0.0002 \ 0.0063 \ 0.0824 \end{pmatrix}$$
(A.8)

EO:

$$\varepsilon_{2\to2}^{ij} = \begin{pmatrix} 0.4243 \ 0.1944 \ 0.0192 \\ 0.1111 \ 0.3181 \ 0.2275 \\ 0.0004 \ 0.0171 \ 0.2746 \\ 0.2083 \ 0.1122 \ 0.0150 \\ 0.0672 \ 0.1505 \ 0.1229 \\ 0.0003 \ 0.0092 \ 0.0957 \\ \end{pmatrix} \qquad \varepsilon_{3\to3}^{ij} = \begin{pmatrix} 0.1772 \ 0.0869 \ 0.0126 \\ 0.0483 \ 0.1277 \ 0.1051 \\ 0.0003 \ 0.0062 \ 0.0916 \\ 0.0224 \ 0.0127 \ 0.0009 \\ 0.0070 \ 0.0169 \ 0.0120 \\ 0.0000 \ 0.0007 \ 0.0085 \\ \end{pmatrix} \qquad (A.9)$$

$$\varepsilon_{3\to3}^{ij} = \begin{pmatrix} 0.4258 \ 0.1958 \ 0.0181 \\ 0.1085 \ 0.3229 \ 0.2266 \\ 0.0004 \ 0.0173 \ 0.2815 \\ \end{pmatrix} \qquad \varepsilon_{3\to2}^{ij} = \begin{pmatrix} 0.1637 \ 0.0793 \ 0.0118 \\ 0.0470 \ 0.1170 \ 0.0975 \\ 0.0001 \ 0.0063 \ 0.0887 \\ \end{pmatrix} \qquad (A.10)$$

$$\varepsilon_{3\to3}^{ij} = \begin{pmatrix} 0.0004 \ 0.0173 \ 0.2815 \end{pmatrix} \\ \begin{pmatrix} 0.0001 \ 0.00063 \ 0.0887 \end{pmatrix} \\ (A) \\ (A)$$

$$\begin{array}{l} \text{PO:} \\ \varepsilon_{2 \rightarrow 2}^{ij} = \begin{pmatrix} 0.4247 & 0.1940 & 0.0181 \\ 0.1075 & 0.3230 & 0.2256 \\ 0.0004 & 0.0169 & 0.2821 \end{pmatrix} \\ \varepsilon_{3 \rightarrow 3}^{ij} = \begin{pmatrix} 0.1564 & 0.0749 & 0.0108 \\ 0.0439 & 0.1106 & 0.0923 \\ 0.0001 & 0.0066 & 0.0857 \\ 0.0263 & 0.0134 & 0.0015 \\ 0.0000 & 0.0009 & 0.0102 \end{pmatrix} \\ \varepsilon_{3 \rightarrow 3}^{ij} = \begin{pmatrix} 0.4283 & 0.1257 & 0.0157 \\ 0.0751 & 0.1713 & 0.1391 \\ 0.0004 & 0.0093 & 0.1122 \end{pmatrix} \\ \varepsilon_{2 \rightarrow 3}^{ij} = \begin{pmatrix} 0.1591 & 0.0840 & 0.0118 \\ 0.0466 & 0.1167 & 0.0902 \\ 0.0001 & 0.0055 & 0.0736 \\ 0.0000 & 0.0009 & 0.0102 \end{pmatrix} \\ \text{D:} \\ \varepsilon_{3 \rightarrow 3}^{ij} = \begin{pmatrix} 0.2214 & 0.1206 & 0.0163 \\ 0.0725 & 0.1535 & 0.1337 \\ 0.0004 & 0.0089 & 0.0983 \end{pmatrix} \\ \varepsilon_{2 \rightarrow 3}^{ij} = \begin{pmatrix} 0.1353 & 0.0646 & 0.0077 \\ 0.0001 & 0.0055 & 0.0736 \\ 0.0004 & 0.0089 & 0.0983 \end{pmatrix} \\ \varepsilon_{2 \rightarrow 3}^{ij} = \begin{pmatrix} 0.1353 & 0.0646 & 0.0077 \\ 0.0381 & 0.0956 & 0.0704 \\ 0.0001 & 0.0058 & 0.0760 \\ 0.0014 & 0.0096 & 0.1301 \end{pmatrix} \\ \varepsilon_{2 \rightarrow 3}^{ij} = \begin{pmatrix} 0.1994 & 0.1033 & 0.0164 \\ 0.0666 & 0.1544 & 0.1334 \\ 0.0006 & 0.0190 & 0.2838 \\ 0.0006 & 0.0190 & 0.2838 \end{pmatrix} \\ \varepsilon_{2 \rightarrow 3}^{ij} = \begin{pmatrix} 0.1994 & 0.1033 & 0.0164 \\ 0.0033 & 0.0164 \\ 0.0033 & 0.0171 & 0.1211 \\ 0.0006 & 0.0190 & 0.2838 \end{pmatrix} \\ \varepsilon_{3 \rightarrow 3}^{ij} = \begin{pmatrix} 0.1994 & 0.1033 & 0.0164 \\ 0.0035 & 0.0190 & 0.0210 \\ 0.0001 & 0.0005 & 0.0101 \end{pmatrix} \\ \text{(A.14)} \end{aligned}$$

$$\begin{split} \varepsilon_{2 \to 2}^{ij} &= \begin{pmatrix} 0.4188 \ 0.1975 \ 0.0212 \\ 0.1285 \ 0.3360 \ 0.2420 \\ 0.0007 \ 0.0189 \ 0.2907 \end{pmatrix} \varepsilon_{2 \to 2}^{ij} &= \begin{pmatrix} 0.1973 \ 0.1018 \ 0.0156 \\ 0.0659 \ 0.1528 \ 0.1248 \\ 0.0008 \ 0.0900 \ 0.1248 \\ 0.0008 \ 0.0900 \ 0.1248 \\ 0.0088 \ 0.0900 \ 0.1248 \\ 0.0088 \ 0.0127 \ 0.0018 \\ 0.0080 \ 0.0187 \ 0.0147 \\ 0.0001 \ 0.0009 \ 0.0125 \end{pmatrix} \end{split} \text{(A.15)}$$

$$\varepsilon_{3 \to 3}^{ij} &= \begin{pmatrix} 0.4198 \ 0.1934 \ 0.0192 \\ 0.1258 \ 0.3348 \ 0.2353 \\ 0.0006 \ 0.0191 \ 0.2949 \\ 0.0005 \ 0.0073 \ 0.1125 \\ 0.0005 \ 0.0073 \ 0.1125 \end{pmatrix} \end{aligned} \qquad (A.16)$$

$$\varepsilon_{3 \to 3}^{ij} &= \begin{pmatrix} 0.2753 \ 0.1412 \ 0.0244 \\ 0.1020 \ 0.2280 \ 0.1899 \\ 0.0012 \ 0.0127 \ 0.1572 \end{pmatrix} \varepsilon_{3 \to 3}^{ij} &= \begin{pmatrix} 0.0236 \ 0.0156 \ 0.0019 \\ 0.0082 \ 0.0168 \ 0.0146 \\ 0.0001 \ 0.0008 \ 0.0117 \end{pmatrix} \end{aligned} \qquad (A.16)$$
PO:

D:

EO:

G: 0.3958 0.1826 0.0160 0.1412 0.0643 0.0105 $\varepsilon_{2\rightarrow2}^{ij} =$ 0.1171 0.3201 0.2183 0.0422 0.1156 0.0819 $\varepsilon_{2\rightarrow 2}^{ij} =$ 0.0006 0.0195 0.2867 0.0001 0.0073 0.0944 (A.19) 0.3107 0.1623 0.0284 0.0461 0.0256 0.0029 $\varepsilon_{3\rightarrow 3}^{ij} =$ 0.1159 0.2516 0.2147 0.0129 0.0296 0.0279 $\varepsilon^{ij}_{3 \rightarrow 3} =$ 0.0013 0.0129 0.1804 0.0002 0.0003 0.0222

Appendix B

Coefficients for the Jet Rates Calculations

y_{cut}	Α	В	С	y_{cut}	Α	В	С
Е				EO			
0.01	37.24	668.4	334.8	0.01	37.24	84.4	334.8
0.015	29.36	661.0	187.6	0.015	29.36	238.9	187.ß
0.02	24.35	615.4	116.6	0.02	24.35	285.8	116.Ei
0.03	18.10	515.6	52.4	0.03	18.10	289.2	52.4:
0.04	14.21	430.9	9 26.0	0.04	14.21	261.7	26.C)
0.05	11.52	362.7	13.4	0.05	11.52	229.8	13.4
0.06	9.53	307.5	6.99	0.06	9.53	199.9	6.99
0.08	6.75	224.9	1.77	0.08	6.75	149.9	1.77
0.10	4.92	167.0	0.343	0.10	4.92	112.0	0.343
0.13	3.12	108.2		0.13	3.12	71.9	
0.16	1.98	69.9		0.16	1.98	45.2	
0.19	1.21	44.3		0.19	1.21	27.4	
0.23	0.57	22.4		0.23	0.57	12.6	
0.28	0.13	7.12		0.28	0.13	3.17	

y_{cut}	Α	В	С	Ycut	Α	В	С
Р				D			
0.01	37.24	-145.5	334.8	0.01	5.71	154.3	83.98
).015	29.36	47.80	187.6).015	2.32	149.1	46.02
0.02	24.35	121.6	116.6	0.02	0.21	137.1	28.02
0.03	18.10	161.0	52.4	0.03	7.61	113.2	12.10
0.04	14.21	156.6	26.0	0.04	6.00	93.63	5.76
0.05	11.52	141.1	13.4	0.05	4.89	78.13	2.86
0.06	9.53	123.5	6.99	0.06	4.06	65.71	1.43
0.08	6.75	91.04	1.77	0.08	2.90	47.31	0.349
0.10	4.92	65.16	0.343	0.10	2.12	34.56	0.0714
0.13	3.12	37.41		0.13	1.36	21.76	
0.16	1.98	19.42		0.16	.863	13.55	
0.19	1.21	8.12		0.19	.530	8.14	
0.23	0.57	0.048		0.23	.246	3.63	
0.28	0.13	-2.84		0.28	.053	0.658	
РО				G			
0.01	37.24	-130.19257	334.8	0.01	0.0	0.0	0.0
1.015	29.36	73.27059	187.6	3.015	0.0	0.0	0.0
0.02	24.35	149.52218	116.6	0.02	4.51	-595.3	750.9
0.03	18.10	188.56036	52.4	0.03	6.76	-244.0	435.3
0.04	14.21	182.77556	26.0	0.04	1.79	-86.65	277.1
0.05	11.52	166.08951	13.4	0.05	8.24	-6.84	185.7
0.06	9.53	147.45415	6.99	0.06	5.52	35.71	128.5
0.08	6.75	113.41357	1.77	0.08	1.61	69.67	64.76
0.10	4.92	86.26788	0.343	0.10	8.89	74.64	33.60
0.13	3.12	56.64114		0.13	6.07	64.39	12.41
0.16	1.98	36.50658		0.16	4.14	48.97	4.21
0.19	1.21	22.70143		0.19	2.75	34.27	1.23
0.23	0.57	10.71281		0.23	1.45	18.47	0.217
0.28	0.13	2.12175		0.28	.390	5.46	

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(C.4)

(C.5)

APPENDIX C. UNFOLDING R_3^J

the vector \vec{A} as

Appendix C

The Mat hemat ical Formulation of Unfolding R_3^j

As a reminder, the equations to be inverted for R_3^j are, from Equation 9.1:

$$n_2^i = \sum_{j=1}^3 (\varepsilon_{(2\to2)}^{ij}(1-R_3^j) + \varepsilon_{(3\to2)}^{ij}R_3^j) f^j N$$
(C.1)

$$n_{3}^{i} = \sum_{j=1}^{3} (\varepsilon_{(3\to3)}^{ij} R_{3}^{j} + \varepsilon_{(2\to3)}^{ij} (1-R_{3}^{j})) f^{j} N.$$
(C.2)

The index *j* refers to the primordial quark flavor; *i* denotes the flavor tag which selected a particular event. The subscripts $n \rightarrow m$ represent "migration" of an event from having n jets at the parton level to having m jets at the detector level due to acceptance and finite detector resolution. This equation must be solved for the R_3^j 's, which are then used to compute α_s^j .

We begin by defining the vector of inputs \vec{Y} by

$$\vec{Y} = \begin{pmatrix} n_1^2 \\ n_2^2 \\ n_3^2 \\ n_3^1 \\ n_3^2 \\ n_3^3 \\ n_3^3 \end{pmatrix}, \qquad (C.3)$$
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 $\vec{A} = \begin{pmatrix} \sum_{j} f_{j} N \varepsilon_{(2 \rightarrow 2)}^{2j} \\ \sum_{j} f_{j} N \varepsilon_{(2 \rightarrow 2)}^{2j} \\ \sum_{j} f_{j} N \varepsilon_{(2 \rightarrow 2)}^{3j} \\ \sum_{j} f_{j} N \varepsilon_{(2 \rightarrow 3)}^{1j} \\ \sum_{j} f_{j} N \varepsilon_{(2 \rightarrow 3)}^{2j} \end{pmatrix}$

and the vector of desired quantities $ec{r}$ by

 $ec{r} = egin{pmatrix} R_3^1 \ R_3^2 \ R_3^3 \ R_3^1 \ R_3^1 \ R_3^2 \ R_3^3 \ R_3^3 \end{pmatrix} \,.$

$$\begin{pmatrix} -(\varepsilon_{(2\to2)}^{11} - \varepsilon_{(3\to2)}^{11})f_1 N & -(\varepsilon_{(2\to2)}^{12} - \varepsilon_{(3\to2)}^{12})f_2 N & -(\varepsilon_{(2\to2)}^{13} - \varepsilon_{(3\to2)}^{13})f_3 N \\ -(\varepsilon_{(2\to2)}^{21} - \varepsilon_{(3\to2)}^{21})f_1 N & -(\varepsilon_{(2\to2)}^{22} - \varepsilon_{(3\to2)}^{22})f_2 N & -(\varepsilon_{(2\to2)}^{23} - \varepsilon_{(3\to2)}^{23})f_3 N \\ -(\varepsilon_{(2\to2)}^{31} - \varepsilon_{(3\to2)}^{31})f_1 N & -(\varepsilon_{(2\to2)}^{22} - \varepsilon_{(3\to2)}^{23})f_2 N & -(\varepsilon_{(2\to2)}^{33} - \varepsilon_{(3\to2)}^{33})f_3 N \\ (\varepsilon_{(1\to2)}^{11} - \varepsilon_{(2\to3)}^{11})f_1 N & (\varepsilon_{(2\to3)}^{12} - \varepsilon_{(2\to3)}^{12})f_2 N & (\varepsilon_{(3\to3)}^{13} - \varepsilon_{(2\to3)}^{13})f_3 N \\ (\varepsilon_{(2\to3)}^{21} - \varepsilon_{(2\to3)}^{21})f_1 N & (\varepsilon_{(3\to3)}^{22} - \varepsilon_{(2\to3)}^{22})f_2 N & (\varepsilon_{(3\to3)}^{23} - \varepsilon_{(2\to3)}^{23})f_3 N \\ (\varepsilon_{(3\to3)}^{31} - \varepsilon_{(2\to3)}^{31})f_1 N & (\varepsilon_{(3\to3)}^{22} - \varepsilon_{(2\to3)}^{22})f_2 N & (\varepsilon_{(3\to3)}^{23} - \varepsilon_{(2\to3)}^{23})f_3 N \\ (\varepsilon_{(3\to3)}^{31} - \varepsilon_{(2\to3)}^{31})f_1 N & (\varepsilon_{(3\to3)}^{22} - \varepsilon_{(2\to3)}^{32})f_2 N & (\varepsilon_{(3\to3)}^{33} - \varepsilon_{(2\to3)}^{33})f_3 N \\ (\varepsilon_{(3\to3)}^{31} - \varepsilon_{(2\to3)}^{31})f_1 N & (\varepsilon_{(3\to3)}^{22} - \varepsilon_{(2\to3)}^{32})f_2 N & (\varepsilon_{(3\to3)}^{33} - \varepsilon_{(2\to3)}^{33})f_3 N \\ (\varepsilon_{(3\to3)}^{31} - \varepsilon_{(2\to3)}^{31})f_1 N & (\varepsilon_{(3\to3)}^{22} - \varepsilon_{(2\to3)}^{32})f_2 N & (\varepsilon_{(3\to3)}^{33} - \varepsilon_{(2\to3)}^{33})f_3 N \\ (\varepsilon_{(3\to3)}^{31} - \varepsilon_{(2\to3)}^{31})f_1 N & (\varepsilon_{(3\to3)}^{32} - \varepsilon_{(2\to3)}^{32})f_2 N & (\varepsilon_{(3\to3)}^{33} - \varepsilon_{(2\to3)}^{33})f_3 N \\ (\varepsilon_{(3\to3)}^{31} - \varepsilon_{(2\to3)}^{31})f_1 N & (\varepsilon_{(3\to3)}^{32} - \varepsilon_{(2\to3)}^{32})f_2 N & (\varepsilon_{(3\to3)}^{33} - \varepsilon_{(2\to3)}^{33})f_3 N \\ (\varepsilon_{(3\to3)}^{31} - \varepsilon_{(2\to3)}^{31})f_1 N & (\varepsilon_{(3\to3)}^{32} - \varepsilon_{(2\to3)}^{32})f_2 N & (\varepsilon_{(3\to3)}^{33} - \varepsilon_{(2\to3)}^{33})f_3 N \\ (\varepsilon_{(3\to3)}^{31} - \varepsilon_{(2\to3)}^{31})f_1 N & (\varepsilon_{(3\to3)}^{32} - \varepsilon_{(2\to3)}^{32})f_2 N & (\varepsilon_{(3\to3)}^{33} - \varepsilon_{(2\to3)}^{33})f_3 N \\ (\varepsilon_{(3\to3)}^{31} - \varepsilon_{(2\to3)}^{31})f_1 N & (\varepsilon_{(3\to3)}^{32} - \varepsilon_{(2\to3)}^{32})f_2 N & (\varepsilon_{(3\to3)}^{33} - \varepsilon_{(2\to3)}^{33})f_3 N \\ (\varepsilon_{(3\to3)}^{33} - \varepsilon_{(2\to3)}^{33})f_3 N & (\varepsilon_{(3\to3)}^{33} - \varepsilon_{(2\to3)}^{33})f_3 N \\ (\varepsilon_{(3\to3)}^{33} - \varepsilon_{(2\to3)}^{33})f_3 N & (\varepsilon_{(3\to3)}^{33} - \varepsilon_{(2\to3)}^{33})f_3 N \\ (\varepsilon_{(3\to3)}^{33} - \varepsilon_{(2\to3)}^{33})f_3 N & (\varepsilon_{(3\to3)}^{33} - \varepsilon_{(2\to3)}^{33})f_3 \end{pmatrix}$$

The expressions in Equation C.l can be rewritten as

$$\vec{Y} = \vec{A} + \mathcal{B} \cdot \vec{r} . \tag{C.7}$$

If we introduce the weight matrix \mathcal{W} ,

$$\mathcal{W}_{ij} = 1/(\sigma_i \sigma_j) , \qquad (C.8)$$

then we can write the fit χ^2

The matrix \mathcal{B} is given by

$$\chi^2 = (\vec{Y} - \vec{A} - \mathcal{B} \cdot \vec{r})^T \mathcal{W} (\vec{Y} - \vec{A} - \mathcal{B} \cdot \vec{r}) .$$
 (C.9)

Minimizing χ^2 requires

$$\frac{\partial \chi^2}{\partial r} = -2\mathcal{B}^T \mathcal{W}(\vec{Y} - \vec{A} - \mathcal{B} \cdot \vec{r}) = 0 . \qquad (C.10)$$

So, \vec{r} can be obtained by

$$\vec{r} = (\mathcal{B}^T \mathcal{W} \mathcal{B})^{-1} \mathcal{B}^T \mathcal{W} (\vec{Y} - A) .$$
(C.11)

The covariance is given by

$$C = \frac{1}{2} \frac{\partial^2 \chi^2}{\partial r^2} = \frac{1}{2} \frac{\partial}{\partial r} \left(-2\mathcal{B}^T \mathcal{W} (\vec{Y} - \vec{A} - \mathcal{B} \cdot \vec{r}) \right)$$
$$= \mathcal{B}^T \mathcal{W} \mathcal{B}$$
(C.12)

Appendix D

Experience with Accelerator Backgrounds at the SLC

A large part of the recent improvement in SLC luminosity has resulted from better beam diagnostics and machine stability, and this in turn has allowed us to arrive at a more detailed understanding of the problems of experimental backgrounds. Part of this success is also due to the recent strategy of running the damping rings with their tunes uncoupled to produce a vertical/horizontal emittance ratio of approximately 1:4. This has allowed the achievement of a significantly smaller vertical beam size throughout the accelerator without much growth in the horizontal size, resulting in higher instantaneous luminosity and lower backgrounds[210]. Currently, upgrades to SLC and SLD are being carried out or planned which will increase the beam angular divergence and possibly decrease the beam pipe radius. A large amount of effort has recently been expended to compare the background models with observations of actual detector backgrounds in order to assess the potential impact of the upgrades. Here, we summarize recent experience in understanding and controlling backgrounds from the accelerator which have an adverse effect on the SLD detector*.

^{&#}x27;This appendix is an updated version of a 1991 memo entitled "You too can tune backgrounds" combined with results presented at the 5th International Workshop on Next-Generation Linear Colliders at SLAC, October 13-21, 1993. This work was done in collaboration with H. Band, D.L. Burke, S. Hertzbach, T. Usher, W. Ash, H. DeStaebler, F.-J. Decker, P. Emma, S. Hedges, J. Huber, T. Markiewicz, T. Maruyama, N. Phinney, G.D. Punkar, D. Su, and J. Yamartino.



Figure D.1: An overview of the SLC collimation system. A symmetric set of collimators exists in the South Arc and South Final Focus,

D.1 The SLC Optics and Collimation System

For completeness, this section presents a detailed introduction to the collimation system for the SLC and its relationship to the beam orbit parameters along its length. These collimators and orbit characteristics will be referred to throughout this paper. Figure D.l shows a schematic view of the end of the linac, the arcs, and the Final Focus, with the major groups of collimators labeled. Since their installation, the collimators in sectors 29 and 30 of the 30 linac sectors have been the primary collimators for the SLC. These collimators are designed to collimate both beams simultaneously, and thus are placed in alternating pairs at high-beta points for the two beams in each plane, as shown in Figure D.2. Effectively, these collimators cut the beam to a box in the x-y plane, establishing the acceptance of the beam transport system. These are the only SLC collimators systematically designed to define the beam phase-space aperture. These jaws are usually used in a configuration where the primary collimation is done with the apertures in Sector 29, with the Sector 30 jaws used to clean up some of the repopulated beam tails and to catch the spray off the upstream collimators. In the 1993 running with "flat" beams of aspect ratios of approximately 1:4 vertical to horizontal, these collimators were typically at 5 beam

sigma in the horizontal plane and 8 to 10 sigma in the vertical. Although collimation to about 5 sigma is viewed as optimal[211], the tightness of the aperture is limited by the precision and stability of two-beam steering in this segment of the linac.

Along the region of maximum horizontal dispersion at the front of the arcs sit the first series of momentum collimators. These allow the removal of any low or highenergy beam tails that are the result of beam acceleration in the linac and are denoted SL-1 and SL-3. Figure D.3 shows the locations of the movable jaws at the front of the SLC arc superimposed upon the dispersion along the beam orbit. Of these jaws, the low-energy jaw is used the most frequently, as there is often a non-negligible fraction of the production bunch that arrives at the end of the linac sufficiently low in energy as to be a potential problem downstream. In addition, it has often proved helpful to collimate away vertical beam tails here, as the cross-plane coupling introduced by the arc rolls will populate all undesirable corners of phase-space by the time the beam reaches the interaction point. To that end, all of the vertical collimators along the arc have been extremely useful. At about one third of the distance down the arc, the beam curvature changes sign at the "Reverse Bend". As shown in Figure D.4, this provides the opportunity for another zero-dispersion collimation point, as well as another (relatively) high-dispersion point for secondary momentum collimation. The zero-dispersion aperture is provided by the jaws SL-9. SL-10 completes the set as the high-dispersion pair to allow removal of smaller tails that were missed by SL-1 and SL-3 far upstream. Beyond SL-10, the beam proceeds to the Final Focus without encountering another set of collimators.

At the entrance to the Final Focus there lie a series of fixed- aperture protection collimators. Their effect on the beam is seen primarily by the rate of Bethe-Heitler production of muons which can penetrate the SLD detector. These collimators cannot be used to remove a significant part of the beam without flooding the entire detector with noise from hundreds of high energy muons (see below.) The set of jaws labeled C1-X and C1-Y are the first movable collimators in the Final Focus. C1-X has a keyhole shape, and can be moved both horizontally and vertically. Before the linac collimators were installed, C1-X had to function **as** the primary collimator for the entire beam-transport system. Following Cl-X/Y are the last useful momentum slits,



Figure D.2: The Sector 30 collimator layout at the end of the SLC linac. The (a) horizontal and (b) vertical β functions for electrons (positrons) are drawn with solid (dashed) lines. The symbols \Box and \exists represent the positions of the electron and positron collimators, respectively.



Figure D.3: The collimators in the 51-Line, at the beginning of the North Arc. The positions of the collimators as well as the horizontal (η_x) and vertical (η_y) dispersion is also shown. The double pairs of lines represent a complete set of vertical and horizontal jaws. SL-4 only has vertical jaws.



Figure D.4: The collimators in the Reverse Bend, approximately one third of the way down the North Arc. The positions of the collimators as well as the horizontal (η_x) and vertical (η_y) dispersion is also shown. Note that the SL-9 jaws sit at a zerodispersion point. The two sets of SL-9 jaws are horizontal and vertical collimators. The SL-10 jaws are horizontal only.



Figure D.5: The layout of the Final Focus collimation system showing the positions of all of the apertures along the beamline. Movable collimators are shown in their fully-inserted positions. The 10σ beam ellipse in the horizontal and vertical planes is shown for optical configuration used in the 1993 SLC running.

PC-12, and the other movable jaws in the Final Focus, PC-10.5 and PC-10, PC-7.5, and PC-3. As discussed below, these collimators *are* used sparingly, as there is a constant need for compromise in how much beam loss one can tolerate here versus how much beam one needs to remove to run the more sensitive parts of the SLD detector. Figure D.5 shows the lo-sigma beam envelope for nominal 1993 emittances at each of the apertures in the Final Focus.

D.2 Types of Backgrounds Seen in the SLD Detector

There are three major types of accelerator-induced backgrounds visible in the SLD detector. A brief introduction to each is provided here. Detailed discussions of the backgrounds follows.

As mentioned above, muons are created by photoproduction from primary beam particles hitting an aperture upstream of the interaction point (IP). As the Final Focus tunnel needs to be large enough for access to the beamline, there is not sufficient material present to absorb all of these muons. Hence, the muons can travel from their production point along the beamline and can penetrate to the SLD detector.

The SLC produces copious synchrotron radiation from the focusing elements along the beamline. A masking scheme has been implemented around the IP in an attempt to shield the sensitive portions of SLD from the intense synchrotron radiation produced by the final dipoles and quadrupole triplet. Particles in the tail of the beam, however, can create synchrotron radiation which strikes the masks directly and can re-scatter into the detector, causing unwanted noise.

In addition to these other backgrounds, we also have evidence that a number of particles derived directly from the primary beam are showering in or near the detector. These showers produce visible secondary particles that are detected by SLD and are a potential source of backgrounds for physics analyses.

D.2.1 Muon Backgrounds

With the initial commissioning of the SLC came the surprise that there are numerous muons being created in the Final Focus as beam tails are scraped off on the collimators. This led to installation of magnetized iron muon toroids placed in such a manner as to deflect most of the muons away from the detector[212]. Figure D.8 shows the layout of the muon toroids relative to the collimators and detector for the north Final Focus. This approach was successful for the running of the Mark II detector for its years at SLC. Since SLD is a much larger detector, however, it presents a much larger



Figure D.6: A view of the SLD WIC detector, showing the paths of muons produced by the SLC beam and traveling essentially parallel to the beam axis. The WIC endcaps have crossed strips in order to obtain the positions of the muons; only one projection can be seen; (The vertical lines are the vertical WIC strips, which are parallel to the plane of the page.)

target for both deflected and undeflected muons. The Warm Iron Calorimeter (WIC) of SLD is a backing calorimeter and muon detector made of steel plates interspersed with Iarocci tubes, Since most of the background muons are traveling parallel to the beam, they do not point back through the IP along tracks and thus are not confusing for muon identification. They do present a constant signal of tracks in the WIC (see Figure D.6), and thus make the creation of a trigger for low- angle muon pairs from Z decays impossible. Also, muons traversing the Liquid Argon Calorimeter (LAC) leave long continuous strips of hit towers (see Figure D.7), which can be removed but occasionally overlap with clusters of energy from particles in Z events, potentially skewing the determination of event shape variables. During the 1993 run, the mean number of muons seen in the calorimeter per Z event was 0.32, down from 0.7 in 1992. This is



Figure D.7: A view of the SLD LAC, showing the paths of muons produced by the SLC beam and traveling essentially parallel to the beam axis. In the LAC barrel, the muons deposit energy in the liquid argon gaps between the lead tiles; longitudinal streaks of adjacent towers signal the presence of these muons. The LAC has been unrolled into a plane for this perspective.



Figure D.8: The layout of the Final Focus collimation system showing the positions of the collimators, several optical elements, and the muon-spoiler toroids. Note the difference in vertical and horizontal scales for this drawing, and the large size of the SLD detector relative to the tunnel aperture.

a relatively small perturbation on a calorimetric analysis and thus does not degrade the performance of the calorimeter by that much. The number of muon tracks in the WIC endcaps is substantially larger, causing approximately 50% more tracks to be called muons by the muon identification system; since muons pass along the barrel axis, the WIC barrel is virtually unaffected. Infrequently, a muon passes through the Central Drift Chamber (CDC) parallel to the sense wires, depositing hundreds of hits in a small number of jet cells. Since these hits are localized to a relatively small area of the tracking volume, they do not significantly effect the tracking efficiency. Studies of the locations and strengths of the muon toroids[213] were undertaken to ascertain whether or not the optimum position calculated in Ref. [212] for running the Mark II detector was also the optimum for SLD. Possibilities for some improvement over the current placement and strengths of the toroids was seen, but these have yet to be subjected to an experimental trial.

D.2.2 Synchrotron Radiation Backgrounds

In the difficult environment of a linear collider, it is necessary to have as detailed as possible an understanding of the potential for backgrounds due to synchrotron radiation (SR). Since there is no opportunity to use scrapers or other such devices to reduce beam tails on multiple turns of the beam around the machine, the interaction region must be designed to accommodate the additional, wider-divergence SR from the beam tails. This makes simulations of the sources and scattering of SR photons and their interaction with the detector essential to minimize the adverse effects of beam tails on detector performance. Indeed, modeling of the synchrotron radiation from the final triplet followed by EGS4[214] simulation of scattering in the detector has illustrated the expected importance of collimation and the sensitivity to beam tails.

The problem is made more severe, however, by the constant effort to minimize the beam spot sizes at the interaction point in order to maximize the luminosity. In general, the linear spot size σ^* at the IP is given by $\sigma^* = \epsilon/\theta^*$, where ϵ is the beam emittance, and θ^* is the angular divergence of the beam core. A larger angular divergence implies a larger beam size in the final quadrupoles that focus the beam, which generates much more synchrotron radiation from any beam tails that are present. At some point, the luminosity cannot be made larger without making the detector useless. As an introduction, Figure D.9 shows the SLD IP masking on a distorted scale; the overall distance along the **beamline** is 3 m, and the diameter of the M2 aperture is 27 mm. Sensitive detector elements are shown in gray. They include the Central Drift Chamber (CDC), the CCD-pixel Vertex Detector (VXD), the Medium Angle Silicon Calorimeter (MASiC), and the silicon calorimeter Luminosity Monitor





Figure D.9: A diagram of the masking around the SLC IP designed to shield the SLD detector components from the primary synchrotron radiation produced in the focussing elements upstream. Also shown are the paths synchrotron photons must take to be scattered into the active detector elements. Primary synchrotron radiation background (which is never allowed to strike the detector) is shown in solid lines. Secondary and tertiary background is in wavy lines.

(LUM). Each detector has a slightly different response to SR backgrounds.

In considering the impact on physics analyses, the two most sensitive systems are the ones used in tracking, the CDC and the VXD. Excessive numbers of synchrotron photons in either detector produce randomly scattered photoelectric or compton electrons, which leave large local depositions of energy. This can result in very large currents on the CDC high-voltage grid, potentially damaging the chamber or increasing the rate of aging of the wires. On a smaller scale, these low energy electrons leave huge amounts of charge in an area equivalent to one CDC sense wire, saturating the digitization ADCs that are necessary to get pulse-shape information. This renders part of the drift length of the drift chamber cell useless for tracking, as the real track hits are obscured by the huge pulse from the secondary electron. Some capacitive coupling between adjacent cells on the CDC also leads to crosstalk from this large pulse on other nearby channels, creating large blotches of hit cells in the CDC. From timing and charge information, this crosstalk is easily recognizable and can be removed in software offline; it does, however, affect the operation of the tracking trigger. Drift chamber noise is always quantified in terms of "occupancy", which is defined as the percent of wires (out of 5120) with hits in a given beam crossing. Sufficiently large occupancy, then, can result in a significant degradation in the detector's ability to find and fit tracks.

In the vertex detector, these local depositions of energy translate into spurious clusters of pixels which could be linked onto the tracks projected inward from the CDC, potentially resulting in mismeasured track parameters. Since the VXD is essentially a two- layer device, the track linking is not sufficiently protected from linking two background hits to an otherwise well-measured track. This has been identified as one of the largest sources of efficiency loss for linking tracks to the VXD hits; the three-layer VXD3[215] will be considerably less vulnerable to this problem.

Synchrotron Radiation Background Simulations

It is useful to define a maximum divergence angle, θ^{MAX} , which is obtained by multiplying the beam core divergence, θ^* , by the aperture of the collimation system in beam sigmas. Typically, θ^{MAX} is about 5 to 8 times θ^* . This represents the maximum

possible excursion of a particle traveling through the final lenses.

If the total angular spread (including tails) of the beam at the IP, θ^{MAX} , is limited by collimation to a small enough range, the only quadrupole synchrotron radiation (QSR) incident on the detector masks strikes the interior of the downbeam M2 mask (M2B) after passing the IP. With the current masking, a photon from this mask cannot reach the beam pipe without additional scattering, and the probability that a photon incident on this mask results in a photon in the CDC is rather small (6 x 10⁻⁷).

If θ^{MAX}_{IS} increased somewhat, there will also be photons incident on the upbeam M2 mask (M2A). Photons originating in this mask also must scatter again in order to reach the beam pipe, but they are typically more energetic than those from the downbeam mask, and have a larger probability of resulting in a photon in the CDC.

If θ^{MAX} is sufficiently large, photons passing through the 27 mm diameter aperture in M2A can hit the annular M4 mask inside the beryllium beam pipe, or the face of the Medium Angle Silicon Calorimeter (MASiC). Photons striking these surfaces directly have probabilities of 18% and 7% respectively of producing a photon in the CDC. Because the flux of photons passing through the M2A aperture can vary rapidly with radius, this situation is likely to result in high and unstable CDC background occupancy and must be avoided.

The SR and EGS4 models can be combined with a simulation of interactions in the CDC to obtain an estimated CDC occupancy for a given beam profile and opaque collimators. We concentrate on the effects of beam tails or halo because a Gaussian beam with nominal θ^* results in negligible CDC occupancy. In order to understand the geometry and the sensitivity of the QSR background to various parts of beam phase space it has been useful to consider a distribution of electrons uniform in θ^*_x and θ^*_y , and to set a reasonable scale we assume this uniform distribution constitutes 1% of the beam at the IP. In the absence of a measurement of the shape or magnitude of beam tails we use the preceding assumptions as a model of a 'flat tail'.

The general behavior of the calculated CDC occupancy due to QSR is illustrated in Fig. D.10 where we show the occupancy for this model as a function of θ^{MAX} . This particular graph is for the case of round beams with $\theta_x^{MAX} = \theta_y^{MAX}$. The occupancy due to scattering from the downbeam mask (M2B) is relatively small and





Figure D.IO: Measured and calculated CDC occupancies. The solid curve is the calculated occupancy in the central drift chamber due to synchrotron radiation from the quadrupoles as a function of the maximum beam divergence angle at the interaction point, assuming round beams and 1% of the beam spread out in a very wide tail. The dashed curves show the contributions from each mask. Typical operating points are indicated by the arrows below the graph, which show approximate values for each beam in June 1993. Also shown is a measured occupancy due to synchrotron radiation, determined from a sample of random trigger events recorded in June and July 1993.

independent of θ^{MAX} in the range of interest, and occupancy due to QSR incident on the upbeam mask (M2A) is significant for θ^{MAX} larger than 1400 µrad. Beyond 1600 µrad scattering from M4 and the MASiC is significant, and quickly dominates as θ^{MAX} increases further.

D.3 Spurious Track Trigger Background

In order to collect events from μ - and r-pair decays that do not deposit large amounts of energy in the calorimeter, SLD has implemented an event trigger based solely on charged tracks in the CDC. For an event to pass the trigger requirements, there must be at least two charged tracks separated by more than 150". Since the trigger's initial commissioning, it has been clear that large numbers of charged particles are being produced in beam crossings that contain no other electroweak physics processes. Occasionally, the rate of hard charged particle production exceeds 0.1 Hz, and the trigger has to be prescaled or rate-limited to avoid excessive dead time in the detector, causing a potential loss of events from ZO decays. We have undertaken an extensive study of these events to understand their source, and some of the results are presented below.

The charged particles in these events have a transverse momentum spectrum with an average of 400 MeV/c, where the mean momentum for the fastest particle is about 1 GeV/c. The average charged multiplicity of these events is approximately 5 tracks per event. The particles with enough transverse momentum to fire the tracking trigger have been shown through a dE/dx analysis to be predominantly protons, and the ratio of positive to negative particles is 7:1. All of these attributes suggest that it is impossible that these tracks are due to interactions from synchrotron radiation, as the SR spectrum typically has very few photons above 5 MeV. The presence of large numbers of protons also implies energetic interactions with the material surrounding the detector. In addition, since the SR spectrum falls off quite rapidly with energy, there would be a huge occupancy in the CDC due to the lower energy photons which would also scatter into the detector. The scattering process does harden the SR spectrum somewhat, but not nearly enough to produce only the high-energy particles seen in these events.

Figure D.ll shows the position in z relative to the nominal IP of reconstructed charged-track vertices from events which fired the tracking trigger. Approximately 60from the two masks M4 located 12 cm to either side of the interaction point. Single-beam studies have shown that the dip angles of the tracks are inclined along



Figure D.ll: The primary vertex positions along the beam axis in low-multiplicity tracking-trigger events. Some number of two-photon events or low-multiplicity Z^0 events which come from the SLC IP are included in the sample, giving the peak at z = 0. However, most of the events come from interactions in the M4 masks which are located at ± 12 cm from the IP.

the direction of the beam that produced them, but no strong correlation is observed between the track direction and whether they hit upstream or downstream of the IP. An analysis of vertices of the tracks in the transverse plane clearly shows that most of the tracks originate in the top and bottom of the beam pipe/masks, as shown in Figure D.12.

There is also often a correlation between the presence of tracks in the CDC and deposits of large clumps of isolated energy in the luminosity monitor. These also occur mostly in the top and bottom of the luminosity monitor, but do not have a strong correlation to the location of the track vertices found in the same event. The isolated showers in the luminosity monitor have a different longitudinal profile than isolated electrons from Bhabha scattering, as the fraction of energy deposited in the first of the two layers is much larger in the case of the background events, implying



Figure D.12: The x - y primary vertex positions in low-multiplicity tracking-trigger events. Some number of two-photon events or low-multiplicity Z^0 events which come from the SLC IP are included in the sample, but most of the events come from interactions at the top and bottom of the beampipe/masks.

that the shower has already started before the particle hits the luminosity monitor. This is supported by numerous tracking trigger events where the innermost layer of the MASiC has a shower which correlates exactly in azimuthal position to the luminosity monitor shower behind it. Geometrically, a particle entering from outside the detector has to hit the MASiC first to deposit energy into the LUM, suggesting that these energetic particles do not originate at the IP. In addition, there are often multiple high-energy isolated showers in events where the total energy deposited in the luminosity monitor is hundreds of GeV.

D.4 SR Background Measurements

SLD routinely writes to tape a sample of events obtained by triggering on random beam crossings. The analysis of a set of random triggers from June and July 1993 resulted in a value of 2.3% CDC occupancy due to synchrotron radiation. Subtracting 0.5% for radiation from the last SLC bend magnet gives 1.8% due to the quadrupoles, which is indicated on Figure D.IO. Typical SLC operating conditions for a 24-hour period in mid-June 1993 are as follows:

Electrons: $\theta_x^{MAX} = 4.5 \times 350 \mu rad = 1575 \mu rad$ $\theta_y^{MAX} = 9.5 \times 175 \mu rad = 1663 \mu rad$ Positrons: $\theta_x^{MAX} = 4.5 \times 350 \mu rad = 1575 \mu rad$ $\theta_y^{MAX} = 8.8 \times 150 \mu rad = 1320 \mu rad$

Figure D.10 indicates these operating conditions. Note that the collimation seems to result in what are effectively "round" beam tails. This is, as mentioned above, because the beam positions at the collimator jaws at the end of the linac is sufficiently unstable as to require a slightly larger aperture to limit the possibility of large beam losses and the even larger backgrounds that this causes. In this case, the electron beam would be expected to produce most of the CDC occupancy. The occupancy predicted by the simple model agrees surprisingly well with the measured value. However, the θ^{MAX} values above must be taken as very approximate, as the collimation in the *x-y* plane at the end of the linac results in a rectangular cut in z-y position space. Particles up to 50% farther away from the core of the beam can thus remain in the beam if they sit in the appropriate location in space. SLC history plots of beam size and angular divergence show significant time variation, but it is common to operate in this situation, with one beam contributing negligible occupancy, and the other on the edge, contributing several percent background CDC occupancy.

Even though the "effective collimation" calculated above for runs during 1993 gives equal θ^{MAX} values for both the x and y planes, the smaller beam emittance in the vertical has had a significant effect on the severity of backgrounds experienced by SLD. Comparing good running conditions from the 1992 run with runs having

similar θ^{MAX} values from 1993, we have seen a reduction of approximately a factor of 3 in the average CDC occupancy due to SR in the 1993 data relative to 1992. This suggests that our assumption of a 1% flat tail as the source of SR backgrounds to be an overestimate of the tail population, or that the composition of the tails primarily responsible has changed as beam conditions have improved. In addition, due to the smaller angular divergence in 1993, the vertical beam size in the Final Focus is also smaller, which may reduce the amount of beam tail generated there. The optical properties of the final quadrupole triplet also may contribute to this reduction in background. From Figure D.5, one can see that the beam has a large vertical β function closer to the interaction point than in the horizontal plane, and that it decreases rapidly to the IP. This violent focusing is expected to generate a larger amount of synchrotron radiation. This problem would be ameliorated if the beam size in the vertical plane were substantially smaller, as there would be fewer particles that would feel the stronger region of magnetic field within the quad.

Late in the 1993 run two shifts were devoted to background studies with single beams. Analysis of this data is continuing, but some preliminary conclusions can be made. The collimation was not optimized for the positron beam, and because of time constraints most of the data were taken with the electron beam.

SLD data were recorded at four values of θ^* for the electron beam: small θ^* , the θ^* at the end of the 1993 run, θ^* expected in 1994, and a larger value. For the reasons mentioned above, positron data was obtained only for the first two values. The CDC occupancy due to synchrotron radiation, determined from random beam crossing data, is shown for electrons and positrons in Figure D.13 as a function of both θ_y^* and θ_y^{MAX} . In the first case the occupancies appear to differ significantly, but when plotted as a function of θ_y^{MAX} the electron and positron occupancies are quite similar, illustrating the importance of collimation on backgrounds. Figure D.13 also shows the occupancy calculated as for Figure D.10, except that the measured values of θ_x^* , θ_y^* and the actual X and Y collimator settings have been used. The agreement between the data and this simple model is reasonable, but not dramatic.

The data taken at the small θ^* point has been used to measure the SR due only to the final bend magnet before the IP, as this is an important baseline number in the



Figure D.13: The solid and dashed curves are measured CDC occupancies, using random triggers in dedicated background runs, for single electron and positron beams respectively. The dotted curve is the calculated CDC occupancy due to synchrotron radiation, assuming 1% of the beam spread out in a very wide tail, and using measured X and Y values for θ^* and collimation. Otherwise the calculation is identical to that used in Figure D.IO. In the upper graph the curves are plotted as a function of the measured . In the lower graph they are plotted as a function of θ_y^{MAX} , the product of θ_y^* and the collimation in units of the beam σ_y at the collimator. The similarity of the measured occupancy curves in the lower graph illustrates the importance of collimation in determining backgrounds.

calculations of the occupancies due to scattering off the masks, as mentioned above. At this small angular divergence, the SR from the final quadrupoles is expected to be small, allowing us to observe the dipole SR directly. The CDC occupancy due to synchrotron radiation from the last SLC bend magnet has been calculated as 0.25%, and the preliminary measured value for a small angular divergence (235 x 164 μ rad) is between 0.25 and 0.50%. This is a partial measure of the ability of the models to predict CDC occupancy due to SR.

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When the angular divergence of the electron beam was varied, the vertex detector occupancy (# of clusters per CCD) and cluster size (# of pixels per cluster) did not vary significantly between current conditions and those expected after the upgrades. The observed increased occupancy in the CDC and vertex detector was not large enough to generate concerns about backgrounds in the 1994 run.

D.5 Tracking Trigger Backgrounds Studies

In an attempt to understand more about the tracking trigger background, we included in the background studies at the end of the 1993 run single- beam running with reduced and near-zero field in the SLD solenoid. At 10% of the nominal .6 T field, the events look very similar to those seen in the original Mark II vertex drift chamber as shown in Ref. [211]. An event picture of the SLD CDC is shown in Figure D.14. This is a random triggered event from this run. Any track that passes through all layers of the CDC has a momentum of at least 20 MeV at this field, and there are typically many of these tracks per event. Virtually all events have some tracks with an energy greater than 20 MeV. A study of the CDC occupancy as a function of radial wire layer while the B-field is varied has shown that the majority of these tracks probably come from photon conversions within the material between the IP and the drift chamber. Also, the z of the track-beam pipe intersection does not show the dramatic, peaked structure seen in Figure D.ll. The overall CDC occupancy during this running was large but stable, implying statistically that the number of particles causing the underlying event must be large enough not to fluctuate to zero very often.



Figure D.14: An event picture showing the CDC for a random trigger during running with the electron beam only at 10% of nominal solenoid field. Tracks that traverse the entire one-meter radius have a transverse momentum of greater than 20 MeV/c.

The above evidence leads us to the conclusion that the tracking trigger events are caused by multiple off-axis, off-energy primary beam particles showering inside the detector as shown in Figure D.15. The pattern of vertical versus horizontal hits is also comparable to what is expected from a study of the transport of off-energy and off axis particles through the final focusing triplet, as off-axis, off-energy particles are much more strongly over-focused in the vertical plane (see Figure D.5.). The additional evidence from single-beam tests with a lowered solenoid field implies that this process occurs on almost every event, as the spectrum of particles produced is

far too energetic to come solely from synchrotron radiation. The mask M2, whose aperture has an inner radius of 1.35 cm at 1.5 meters from the IP, is the likely source of these scatters. We have also observed events in running single beams where the entire detector is filled with noise, including the upstream luminosity monitor, which implies that the shower must have originated before the beam has arrived at the IP. A shower near M2 is the only thing that could produce enough local energy to generate this signal. EGS studies are underway to confirm that enough energy can be generated in these showers to produce the observed tracks. In the Mark II studies discussed in Ref. [211], the observed backgrounds could be simulated by assuming that between 10 and 100 50 GeV electrons hit a mask upstream of the final quadrupole triplet on every beam crossing. However, the SLD masking scheme is substantially different near the IP, so it is not clear if this will remain as an accepted explanation.

It remains to be seen whether or not the actual source of these off-energy, off-axis particles can be found and eliminated. The inner radius of M2, though it is at 1.35 cm, is still more than 30 beam sigma from the core of the beam. The momentum bandwidth of the Final Focus beam transport system is narrower than this, which implies that these particles must be generated somewhere in the Final Focus itself. A ray-tracing study should probably be done to see if any potential sources can be identified. Interestingly enough, the rate of these events decreased markedly when the angular divergence of the beam was lowered. This points strongly at some aperture in the Final Focus as a cause of these tails, as the beam in this case is much smaller as it travels along the beamline, and thus will be more apt to miss the offending obstruction.

D.6 Conclusions

We have undertaken an extensive study of the background conditions at the SLC using the SLD detector in an attempt to ameliorate trigger problems and assess the affect of potential upgrades. Although the results are preliminary, good agreement is found between models of synchrotron radiation interactions and actual data taken from the detector. Many unsolved questions, especially as to the source of the tracking trigger



Figure D.15: Schematic showing the expected source of the track triggers. The masks are shown on the conventional distorted scale. Primary, but degraded-energy, electrons strike the mask with shower debris accounting for the large energy in the low-angle calorimetry and the production of low-energy particles at large angle.

APPENDIX D. ACCELERATOR BACKGROUNDS AT THE SLC

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background, still remain, however, and will have to be pursued during the upcoming runs.

Appendix E

The SLD Collaboration

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