PRODUCTION AND DECAY OF CHARMED MESONS IN THE UPSILON ENERGY REGION

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PRODUCTION AND DECAY OF CHARMED MESONS IN THE UPSILON ENERGY REGION

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We study the properties of charmed, nonstrange, D mesons produced in continuum $e^+e^-$ annihilations. The $e^+e^-$ collisions were generated in the energy region of the Υ(3S) and Υ(4S) resonances by the Cornell Electron Storage Ring, and were analyzed by the CLEO detector. We make extensive study of the decay topology $D^0, D^+ \rightarrow \overline{K}^0X$, where the $\overline{K}^0$ is observed in the final state $K_S^0 \rightarrow \pi^+\pi^-$. The decay $D^+ \rightarrow K^-\pi^+\pi^+$ is also considered. We analyze the hadronization process through which charmed quarks become charmed hadrons. We measure the probability that a charmed meson will be produced in a state of non zero angular momentum. A comparison of the $D^0$ and $D^+$ fragmentation distributions is made to different fragmentation models, and to other charmed hadrons. A measurement of the relative production of several $D^+$ decay modes is made, and the total $D^0$ and $D^+$ cross sections are estimated. Techniques are presented for the analysis of satellite mass peaks and $D^+ \leftrightarrow D_s^+$ reflections. We search for nonspectator effects in the weak decays of charmed mesons, and for evidence of hadronic final state interactions in $D^0$ decay.
BIOGRAPHICAL SKETCH

David Russell Perticone, son of Frank and Lena B. Perticone, was born and raised on America’s North Coast in the city of Rochester, New York. He was educated in the West Irondequoit public school system. During the summer which followed his junior year at Irondequoit high school, he spent six weeks traveling through Europe and the Soviet Union as a student ambassador on the People to People program. In his senior year, he was elected President of the Robert Fulton Society. He also was a four year participant in the interscholastic tennis program. In the fall of 1979 he entered Rensselaer Polytechnic Institute as a chemical engineering major, only to emerge four years later with a degree in physics. During the summer of 1981 he was selected to participate in the Summer Science Program of the Stanford Linear Accelerator Center. It was there that he received his introduction to experimental high energy physics and drift chamber technology. He returned to SLAC the following summer. In June of 1983, shortly after the completion of his B.S., he began graduate study at Cornell University. At that time he joined the Laboratory of Nuclear Studies. In his first year of graduate school he was supported by a full graduate research assistant-ship, the only Cornell physics graduate student in recent history to hold such an appointment. He remained a G.R.A. until completion of his degree, with a brief hiatus as a half time teaching assistant during the spring of his third year. An avid reader, his recreational diversions include ice hockey, croquet, and bicycling. Upon completion of his Ph.D. program he has accepted a Research Associate position with the University of Minnesota.
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It is in many ways ironic that an experimental high energy physics paper bears the names of a hundred or so physicists, while a thesis from the same experiment bears the name of only one. Many people have contributed to this work, here I shall offer an incomplete list.

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CHAPTER 1
INTRODUCTION

After all a man cant only do what he has to do, with what he has
to do it with, with what he has learned, to the best of his judgment.
And I reckon a hog is still a hog, no matter what it looks like. So here
goes,
William Faulkner, Old Man

Man’s desire to understand the behavior of nature at its most fundamental
level has existed for several millennia. As technology has advanced, ever finer
distance scales have been revealed. Accompanying these revelations have been
revolutions in what has been considered fundamental, as displayed in Figure
1.1. Revolutions may be characterized by the observation of a plethora of states
considered to be fundamental, which are ultimately explained as being com-
posed of a newer set of constituents. The theory currently in vogue is referred
to as the $SU(3) \times SU(2) \times SU(1)$ Standard Model of the Strong and Electroweak
interactions. It describes the interactions of twelve fundamental fermions (mat-
ter fields) via the exchange of gauge bosons belonging to three distinct forces
(strong, electromagnetic, weak). It is an anomaly free, fully renormalizable the-
ory, with an impressive record of performance. Yet, in its minimal form it has
a total of 18 parameters; three gauge couplings, two parameters for the Higgs
sector, nine masses for the fundamental fermions (three are expected to be mass-
less), and three angles plus a phase which describe quark mixing in the weak
interaction. The centerpiece of the standard model is the Higgs boson, which
breaks the electroweak symmetry and provides the mass of the weak interme-
diate vector bosons and the fundamental fermions. The Higgs remains unob-
Figure 1.1: Illustrative sketch of the approximate number of elementary particles versus time.

...erved experimentally, and many of the parameters are not predicted by the theory. These must be determined experimentally (the Higgs mass is among them). We briefly recount the salient features of the fundamental fermions, and the three forces that govern their dynamics. Detailed expositions on the standard model can be found in several textbooks [1, 2].
Table 1.1: Quark and Lepton Doublets

| Generation | Quarks | Leptons |
|------------|--------|---------|
| 1          | \(u_L\) | \(\nu_e\)  |
| 2          | \(d_L\) | \(\nu_\mu\) |
| 3          | \(c_L\) | \(\nu_\tau\) |
| 3          | \(s_L\) | \(\nu_\mu\) |
| 3          | \(t_L\) | \(\nu_\tau\) |

1.1 Matter Fields

Symmetry pervades our classification of the fundamental fermions. The twelve are broken into two groups of 6, being the leptons and the quarks. Of the six in each family there are two types of particles which are distinguished by electric charge (-1,0 for leptons, \(\frac{2}{3}\), \(-\frac{1}{3}\) for quarks). It is then convenient to group the \((\frac{-1}{0}), (\frac{2}{3})\) pairs into generations. For the left-handed states each generation forms a natural isodoublet. This hierarchy is summarized in Table 1.1. The electrically neutral leptons are called neutrinos. Since neutrinos are thought to be massless, right handed states are only allowed for anti-neutrinos. Thus the charged leptons form right handed singlets, as do both types of electrically charged quarks. Despite the topological similarity, two rather profound differences exist between the quarks and leptons:

1. The first great distinction between the two families is that quarks carry color charge and are thus allowed to participate in the strong interaction. Since there are three types of color charge, there are in principle three times as many quarks as leptons.

2. The second is that each lepton generation possess a distinctive identity, such that the total number of leptons \((L_i = 1)\) and anti-leptons \((L_i = -1)\) from each doublet \((i = 1, 2, 3)\) are separately conserved for all known interactions. Quarks have a much weaker identity. The first generation forms
an isospin doublet, while each of the remaining four flavors each carry their own quantum number. The quark flavor quantum numbers are only preserved in the strong interaction.

We shall now examine the ramifications of these differences.

### 1.2 Gauge Fields

| Force    | Symmetry | Quanta          | Mass      | Coupling | Charge |
|----------|----------|-----------------|-----------|----------|--------|
| Strong   | $SU(3)$  | 8 gluons (g)    | No        | $\alpha_s(Q^2)$ | color  |
| Electroweak | $U(1)$   | Photon ($\gamma$) | No        | $e$      | No     |
| Electroweak | $SU(2)$  | $Z$             | 92.6 GeV  | $g = e/\sin \Theta_W$ | No     |
| Electroweak | $SU(2)$  | $2 W^\pm$       | 81.8 GeV  |          | electric |

Identical fermions are prevented from occupying the same spacetime point, integer spin (bosons) particles suffer no such restriction and are thus naturally associated with fields. Particles are subjected to a given force by exchanging the field quanta of that particular force. Table 1.2 presents the properties of the Gauge bosons, and Figure 1.2 collects the Feynman rules for these interactions. The coupling constants for these forces are very different, obeying the approximate relationship

$$\text{Strong} : \text{EM} : \text{Weak} \approx 1 : 137 : 10^{-5}$$

The dominant theme in the Standard Model is local gauge (or phase) invariance. The Lagrangian for any interaction is the particle physics analog of DNA. It contains all the information and possible reactions for a given interaction. All Lagrangians are required to be invariant under the transformation $e^{-ia(x)}$. It is
precisely the terms which arise in the Lagrangian to preserve local gauge invariance which insure renormalizability and give each force its distinctive character.

### 1.2.1 Strong Interaction

Gluons themselves possess color charge, thus not only transmit the strong force between colored objects, but are capable of interacting with themselves. The strong interaction is Non-Abelian, and part of this is manifested in that the coupling constant of the strong force is not constant, but changes with energy or $1/distance$. To lowest order the coupling constant is:

$$
\alpha_s(Q^2) = \frac{12\pi}{(33 - 2n_f) \ln(Q^2/\Lambda^2)}
$$

where $n_f$ is the number of quark flavors and $\Lambda$ is the QCD scale parameter. This causes an effect referred to as asymptotic freedom. When quarks are probed at very short distance they appear almost free, however as quarks are separated
the force becomes so great that it is energetically favorable to pop another $q\bar{q}$ pair from the vacuum. Colored objects (i.e. quarks) have never been directly observed. They always combine in two’s ($q\bar{q}$) and three’s ($qqq$) to produce colorless final states. These hadrons are very complex structures, consisting of the valence quarks, the gluons which bind them, and virtual ($q\bar{q}$) pairs (or sea quarks). Deep inelastic scattering experiments, which measure the momentum fractions of composite objects, found that about 50% of a proton’s momentum is carried by neutral particles, substantiating the role of gluons as the hadronic “glue.” Figure 1.3 shows a calculation of the momentum fraction of the com-

Figure 1.3: Momentum fraction of proton components versus $Q^2$ of the probe.

ponents of a proton versus $Q^2$ of the probe, done by Eichten [3]. The theory of color interactions is known as quantum chromodynamics (QCD).
1.2.2 Electroweak Interaction

In a master stroke, the weak and electromagnetic forces were unified in the same fashion as the electric and magnetic forces were unified by Maxwell. This is the so called Weinberg-Salam $SU(2)_L \times U(1)_Y$ theory, which merited a Nobel prize. Here the four electroweak vector bosons are split, where the three weak vector bosons (two charged and one neutral) acquire a mass on order 100 GeV, and the photon remains massless. Because the weak bosons are massive, the range over which the weak force can be transmitted is greatly reduced.

The exchange of the charged $W^\pm$ bosons governs the transmutations of the quarks and leptons. The most general transformation consists of a quark undergoing a flavor and charge changing transition by emitting a virtual $W$ boson. The $W$ then fragments into either a particle antiparticle pair from a lepton doublet, or quark pair. The pivotal factor is that no constraint such as lepton number exists for weak transitions of quarks. Thus a $\frac{+2}{3}$ quark could transform into any lighter $-\frac{1}{3}$ quark. Similarly, the $W$ may also fragment into quarks which do not belong to the same doublet. The quark “mass” eigenstates are not the eigenstates of the weak interaction. The weak eigenstates are obtained by the rotation of the $-\frac{1}{3}$ quarks

$$
\begin{pmatrix}
    d' \\
    s' \\
    t'
\end{pmatrix} =
\begin{pmatrix}
    V_{ud} & V_{us} & V_{ub} \\
    V_{cd} & V_{cs} & V_{cb} \\
    V_{id} & V_{is} & V_{ib}
\end{pmatrix}
\begin{pmatrix}
    d \\
    s \\
    t
\end{pmatrix}
$$

$V$ is a $3 \times 3$ unitary matrix often referred to as the K-M matrix, after Kobayashi and Maskawa [4] who first generalized the quark mixing matrix to three generations. It contains three angles and a phase for three generations of quarks. The parameterization of Chau and Keung [5] is shown in Figure [1,4]. It has the
Figure 1.4: Quark mixing matrix in the parameterization of Chau and Keung. The three angles are x, y, and z. $s_x = \sin(x)$ and $c_x = \cos(x)$.

approximate value

$$
V_{ij} = 
\begin{pmatrix}
1 & s & s^3 \\
-s & 1 & s^2 \\
 s^3 & -s^2 & 1
\end{pmatrix}
$$

$s \sim 0.23$

Off diagonal transitions occur much less frequently, and are historically referred to as Cabbibo suppressed. This fortunate feature of weak decays lends a great deal of richness to the decays of heavy quarks. Flavor changing neutral currents, however, are excluded by the unitarity of the K-M matrix.

1.3 The Role of Charm

The prediction of charm and the four quark mixing (Cabbibo) matrix were important progenitors to the three generation picture of the standard model. Initially, three quarks $u, d, s$ formed an approximate flavor symmetry, and were used to classify the known states of the day. At that time the first two lepton generations were known, and the lepton quark asymmetry was both technically and aesthetically displeasing to theorists. It became apparent that certain kaon decays, such as $K_L^0 \to \mu^+\mu^-$ were notoriously absent. Unitarity of the quark mixing matrix was first prognosticated by Glashow, Iliopolos, and Maiani.\footnote{[6]}
who completed the Cabbibo mixing matrix, casting it in the form:

\[
\begin{pmatrix}
\tilde{d}' \\
\tilde{s}'
\end{pmatrix} =
\begin{pmatrix}
\cos \theta_c & \sin \theta_c \\
-\sin \theta_c & \cos \theta_c
\end{pmatrix}
\begin{pmatrix}
d \\
s
\end{pmatrix}
\]

This provided a clean mechanism to remove \( K^0_L \rightarrow \mu^+ \mu^- \) (Figure 1.5). This prediction was confirmed some four years later by the observation of a \( c \bar{c} \) bound state by groups at SLAC and Brookhaven. This represented the first major step in proliferation of the heavy quarks, which now totals two and a third is expected. The observation of the charmed quark was an event of such magnitude that the experimentalists who made the discovery were awarded a Nobel prize. This sparked off a flurry of activity in the study of charmed particles, which still continues briskly today.

Figure 1.5: Cancellation of the decay \( K^0_L \rightarrow \mu^+ \mu^- \) via the GIM mechanism.
CHAPTER 2
THEORY OF THE PRODUCTION AND DECAY OF CHARMED MESONS

Here we review the progress made in understanding how charmed quarks hadronize and eventually decay. These are two entirely different mechanisms, and this is reflected in the paths taken in our theoretical understanding of these processes. Quark hadronization is theoretically intractable, and relies on extensive computer modeling and phenomenology. The theory of the decays of charmed mesons enjoys a (limited) quantitative stature, however the predictive power of these theories has been poor. In both instances, theoretical advancement has relied heavily on experimental analysis.

2.1 Fragmentation Fundamentals

Experimentally, we have not observed free quarks. Fragmentation theory aims to describe the process through which an uncombined quark evolves into the hadrons which we observe in our detectors. This section aims to address the theoretical progress that has been made in understanding heavy quark fragmentation. It will emphasize the physics of charm quark fragmentation to charmed mesons in $e^+e^-$ collisions.

One of the differences among the flavors is the quark mass. While this is in general an ill defined concept since quarks cannot be observed in isolation, qualitative trends still exist. The estimated constituent (quark + surrounding gluons) quark masses obey the approximate proportions:

$$m_u : m_d : m_s : m_c : m_b \rightarrow 1 : 1.4 : 5.1 : 14.9$$

The first three flavors are referred to as ‘light’ quarks or $q$, and the rest as ‘heavy’
quarks or $Q$. The most important effect of mass occurs during quark anti-quark production in the color field. When examined from a “thermodynamic” standpoint [7] quark pair production is governed by the expression

$$\text{rate} \propto \exp \left( -\frac{2\mu c^2}{kT} \right)$$

where $\mu$ is the quark mass and $T$ is the universal temperature, corresponding to approximately 160 MeV. This leads to the hadron production ratios [8]

$$\pi : K : D : B \simeq 1 : 0.04 : 10^{-6} : 10^{-15}$$

This suggests that heavy $Q\bar{Q}$ pairs are almost never predicted in the hadronization process. Our experience with the electromagnetic production of quark pairs points toward more of a threshold effect, where quark electric charge and not mass is the dominant factor in the fragmentation of a virtual photon. Gluons hadronize independently of quark flavor, and a sufficiently hard gluon spectrum could be a prominent source of heavy quark production. Charmed particle cross sections in $p\bar{p}$ collisions have been much larger than originally anticipated. The inclusive charmed particle cross section [8] $\sigma_{\text{inc}}$ ranges from $\simeq 5 \times 10^{-1}$ mb at $s = 2 \times 10^2$ GeV$^2$ to $\simeq 1$ mb at $s = 1 \times 10^4$ GeV$^2$. This is to be compared with a pion cross section that remains essentially constant at $10^2$ mb. While this effect is not fully understood, QCD flavor-annihilation and flavor-excitation processes are the most likely explanation. At CESR energies, it is highly unlikely that radiative gluons are capable of producing $c\bar{c}$ pairs, and it is a physically reasonable to assume that charmed quarks are only produced electromagnetically.

The variation in mass will also introduce kinematical differences between light and heavy quark fragmentation. Attaching a light quark to a moving heavy quark will not decelerate it very much. Thus hadrons formed from a primary heavy quark are expected to be ‘hard’, retaining a large fraction of the
original quark momentum. This is in contrast to light quark fragmentation, which is ‘soft.’

A particularly convenient way to observe hadronization is the reaction

\[ e^+ e^- \rightarrow \gamma^*, Z^0 \rightarrow f \bar{f} \]

where \( f \) is a fermion. Electromagnetically, q\( \bar{q} \) pairs are produced relative to muon pairs by \( 3e_q^2 \) where \( e_q \) is the quark charge. Thus heavy quarks are produced, above threshold, as often as light quarks with the same charge.

If heavy quarks are not produced during hadronization, a reconstructed heavy hadron can be directly compared to the primary Q that may have caused it. One exception to this, of course, would be heavy quarks produced in the cascade decays of even heavier quarks. The kinematics of these decays are well known and the contribution to the fragmentation data of these hadrons can be sorted from those of primary Q\( \bar{Q} \) pairs. Heavy quark production is not limited to \( e^+ e^- \) collisions. Important early results in charm fragmentation were gleaned from \( \nu N \rightarrow \text{charm} \) reactions and lepton production \[9\] but these will not be discussed here.

Fragmentation of quarks into hadrons is not well understood. The principal reason for this is that the process is very violent, and takes place where \( \alpha_s(Q^2) \) is large. We are therefore unable to apply perturbative QCD to do any calculations. We have circumvented this by proposing a probability function \( D_{qH}(z) \), which describes the probability that a hadron H will be found in a jet formed by the fragmentation of q with a fraction z of the original quarks energy-momentum. These are the so-called fragmentation functions. A number of different forms for these fragmentation functions have been put forth, based on a variety of ideas. This section will attempt to illuminate the various approaches that have used to
describe fragmentation. Some of the fundamental tools needed to understand fragmentation will also be described.

There are three scenarios which have been put forward to describe fragmentation. These are Independent Fragmentation (IF), String Fragmentation (SF), and, more recently, Cluster Fragmentation (CF). IF was proposed by Field and Feynman [10] as a first guess model of jet formation. Its central theme is that all partons fragment independently of each other. The SF model, based on an idealization of the color field, was due initially to Artu and Messier [11]. Their work was highly embellished by the Lund [12] group and provides the foundation for many of the Monte Carlos currently employed in high energy physics. A promising development, CF is based on leading log QCD. Fragmentation takes place as a QCD shower generated as an off shell parton comes on shell.

2.1.1 Independent Fragmentation

IF is a relatively uncomplicated idea. It is described by the diagram shown in Figure 2.1. A $q_1\bar{q}_1$ pair is produced in the color field of uncombined quark $q_0$ moving through spacetime. The initial quark combines with the appropriate anti quark to form a meson ($q_0\bar{q}_1$) while the remaining quark $q_1$ is left uncombined to continue the fragmentation process. A serious problem with this concept is that there is eventually one quark left over, which creates difficulties in terms of flavor and four momentum conservation. This is dealt with in Monte Carlo at the end of event generation by globally imposing energy and momentum conservation throughout the entire event. Despite its counter intuitive nature, IF was unsophisticated, available and became popular. A number of its
deficiencies, particularly the way it handles gluons, have made it less popular today in light of the better models available.

### 2.1.2 String Fragmentation

Because the field quanta of QCD, the gluons, carry color charge, the field lines of the color field collapse into a flux tube or ‘string’ between quarks. String models are defined in 1+1 dimensions \( i.e. \) 1 time and 1 space or 1 energy and 1 momentum. The primary parameter is the string tension constant \( \kappa \), which has the approximate value \( \kappa \approx 0.2 \text{ GeV}^2 \approx 1 \left( \frac{\text{GeV}}{\text{fermi}} \right) \). The key equations of SF are

\[
\frac{dp}{dt} = \pm \kappa
\]

\[
\frac{dE}{dx} = \kappa
\]

where the + – refer to the string pulling the parton to the right and left respectively. These two equations can be integrated and combined with the Lorentz
invariant expression \( E^2 = p^2 + m^2 \) to yield
\[
(x - x_0)^2 - (t - t_0)^2 = \left(\frac{\mu}{k}\right)^2
\]
(1)

Thus, massive quarks move along hyperbolas in spacetime with light cone asymptotes. ‘Massless’ quarks would move along the light cone. There are several important questions that need to be addressed: what determines how and when the string breaks, and is there any relation between the color field and the mass of the hadron formed? The first question is a pivotal issue with SF, and we will see later how different approaches to this question lead to slightly different fragmentation functions.

### 2.1.3 Cluster Fragmentation

A third and relatively new approach to fragmentation is the called Cluster Fragmentation. It is based on leading log QCD and the Altareli-Parisi formalism. Each branching is considered to be an independent event, and the parton comes more on shell at each branch.

The classical branching probabilities are given by
\[
P_{q \to qg} = \frac{4}{3} \left( \frac{1 + x^2}{1 - x} \right)
\]
\[
P_{g \to gg} = 3 \left[ \frac{1 - x}{x} + \frac{x}{1 - x} + x (1 - x) \right]
\]
\[
P_{g \to qq} = \left[ x^2 + (1 - x^2) \right]
\]

where \( P_{a \to bc} \) represents the probability that \( b \) will retain a fraction \( x \) of \( a \)'s momentum. Work has been done in this area by Field and Wolfram \[13\], Gottschalk \[14\] and the most popular by Webber \[15\]. The shower is terminated at the
point which the parton reaches the appropriate quark mass or cutoff $Q_o$ for gluons (to prevent infrared divergences). At this point all the gluons are split into $q\bar{q}$ pairs and color flow is regulated to produce colorless clusters which decay through phase space. The Webber model has the advantage of requiring only a few parameters: $\Lambda_{QCD}, M_{max}, m_u, m_d, m_s, Q_o$. Additional features of the Webber model include successively reducing the opening angles of parton emission to account for leading interference effects, and decaying clusters exceeding some $M_{max}$ as a string. Initially, $c$ and $b$ quarks decayed weakly before the clusters were formed. Now $c$ quarks are kept in the shower and form charmed clusters which decay into $D^*$ and a $\pi$ or $K$. While this is not directly suited to predicting charm yields, charm fragmentation is generated directly via the cluster decay algorithm. While Cluster Fragmentation is still in the developmental phase, it has shown itself able to compete with SF and IF. The purpose of this cursory treatment of CF is to demonstrate that we may be finally approaching the point where fragmentation is able to proceed naturally out of our Monte Carlo models without the imposition of phenomenological kludges.

2.1.4 Reciprocity

The reciprocity relation is often invoked as a check on fragmentation functions. It is essentially a boundary condition that proposes

$$D_q^H(z \rightarrow 1) \Rightarrow f_H^{D}(x)_{z \rightarrow \frac{1}{x} \rightarrow 1}$$

What this means is that the fragmentation function of a hadron containing virtually all of the original quarks momentum should be equal to the structure function of a hadron where that quark posses a momentum fraction that approaches
one. This can be understood in that as a quark obtains a large momentum fraction, the remaining quarks in the hadron must compensate for this by giving up their momentum. As \( x \to 1 \), the other valence quarks must be almost at rest. If this reaction is reversed in time, one obtains fragmentation.

If reciprocity holds, the appropriate question is how the structure functions behave as \( x \to 1 \)? Structure functions are on a somewhat better theoretical footing than fragmentation functions, but the answer is still uncertain. The ‘standard’ textbook formula \([2][p. 200]\) is

\[
f_H^q(x)_{x \to 1} = (1 - x)^{2n_s - 1}
\]  

where \( n_s \) is the number of spectator valence quarks. This relation predicts \((1 - x)\) for mesons and \((1 - x)^3\) for baryons. An alternate \([16]\) form is

\[
f_H^q(x)_{x \to 1} = (1 - x)^{2n_s - 1 + 2\Delta S}
\]  

The \( \Delta S \) term is the absolute difference between the initial hadron spin and the quark spin \( (\frac{1}{2}) \). This conflicts with the above expression in the prediction for mesons, yielding \((1 - x)^2\). As we shall see, almost all the fragmentation functions show a \((1 - z)\) behavior as \( z \to 1 \), agreeing with the first prediction.

### 2.1.5 Odds and Ends

Finally, we present three items that will be needed in order to coherently proceed to the next section. These are the meaning of \( z \), the light-cone variables, and some elementary relations of fragmentation functions.

Fragmentation functions are parametrized in terms of ‘\( z \)’, but what exactly is \( z \)? It is meant to describe the fraction of energy-momentum or momentum
of the primary quark retained by the hadron in the fragmentation process. Two definitions that appear in the literature are

\[ z^+ \equiv \frac{(E + p_\parallel)_{\text{had}}}{(E + p)_{\text{quark}}} \]

and

\[ z_p = \left( \frac{P_{\text{had}}}{P_{\text{quark}}} \right) \]

The first definition is Lorentz invariant for boosts along the quark direction, hence is more desirable theoretically. This is true because \( L(E \pm p) = f^\pm(\beta)(E \pm p) \), where \( f^\pm(\beta) \) is a constant that is a function of the boost parameter \( \beta \). These constants trivially cancel in numerator and denominator. We adopt the convention that \( z \) will refer to the first definition.

Unfortunately, the kinematical variables \( E \) and \( p \) of the quark are unavailable to the experimentalist. Fragmentation functions are measured in terms of \( x \), where \( x \) also has two definitions

\[ x_E \equiv \left( \frac{E_{\text{had}}}{E_{\text{beam}}} \right) \]

\[ x_p \equiv \left( \frac{P_{\text{had}}}{\sqrt{E_{\text{beam}}^2 - m_{\text{had}}^2}} \right) \]

Both \( x_E \) and \( x_p \) have been used by experimentalists, although \( x_p \) has the advantage of ranging from 0 to 1 for all experiments. The variables \( x \) and \( z \) are not equivalent. In general \( z \geq x \) because perturbative QCD gluon radiation and initial state photon radiation tend to make \( E_{\text{quark}} \leq E_{\text{beam}} \).

The light cone formalism is a handy tool when working with fragmentation.

\[ x^\pm \equiv x^0 \pm x^3 \]

\[ p^\pm \equiv p^0 \pm p^3 \]
$p_\perp$ is absorbed in the mass term. As mentioned earlier these combinations transform trivially under Lorentz transformations, making frame transformations straightforward. They have the added convenience that

$$z \equiv \frac{(E + p_\parallel)_{\text{had}}}{(E + p)_{\text{quark}}} \equiv \left( \frac{p_{\text{had}}^+}{p_{\text{quark}}^+} \right)$$

when $p_\perp$ is neglected. Thus, the light cone variables lend themselves to a more natural expression of $z$. One is still faced with a dilemma as to what the denominator should be. It has been suggested [17] that

$$p_{\text{quark}}^+ \approx E_{\text{max}} + p_{\text{max}} = E_{\text{beam}} + \sqrt{E_{\text{beam}}^2 - \mu^2}$$

however it is not obvious that this is a physically reasonable assumption.

Measurements of fragmentation functions often appear as plots of $s \left( \frac{d\sigma}{dz} \right)$ vs. $(X_E$ or $X_p)$. The master equation for fragmentation in $e^+ e^-$ collisions is [18]

$$\left( \frac{1}{\sigma_{\text{had}}} \right) \frac{d\sigma (e^+ e^- \to HX)}{dz} = \frac{\Sigma_i e_i^2 \left[ D^H_i(z) + D^{\bar{H}}_i(z) \right]}{\Sigma_i e_i^2}$$

where $i$ sums over the quark flavors participating in the reaction. The fragmentation functions of both $i$ and $\bar{i}$ are summed over if the hadron H could be produced in the jets of $q_i$ and $\bar{q}_i$. Applying $\sigma_{\text{had}} = \sigma_{\mu\mu} \left( 3\Sigma_i e_i^2 \right)$ and $\sigma_{\mu\mu} = \left( \frac{4\pi\alpha^2}{3s} \right)$, this can be recast into

$$s \left( \frac{d\sigma}{dz} \right) = 4\pi\alpha^2 \Sigma_i \left[ D^H_i(z) + D^{\bar{H}}_i(z) \right]$$

which explicitly shows the proportionality of $s \left( \frac{d\sigma}{dz} \right)$ to the fragmentation functions. For the production of heavy mesons $H_Q = (Q\bar{q})$ the form is slightly different. $H_Q$ can only be only be found in the debris of Q since heavy quarks are not produced in the color field. The first equation becomes

$$\left( \frac{1}{\sigma_{\text{had}}} \right) \frac{d\sigma (e^+ e^- \to H_Q X)}{dz} = \frac{3e_Q^2 D^H_Q(z)}{\Sigma_i e_i^2}$$
2.2 Fragmentation Functions

FF arise out of our basic inability to calculate hadronization from perturbative QCD. None of them represent a great tribute to theoretical physics, and actually only one has received considerable attention from experimentalists. We analyze five different functions; two based on IF, two on SF, and one derived from the reciprocity relation. An important fact to bear in mind is that these functions describe the initial fragmentation of a primary heavy quark Q into a heavy hadron $H_Q (Q \bar{q})$ containing Q.

The gross qualitative features of heavy quark fragmentation were predicted by several theorists in the late 70’s. Bjorken [19] on the basis of simple kinematics, arrived at

$$< z_p > \sim 1 - \left( \frac{1 \text{GeV}}{\mu} \right)$$

where $\mu$ is the quark mass and 1 GeV is a constant chosen for ‘didactic convenience.’ His argument was that ordinary hadrons were produced with $\gamma$ less than or the order of that of the primary Q. Since $p = \gamma mv$, the heavier objects would take a larger fraction of the available momentum.

2.2.1 Functions Based on Reciprocity

One of the first charm fragmentation functions was that of Kartvelishvili [20]. Their efforts were essentially directed toward computing the structure function of a heavy charmed meson using the Kuti-Weisskopf [21] model, and connecting that to the fragmentation function via reciprocity. They arrived at the structure
function:

\[ f_{H_c}^c(z) = \frac{\Gamma(2 + \gamma - \alpha_c - \alpha_q)z^{-\alpha_c}(1 - z)^{\gamma - \alpha_q}}{\Gamma(1 - \alpha_c)\Gamma(1 + \gamma - \alpha_q)} \]

Where \(\gamma\) is a constant that equals \(\frac{3}{2}\) and \(\alpha_c\) is the intercept of the Regge trajectory for the c quark. \(\alpha_q\) is the intercept of the light quark Regge trajectory, set to \(\frac{1}{2}\) based on the \(\rho, \omega, A_2\), and \(f\). This expression can be used to calculate the average momentum carried by the valence quarks

\[ <z_c> = \frac{1 - \alpha_c}{2 + \gamma - \alpha_c - \alpha_q} \]

Their rational was that since the structure function peaks at high \(z\), the structure function can be set equal to the fragmentation function at all \(z\). As can be easily seen from the expression for the structure function, the fragmentation function becomes

\[ D_{H}^c(z) = Nz^{\alpha_q}(1 - z) \]

The parameter \(\alpha_c\) was unknown, the initial guess was -3. In a later and often referenced paper \cite{22} they choose \(13.44z^{2.2}(1 - z)\) without explanation. It certainly follows, though, from the logic of their earlier work. Claiming that reciprocity is valid at all \(z\) (or at least from .6 to 1) is certainly dubious and must raise doubts as to the credibility of the Kartvelishvili function. The Kartvelishvili function is based on the structure functions of charmed mesons. Differences between charmed meson and baryon structure functions would have to be accounted for before it could properly be applied to charmed baryons.

### 2.2.2 Functions Based on Independent Fragmentation

The most celebrated heavy quark FF is that of Peterson \cite{23} et al. Its success is based on the simplicity of its form and its flexibility. The function is based
on quantum mechanics and IF. Basically a q̅q pair is produced in the color field of a heavy quark Q. Q then combines with the appropriate anti-quark, and one quark is left uncombined to continue the fragmentation process. The rate goes as $\Delta E^{-2}$, the difference in energy between the initial heavy quark state and the final state hadron + q,

$$\Delta E = \sqrt{m_H^2 + z^2 p^2} + \sqrt{m_q^2 + (1-z)^2 p^2} - \sqrt{m_Q^2 + p^2}$$

Approximating $m_H = m_Q$ and after a binomial expansion we obtain

$$\Delta E \propto 1 - \frac{1}{z} - \frac{\epsilon_Q}{1-z}$$

Squaring, and tacking on a factor of $\frac{1}{z}$ for longitudinal phase space produces the well known result

$$D^{H}_{Q}(z) = N \left( \frac{1}{z} \left[ 1 - \frac{1}{z} - \frac{\epsilon_Q}{(1-z)} \right]^2 \right)^{-1}$$

The parameter $\epsilon_Q$ is proportional to $m^2_{q\perp}/m^2_{Q\perp}$, and the transverse mass is defined as $m^2_{\perp} \equiv m^2 + p^2_{\perp}$. As defined, $z$ is the fraction of the momentum of the heavy quark Q retained by the hadron. The binomial expansion is a delicate issue, especially in the region $z \to 0$. It is mentioned in the paper that the expansion is strictly valid in the $p \to \infty$ limit. The function peaks at

$$z_{\text{peak}} \approx 1 - 2\epsilon_Q$$

with width $\approx \epsilon_Q$.

The function was not initially derived in terms of the light cone variables, but it is pointed out that light cone variables may be more appropriate at finite energies, and the above expression may be carried over provided it is cut off below $p^+_{\text{min}}$.

It is reasonable to assume that a di-quark pair could have been popped and

$$m_{q\perp}^2 \to m_{d\perp}^2$$

in the formula for $\epsilon$. Both the Peterson and Andersson functions
contain a term that depends on $m_\perp$. Often the $p_\perp$ is neglected and $m_\perp \rightarrow m$. A consistent interpretation needs to be arrived at.

The Peterson function goes as $(1-z)^2$ as $z \rightarrow 1$. This is in disagreement with reciprocity using Eq. 2. This potentially disquieting fact motivated some recent work by Collins and Spiller [24] who set out to produce a fragmentation function that did not ‘violate’ reciprocity and dimensional counting rules. They calculate the cross section $\sigma (e^+e^- \rightarrow HX)$ which depends on the fragmentation function $D_H^Q(z)$. They use Independent Fragmentation and compute the vertex $Q \rightarrow Hq$. This depends on the transverse momentum distributions of H and q with respect to the Q direction, and the longitudinal momentum distributions. The longitudinal momentum distribution, they claim, happens to be equal to the function which describes the momentum distributions of the two valence quarks in a meson. The final result with the approximations $m_H = m_Q \gg \langle k_T^2 \rangle$

$$D_H^Q(z^+) \approx N \left( \frac{1-z^+}{z^+} + \frac{2-z^+}{1-z^+} e_Q \right) \left( 1 - \frac{1}{z^+} - \frac{e_Q}{1-z^+} \right)^2$$

We have elected to use $z^+$ to indicate functions explicitly derived in terms of light cone variables. The parameter $e_Q$ is defined $\equiv \left( \frac{\langle k^2_T \rangle}{m_Q^2} \right) < k_T^2 > = (0.45 \text{ GeV})^2$ ‘represents the size of the hadron in momentum space’. A Peterson like term is apparent, and the two functions are noticeably similar. This function, in contrast to the Peterson function, exhibits a $(1-z^+)$ behavior as $z^+ \rightarrow 1$. It is uncertain how this function would differ for D and $D_s^+$ mesons, since it is unclear if $k_T$ is a constant or depends on $u$, $d$ or $s$ quarks. This function would not be directly suitable for baryons. Like the Kartvelishvili function, it is partially derived from the structure function of mesons.

They have also proposed a kernel for vector $\Rightarrow$ pseudoscalar decays, to allow calculation of inclusive fragmentation functions. This is an important problem
when attempting to determine $D^0$ or $D^+$ fragmentation, since the feed down from $D^*$ must be taken into account. They calculate the fragmentation function for the decay $H_Q^* \rightarrow H_Q \pi$, $z$ being the fraction of momentum retained by $H_Q$

$$P_Q(z) = \frac{N}{z(1-z)\left( \frac{m_{H_Q}^2 - <k_H^2>}{z} - \frac{m_{H_Q}^2 + <k_H^2>}{1-z} \right)^2}$$

The inclusive fragmentation is the given by

$$D_Q^{H_Q}(z) = \delta D_Q^{H_Q}(z) + (1-\delta) \int_{z}^{1} \frac{dy}{y} D_Q^{H_Q}(y)P_Q\left(\frac{y}{z}\right)$$

$0 \leq \delta \leq 1$ is the fraction of direct pseudoscalar production. The first term represents direct pseudoscalar production and the second feed down from vector mesons.

### 2.2.3 Functions Based on String Fragmentation

The function of Andersson [25] et al. is based on the symmetric Lund String Fragmentation model. The symmetric Lund function assures that starting at either the $q$ or $\bar{q}$ end and iterating will produce the same result. The left right symmetry is a non-trivial property of an iteratively generated cascade. The Andersson function is derived by solving a set of coupled equations that demand this property. The resulting form is

$$D_Q^{H}(z^+) = \frac{N}{z^+}(1-z^+)^a exp\left( -\frac{b \cdot m_H}{z^+} \right)$$

Where $a$ and $b$ are constants related to the production of $q\bar{q}$ vertices in space-time. The constant $b$ is flavor independent, while $a$ may depend on flavor. The above expression assumes all $a$’s are equal. The constant $m_{\perp}$ is the transverse mass of the hadron produced. It is anticipated that $0 \leq a \leq 2$ and $b \geq 0$, with the
current values set at $a \sim 1$ and $b \sim (\frac{1}{2.25})$. The mean is predicted to be

$$<z^+> \simeq 1 - \frac{(a + 1)}{bM^2}$$

$M$ is the mass of the meson produced.

The Andersson expression was derived using massless quarks moving along light cones. It is expected to carry over to an initial heavy $QQ$ pair, which move along hyperbolas with light cone asymptotes. A definite difficulty with this function is that it differentiates light and heavy quark fragmentation only by the mass of the hadron formed. Although it is derived in the full glory of the string model, it has to be recognized that symmetry was the main impetus for this function.

An event generated using SF has the following structure (Figure 2.2). $q\overline{q}$ vertices appear in spacetime in the light cone of the initial $q_0\overline{q}_0$ pair. When quark world lines cross a hadron is formed. The hadron mass is proportional to the area of the color field between the quarks (hatched region). We know from Eq. 1 that massive objects must move along hyperbolas in spacetime. In order to produce a hadron of mass $m$, the secondary $q\overline{q}$ pair must lie on this hyperbola. The heart of the Lund model is that when a string breaks to produce a hadron of mass $m$, vertices of the secondary $q\overline{q}$ will be uniformly distributed along the production hyperbola. In this scheme, quarks fragment into hadrons with a discrete mass spectrum. An alternative point of view was taken by Bowler [26] who choose to break the string in a more classical way. The probability $dP$ of a break occurring at spacetime coordinates $(x + dx, t + dt)$ is

$$dP = \Pi dx dt$$

$\Pi$ is a constant per unit length of string and unit time. The function is derived
Figure 2.2: Models of string fragmentation. The top diagram (a) shows a model where the energy stored in the quark flux tube becomes so great that additional quark pairs are produced. The bottom (b) diagram shows a light cone view.
utilizing the basic equations of the string model. The primary $Q\bar{Q}$ pair is massive and the secondary quark pairs are massless. Two sets of coordinates are considered in the fragmentation process, $(x_1, t_1)$ the point where secondary q$\bar{q}$ pair is produced, and $(x_m, t_m)$ the coordinates where the world lines cross and the hadron is formed. The derivation is straightforward, applying

$$dP = \Pi dA \exp(-\Pi A)$$

the probability of creating the secondary pair within $dA$ where $A$ is the area of the color field in the absolute past of the point $(x_m, t_m)$ gives

$$D_{Q}^{H}(z) = \left(\frac{B}{z}\right) \exp\left(-B m_{Q}^{2} \left(\frac{m_{H}^{2}}{m_{Q}^{2} z} - 1 - \ln\left(\frac{m_{H}^{2}}{m_{Q}^{2} z}\right)\right)\right)$$

$B$ equals $(\Pi/2k^2)$. $\Pi$ being the constant probability the string will break, and $k$ is the string constant.

The function has recently been modified [17] to include light cone formalism and be made more suitable for CESR energies. Bowler also later suggested the addition of the term $(1 - z)\beta$ to account for the fact that the string is not ‘straight’ due to soft gluon emission. The modified Bowler function is:

$$D_{Q}^{H}(z^+) = (1 - z^+)\beta \left(\frac{B}{z^+}\right) \exp\left(-B m_{Q}^{2} \left(\frac{m_{H}^{2}}{m_{Q}^{2} z^+} - 1 - \ln\left(\frac{m_{H}^{2}}{m_{Q}^{2} z^+}\right)\right)\right)$$

It is expected that $\beta$ will be close to 1. One criticism of this approach is that the Bowler function is singular as $m_H \to 0$ and would highly favor 0 mass mesons unless a low mass cutoff is imposed.

### 2.2.4 Comparison of the Functions

Table 2.1 presents the collected features of the various fragmentation functions. The two functions derived from IF are similar, as are the two based on SF. There
Table 2.1: Properties of Fragmentation Functions

| Function         | Model | Reciprocity | Variable | Baryons | \( z \to 1 \)  |
|------------------|-------|-------------|----------|---------|-----------------|
| Kartvelishvili   | -     | Yes         | CT       | No      | \((1 - z)\)     |
| Peterson         | IF    | No          | CT,LC    | Possible| \((1 - z)^2\)   |
| Collins          | IF    | Yes         | LC       | No      | \((1 - z)\)     |
| Andersson        | SF    | No          | LC       | ?       | \~\((1 - z)\)   |
| Modified Bowler  | SF    | No          | LC       | ?       | \~\((1 - z)\)   |

LC = Light Cone  CT = Cartesian

are differences between the two models. The most distinct difference is the exponential term in the SF functions. This causes the fragmentation function to go to 0 prematurely. SF also appears to be flatter, but this feature strongly depends on \( \epsilon \) for the IF functions. We show the Andersson and Bowler functions in Figure 2.3, and the Collins and Peterson forms in Figure 2.4. Another difference is the \( z \to 1 \) behavior, where reciprocity can be used as a guide. All the functions, save the Peterson, lean toward a \((1 - z)\) behavior. The Peterson function has a \((1 - z)^2\) as \( z \to 1 \). From an experimental perspective, the low \( z \) region is difficult to measure because of poor efficiencies of low momentum tracks. For the high \( z \) end, a large amount of statistics would be necessary to accurately determine the power dependence of the curves. So experimental clarification of these issues will only come with great effort. A detailed comparison of these five functions to CLEO’s \( D^{*+} \) fragmentation distribution can be found elsewhere \[27\]. One of the more serious shortcomings of IF stems from the so-called string effect (Figure 2.5). In IF models all the partons fragment independently, while in SF models a color “web” stretches from the quarks to the gluons. This acts to increase the density of hadrons in the vicinity of the gluon direction. This has been observed experimentally \[28\]. While the exact cause of this effect in SF Monte Carlos is unclear, it cannot be accommodated in IF schemes. This, along with other undesirable features has all but terminated interest in IF as a viable
2.3 Theoretical Approaches to Charm Decay

Here we briefly recount the evolution of our understanding of the weak decay of charmed mesons. The weak force has supplied the physics community with a great deal of beauty and bedazzlement, with the charm sector being no exception. This field also attests to the vital interplay of experimental and the-
Figure 2.4: Independent fragmentation inspired functions. Collins function (dash-dotted line), and the Peterson function (solid line).

theoretical physics, as theoretical predictions for charm decays have been consistently defied by experimental observations. Here we trace our understanding of charm decays from the failings of the simple spectator model to more advanced approaches. We also detail outstanding conflicts with experimental measurements. Voluminous amounts of material has previously been written on this subject. The reader may find of particular use the experimental summaries of Hitlin [29], Schindler [30], and the theoretical treatments of Rückl [31], Shifman [32] and the recent review by Bigi [33] to append the rather quick distillation presented here.
2.3.1 The Trivial Spectator Model

The lowest level of understanding of charm decays utilizes the method of quark diagrams. The first order processes are collected in Figure 2.6.

They may be distinctly be separated topologically into two classes; the “spectator class” (a - b), and the nonspectator or “annihilation class” (c - d). The primary difference is in the location of the W vertices. In spectator processes the
W vertex occurs only on the c quark world line, it then hadronizes either independently (spectator) or in conjunction with the light quark (color suppressed). In annihilation processes the W vertex touches both quark lines of the initial meson. The trivial spectator model makes the following predictions:

1. The dominant process in the weak decay of charmed mesons is the spectator diagram. This approximately corresponds to the beta decay of the charmed quark. The rate for charm decay \( \Gamma_0 \) can be obtained (to first order) by scaling the muon lifetime \( \tau_\mu = \frac{192\pi^3}{G_F m_\mu^5} \approx 2.2 \times 10^{-6}\text{sec} \) by a factor of \( \frac{1}{3} \left( \frac{m_\mu}{m_c} \right)^5 \approx 7 \times 10^{-13}\text{sec} \), to account for the charmed quark mass and the additional number of decay channels available to the W. This form suffers from an exceptional (fifth power) dependence on the charmed quark mass which is an undefined quantity. It is generally approximated to be in the range of \( 1.5 - 1.6 \text{ GeV} \).

2. The color suppressed spectator diagram occurs at a rate reduced by a factor of \( \approx 10 \). The is because the W has to hadronize a quark pair with the right color combination to make color singlet hadrons.

3. The exchange diagram, which is contributes to D\(^0\) decay on the Cabbibo allowed level is expected to be strongly suppressed by a small wave function overlap |\( \psi(0) \)|\(^2 \propto \frac{f_D^2}{M_c^2} \) where \( f_D \sim 0.15 \text{ GeV} \). Similarly the W annihilation is expected to suffer helicity suppression (as in pion decay) by a factor \( \frac{m_D}{M_c} \).

4. Taking these factors into account, the D\(^0\) and D\(^+\) are expected to have the same lifetimes and semi-leptonic branching ratios.
Figure 2.6: Lowest order decay diagrams for charmed mesons. a) the spectator diagram, b) color matched or color suppressed spectator diagram, c) W exchange diagram, d) W annihilation.

| Group   | $D^+ \times 10^{-13}$sec | $D^0 \times 10^{-13}$sec | $D_s^+ \times 10^{-13}$sec |
|---------|--------------------------|--------------------------|---------------------------|
| CLEO    | 11.4 ± 1.6 ± 0.7         | 5.0 ± 0.7 ± 0.4          | 4.7 ± 2.2 ± 0.5           |
| World Average | 10.45$^{+0.31}_{-0.29}$ | 4.27 ± 0.10              | 4.31$^{+0.36}_{-0.32}$    |

2.3.2 Experimental Controversies

After the experimental picture of charm decays began to unfold, several failings of the trivial Spectator model were revealed.

1. The $D^+$ and $D^0$ mesons were found to have very different decay rates.
This has been determined by measuring the lifetimes of the D⁰ and D⁺, as collected in Table 2.2. We have presented the measurements of CLEO [34] and the current world average [29]. Empirically, we measure $\frac{\tau(D^+)}{\tau(D^0)} = 2.45 \pm 0.09$. This difference has also been established by a large discrepancy in the semileptonic branching ratios of the D⁰ and D⁺. A recent Measurement by the MARK III [35] has determined $\frac{B(D^+ \rightarrow e^+X)}{B(D^0 \rightarrow e^+X)} = 2.3^{+0.5+0.1}_{-0.4-0.1}$, which is in good agreement with the measured lifetime difference.

2. Several color suppressed D decays have been observed to occur with a healthy rate. Among them $D^0 \rightarrow \bar{K}^0\pi^0$ has been found to occur at a rate of $0.45 \pm 0.9$ times that of $D^0 \rightarrow K^-\pi^+$. CLEO [36] and ARGUS [37] have also observed a color suppressed decay of the B meson, $B(B \rightarrow \psi X) \sim 1.2\%$.

3. Evidence for annihilation processes appears to exist. The first was the observation of $D^0 \rightarrow \phi \bar{K}^0$. This mode was first observed by ARGUS [38] and later by CLEO and MARK III. These measurements have confirmed a branching fraction of $\approx 1\%$. Other clean annihilation class signatures have been sought in $D^+_s$ decay. The E-691 experiment [39] has placed a stringent limit on the $D^+_s \rightarrow \rho^0\pi^+$ mode, finding $\frac{B(D^+_s \rightarrow \rho^0\pi^+)}{B(D^+_s \rightarrow \phi\pi^+)} = < 0.08$ at the 90 \% CL. This same group has, however, has observed the annihilation decay candidate $\frac{B(D^+_s \rightarrow \pi^+\pi^-\pi^0)}{B(D^+_s \rightarrow \phi\pi^+)} = 0.29 \pm 0.07 \pm 0.05$. 

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2.3.3 Patching up the Spectator Model using QCD

In light of these difficulties, reexamination of the spectator model was in order. The charged weak current has as its hadronic component:

\[ J^\mu = (J^\mu)^i = (\bar{u}, \bar{s}, \bar{t}) \gamma^\mu \left(1 - \gamma^5\right) \begin{pmatrix} V_{ud} & V_{us} & V_{ub} \\ V_{cd} & V_{cs} & V_{cb} \\ V_{td} & V_{ts} & V_{tb} \end{pmatrix} \begin{pmatrix} d \\ s \\ t \end{pmatrix} \]

Where the \( V_{xy} \) are the familiar K-M matrix elements. The current-current approximation \((q^2 << M^2_W)\) of the weak Hamiltonian leads to the “Bare” Hamiltonian for hadronic charm changing interactions of

\[ H^{(0)}_W(\Delta c = -1) = \frac{G_F}{\sqrt{2}} V_{cs} V_{ud}^* \left[(\bar{s}c)_L(\bar{d}u)_L\right] \]

where the notation \((\bar{a}b)_L\) implies the canonical V-A structure \(\bar{a}\gamma^\mu(1 - \gamma^5)b\). Since

![Figure 2.7: One loop gluon corrections at a four quark vertex.](image)
a $W$ vertex can be thought of as a four fermion interaction, Rückl [31] first studied the effects of 1 loop gluon exchange at a $W$ vertex. The diagrams for these processes are shown in Figure 2.7. The diagrams of a) are absorbed in the renormalization of $G_F$, while those of b) introduce new effects. We note that none of the diagrams in b) are possible in semileptonic decay. Rückl predicted the following ramifications of the hard gluon exchanges:

1. weak couplings would be renormalized.
2. distortions to the color structure of the interaction from octet currents
3. possibility of new chiral structures, $(V+A)$ for example
4. possibility of new Lorentz structures (scalar or tensor)

The last two effects should be small, but the first two are not. He calculated the first order correction to the Hamiltonian to be

$$H^{(1)} = \frac{G_F}{\sqrt{2}} V_{cs} V^*_{ud} \frac{3 \alpha_s}{8 \pi} \log \left( \frac{M_W^2}{\mu^2} \right) (\bar{s} \lambda^a c)_L (\bar{d} \lambda^a u)_L$$

The effective weak Hamiltonian then becomes

$$H_{eff}(\Delta c = -1) = \frac{G_F}{\sqrt{2}} V_{cs} V^*_{ud} [c_+ O_+ + c_- O_-]$$

where

$$O_\pm = \frac{1}{2} \left[ (\bar{s} c)_L (\bar{u} d)_L + (\bar{s} d)_L (\bar{u} c)_L \right]$$

and to lowest order

$$c_\pm = 1 \mp \frac{\alpha_s}{2 \pi} \log \left( \frac{M_W^2}{\mu^2} \right)$$

Where the $c_\pm$ coefficients obey the relation $c_- = \frac{1}{\sqrt{c_+}}$ and are approximately 2.0 and 0.7, respectively at $q^2 = 1.5$ GeV. A plot of $c_\pm$ verses mass scale for leading log (LL) and next to leading log (NLL) is displayed in Figure 2.8. Making one
The final transformation allows for a convenient representation of the Hamiltonian.

We define

\[ c_1 = \frac{c_+ + c_-}{2} \]
\[ c_2 = \frac{c_+ - c_-}{2} \]

The Hamiltonian then becomes

\[ H_{eff}(\Delta c = -1) = \frac{G_F}{\sqrt{2}} V_{cs} V^*_{ud} [c_1 (\bar{s}c)_L (\bar{u}d)_L + c_2 (\bar{s}d)_L (\bar{u}c)_L] \]

The resulting \( c_1 \) and \( c_2 \) processes are displayed in Figure 2.9

Several important relationships can be derived based on the \( c_\pm \) coefficients.
The non-leptonic width becomes enhanced by a factor

\[ \Gamma_{NL} = (2c_+^2 + c_2^2)\Gamma_0 \]

where we would naively expect three. The semi leptonic branching ratio is also modified

\[ B(c \rightarrow eX) = \frac{1}{(2c_+^2 + c_2^2 + 2)} \approx 15\% \]

Both of these are clearly pushing things in the correct direction. An induced
neutral current ($c_2$ process) has also arisen which has precisely the form of the
color suppressed spectator diagram, and may neatly account for such processes. In spite of these successes, the spectator model still does not predict the large difference in the $D^+$ and $D^0$ lifetimes which has been confirmed experimentally, nor the annihilation type processes. To this end we must either consider the addition of nonspectator processes or resort to the alternate approach of final state interactions.

### 2.3.4 Final State Interactions

We examine final state interactions on the quark level and the hadron level.

Several $D^+$ decay modes contain identical quarks in the final state. An example is the decay mode $D^+ \rightarrow \bar{K}^0 \pi^+$ depicted in Figure 2.10. Since there are two identical fermions in the final state, the Pauli principle demands that the wave function be suitably altered. This induces an additional term to the $D^+$ rate

$$\Gamma_{\text{pauli}} = (2c_+^2 - c_-^2) \frac{12\pi^2}{M_D^2} f_D^2 \Gamma_0$$

Since this term is linked to $f_D^2$ which depends on the wave function overlap, uncertainty is introduced into the magnitude of this effect. General considerations predict that the size of the interference term compared to the spectator term will be about 20%. We note that for $c_- >> c_+$ the overall term is negative and acts to increase the $D^+$ lifetime, as expected.

It is also possible that in the case of two body decays, the two final state hadrons may re-scatter into a different final state (post hadronization). This has been proposed [40] as a mechanism by which ordinary spectator decays
Figure 2.10: The decay mode $D^+ \rightarrow \bar{K}^0\pi^+$. Both color clustered terms contain a $\bar{d}$ quark.

may produce more exotic nonspectator final states. Hadron re-scattering is also required by unitarity.

2.3.5 **Summary**

An easy lesson to be learned from the previous discussion is that simple quark diagrams are meaningless without fully considering the role of gluons. Their presence substantially alters the weak Hamiltonian, and may well act to catalyze the nonspectator diagrams (by removing helicity suppression in $W$ ex-
change, for example). While the spectator model with QCD improvements is a more reliable theory, it is still incomplete. Both the nonspectator processes along with pre- and post- hadronization final state interactions seem to have established a foothold in the picture of charm decay. Some of the more modern theoretical systems have been successful at predicting the coarse features of relative D decays, and these have been limited to the two body case (fortunately a large number of D decays are two body or quasi-two body). Stumbling blocks persist however. Implementation of nonspectator processes will require some effort. The Bauer-Stech-Wirbel [41] group has typically neglected annihilation processes, and thus has a difficult time explaining $\phi K^0$. The 1/N expansion method of Buras-Gerard-Ruckl [42] includes such processes, however they are often of higher order and are then dropped out. It is also unclear whether final state hadron scattering can ever be integrated into calculations in a systematic way. Although the $D^0 \leftrightarrow D^+$ lifetime difference has been known for almost a decade, we still await a quantified theoretical answer.
CHAPTER 3
APPARATUS

In order to perform high energy physics experiments, one needs to generate and record high energy interactions. Two general methods exist for producing interactions, accelerating a particle and firing it into a target (fixed target), or accelerating two particles (usually particle-antiparticle) and colliding them. $e^+e^-$ colliders are useful as both particles completely annihilate into a state of pure energy. The non resonant cross section for production of a fermion-antifermion pair in an $e^+e^-$ collision is approximately:

$$\sum_{nf} \frac{4\pi \alpha^2 e_f^2}{3s}$$

where $nf$ is the number of species of the particular fermion, $\alpha$ the fine structure constant, and $s$ the center-of-mass energy squared (GeV$^2$). In addition the $e^+e^-$ annihilations, $e^+e^-$ scattering (Bhabha) and two $\gamma$ process also occur. Of greater importance to the experimentalist are the several resonant enhancements to the $e^+e^-$ cross section. These are shown in Figure 3.1.

Reactions from proton colliders are extremely rich and dense, since the protons themselves are composite objects. Because the interaction rates for proton colliders are substantially larger, and the physics more difficult to disentangle, they have assumed the role of discovering phenomena, while $e^+e^-$ machines excel at refined measurements.

3.1 CESR

The data for this research project was gathered at the Cornell Electron Storage Ring (CESR), located at Wilson Synchrotron Laboratory on the campus of Cor-
Figure 3.1: Electron-positron cross section as a function of center of mass energy. The cross section has been measured up to 50 GeV. Above this energy only the Z resonance is established.
nell University in Ithaca, New York. It is capable of producing $e^+e^-$ collisions at a center-of-mass energy from 7 to 14 GeV (here, and throughout this thesis $c = 1$ unless explicitly stated otherwise). Electrons and positrons are produced in a linear accelerator and injected into a synchrotron where they are accelerated to near collision energy. Finally the counter-rotating beams are transferred into a storage ring 770 m in circumference where they are made to collide in two interaction regions at multiples of a fundamental frequency of 390 KhZ. The layout of this facility is shown in Figure 3.2. Individual data runs last an average of 1-2 hours. The energy resolution, which is largely a function of the bending radius of the machine is of order 3 MeV.

Storage rings maximize the opportunity for obtaining collisions from a generated group of electrons and positrons. This is offset by synchrotron radiation, the energy lost by the beams as they are accelerated so as to travel in a (near) circular orbit. The power lost per electron per turn is:

$$P_\gamma = \frac{2}{3} \frac{c r_e E^4}{(mc^2)^3 \rho^2}$$

$r_e$ is the classical electron radius, $E$ the beam energy, $m$ the mass of the particles in the beam, $\rho$ the radius of curvature. At CESR energies this amounts to $\approx 1$ MeV. To maintain a constant beam energy, this energy must be constantly replenished by boosts through R.F. cavities. The R.F. power costs makes a large contribution to the operating expenses of a storage ring. The cost of storage ring is predicted by the Richter relation:

$$S = \alpha R + \beta \frac{E^4}{R}$$

which the first term accounts for magnets, tunnels, etc., and the second for R.F. expenditures. This predicts both the size and the cost of a storage ring increase as the energy squared.
Figure 3.2: The Cornell electron storage ring and associated components.
The critical performance parameter in any collider facility is the luminosity (L). The number of events observed $N_{\text{obs}}$ is given by:

$$N_{\text{obs}} = \sigma B \epsilon \cdot \int L$$

The integrated luminosity enters linearly in the rate, along with the cross section, branching ratio and reconstruction efficiency. Since the only variables available to the experimenter are the luminosity and efficiency (which is largely determined by detector design) it is desirable to accumulate as much luminosity as possible. The instantaneous luminosity, for a storage ring such as CESR, is governed by:

$$L = \frac{f N s \xi_v}{2 r \beta^*}$$

where $f$ is the frequency of the collisions, $N$ the number of particles per beam, $\xi_v$ the vertical tune shift, and $\beta^*$ is a parameter which represents how tightly the two beams can be focused at the interaction point. During this data taking the CESR group began a program to increase the luminosity by simultaneously circulating several bunches of electrons and positrons. This project was successful, although the strain on the R.F. system caused, at first, the machine to operate less reliably. The average luminosity during this running period was roughly $3 \times 10^{31}$ cm$^{-2}$sec$^{-1}$.

### 3.1.1 Data Sample

The data sample was acquired from August 1984 through February 1986 with the CLEO detector, operating in what shall be referred to as the 85VD configuration. Table 3.1 collects the operating conditions for collection of the data sample. The various states of timing and bunches reflect the multibunch upgrade pro-
gram. The luminosity reflects the runs to first order which are considered good, the data analysis utilizes a subset of the above luminosity. The major impetus for CLEO physics running is to study the decays of B mesons which are produced on the Υ(4S) resonance. To efficiently separate processes connected with the Υ(4S) from those of the surrounding continuum, the Υ(4S) running time is divided between on resonance and a center-of-mass energy approximately 60 MeV below resonance in a 2 : 1 ratio. Also, running on the Υ(1S) is partially motivated to study lepton faking backgrounds.

### Table 3.1: 85VD Data Sample

| Resonance region | Machine timing | # bunches | ∫L (pb⁻¹) |
|------------------|----------------|-----------|-----------|
| Υ(1S)            | 3              | 3         | 17.0      |
| Υ(4S)            | 7              | 7         | 10.5      |
| Υ(4S)            | 7              | 3         | 63.3      |
| Υ(1S)            | 7              | 3         | 3.4       |
| Υ(4S)            | 3              | 3         | 43.8      |
| Υ(3S)            | 3              | 3         | 33.3      |

3.2 The CLEO Detector

The results of the $e^+e^-$ collisions were recorded by the CLEO [43] detector, a large magnetic spectrometer with excellent charged particle tracking capabilities. Operational since 1979, in the summer of 1984 the first stages of a major detector upgrade [44] were implemented. This consisted of the installation of a precision vertex detector (VD) and instrumentation of the central drift chamber (DR1.5) to simultaneously perform drift time and specific ionization (dE/dx) measurements in all 17 layers. The CLEO detector so configured shall be referred to the 85VD configuration.
The CLEO detector, not unlike a high fidelity stereo system, consists of several distinct detector elements, each of which performs a specific function in the event reconstruction. The single most important property that determines the detector response is whether or not the particle possesses electric charge. The methods of detecting and analyzing charged and neutral particles differ such that detectors of CLEO’s generation were often polarized to one extreme. In brief, the CLEO consists of an inner detector dedicated to charged particle tracking and an outer detector dedicated to the identification of both charged and neutral particles. The natural boundary between the inner and outer detectors is a 1.0 Tesla superconducting solenoid magnet. The inner detector consists of the two drift chamber devices mentioned above, as well as shower counters (ES) on either end of the drift chamber face. The outer detector contains a dedicated dE/dx system (DX) for charged particle identification, a time of flight scintillator detector (TF), also for particle identification, and an electromagnetic shower detector (RS) for γ’s, π₀’s and electron identification. These three devices are arranged in eight identical octant sections around the coil. Surrounding the entire detector is a steel-drift chamber sandwich used to identify muons (MU). Several other detector elements exist from previous detector configurations, they do not affect this analysis and will not be discussed here. A beams eye view of the CLEO detector is displayed in Figure 3.3. The octant structure is readily apparent, and a side view is provided in Figure 3.4. We now turn to a more involved discussion of the detector function.
3.2.1 Tracking Fundamentals

The inner detector is globally cylindrical in geometry, consisting of fine wires strung parallel to the beam direction, on circles of constantly increasing radius. Drift chamber detectors are gaseous detectors which detect the passage of charged particles through ionization. The ionized electrons drift toward sense wires held at high electrostatic potential. The extent to which the sense wires are isolated and the way the geometry of the electric field is defined determines the intrinsic performance characteristics of the chamber. As the electrons are ac-
Figure 3.4: Side view of the CLEO detector.
celerated toward the sense wire, they begin to acquire enough energy to liberate other electrons. This develops into a process referred to as avalanche multiplication, which the electron gain is on order $10^4$. The gas chosen is typically argon, which has high specific ionization, good gas gain, and undergoes roughly 30 ionizing collisions/cm at STP. Since argon is in the same chemical family as neon, the high voltage conditions could result in breakdown or self sustaining discharge. This is obviated by the addition of an organic vapor. Nobel gases can only be excited by the emission or absorption of photons, while organics posses a myriad of rotational and vibrational states. This leads to a substantial amount of energy dissipation is radiation-less transitions. Organics also tend to increase the drift velocity of the gas, thereby decreasing diffusion effects. CLEO operates its wire chambers in a 50-50 argon-ethane mixture, which has a mean drift velocity of 50 $\mu$m/nsec.

The central operating principle of a drift chamber is drift velocity saturation. As the electric field is increased, the drift velocity plateaus or saturates. Insuring a constant drift velocity across the cell allows for a theoretically simple (in practice the drift-time relationship often represents the most difficult aspect of drift chamber calibration) method for reconstructing the trajectory through the cell by measuring the time difference between entry into the cell and a hit on the wire. The accuracy with which trajectories can be reconstructed is in the range of hundreds of microns, and requires timing precision on the nanosecond level.

Immersing the drift chamber in a solenoidal magnetic field ($B$) allows for the measurement of two vital quantities for a charged track, the sign of the electric charge and the momentum. The path of a charged particle moving in a magnetic field is a helix. The CLEO reference coordinate system is depicted in
CLEO uses a track coordinate system that consists of three $r - \phi$ parameters; CUDR is one half times the reciprocal of the radius of curvature and is signed, FIDR is measured from the distance of closest approach to the beam line, DADR is the signed impact parameter, and two additional parameters CTDR, the cotangent of the angle between the track and the beam line (polar angle), and Z0DR, the z coordinate of the distance of closest approach. The transverse momentum $p_\perp$ can then be calculated by:

$$p_\perp = 0.015 \frac{B}{|\text{CUDR}|}$$

where the dimensions are GeV, KiloGauss, and meters. Once the transverse momentum is determined the total momentum is trivially extracted from the measurement of the polar angle.

Drift chamber performance is gauged by the spatial resolution in the $r - \phi$ plane $\sigma_{r\phi}$ and the momentum resolution $\frac{\delta p}{p}$. The two are linked by the relation:

$$\left( \frac{\delta p_\perp}{p_\perp} \right)_{\text{res}} = \frac{p_\perp \sigma_{r\phi}}{(0.03) L^2 B \sqrt{750 N + 5}}$$
for a drift chamber with N equally spaced measurements over a lever arm L in an axial magnetic field B (the units are GeV, meters, and kilogauss). The other dominant term comes from multiple scattering:
\[
\frac{\delta p_{\perp}}{p_{\perp}}_{ms} = 0.5 \frac{L}{LB} \sqrt{1.45 \frac{L}{X_0}}
\]
here \(X_0\) is the average radiation length of the detector in meters. The resolution dominated term is prominent for high momentum tracks whilst the multiple scattering limits the resolution at low momentum. The CLEO VD+DR1.5 tracking system achieves a momentum resolution of \(\left(\frac{\delta p}{p}\right)^2 = (0.007p)^2 + (0.006)^2\) (\(p\) in GeV/c).

### 3.2.2 Vertex Detector

The CLEO vertex detector is a high precision drift chamber which forms the innermost detector element. 800 cells are divided among 10 axial layers, which range in radius from 8 to 16 cm (see Figure 3.6). The inner 5 layers have 64 cells per layer, and the outer 5 layers each have 96 cells. The cells are hexagonal in shape, and expand in size so as to subtend a constant angle in \(\phi\). The \(\sigma_{r\phi}\) resolution for the VD [45] on average is 100 microns.

The active z length is 70 cm. Two methods are provided for z measurements. A resistive sense wire is read out at both ends allowing a measurement by charge division. The resolution for this method is in the realm of 9 mm. More precise information is obtained from the conducting cathode strips on the inner and outer support tubes. The strips are segmented into 8 \(\phi\) regions which are further divided into 64 (96) sections on the inner and outer tubes. The cathode strips z measurements have a resolution of 750 microns.
Figure 3.6: Vertex detector cell arrangement.
The device derives its name from the design purpose to extrapolate tracks back to their production vertex. The instrument is separated from the interaction point by a silver coated beryllium beam pipe and a carbon inner support tube. The interceding material amounts to 0.1 % of a radiation length. The extrapolation error is given by:

$$\sigma_{ext} = \sqrt{100^2 + \frac{115}{\sin \Theta(p\beta)^2}}$$

$\Theta$ is the polar angle and $\beta = v/c$. $p$ is in GeV and the result has the units of microns.

A plethora of benefits accompanied the installation of the VD. The tracking momentum resolution was greatly improved, and the momentum range for low momentum track reconstruction was extended. Beam wall and beam gas events were efficiently removed, and the capabilities of the experimental trigger were extended.

### 3.2.3 Drift Chamber

In contrast to the VD, the design of the drift chamber relies on the repetition of a single cell geometry. The unit cell consists of a 20 $\mu$m diameter gold plated tungsten sense wire responsible for a radial region of 11.3 mm. The cell boundaries are formed by three 115 micron diameter silver plated beryllium copper field wires located on either side of the sense wire, and arranged in a straight line perpendicular to the radius of the cylinder over a region of 10 mm centered about the sense wire. The inner and outer boundaries of the device are kept at ground, which coupled with the open rectangular cell geometry leads to certain peculiarities in the function of the detector which shall be addressed.
The arrangement of cells is rigidly described by the prescription $\text{cells/layer} = 24 \times (N + 4)$ which are located on nearly concentric cylinders of radius $R_N = 42.49 \times (4 + N)$ (mm) and $N$ ranges from 1 to 17. The lever arm of the drift chamber ranges from 212.5 to 892.5 mm. A schematic illustration of the drift chamber can be found in Figure 3.7. Also distinct from the VD is the method of $z$

Figure 3.7: The CLEO central drift chamber.

measurement. Starting with the second cylinder, the 8 even numbered layers are
strung in a “stereo” mode at alternating ± angles of 2.9° to the beam axis. The odd numbered layers are all axial, strung parallel to the beam line. High voltage is distributed through the field wires, with the sense wires kept at ground for ease of electronic readout. During August 1984, the readout electronics for the drift chamber was replaced, with the new system allowing both timing and pulse height measurements. With the addition of the vertex detector, the drift chamber underwent its most thorough calibration, revealing several biases. It is this mode of operation which shall be recounted here.

The chamber was operated at 2050V, corresponding to an approximate electric field strength of 3630 v/cm and gas gain $\approx 2 \times 10^4$. Performance tests [46] of the new electronics were done by T. Copié and the author using a 9 layer, 220 channel $\phi$ slice mock-up of the drift chamber using cosmic ray muons. Plots of the timing and pulse height performance are shown in Figure 3.8 and Figure 3.9. Because of the open cell geometry, a symbiotic relationship exists among the gains of the cylinders. Shown in Table 3.2 are the effects on the average measured pulse height in layer 16 subject to variations in the outermost layer 17. Layer 15 was held at 2050 V. Since the inner and outermost cylinders have one high voltage layer and ground as their two radial nearest neighbors, one can deduce from the above table that they will also suffer from reduced gain. To compensate for this problem, these two cylinders were operated at 100V higher than the nominal operating voltage of the chamber. A more deleterious effect from the outer ground planes was that the field would “leak” out

| Layer 16 ($PH$) | Layer 16 Voltage (V) | Layer 17 Voltage (V) |
|----------------|----------------------|----------------------|
| 252            | 2050                 | 2150                 |
| 148            | 2050                 | 0                    |
| 186            | 2100                 | 0                    |
of the cell creating asymmetries in the drift-time relationship on the right and left sides of the cell. The measured versus fit drift distance for the right and left sides of a wire are shown for an outer layer (Figure 3.10) and an inner layer (Figure 3.11). Note the asymmetry in the outer layer. An analogous effect was observed in the vertex detector, which has a similar grounding arrangement. This problem was partially calibrated away by having separate drift-time relationships with right-left asymmetries for the inner and outer layers. A plot of the \( r - \phi \) residuals from hadronic events is shown in Figure 3.12. Fitting the central peak to a Gaussian yields \( \sigma_{r\phi} = 160 \pm 10 \) microns.

Figure 3.8: Measured test lab spatial resolution \( \sigma_{r\phi} \) verses chamber voltage.
Figure 3.9: Test lab pulse height resolution $\frac{\sigma}{\text{peak}}$ versus chamber voltage.

Since the primary interaction of charged particles and a drift chamber detector is through ionization, it becomes possible to use the measured pulse heights to perform particle identification based on energy loss ($dE/dx$). The amount of energy lost for a relativistic charged particle traversing a medium is predicted from the Bethe-Bloch equation:

$$\frac{dE}{dx} = \frac{A}{\sin \Theta} \cdot \frac{1}{\beta^2} \cdot \left[ \ln \left( \frac{2m_e \beta^2 \gamma^2}{I_0} \right) - \beta^2 \right]$$

where $A$ is a term related to the medium traversed, $\frac{1}{\sin \Theta}$ the path length, $I_0$ the ionization potential, $\beta, \gamma$ the canonical relativistic variables. To engineer such measurements a precarious balance must be struck between preserving
Figure 3.10: Measured vs fit drift distance for the left (cross) and right (diamonds) side of the cell in layer 17. The units are in meters.
Figure 3.11: Drift time relations for left and right sides of a normal layer.
the spatial resolution which is improved with increasing gas gain, and the dE/dx measurements which deteriorate with increasing gas gain. Timing circuits are equipped with discriminator circuits which trigger once a pulse amplitude greater than a specified size is registered. It is desirable to have the circuit respond as quickly as possible to a given pulse in order to obtain the best resolution. The path of least resistance is to make the gas gain as large as possible, providing hefty pulses from the chamber which are easy to trigger on. This approach was the initial operational mode of the CLEO drift chamber, where
only timing information was collected. To also make ionization measurements, the gas gain must be lowered. Since the incoming pulses are smaller the system must be capable of operating at a reduced threshold. The major obstacle is that lowering the sensitivity means that noise hits may penetrate the system, and then real information can be washed out by the background. Noise immunity was enhanced by mounting preamplifiers, 24 channels to a preamplifier card, directly on to the drift chamber face. The entire drift chamber was read out one end face. The boosted signals were transmitted via a symmetric differential transmission system on a 7.5m flat twisted pair cable, to receivers on 48 channel data cards, located in standard CLEO readout crates outside the detector. Each data card channel contained timing and pulse height circuits which analyzed the same incoming pulse. The operating sensitivity was 300 nanoamps, and the resulting dE/dx resolution was in the range of 10-14%. Part of the difficulty in performing dE/dx measurements was that the open rectangular cell geometry was not ideally suited for charge collection. The system was engineered and implemented by John Dobbins and Don Hartill. The calibration and analysis of the dE/dx information was primarily done by Thierry Copie and Tom Ferguson, while the timing performance system was the responsibility of Paul Avery and the author. Although the dE/dx resolution was somewhat inferior to the outer detector, simply being able to have information associated with every track (especially at low momentum) dramatically improved the particle identification performance of the detector.
3.2.4 End Cap Shower Counters

Beyond the drift chamber face on both sides of the detector are the end cap shower counters. This detector is formed from aluminum proportional tubes and lead, with 21 layers of the device oriented in 3 groups oriented at 120° to each other. The energy resolution is $\sigma_E/E = 0.39/\sqrt{E}$. The primary function of the device is to provide a measurement of the luminosity from Bhabha scattering.

3.2.5 Superconducting Coil

Since in the initial CLEO conceptual design the particle identification detectors were located beyond the coil, it was necessary to produce the field using as little material as possible, hence the decision was made to use a superconducting coil. The coil is 3.1 m in length and 2 m in diameter. The winding is made of Nb-45% Ti in a copper matrix. The standard running conditions mandate a 1 Tesla field which requires a current of 1500 amps. Measurements of the field have determined uniformity to within 2%. The net material, including the cryostat is 0.7 radiation lengths. This material ranges out pions with momenta less than 150 MeV, kaons with momenta less than 400 MeV, and protons with momenta less than 600 MeV.

3.2.6 dE/dx

Immediately outside the coil is the dedicated dE/dx system. Each octant contains 124 modules, each of which contains 117 wires oriented perpendicular to
the beam line. The large number of wires was desirable to obtain a high statistics measurement of the average pulse height, which CLEO calculates from the lowest 50% of the measurements to avoid the long tail associated with the Landau distribution. The detector was operated on a gas mixture of 91% argon and 9% methane at a regulated pressure of 45 psia. The resolution, as measured by the peak divided by the standard deviation of the distribution is 5.8% for hadrons. Despite this fine performance, the inability to distinguish multiple tracks and the loss of low momentum tracks where hadrons are most easily separated seriously undermined the usefulness of this device.

3.2.7 Time of Flight

Occupying the next radial position in the octants is the time of flight system. Each octant contains twelve $2.03 \times 0.312 \times 0.0023$ m scintillation counters. Each scintillator is read out on one end by an Amperex XP-2020 photomultiplier tube. The TF geometry per octant is shown in Figure 5.13. The counters operate using two discriminators sensitive to pulse height from the photomultiplier tube. The time for a given hit is then determined by extrapolating to zero pulse height and correcting for travel time in the scintillator medium, as well as measured pulse height. The rms resolution of the detector is 350 psec.

TF hits are matched to drift chamber tracks, and the flight time coupled with the tracks momentum is used to determine the mass of the track. In addition the TF plays a vital role in the experimental trigger.
3.2.8 Octant Shower Counters

The last device in the octant is the octant shower counter usually referred to as the RS. It is a twelve radiation length thick lead proportional tube sandwich, operated on a gas mixture of 91% argon and 9% methane. It covers a solid angle of $\Omega/4\pi = 0.47$ and the device performance can be characterized by $\sigma_E/E = 0.17/\sqrt{E}$ with E in GeV. Neural particle generate characteristic electromagnetic showers and can thus be identified. Charged particles also are detected, though the response is much weaker, allowing electrons to be distinguished from other charged particles. The instrument is also used to calibrate the luminosity using large angle Bhabha events, and serves in the experimental
3.2.9 Muon Detectors

The muon system is the outermost element and surrounds the CLEO detector in a box-like geometry. Beyond a steel hadron absorber of 6 to 12 hadron interaction lengths are two orthogonal sets of drift chamber planes with cell widths of 10 cm, operated in a 50-50 argon ethane gas mixture. The solid angle coverage is $\Omega/4\pi = 0.72$. The hadronic faking for tracks in the 1 GeV range is 1% while the muon detection efficiency is $\sim 30\%$.

3.2.10 Luminosity Monitors

CLEO uses 3 devices to measure the luminosity, the two shower counters described above, and a of set two scintillation-lead shower counter units (LM) to detect small angle bhabha scattering. Each unit contains 4 sections, and are mounted at small angles to the beam trajectory. The device is triggered by hits in any of sections diametrically opposite from the beam spot. This device is useful for measuring relative luminosity, since the orientation of the two units cannot be measured reliably enough to match the rapidly varying Bhabha cross section at those angles. The absolute measurement of the luminosity comes from the two shower counters, and is determined with a systematic error of 2%.
3.2.11 Experimental Control

Information from the individual channels must be sampled, organized, and stored in a coherent fashion for each accepted collision. In the CLEO system each associated detector circuit occurs redundantly in groups of 24-60 on discrete data cards housed in water cooled “crates.” The crate supplies power for the electronics, and contains a “controller” which directs the readout of the digitizing electronics. Data from the channels is converted to a voltage and stored on buffered capacitors which are digitized by the 12 bit ADC of the controller. All crates are connected to a 16 bit data bus known as the y-bus, which is interfaced to the crate via the controller. The y-bus was driven by a VAX-11/750 computer during this run period. The 750 collects the data and writes it to 6250 bpi magnetic tape. An 8 bit data line (x-bus), also driven by the 750, is used to control and monitor detector functions such as high voltage and calibration pulsing. In all, nearly 80 crates comprise the data acquisition system. Virtually all of the readout electronics and crate controller system were designed by members of the CLEO collaboration. A block diagram of the control system is shown in Figure 3.14.

3.3 Data Selection

In the vast majority of beam crossings no interesting interaction takes place. A fast trigger must be able to sample the initial response of the detector and decide to pursue the event before the next collision takes place. Accepted events are recorded merely as addresses and values, these must be converted into a format suitable for physics analysis. This “arduous” process is collectively re-
Figure 3.14: CLEO control system.
ferred to as compress. Detectors must be calibrated and highly complex algorithms sift through the data to reconstruct the interactions of the decay products with the detector. At this stage the data is analyzable, however since the trigger requirements are fairly loose, some of the events at this stage are still uninteresting. Selective event classification schemes are activated to filter out the types of events desired for particular analysis needs. These facets of the analysis shall be examined here.

3.3.1 Experimental Trigger

The experimental trigger allows for reduction of the 1-3 MHz crossing rate to a raw data recording rate of 1-2 Hz, with an average live time fraction of 0.9. The detector elements which play a pivotal role in the trigger system are the VD, DR1.5, TF and the RS. To a minor extent MU and ES are also involved. Channels from the various detector elements are ganged together to provide a course overview of the detector response. Because of accelerator development work that transpired concurrently with the data acquisition program, two trigger modes were used in this data set. They are equivalent from a physics standpoint, differing primarily with the speed in which the decisions had to be made. In 7 (3) bunch mode the crossings occurred every 360 (854) nsec.

In the tracking chambers, four wires are OR-ed together to form a fast-or bit. The patterns of these fast-or bits are correlated in a track segment processor, which selects acceptable topologies. In the vertex detector, the inner and outer five layers are each segmented into blocks, where a block consists of the fast-or bits of five vertically adjacent layers. The five layers in a block are or-ed,
the block being on if four of the fast-or bits are on. A physical combination of an inner and outer block turns on a VD bit. In the drift chamber blocks are formed from 3 vertically adjacent axial layers, containing all the fast-or bits in an approximate $24^\circ$ phi slice for each layer. A DR bit is set if 2 of the 3 layers in a block are hit. The nine axial DR layers contribute to a total of 3 possible drift chamber bits. The VD and DR bits are correlated in a loose road in $\phi$ to form track track segments. A “medium track” consists of a VD bit and the first two DR bits, and a long track is defined by all four bits being on.

Two signals are sent from each time of flight octant, which are the or’s of groups of six scintillators which share the same z side of the octant. The TF bit pattern is analyzed for acceptable patterns, such as TFNADJ (two TF hits in nonadjacent octants) or BBTF (TF hits in diametrically opposite octant segments). Trigger information is also examined from the shower detector. Pulse heights are accumulated with an analog or of all 24 channels on a data card, these are likewise summed on the crate level and then on the octant level (two crates). The octant energy threshold (OCT) was approximately 1 GeV. Due to the analog summation, the input to the trigger discriminator was susceptible to noise. In particular an odious glitch in the baseline of the discriminators occurred after analog reset. This partially contributed to the loss of 2 (1) crossings in 7 (3) bunch mode to allow for settling time. Analog reset was fired every 19 (16) crossing during this running period. This effect coupled with other engineering instabilities prompted a sweeping renovation of the energy trigger in the fall of 1986. This upgrade was developed by John Dobbins with assistance from the author. Preamplifiers (gain 25) were placed at the octant sum junction making the transmitted signals less susceptible to noise. The boosted signals were differentiated with a time constant of 560 nsec at the discriminator input.
to smooth out the $\overline{AR}$ glitch. The improvement system has proved to be significantly more robust than its predecessor. Similar information from the end cap shower detectors (ES) was supplied to the trigger.

The trigger logic analyzes the input from the detectors in a two tiered approach. The level 1 trigger is enabled by a hit in the TF, OCT or ES lines. The singles rate with beam in the machine is about 5 KhZ for the ES and OCT lines. Level 1 initiates a search through all of the possible level 2 topologies, which are much more restrictive. The CLEO trigger logic implemented during this data taking is organized in Table 3.3. The level 2 lines can be physically grouped into a hadronic element which fires the tracking devices and some of the outer detectors (to avoid beam-wall, beam gas triggers) and a two track trigger for QED processes. Because of higher rates, some lines are prescaled. Only 1 of 64 2ES triggers in accepted, while the MU2 $1 + TF$ line is accepted $1/64$ th ($16/64$) of the time in 7 (3) bunch mode. Each level 2 search takes roughly $2.5 \mu$sec, and a successful $L1 \cdot L2$ trigger initiates a digitization and data readout sequence lasting 20 msec.

### Table 3.3: CLEO Trigger Logic

| Level | option1 | option2 | option3 | LOGIC | prescale |
|-------|---------|---------|---------|-------|----------|
| 1     | TF      | ES      | OCT     | OR    | N/A      |
| 2     | 1M2L    | TFNADJ  | AND     | N/A   |
| 2     | 2M      | BBTF    | AND     | N/A   |
| 2     | 1M1L    | 1 + TF  | 1 + OCT | AND   | N/A      |
| 2     | OCTOPP  |         | AND     | N/A   |
| 2     | 2ES     |         | AND     | 1/64  |
| 2     | MU2     | 1 + TF  | AND     | 1/64 (16/64) |
3.3.2 Compress

The process of turning the raw data into analyzed events is a complex, iterative, feedback system. A diagram representing the compress system used to analyze this data set is shown in Figure 3.15. Data runs are taken in two modes, physics running and calibration. As seen from the diagram, the full calibration of the detector uses both calibration and physics data. Calibration runs for the drift chamber were taken every 2 to 3 days, and consisted of pulsing the chamber to determine the pedestals and gains for the timing channels and the pedestals for the pulse height channels.

As the physics data was gathered, it was subjected to a “first pass” track finding algorithm (SOLO [47]) to find seed tracks. Those events which were judged interesting by the track finder were then segregated into QED and hadronic event candidates. For the drift chamber, information from hadronic events at this stage was used to calibrate out global changes (such as machine timing) and systematic shifts in the electronics or performance of the drift chamber (non standard operating conditions). After completion of this stage all the devices have been fully tuned, and the final analysis algorithms are run. A second, more rigorous, track finding algorithm (DUET [48]) uses the seed tracks as initial conditions and systematically searches the data for additional tracks. At this stage information from the vertex detector is applied, and the intervening material between the vertex detector and drift chamber is compensated for by allowing a “kink” in the track.

Those fully analyzed events which are classified as hadronic are compacted into a format called ROAR [49]. This format strips the event down to the minimal amount of information necessary for analysis, and computes and stores
Figure 3.15: CLEO compress data flow for hadronic type events.
a variety of invariant mass combinations for quick analysis. The $\Upsilon(4S)$ data sample of 113.7 pb$^{-1}$ produced 113 2400 foot, 6250 bpi tapes of hadronic events. After being ROARed, the set fit on to five such tapes and could be analyzed in 1-2 VAX 8600 cpu hours.

3.3.3 Hadronic Event Selection

All events used in this analysis have been classified as hadronic in nature. The primary selection criteria evolve from charged particle tracking in the sequence single tracks, event vertex, and event shape. To eliminate beam gas, beam wall, cosmic ray showers, and radiative bhabhas, the energy response of the octant shower detectors is also employed.

To begin, the charged multiplicity of the event is examined. A charged track is considered good if no more than one following tests is failed:

1. The z value at the distance of closest approach to the beam line (DOCA) is less than 50 mm from the origin.
2. The track’s DOCA with respect to the nominal beam spot in the r-$\phi$ plane is less than 5 mm.
3. The average residual from the 5 parameter helical fit must be less than 700 microns.
4. The track is defined using at least 8 drift chamber layers.

The next level of selection is a crude classification of a hadronic event:

1. Tracking.
(a) At least 3 good charged tracks, which may come from the primary or a secondary vertex, excluding those consistent with originating from a converted photon.

(b) At least one of the tracks must originate from the primary vertex with a DOCA less than 5 mm.

2. Vertex.

(a) The x and y position of the vertex must be within 20 mm of the nominal x-y beam position.

(b) The z position of the vertex must be within 50 mm of the origin.

3. Energy.

(a) The sum of the energies of the charged and neutral tracks in the event must exceed 15% of the center-of-mass energy.

All events that have survived to this stage are finally subjected to a stringent, well developed set of cuts known as morcut \[50\]

1. Tracking.

(a) At least 15% of the hits in the drift chamber must be associated with tracks.

(b) The fraction of bad tracks to good tracks must be less than 1.15

2. Energy.

(a) The total energy of the charged and neutral tracks in the event must exceed 30% of the center-of-mass energy.
(b) The energy measured in the octant shower counters must exceed 250 MeV.

3. Topology.

(a) The event must not be consistent with a radiative bhabha.

(b) The event must not be classified as a beam wall collision.

(c) The Fox-Wolfram parameter $H_1/H_0$ which measures momentum imbalance must be less than 0.4.

(d) The Fox-Wolfram parameter $H_2/H_0$ which measures event shape (0 = spherical, 1 = back to back 2 track event) must be less than 0.98.

The non hadronic background for events which have passed morcut is only a few percent, while a Monte Carlo simulation of D decays shows that 98 percent of the events that are successfully reconstructed pass morcut.
CHAPTER 4
TECHNICS

After the data has been collected and processed, additional “second order” analysis systems need to be developed to fully exploit the underlying physics. Here we detail the other relevant tools essential to this analysis. They are: reconstruction of secondary vertices, charged particle momentum correction, hadron identification, and Monte Carlo simulation. Also presented is a discussion of particle decay kinematics, which forms the basis for a number of analysis procedures used in the following chapters. To conclude, a presentation of the analysis architecture used to derive the results of this thesis shall be crystallized.

4.1 Reconstruction of Secondary Vertices

As discussed in chapter 1, the tremendous diversity in the strength of the four forces causes a difference in the decay times associated with each force of several orders of magnitude. At CESR energies, strong and electromagnetic processes evolve at such a rate that they cannot be distinguished from the primary vertex. Particles that decay weakly, in contrast, can move from order 100 microns (charm, $\tau$ decays) to tens of millimeters ($K^0_S$, $\Lambda$ decays). At CLEO, these particles are reconstructed from their decays into purely charged modes. For the first class of decays, the decay length is in the realm of the extrapolation error of the individual tracks. Evidence for these vertices comes from the observation of systematic offsets from zero in distributions of extrapolated decay lengths and impact parameters. The vertices of the second group are often clearly visible with decent event display graphics. Decays of $K^0_S$’s and $\Lambda$’s have been historically referred to as “vees” because their 2 prong decay topology resembles the
Since the track finding algorithm performs searches for individual tracks, an additional program is required to isolate vees. The vee finder examines all good tracks that have a full 3-D reconstruction, with the two fold mission of finding pairs of tracks which come from a vertex distinct from the common event vertex, and properly determining the momentum of the original neutral particle. Vectors of the two daughter particles must be re-evaluated, since they were determined at the point of closest approach to the drift chamber origin. The algorithm used in this analysis was developed by M. Ogg and M. Mestayer. The search is initiated by parametrizing the tracks as circles in the r-φ plane. This is a preferred starting point since the r-φ tracking resolution is about a factor of five better than that in z. Vee candidates are formed from pairs of oppositely charged tracks where the sum of the absolute values of the track DOCAs exceeds 1mm. The two circles are tested for intersection, and only those candidates with 2 intersection points are considered. If either solution is consistent with originating from the primary vertex (z of the vertex less than 3 cm from the origin, and radius of the vertex in the r-φ plane with respect to the average beam spot (rv) less than 0.3 cm ) the candidate is rejected. Each solution is required to have rv in the range of 0.5 cm to 50 cm, and the reconstructed vee vector must point away from the origin, as determined by requiring the normalized dot product of the vee vector and the vector drawn from the origin to the secondary vertex be positive. The selection procedure is completed with the definition of a vee quality factor:

\[ \chi^2_v = \left( \frac{\Delta z}{\sigma_z} \right)^2 + \left( \frac{b_v}{\sigma_b} \right)^2 \]

where the parameter Δz is the z difference of the two candidate charged tracks at their r-φ intersection point, and bv represents the impact parameter of the fit
secondary vertex with respect to the run by run average x-y position of the beam spot. The nominal rms errors $\sigma_z$ and $\sigma_b$ were chosen from Monte Carlo studies to be 10 mm and 2 mm, respectively. Application of a cut on $\chi^2_v$ improves the purity of the $K^0_S$ sample by rejecting poorly determined and incorrect vertices. Vee candidates are subject to a loose cut of $\chi^2_v \leq 12.0$ in order to provide a maximal vee base available to the user. In practice, the author uses $\chi^2_v$ cuts of 2-3 for $K^0_S$'s. Should both solutions pass all cuts, the solution with smaller $\Delta z$ is selected. The variables involved in the vee finder as discussed above are visually presented in Figure 4.1. Once the position of the secondary vertex has been determined, the momenta of the two daughter tracks are redetermined with respect to the new vertex, and added to form the vee three vector.

Two measures of the robustness of the vee finder are the reconstruction efficiency and FWHM (width) of the vee invariant mass peak. We examine the performance with respect to the decay $K^0_S \rightarrow \pi^+\pi^-$, which is vital to this analysis. Details on $\Lambda$ reconstruction can be found elsewhere [52]. Figure 4.2 contains a plot of the efficiency for reconstructing $K^0_S \rightarrow \pi^+\pi^-$ versus $K^0_S$ momentum, as determined form Monte Carlo simulation. The efficiency folds in hadronic event selection, single particle tracking, and secondary vertex finding. The efficiency is seen to plateau at $\approx 50\%$ from 1 to 3 GeV. The momentum dependence of the mean and width of the $K^0_S$ observed in the data are shown in Figure 4.3 and Figure 4.4. The invariant mass spectrum is fit to a Gaussian signal and a polynomial background. A systematic deviation from the known $K^0_S$ mass versus momentum is prominent at at low $K^0_S$ momentum. This is due to energy loss of the $K^0_S$ daughters. The correction for this effect will be presented in the next section. The width obeys an approximately linear relationship as a function of momentum. Convolving this with the observed $K^0_S$ momentum spectrum yields
Figure 4.1: The $r$-$\phi$ intersection of two helical tracks. The intersection occurs at the point $(x_V, y_V)$ at radius $r_V$. The reconstructed vector has impact parameter $b_V$.

an average value for the width of the $K_S^0$ mass peak of 9 MeV. The behavior of both the mean and width are in agreement with Monte Carlo simulation of this decay in the CLEO detector.
4.2 Corrections to Charged Particle Momenta

Tracks in the vertex detector appear largely as straight line segments, making track curvature (hence momentum) measurements the sole responsibility of the drift chamber. Energy lost in interactions with material before the drift chamber will manifest itself in systematically lower invariant masses for reconstructed particles. The material preceding the drift chamber is collected in Table 4.1. The material corresponds to the beam pipe, inner and outer VD supports, and drift
Figure 4.3: Observed mean of the $K_S^0$ mass peak versus $K_S^0$ momentum.

Table 4.1: Material Preceding the Drift Chamber Detector

| outer radius (cm) | Material                        | $dE/dx$ (MeV) |
|-------------------|---------------------------------|---------------|
| 7.57              | silver coated beryllium         | 0.246         |
| 8.02              | Carbon filament tube            | 0.322         |
| 16.44             | VD gas + Carbon filament tube   | 0.208         |
| 17.44             | Carbon filament tube            | 0.270         |

As described earlier, we choose the Bethe-Bloch form $\frac{dE}{dx} = \frac{A}{\sin \theta} \cdot \frac{1}{\beta^2} \cdot \left[ \ln(2m_e^2 \gamma^2 / I_0) - \beta^2 \right]$ to model the expected energy loss. The algorithm used was developed by P Avery [53]. The average momentum loss $\frac{dp}{dx} = \frac{1}{\beta} \cdot \frac{dE}{dx}$ is calculated based on the track’s momentum and cell entrance angle, for several different mass hypotheses. The momentum loss is added to
the measured track in such a way that only the magnitude, not the direction
is altered. The average corrections for pions, kaons, and protons are shown in
Figure 4.5. The net result when calculating the invariant mass of several charged
tracks is to shift the peak of the mass distribution upward 1-2 MeV, with the
FWHM of the distribution largely unaffected. Monte Carlo simulations mirror
the effects of the correction observed in the data.

We can now re-address the properties of $K_S^0$'s after energy loss corrections
are applied. The observed $K_S^0$ mass (Figure 4.6) is found to be $497.8 \pm 0.1$ MeV,
in good agreement with the world average of $497.72 \pm 0.07$ MeV. The $K_S^0$ mass as
Figure 4.5: Average energy loss corrections for tracks as a function of momentum and mass hypothesis.

a function of momentum (Figure 4.7) is now centered about the correct mass to within a few tenths of an MeV over the entire momentum range.
Figure 4.6: Invariant mass of $K^0_S$ candidates with $\chi^2 \leq 3.0$ from data.

4.3 Hadron Identification

Hadron ($\pi$, $K$, $p$) identification is accomplished by coherently combining information from the three detector elements with this capability; DR1.5, $dE/dX$, and TF. Complete details of the hadron identification system can be found in [52]. The hadron identification algorithm is unsophisticated in nature yet powerful in performance. For each device, a “probability” is calculated for each of the three mass hypothesis $p_{\text{HYPi}}$ ($1 = \pi$, $2 = K$, $3 = p$) defined by:

$$p_{\text{HYPi}} = \exp^{-\frac{1}{2}} \left( \frac{M_{\text{DEV}} - M_{\text{HYPi}}}{\sigma_{\text{HYPi}}} \right)^2$$
Figure 4.7: The $K_S^0$ mass as a function of momentum, after application of the energy loss correction.

where $M_{\text{DEV}}$ are the measured values (mean of lowest 50% pulse height for DR1.5 and dE/dx, time for the TF), $E_{\text{HYPi}}, \sigma_{\text{HYPi}}$ are the expected measurement and resolution for each device, respectively. The distributions of $M_{\text{DEV}}$ for each of the three devices (Figures 4.8 - 4.10) all show distinctive $\pi$, K, and p bands. $E_{\text{HYPi}}$ and $\sigma_{\text{HYPi}}$ are the expected detector responses for the mean and width of three species, as determined from non-trivial measurements from the data. Since dE/dx and the TF detectors are separated from the drift chamber by the superconducting coil, these devices can only be used if successful matches are made to drift chamber tracks. The overall hypothesis is calculated from the
Figure 4.8: Mean of the lowest 50% of the pulse heights from the drift chamber detector versus track momentum.

Product of the probabilities for each functional device:

\[ P_{\text{HYPi}} = \prod P_{\text{HYPi}} \]

\( P_{\text{HYPi}} \) is set to 1 for each hypothesis if there was no information available. To discern among the species, we define a normalized weight:

\[ W_{\text{HYPj}} = \frac{P_{\text{HYPj}}}{\sum_{i=1}^{3} P_{\text{HYPi}}} \]
Figure 4.9: Mean of lowest 50% of the pulse heights from the dE/dx chamber detector versus track momentum.

The normalized weight is set to zero if there was no information available \( \left( \sum_{i=1}^{3} P_{\text{HYP}i} = 3 \right) \) or the information available favored none of the three hypotheses \( \left( \sum_{i=1}^{3} P_{\text{HYP}i} \leq 0.001 \right) \)

The utility of this approach is pragmatically illustrated for the case of \( D^+ \rightarrow K^-\pi^+\pi^+ \). The \( D^+ \) candidates have momentum greater than 2.5 GeV, with the momentum of pions greater than 0.3 GeV.
Figure 4.10: Measured velocity in the Time of Flight detector versus track momentum measured in the drift chamber.

Figure 4.11 a) shows the invariant mass plot where the kaon candidate has $W_K \geq 0.1$, while in b) the kaon candidate is tightly identified $W_K \geq 0.7$ and the pion candidates are loosely identified $W_\pi \geq 0.2$. The reduction in the background is startling. Obtaining a $D^+ \rightarrow K^-\pi^+\pi^+$ mass peak with a signal to noise of 1 : 1 was a significant experimental achievement. This has allowed CLEO to measure the $D^+$ lifetime with the world’s second largest $D^+$ sample.
Figure 4.11: $D^+ \rightarrow K^-\pi^+\pi^+$ candidates with momentum greater than 2.5 GeV. a) $W_K \geq 0.1$. b) $W_K \geq 0.7$ and $W_{\pi} \geq 0.2$. 
4.4 Monte Carlo Simulation

In order to complete physics measurements, computer modeling of the detector and the 10 GeV $e^+e^-$-environment is required. For this analysis, Monte Carlo methods are used to determine the detector acceptance for specific D decay modes, along with backgrounds from $D_s^+$ decays. We simulate the byproducts of $e^+e^-$ annihilations using a modified \cite{54} version of the LUND \cite{55} QQJET generator. For this analysis all Monte Carlo events were generated using $c\bar{c}$ jets, including gluon radiation. The center-of-mass energy was defined to be 10.56 GeV, approximately the average run energy of this data set. Effects of initial state radiation of the beams were not compensated. The strategy for studying a particular decay mode was to force the particle ($D^0, D^+$) to decay into the mode in question, while the anti-particle ($\bar{D}^0, D^-$) was allowed to decay freely. Each $c\bar{c}$ event thrown was selected for further analysis only if it contained the particle under study. Since much of this analysis involves neutral kaons, when the D meson decayed it was forced only to decay to a $K^0$ or $\bar{K}^0$, the event being selected only if the final state $K^0_{S}\rightarrow\pi^+\pi^-$ was thrown. In this way the Monte Carlo was free from bias from either an excessive number of $K^0_S$’s or charged pions. All efficiencies are calculated with respect to the final state $K^0_{S}\rightarrow\pi^+\pi^-$, and later rescaled for the branching ratio $B(\bar{K}^0\rightarrow K^0_{S}) \cdot B(K^0_{S}\rightarrow\pi^+\pi^-)$

D mesons are fragmented according to the Peterson recipe. Because of the symmetric LUND fragmentation scheme, when two objects are produced in the fragmentation process each with a mass of order the beam energy, the first object tends to receive a larger share of the available energy-momentum. This distorts the fragmentation distribution. This effect is mitigated by calculating efficiencies over a small momentum range and performing a summation.
This method also makes calculation of Monte Carlo parameters insensitive to the fragmentation model.

The collection of vectors and vertices for each selected event are then propagated through a detector simulation to mimic CLEO raw data. Detector response in terms of effective efficiency and resolution are folded into the process. The “false data” is passed through the CLEO compress system, and is subjected to the final tracking algorithm DUET. The data at this stage is identical to real data, and can be analyzed by the same program. The entire process requires \( \approx 1.5 \) VAX 8600 cpu hours for about 1000 events, making Monte Carlo generation one of the most cpu intensive tasks in an analysis project.

The detector resolution in most cases dominates the width of an observed mass peak. Monte Carlo predicted widths have been in respectable agreement with the data, indicating a proper simulation of the tracking. Information for neutral particles is approximated to first order based on their trajectories and shower counter efficiencies. The hadronic event selection criteria requires 250 MeV deposited in the octant shower counters. Thorough modeling of this detector would require the cpu intensive EGS shower simulation, so this constraint was relaxed for Monte Carlo events. While charged particles lose energy as they traverse the detector, response of the devices that measure this energy loss is not simulated. The efficiency for identifying hadrons is measured \([52]\) directly from the data (Figure 4.12) \(|\cos \Theta| > 0.6\) only information from the drift chamber using “pure” samples of pions \((K_S^0 \rightarrow \pi^+\pi^-)\), kaons \((\phi \rightarrow K^+K^-\) and \(D^+ \rightarrow D^0\pi^+, D^0 \rightarrow K^+\pi^-)\), and protons \((\Lambda \rightarrow p\pi)\). These efficiencies are injected into the Monte Carlo by simply throwing a random number between 0 and 1, and comparing it to the expected efficiency. While this empirical approach assures to first or-
Figure 4.12: Efficiencies for detecting kaons with weights \( W_K \geq 0.1 \) (squares), \( W_K \geq 0.4 \) (crosses), and \( W_K \geq 0.7 \) (circles). Also shown are the efficiencies for measuring pions with the same kaon weights. For tracks with \( |\cos \theta| > 0.6 \) only information from the drift chamber is available.
der proper detector response, the selection of the modes to produce the “pure sample” might possibly introduce systematic effects.

4.5 Two Body Decay Kinematics

Important insight into the decay mechanism of a parent particle can be gleaned by examining the angular distribution of the decay products. Observation of an angular distribution consistent with an anticipated polarization (or lack thereof) can help substantiate a particular decay hypothesis. Also, an anisotropy of the angular distribution of the background can be used to enhance the signal. A well quantified formalism \[56\] exists for analyzing decays which are either two body or consist of a series of sequential two body decays. Termed the helicity formalism, the following relation can be derived for the angular distribution of a two body decay in the rest frame frame of the decaying particle:

\[
\frac{d\sigma}{d\Omega_f}(\theta_f, \phi_f) = \sum_{\lambda_1, \lambda_2} \frac{1}{4\pi} \left( \frac{2J + 1}{4\pi} \right)^{\frac{1}{2}} D_{M1}^{J*}(\phi_f, \theta_f, \phi_f) A_{\lambda_1, \lambda_2} \left| \frac{D_{Jm_1}^{J*}(\alpha, \beta, \gamma)}{e^{-i\lambda \gamma} \mathcal{A}_{m_1}^{m_2}(\beta)e^{-im\phi}} \right|^2
\]

Here \(\theta_f, \phi_f\) are the angles of the two body decay axis with respect to the spin quantization (z) axis of the parent. J and M are the spin and z projection of the parent particle. \(\lambda_1, \lambda_2\) are the helicities of the daughters, as defined by \((\lambda_i = \vec{s}_i \cdot \hat{p}_i)\), and \(\lambda\) is the overall helicity \(\lambda_1 - \lambda_2\), which is subject to the constraint \(J \geq |\lambda|\). The \(D_{Jm}^{J}\) functions have the definition:

\[
D_{Jm}^{J}(\alpha, \beta, \gamma) = e^{-i\lambda \gamma} \mathcal{A}_{m_1}^{m_2}(\beta)e^{-im\phi}
\]

where \(d_{m_1}^{J}\) are the “D” functions \[57\]. The “D” functions are connected to the spherical harmonics through the relation:

\[
d_{m_0}^{J} \propto Y_{m_0}^{J}(\theta, \phi)e^{-im\phi}
\]
The factor $A_{\lambda_1,\lambda_2}$ is related to the decay matrix element, but does not contain any angular information. We can neglect the overall normalization (and in most instances the $\phi$ dependence) to get at the the general character of the decay:

$$\frac{d\sigma}{d\Omega_f}(\theta_f, \phi_f) \propto \sum_{\lambda_1,\lambda_2} |d^j_{m,1}|^2$$

The first case to consider is the frequently encountered decay of a pseudoscalar decaying into two pseudoscalars ($P \rightarrow PP$). Since everything in this decay is in a state of zero angular momentum, $J=M=\lambda=0$ and

$$\frac{d\sigma}{d\Omega_{PP}}(\theta_{PP}, \phi_{PP}) \propto |Y_0^0|^2 = 1$$

hence the decay axis (PP) is isotropically distributed. Since there exists no particular quantization axis, we choose as the reference direction to be that of the momentum of the parent particle in the laboratory frame. We define the normalized dot product of the decay axis and the reference vector to be the rest frame decay angle (RFDA).

The second angle of interest results from a polarization in the decay chain pseudoscalar $\rightarrow$ pseudoscalar, vector ($P \rightarrow PV$), where the vector particle subsequently decays into two pseudoscalars ($V \rightarrow PP''$). Because the parent has $J=0$, the value for the difference of the helicity state of the daughters $\lambda_1 - \lambda_2 = \lambda \equiv 0$. This polarizes the vector particle into the helicity zero state. The PV decay axis is also isotropically distributed.

$$\frac{d\sigma}{d\Omega_{PV}}(\theta_{PV}, \phi_{PV}) \propto |Y_0^0|^2 = 1$$

We now examine the decay of the vector particle in its rest frame. The quantization direction has already been defined as the PV decay axis. The two vector daughters $PP''$ are then boosted into the vector rest frame. The vector particle
Figure 4.13: a) Rest frame decay angle (RFDA) for a $P \rightarrow PP$ decay. b) polarization angle $\cos \Theta_{p'p''}$ for $P \rightarrow PV, V \rightarrow PP$. 

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has \( J = 1, M = 0 \), and both daughter particles have helicity zero, by definition. The angular distribution of the \( P'P'' \) decay axis in the \( V \) rest frame with respect to the \( PV \) decay axis (Figure 4.13) is

\[
\frac{d\sigma}{d\Omega_{P'P''}}(\theta_{P'P''}, \phi_{P'P''}) \propto |Y_{10}|^2 = \cos^2 \theta_{P'P''}
\]

### 4.5.1 Satellite Mass Peaks

The ramifications of this polarization are profound. In the laboratory frame, the two vector daughters will have diametrically opposite values for their total momentum, with one being produced almost at rest. The invariant mass formed from 2 of the 3 final state particles, neglecting the particle produced with very small momentum, will differ in mass from the parent mass largely by the mass of the neglected daughter. Such a distribution has a distinctly non Gaussian shape, and is deemed a “satellite” mass peak. Figure 4.14 shows a Monte Carlo simulation of the decay \( D^0 \rightarrow K^- \rho^+, \rho^+ \rightarrow \pi^+ \pi^0 \), where the invariant mass is formed from the \( K^- \) and \( \pi^+ \). The two lobed structure corresponds to the cases when the \( \pi^+ \) is the slow and fast particle from the \( \rho^+ \) polarization. Because the low mass lobe is much broader and in a region of extremely high background, it is not visible in the data. Information can be extracted from satellite peaks, provided a proper form can be determined to fit the peak. Extensive study by the author has demonstrated that a form termed the Bifurcated Gaussian provides very agreeable fits to satellite peaks. This curve is composed of two Gaussian peaks that share a mean, but have different areas and widths on either side of the mean. Continuity of the curve is guaranteed by forcing the constraint

\[
\frac{A_1}{F_1} = \frac{A_2}{F_2}
\]

Figure 4.15 shows the invariant mass distribution for of \( D^0 \rightarrow K^- \pi^+ \). A satellite mass peak is clearly evident in the 1.4-1.6 GeV region. The \( D^0 \) peak is
Figure 4.14: Monte Carlo simulation of the decay $D^0 \rightarrow K^- \rho^+, \rho^+ \rightarrow \pi^+ \pi^0$. The Invariant mass spectrum is formed from the $K^-$ and $\pi^+$. 

fit to a Gaussian peak, while the satellite peak is fit to a Bifurcated Gaussian. Fitting a satellite peak to a Gaussian shape underestimates the area by 10-15%.
Figure 4.15: Invariant mass spectrum for the decay $D^0 \rightarrow K^- \pi^+$. A satellite peak is observed in the 1.4 - 1.6 GeV region.

4.6 Analysis Architecture

Having collected all the techniques and machinery employed in this project, let us examine how they are used in concert to derive results. Since most of this analysis involves the reconstruction of exclusive decays using charged particle tracking and the invariant mass technique, the first step required is to prepare a combinatoric driver. This code generates suitable combinations of charged tracks and secondary vertices specific for each decay mode. The most diffi-
cult task in using the invariant mass technique is to isolate a signal. This is accomplished by developing a set of physically motivated cuts to simultaneously maximize the observed signal and provide the highest possible rejection of spurious combinations which form the background. For this analysis, the cut development was accomplished by writing out a binary word for each event containing all the relevant information into a disk file, which ranged in length from 2.5 - 100 K. This facilitated a decrease in cut development by a large factor, since after passing through the data once to create the file, the entire set could be reanalyzed in a matter of minutes. More importantly, this made the cut development largely an interactive process.

After a signal has been isolated, the detector acceptance must be determined from Monte Carlo. Two methods are used to extract the Monte Carlo parameters. This redundancy assures that these parameters are properly determined, as well as demonstrating the correctness of the combinatoric driver. In the first method the Monte Carlo events are subjected to the full data analysis system. Since the Monte Carlo was generated with D₀'s and D⁺'s decaying into the desired final state, decays of the antiparticles into analogous final states are eliminated. The second (TAGGER) method involves using additional stored information on the Monte Carlo data which “tags” the original Monte Carlo tracks to the tracks eventually found in the simulated event by the track finder. The decay products in the event are traced down to particles observable in the detector, these trajectories are compared to the found tracks, and the Monte Carlo trajectory associated with the most hits on the track is “assigned” to the track. For each event the initial D⁺ or D₀ Monte Carlo track is found, and its generated momentum is calculated. The decay daughters are found, and tested for tagged matches to drift chamber tracks. If all daughter tracks are matched, all stan-
standard analysis parameters for that group of daughter tracks are calculated. The Monte Carlo signals are fit to a Gaussian signal and a polynomial background, and the Monte Carlo parameters are calculated from a weighted average of the two methods. The agreement between the two methods is excellent, as evi-

denced from the example listed in Table 4.2 which compares the two methods in calculating the parameters for the decay $D^+ \rightarrow K^0_S \pi^+$. Although we believe the extracted parameters to be correct for a given Monte Carlo, we add an additional 5% systematic error for uncertainties in the generator and the simulation of the detector. When determining fragmentation distributions, because of limited statistics, when performing fits to the data we choose to constrain the mean and the FWHM of the Gaussian signal to be within $\pm 3$ standard deviations of the Monte Carlo parameters. This is a more physical approach than an exactly constrained fit (with or without smoothing), as it can adjust for systematic errors in the Monte Carlo simulation. The analysis structure is summarized in Figure 4.16.

| $x$ range  | Method 1     | Method 2 (TAGGER) |
|------------|--------------|------------------|
| $0 \leq x < 0.375$ | $0.381 \pm 0.009$ | $0.373 \pm 0.009$ |
| $0.375 \leq x < 0.51$ | $0.346 \pm 0.009$ | $0.347 \pm 0.009$ |
| $0.51 \leq x < 0.625$ | $0.358 \pm 0.009$ | $0.352 \pm 0.009$ |
| $0.625 \leq x < 0.750$ | $0.371 \pm 0.009$ | $0.365 \pm 0.009$ |
| $0.750 \leq x < 0.875$ | $0.363 \pm 0.010$ | $0.357 \pm 0.010$ |
| $0.875 \leq x \leq 1.0$ | $0.369 \pm 0.020$ | $0.365 \pm 0.020$ |
Figure 4.16: Data analysis flow diagram.
The \( D^+ \) meson is the spin 0 charged ground state of the charmed, nonstrange meson. Its mass is measured to be \( 1.8693 \pm 0.0006 \) GeV. Initially, the properties of \( D \) mesons were studied in \( e^+e^- \) experiments running on the \( \psi'' \) resonance. This was an optimal running location since the \( \psi'' \) decays into a \( D\bar{D} \) pair produced almost at rest. Experiments performed at this energy were able to accumulate large samples of \( D \) mesons, and successfully reconstruct a vast number of \( D \) decay modes. Not accessible to these experiments was the opportunity to study the charm hadronization process or to measure the lifetimes of the \( D \) mesons.

In \( e^+e^- \) annihilations at higher center of mass energies, charm production accounts for \( \frac{4}{10} \) of the total hadronic cross section, allowing for a fairly copious production of charmed particles. This is offset by a superabundant combinatorial background, which impedes the isolation of charm decay signals. For the \( D^0 \) and \( D^{*+} \) mesons, this problem is obviated by exploiting the well known mass difference of the two states using the decay chain \( D^{*+} \to D^0\pi^+ \), \( D^0 \to X \). Using this technique, \( e^+e^- \) experiments in the energy range of \( 10 - 50 \) GeV have been able to study the properties of these two mesons. No such kinematical artifice exists for \( D^+ \) mesons, as the mass difference between the \( D^{*0} \) and the \( D^+ \) forbids charged pion transitions. Additional difficulties complicate the reconstruction of \( D^+ \) hadronic decays. The \( D^+ \) has a large branching ratio into semileptonic modes (\( B(D^+ \to e^+X) = 17.0 \pm 1.9 \pm 0.7 \% \)) \[58\], and the mass splitting of the \( D \) and \( D^* \) mesons allows cascades from the \( D^{*+} \)’s to \( D^+ \)’s to occur only one third of the time. Since \( D^* \) mesons are favored in the charm hadronization process \( \left( \frac{e^+e^- \to D^*}{e^+e^- \to D_+D^-} \right) \approx 0.7 \pm 0.2 \), this manifests itself in an inclu-
sive $D^0$ cross section which is about a factor of two larger than the $D^+$. Thus, the lack of a convenient reconstruction method and low production rates have inhibited experimental study of $D^+$ mesons.

The excellent charged particle tracking system coupled with hadron identification capabilities make the CLEO experiment a highly competitive facility for studying charmed particles, including $D^+$ mesons. The sample of reconstructed $D^+$ mesons used in this analysis, on a mode by mode basis, is larger than that of the MARK III experiment. The MARK III group has been a leading source of information on D decays, and have measured the greatest number of $D$ meson decay modes. Here we will detail the measurement of three Cabibbo allowed, hadronic decays of the $D^+$. They are $D^+ \rightarrow \bar{K}^0\pi^+$, $D^+ \rightarrow \bar{K}^0\pi^+\pi^+\pi^-$, and $D^+ \rightarrow K^-\pi^+\pi^+$. For comparison and use in future chapters, we also present a measurement of $D^0 \rightarrow \bar{K}^0\pi^+\pi^-$. Signal isolation techniques and corrections for physics backgrounds will be discussed. Fragmentation distributions and relative branching ratios will be compared, and total cross sections for $e^+e^-$ continuum production will be estimated. Information of the CLEO measurement of the $D^+$ lifetime can be found elsewhere [59]. The measurements of the $D^+ \rightarrow \bar{K}^0$ modes are the first to be done at an $e^+e^-$ experiment at an energy greater than the $\psi''$, and the measurements of the fragmentation distributions of these decay modes are the first such measurements.

5.1 Preliminaries

All events under consideration here have passed the hadronic selection criteria. The invariant mass for D candidates is formed only from “good tracks” which
Table 5.1: MARK III D Meson Branching Fractions

| Decay Mode       | Branching Fraction (%) |
|------------------|------------------------|
| $D^+ \rightarrow K^0\pi^+$ | 3.2 ± 0.5 ± 0.2 |
| $D^+ \rightarrow K^0\pi^+\pi^-\pi^+$ | 6.6 ± 1.5 ± 0.5 |
| $D^+ \rightarrow K^-\pi^+\pi^+$ | 9.1 ± 1.3 ± 0.4 |
| $D^0 \rightarrow K^0\pi^+$ | 6.4 ± 0.5 ± 1.0 |
| $D^0 \rightarrow K^-K^+$ | 0.51 ± 0.09 ± 0.07 |
| $D^0 \rightarrow K^-\pi^+$ | 4.2 ± 0.4 ± 0.4 |

have been corrected for energy loss in the material preceding the drift chamber detector. For calculation of continuum cross sections, we used only data taken in the region of the $\Upsilon(4S)$ (77.7 pb$^{-1}$ on resonance, 36 pb$^{-1}$ below resonance). The 33 pb$^{-1}$ of $\Upsilon(3S)$ data is also used for specific measurements, we choose not to use this data to calculate cross sections because of the high combinatorial backgrounds in this region and the possibility of contamination from the process $\Upsilon(3S) \rightarrow c\bar{c}$.

Since $D^+$ mesons can also be produced in the decay of B mesons, this $D^+$ production mechanism must be excluded. A Monte Carlo study of 11.5 K $B\bar{B}$ decays where there was at least 1 $D^+$ in the event found no $B \rightarrow D^+X$ decays where the $D^+$ had a momentum greater than 2.5 GeV. We prudently select momentum cut-off for $D^+$ candidates of $p \geq 2.52$GeV ($x = \frac{p}{p_{\text{max}}} \geq 0.51$) to use the on resonance data. To calculate the differential cross section we elect to use the kinematical variable $x = \frac{p}{p_{\text{max}}} = \frac{p}{\sqrt{E_{\text{beam}}^2 - m_{had}^2}}$. While the approximate “light-cone” variable $x^+ = \frac{(E+p)}{(E+P)_{\text{max}}}$ is described as being the most suitable for comparing fragmentation distributions at different energies, this is mitigated by radiative effects. This variable also suffers from systematic distortions at low $x^+$. The $x$ variable ranges from 0 to 1 for all experiments and is much more useful in fitting and visualizing the data.
Finally, for comparison of the D$^+$ meson decay rates, and to convert those measurements into cross sections we collect the most recent MARK III D meson branching ratios \[60\] in Table 5.1 and the measurements of resonant substructure of three body D meson decays \[61\] in Table 5.2.

| Decay Mode         | Fraction | Branching Fraction (%) |
|--------------------|----------|------------------------|
| $D^+ \to K^-\pi^+\pi^+$ |          |                        |
| $K^{0}\pi^+$       | $13 \pm 1 \pm 7$ | $1.8 \pm 0.2 \pm 1.0$ |
| non-resonant       | $79 \pm 7 \pm 15$ | $7.2 \pm 0.6 \pm 1.8$ |
| $D^0 \to K^{0}\pi^+\pi^-$ |       |                        |
| $K^{0}\rho^0$      | $12 \pm 1 \pm 7$ | $0.8 \pm 0.1 \pm 0.5$ |
| $K^{*-}\pi^+$      | $56 \pm 4 \pm 5$  | $5.3 \pm 0.4 \pm 1.0$  |
| non-resonant       | $33 \pm 5 \pm 10$ | $2.2 \pm 0.3 \pm 0.7$  |

### 5.2 $D^+ \to K^{0}\pi^+$

#### 5.2.1 Signal Isolation

The first approach used in isolating a D$^+$ signal is to use decay modes which contain a secondary vertex, in this circumstance a $K^*_S$. Since approximately 10% of the events contain a $K^*_S$ candidate, this dramatically reduces the combinatorial background. We observe the $K^*_S$ through its decay into two charged pions. To enhance the $K^*_S$ signal we require that the secondary vertex quality factor, $\chi^2_V$, be less than 2.0, and that the invariant mass of the $K^*_S$ formed from its two daughter tracks be within 30 MeV of the nominal $K^*_S$ mass. The effect of the $\chi^2_V$ cut was tested by calculating the efficiency corrected number of D$^+$ candidates for several different values of $\chi^2_V$. The numbers were found to be completely
consistent. To test the loose mass cut, we formed \( D^+ \) candidates where the \( K_S^0 \)
candidates were selected from a similar mass band centered at 400 MeV (Figure 5.1), no enhancement is evident. Since this is a two body decay, restrictive cuts

![Image](image.png)

**Figure 5.1:** \( D^+ \to K_S^0 \pi^+ \) candidates where the \( K_S^0 \) mass was selected from a side band region centered at 400 MeV.

are placed on the on the momentum of the track not associated with the \( K_S^0 \).
We require that that the single pion track have a momentum greater than 0.4 GeV. This track was also required not to be consistent with originating from a
secondary vertex.

The D\(^+\) candidate four momentum is calculated by adding the four momenta of the K\(_S^0\) and \(\pi^+\) candidates, where the K\(_S^0\) four momentum was calculated from the three-momentum of the K\(_S^0\) as determined from the secondary vertex finding algorithm, and defining the mass to be that of the K\(_S^0\). A plot of the invariant mass of D\(^+\) candidates with \(x \geq 0.51\) is shown in Figure 5.2.

We fit the mass spectrum to a Gaussian signal plus a third order polynomial background. The signal was found at a mass of 1871.4 ± 3.1 MeV/c\(^2\) and a FWHM of 48.3 ± 8.0 MeV/c\(^2\) which are consistent with a Monte Carlo simulation of this decay in the CLEO detector.

\[ \text{5.2.2 Background} \]

In addition to the combinatorial background, which is a smoothly varying function of D\(^+\) candidate momentum, certain physics processes cause anomalous enhancements to the combinatorial background. The final state K\(_S^0\)\(\pi^+\) is also a subset of the decays D\(^+\) → \(\overline{K}^0\rho^+\), D\(^0\) → K\(^-\)\(\pi^+\). Because of the polarization of the vector particle, combining the pseudoscalar daughter of the D with only one of the daughters of the vector particle will produce a satellite peak. Evidence for a satellite peak appears in the region of 1.4 to 1.7 GeV/c\(^2\) and was excluded from the fit.

A particularly difficult situation arises when the enhancement occurs in the signal region. This is a known problem for the D\(^+\), which is plagued by reflections from the D\(_s^+\). A D\(_s^+\) final state which is identical to a D\(^+\) final state with the
Figure 5.2: Mass spectrum for $D^+ \rightarrow K^0_L \pi^+$ candidates with $x \geq 0.51$. 
exception of one pion in the D⁺ state replaced with a kaon in the D⁺ s final state can produce an enhancement in the D⁺ region if the D⁺ kaon is misconstrued as a pion. Possibilities also exist for a similar confusion of p from a Λ𝑐 decay, however the reflection from this decay does not significantly overlap with the D⁺ region.

The author has developed a technique for quantifying D⁺ ⇒ D⁺ s reflections for two body decays. While examining the properties of the invariant mass D⁺ s → K⁺ K⁺ when the K⁺ was misidentified as a π⁺, it was noticed that at high momentum (p ≥ 2.5 GeV) the reflected mass had a strong dependence on the Rest Frame Decay Angle (RFDA - section 4.5), which is defined here as the cosine of the angle θ of the ‘pion’ in the ‘D⁺’ center-of-mass frame with respect to the ‘D⁺’ direction in the laboratory frame. A Monte Carlo prediction for this dependence is shown in Figure 5.3. In a certain sector of RFDA (RFDA ≤ −0.2) the D⁺ s reflection does not contaminate the D⁺ region. Thus it becomes possible to decompose the D⁺ s → K⁺ K⁺ reflection into parts which do (contaminated region) and do not (pure region) populate the D⁺ region. Figure 5.4 illustrates a Monte Carlo simulation of the distribution of reflected D⁺ s mass when segregated by RFDA. In addition, the broad peak observed in the contaminated region would change the width of the observed D⁺ peak in that region depending on the D⁺ : D⁺ s ratio. Table 5.3 contains a Monte Carlo study of the properties of the signal at the D⁺ mass for various ratios of D⁺ : D⁺ s production. The presence of a reflection from the decay D⁺ s → K⁺ K⁺ appears in the broadening of the D⁺ sig-

| D⁺ : D⁺ s | P = N_{contaminant} / N_{pure} | pure FWHM (MeV) | cont FWHM (MeV) |
|------------|-------------------------------|----------------|-----------------|
| Pure D⁺    | 1.59 ± 0.05                  | 40 ± 1         | 42 ± 1          |
| 4 : 1      | 2.01 ± 0.06                  | 39 ± 1         | 47 ± 1          |
| 2 : 1      | 2.47 ± 0.08                  | 39 ± 1         | 52 ± 1          |
Figure 5.3: Reflected mass of $D_s^+ \rightarrow K_S^0 K^+$ candidates with momentum $\geq 2.5$ GeV versus the RFDA (see text).

...nal width in the contaminated region and an increase in the ratio of events in the contaminated to pure region. To perform this measurement in the data we include the 33 pb$^{-1}$ of $\Upsilon(3S)$ data. Figure 5.5 shows the $K_S^0 \pi^+$ signal decomposed into pure and contaminated regions. We fit the distributions to the same form as the other $K_S^0 \pi^+$ signals, however we make no constraints on the mean and width of the signal since these properties could be altered by $D_s^+$ reflections. The effi-
Figure 5.4: Monte Carlo Simulation of the decay $D_s^+ \rightarrow K^0 S K^+$ where the invariant mass has been calculated calling the $K^+$ a $\pi^+$. The invariant mass is plotted for two regions of RFDA; a) RFDA < −0.2 (pure region) and b) RFDA ≥ −0.2 (contaminated region). Only the events in b) significantly overlap the $D^+$ region.
Figure 5.5: The $D^+ \rightarrow K^0\pi^*$ split into two regions on the basis of RFDA. a) Pure region (RFDA < -0.2) b) contaminated region (RFDA ≥ -0.2).
ciency corrected ratio of events found in the contaminated region to those found in the pure region is $1.6 \pm 0.4$. The large error, which is due to the combinatorial background, unfortunately prevents a quantitative statement being made about the amount of $D_s^+$ contamination. The width of the invariant mass peak is found to be $43.7 \pm 9.4$ ($47.4 \pm 8.0$) in the pure (contaminated) regions, and are also consistent. While we are prevented from making a strong statistical statement, it would seem unlikely that $D_s^+ \rightarrow K_0^- K^+$ contaminates $D^+ \rightarrow K_0^0 \pi^+$ at more than a $\approx 20\%$ level.

From a theoretical perspective $D_s^+ \rightarrow \bar{K}^0 K^+$, which proceeds through non spectator diagrams, is not anticipated to be large. Kamal [62] predicts $B(D_s^+ \rightarrow \bar{K}^0 K^+)/B(D^+ \rightarrow \bar{K}^0 \pi^+) = 0.04$ to 0.08. Approximating that there are nearly 3 times as many $D^+$’s as $D_s^+$’s, it is expected that the ratio $\sigma(e^+e^- \rightarrow D_s^+)B(D_s^+ \rightarrow \bar{K}^0 K^+)/\sigma(e^+e^- \rightarrow D^+)B(D^+ \rightarrow \bar{K}^0 \pi^+)$ would be only a few percent. CLEO has searched [63] for the decay $D_s^+ \rightarrow K_0^- K^-$ and determined the upper limit the upper limit $B(D_s^+ \rightarrow \bar{K}^0 K^+)/B(D_s^+ \rightarrow \phi \pi^+) < 0.55$. Only one group has seen a signal for this decay mode. MARK III [64] finds the ratio of the above two $D_s^+$ branching ratios to be $0.88 \pm 0.50$. This number, however, has continued to enjoy preliminary status. The amount of contamination from $D_s^+$ reflection for this decay mode can be described by

$$C = \left[ \sigma_{D_s^+} \cdot B(D_s^+ \rightarrow \bar{K}^0 K^+) \right] K_R$$

where $K_R$ is a kinematical factor relating how much of the $D_s^+$ signal actually reflects into the $D^+$ region and is approximately 0.5 for this case. Based on CLEO’s measurement [27] of $\sigma_{D_s^+} \cdot B(D_s^+ \rightarrow \phi \pi^+) = 5.8 \pm 1.0$ ($x \geq 0.5$) we estimate a systematic error of 1.3 pb$^{-1}$ for potential contamination of the $D^+ \rightarrow K_0^0 \pi^+$ cross section from $D_s^+ \rightarrow \bar{K}^0 K^+$ reflections.
5.2.3 Fragmentation Distribution

Table 5.4: Fragmentation Distribution for $D^+ \to \bar{K}^0\pi^+$

| x range     | $N_{\text{obs}}$ | efficiency | $d\sigma/dx \cdot B$ (pb) |
|-------------|------------------|------------|--------------------------|
| 0.000 - 0.375 | 43 ± 32          | 0.367 ± 0.020 | 25 ± 18                  |
| 0.375 - 0.510 | 17 ± 12          | 0.347 ± 0.019 | 30 ± 21                  |
| 0.510 - 0.625 | 73 ± 17          | 0.355 ± 0.019 | 46 ± 11                  |
| 0.625 - 0.750 | 73 ± 14          | 0.368 ± 0.019 | 41 ± 9                   |
| 0.750 - 0.825 | 48 ± 10          | 0.360 ± 0.019 | 27 ± 6                   |
| 0.825 - 1.000 | 14 ± 5           | 0.367 ± 0.024 | 8 ± 2                    |

The results of the fits for the differential cross sections for the $\bar{K}^0\pi^+$ mode are collected in Table 5.4. When quoting fragmentation distributions for modes which contain a $K^0_S$, we adopt the convention that $N_{\text{obs}}$ is number observed in the $K^0_S$ mode, where the error is statistical only. The quoted efficiency ($\epsilon_r$) is the reconstruction efficiency which accounts for geometrical acceptance, tracking efficiency, and cuts for this decay mode where the $\bar{K}^0$ has decayed via the chain $\bar{K}^0 \to K^0_S, K^0_S \to \pi^+\pi^-$. The error on the efficiency includes both statistical and systematic effects combined in quadrature. The value $d\sigma/dx \cdot B$ is defined as $(2.91)N_{\text{obs}}/\epsilon_r \Delta x L$. Here $L$ is the integrated luminosity, $\Delta x$ is the width of the bin, and the factor of 2.91 accounts for $B(\bar{K}^0 \to K^0_S) \cdot B(K^0_S \to \pi^+\pi^-)$. Differential cross sections will always be referenced to the original state of a $\bar{K}^0$. Because of the greatly reduced luminosity and immense combinatorial background, the measurements below $x = 0.51$ have disproportionately large errors. The most statistically significant information about the production cross section comes from a summation of the points above $x = 0.51$. Performing this summation yields $\sigma_{D^+} \cdot B(D^+ \to \bar{K}^0\pi^+) = 14.8 \pm 1.7 \pm 0.6 \pm 1.3$ pb ($x \geq 0.51$), where the errors are statistical, systematic (fitting procedure and Monte Carlo simulation), and an estimate for $D^+_s$ reflection contamination.
5.3 $D^+ \rightarrow \overline{K}^0\pi^+\pi^+\pi^-$

The approach for isolating this decay is quite similar to $K^0_S\pi^+$. We again rely on the $\chi^2_V$ cut, though it is relaxed to 2.5, since in this case the $K^0_S$ is slower and more difficult to reconstruct. All tracks not associated with the $K^0_S$ were required to have a momentum greater than 200 MeV. The mass spectrum for $D^+$ candidates with $x$ greater than 0.51, displayed in Figure 5.6, was fit to a Gaussian signal plus a fourth order polynomial background. A signal from the decay $D^{++} \rightarrow D^0\pi^+$, $D^0 \rightarrow K^0_S\pi^+\pi^-$ appears in the mass spectrum in the high $x$ bins and is excluded from the fit. The signal was found at a mass of $1876.0 \pm 2.6$ MeV/$c^2$ and a FWHM of $26.3 \pm 6.3$ MeV/$c^2$. The FWHM is consistent with Monte Carlo simulation; however the mass is 2.5 standard deviations away from the expected mass. This is caused by an upward fluctuation in the mass spectrum in the region $0.51 < x < 0.625$. Fitting the mass spectrum for $x$ greater than 0.625 yields a $D^+$ mass of $1869.9 \pm 2.7$ MeV/$c^2$, as expected. Because this is a four body decay, $D^+_s$ reflections are extremely broad and do not significantly overlap with the $D^+$ region.

5.3.1 Fragmentation Distribution

| x range     | $N_{\text{observed}}$ | efficiency  | $d\sigma/dx \cdot B$ (pb) |
|-------------|------------------------|-------------|----------------------------|
| 0.510 - 0.625 | 99 ± 33                | 0.183 ± 0.010 | 120 ± 41                   |
| 0.625 - 0.750 | 98 ± 28                | 0.206 ± 0.011 | 98 ± 29                    |
| 0.750 - 0.825 | 44 ± 15                | 0.236 ± 0.013 | 38 ± 13                    |
| 0.825 - 1.000 | 28 ± 8                 | 0.275 ± 0.014 | 21 ± 6                     |

The fragmentation distribution for the $K^0_S\pi^+\pi^+\pi^-$ mode is collected in Ta-
Figure 5.6: Mass spectrum for $D^+ \rightarrow K_S^0 \pi^+ \pi^+ \pi^-$ candidates having $x \geq 0.51$.  

Events / (10 MeV/c^2)
ble 5.5, where the definitions are the same as the previous section. \[ \sigma_{D^+} \cdot B(D^+ \rightarrow K^0\pi^+\pi^+\pi^-) = 33.4 \pm 6.0 \pm 1.4 \text{ pb, (} x \geq 0.51) \]. The errors are statistical and systematic, respectively.

### 5.4 \( D^+ \rightarrow K^-\pi^+\pi^+ \)

The isolation of \( K^-\pi^+\pi^+ \) requires a completely different approach. Here we rely primarily on the hadron identification system to produce a discernible signal. An additional advantage of this decay is that the kaon is produced with the opposite sign of the two pions. This reduces the combinatorial background along with providing signal validation and reflection rejection through the use of right and wrong sign combinations.

To enhance the signal we require that the kaon candidate have a weight \( W_K \geq 0.1 \). In addition we require that the kaon candidate posses at least 200 MeV of momentum. The motivation for this cut is simply that at low momentum it is difficult to produce a pure sample of kaons to measure the efficiency. Since a low momentum a particle’s energy loss is maximal, saturation effects in the electronics may produce unusual responses. Since these effects may not be properly reflected in the measured efficiency, it is best to avoid this particular group. We also make the usual cuts on the momenta of the pions (200 MeV). To further reduce the combinatorial background we require that all tracks have their DOCA’s less than 8 mm. The invariant mass spectrum for \( D^+ \rightarrow K^-\pi^+\pi^+ \) candidates with \( x \geq 0.51 \) is presented in Figure 5.7. The mean of the \( D^+ \) peak is measured to be \( 1.8693 \pm .0013 \text{ GeV} \), which is in excellent agreement with the world average of \( 1.8693 \pm 0.0006 \text{ GeV} \).
5.4.1 Background

This decay mode can be contaminated by the $D^+_s$ final state $K^+K^−\pi^+$. If either (but only one) of the kaons are misidentified a reflection may occur. The misidentification $K^−\pi^+\pi^+$ could count as a valid $D^+$ candidate but $K^+\pi^−\pi^+$ would not. We are therefore able look for the effects of $D^+_s$ reflections by examining wrong sign combinations. Figure 5.8 contains a plot of wrong sign combinations. No enhancements to the background are observed. We also note that because this is a three body decay the reflected peak is much broader and peaks well below the
D\(^+\) region. For this case \(K_R \approx 0.1\) as opposed to the \(K^0_S\pi^+\) mode where \(K_R \sim 0.5\). Enhancement to the background also occurs from the decay chain \(D^* \rightarrow D^0\pi^+\), \(D^0 \rightarrow K^-\pi^+\), which shares the same final state. This is simply removed by excluding the \(D^*\) region from the fit.

### 5.4.2 Fragmentation Distribution

The Monte Carlo simulation of this decay consisted of the nonresonant \(K^-\pi^+\pi^+\) and \(K^{*0}\pi^+ \rightarrow K^-\pi^+\pi^+\) in a ratio of 6.1 : 1 in accord with Table 5.2. Since the kaon efficiencies are not determined from Monte Carlo simulation, we include an
Table 5.6: Fragmentation Distribution for D$^+ \rightarrow K^-\pi^+\pi^+$

| x range     | N$_{\text{obs}}$ | efficiency   | $d\sigma/dx \cdot B$ (pb) |
|-------------|------------------|--------------|---------------------------|
| 0.000 - 0.375 | -201 ± 149       | 0.405 ± 0.046 | -36 ± 30                  |
| 0.375 - 0.510 | 221 ± 67         | 0.421 ± 0.047 | 107 ± 39                  |
| 0.510 - 0.625 | 702 ± 96         | 0.416 ± 0.047 | 129 ± 23                  |
| 0.625 - 0.750 | 609 ± 68         | 0.461 ± 0.052 | 93 ± 15                   |
| 0.750 - 0.825 | 384 ± 49         | 0.475 ± 0.052 | 57 ± 10                   |
| 0.825 - 1.000 | 119 ± 24         | 0.493 ± 0.059 | 17 ± 4                    |

additional term which is combined in quadrature with the error on the Monte Carlo efficiency and the error from the fitting procedure to estimate the overall systematic error. The summary of the differential cross section for the $K^-\pi^+\pi^+$ mode is contained in Table 5.6. The data are fit to a Gaussian signal and a third order polynomial. Performing a summation of the points with $x \geq 0.51$ We find a differential cross section $\sigma_{D^+} \cdot B(D^+ \rightarrow K^-\pi^+\pi^+) = 35.6 \pm 2.7 \pm 2.4$ pb, ($x \geq 0.51$).

5.5 $D^0 \rightarrow \bar{K}^0\pi^+\pi^-$

Table 5.7: Fragmentation Distribution for $D^0 \rightarrow \bar{K}^0\pi^+\pi^-$

| x range     | N$_{\text{obs}}$ | efficiency   | $d\sigma/dx \cdot B$ (pb) |
|-------------|------------------|--------------|---------------------------|
| 0.000 - 0.375 | -50 ± 44         | 0.175 ± 0.010 | -61 ± 62                  |
| 0.375 - 0.510 | 53 ± 21          | 0.221 ± 0.012 | 143 ± 63                  |
| 0.510 - 0.625 | 239 ± 32         | 0.249 ± 0.014 | 213 ± 33                  |
| 0.625 - 0.750 | 212 ± 25         | 0.275 ± 0.016 | 159 ± 21                  |
| 0.750 - 0.825 | 134 ± 17         | 0.284 ± 0.017 | 96 ± 14                   |
| 0.825 - 1.000 | 39 ± 8           | 0.286 ± 0.025 | 28 ± 7.0                  |

For comparison with the $D^+$ measurements and for use in the following chapter we document the properties of the $D^0 \rightarrow \bar{K}^0\pi^+\pi^-$. To detect this signal we make identical cuts on the $K^0_S$ and associated pions that we use in reconstructing $D^+ \rightarrow \bar{K}^0\pi^+\pi^+\pi^-$. Our simulation of this decay included $K^*-\pi^+$, $\bar{K}^0\rho^0$, and non resonant $\bar{K}^0\pi^+\pi^-$ in the ratio as prescribed by Table 5.2. This
decay mode is free from reflections of any kind. Figure 5.9 displays the in-

variant mass for $K^0_S\pi^+\pi^-$ candidates. The curve is a second order polynomial with a Gaussian signal. Table 5.7 collects the measurements of the differential cross section for this mode. Performing the canonical summation, we find

$$\sigma_{D^0} \cdot B(D^0 \to K^0_S\pi^+\pi^-) = 59.7 \pm 4.4 \pm 2.3 \text{ pb}, (x \geq 0.51).$$

Figure 5.9: Mass spectrum for $D^0 \to K^0_S\pi^+\pi^-$ candidates having $x \geq 0.51$. 

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Table 5.8: $D^+$ Cross Section Measurements

| Mode                                      | Experiment | $\sigma \cdot B$          |
|------------------------------------------|------------|---------------------------|
| $D^+ \to K^0\pi^+$                       | LGW        | 0.14 ± 0.05 (nb)          |
|                                           | MARK II    | 0.14 ± 0.03 (nb)          |
|                                           | MARK III   | 0.135 ± 0.012 ± 0.010 (nb)| 14.8 ± 1.7 ± 0.6 ± 1.3 (pb) |
|                                           | CLEO       |                           |
| $D^+ \to K^0\pi^+\pi^+\pi^-$            | MARK II    | 0.51 ± 0.18 (nb)          |
|                                           | MARK III   | 0.305 ± 0.031 ± 0.030 (nb)| 33.4 ± 6.0 ± 1.4 (pb)      |
|                                           | CLEO       |                           |
| $D^+ \to K^-\pi^+\pi^+$ (inclusive)     | LGW        | 0.36 ± 0.06 (nb)          |
|                                           | MARK II    | 0.38 ± 0.05 (nb)          |
|                                           | MARK III   | 0.388 ± 0.013 ± 0.029 (nb)| 35.6 ± 2.7 ± 2.4 (pb)      |
|                                           | CLEO       |                           |

Table 5.9: $B(D^+ \to K^0\pi^+\pi^+\pi^-) / B(D^+ \to K^0\pi^+)$

| Experiment | Ratio         |
|------------|---------------|
| MARK II    | 3.6 ± 1.5     |
| MARK III   | 2.3 ± 0.3     |
| MARK III (2 * T) | 2.1 ± 0.6 |
| CLEO       | 2.3 ± 0.5     |

Table 5.10: $B(D^+ \to K^-\pi^+\pi^+) / B(D^+ \to K^0\pi^+)$

| Experiment | Ratio         |
|------------|---------------|
| LGW        | 2.6 ± 1.0     |
| MARK II    | 2.7 ± 0.7     |
| MARK III   | 2.9 ± 0.4     |
| MARK III (2 * T) | 2.8 ± 0.6  |
| CLEO       | 2.4 ± 0.4     |

5.6 $D^+$ Relative Production Rates

The production measurements of $D^+$ mesons as observed in $e^+e^-$ annihilations are arranged in Table 5.8. The CLEO cross sections are only partial ($x \geq 0.51$), they represent the most statistically significant information available, and are perfectly acceptable for calculating relative rates. We have evaluated the relative rates from the various groups for $\left( \frac{B(D^+ \to K^0\pi^+\pi^+\pi^-)}{B(D^+ \to K^0\pi^+)} \right)$ (Table 5.9) and $\left( \frac{B(D^+ \to K^-\pi^+\pi^+)}{B(D^+ \to K^0\pi^+)} \right)$.
For the MARK III group, we also compare the branching ratios derived from the global fit double tag method \((2 \ast T)\). This analysis determines

\[
\left( \frac{B(D^+ \to \overline{K}^0\pi^+\pi^-)}{B(D^+ \to \overline{K}^0\pi^+)} \right) = 2.3 \pm 0.5
\]

and

\[
\left( \frac{B(D^+ \to K^-\pi^+\pi^+)}{B(D^+ \to \overline{K}^0\pi^+)} \right) = 2.4 \pm 0.4
\]

We find these results to agree well with previous measurements.

### 5.7 Analysis of Fragmentation Distributions

| Mode | \( \epsilon_0 \) | \( \chi^2/\text{d.o.f.} \) |
|------|------------------|------------------|
| \( D^+ \to K^-\pi^+\pi^+ \) | 0.21\( ^{+6}_{-5} \) | 10.7/4 |
| \( D^+ \to K^0\pi^+ \) | 0.17\( ^{+7}_{-5} \) | 3.7/4 |
| \( D^0 \to K^0\pi^+\pi^- \) | 0.22\( ^{+6}_{-5} \) | 9.0/4 |

The precise shape of the fragmentation distribution reveals distinctive features of the hadronization process. Charmed hadrons make an excellent laboratory to study such phenomenon, since they can be abundantly isolated in a variety of forms, and at different center of mass energies. Experimentally, measurements of fragmentation distribution distributions are difficult to execute. This is largely because fragmentation models distinguish themselves most at low momentum, which for reasons of acceptance, feed down, and or lack of
data tends not to be measured well. None the less qualitative still exist trends exist.

We choose to compare two fragmentation functions to our fragmentation data, one from each prevalent philosophy. We elect not to use models which explicitly include meson structure functions in their derivation (Kartvelishvili and Collins). The Peterson function

\[ \Delta^H_Q(z) = N \left( z \left[ 1 - \frac{1}{z} - \frac{\epsilon_Q}{(1 - z)} \right] \right)^2 \]

of independent fragmentation is used as is the Andersson form

\[ \Delta^H_Q(z^+) = \frac{N}{z^+} (1 - z^+)^a \exp\left( -b \frac{m_{H^+}^2}{z^+} \right) \]

representing string fragmentation. These functions happen to also be the most common in the literature. We note that the Andersson function is derived in terms of the light cone variable, analysis in terms of the x variable may not be optimal, but is suitable here for our mainly didactic purposes. The values of A and B are only for comparison with similar fragmentation distributions binned in x, and should not be compared with theoretical expectations. When fitting to the Andersson form, we fix the parameter \( m_{H^+} \) to be the known hadron mass. The results for the fits are exhibited in Table 5.11 and Table 5.12. Fits of both forms to the \( K^- \pi^+ \pi^+ \) and \( \bar{K}^0 \pi^+ \pi^- \) modes are shown in Figures 5.10 and 5.11. We have elected not to analyze the fragmentation distribution of the \( \bar{K}^0 \pi^+ \pi^+ \pi^- \) mode because of the much weaker statistical significance. For comparison Figure 5.12 displays the fragmentation distributions for the two \( D^+ \rightarrow \bar{K}^0 X \) modes, scaled by their relative production rates. Collectively we note that the shapes of the fragmentation distributions are consistent, as endorsed by the similarity of the parameters from the various fits.
Figure 5.10: Fits to the Andersson form (solid line) and Peterson form (dashed line) for $D^+ \rightarrow K^-\pi^+\pi^+$.  

In interpreting the information from the experimental observed fragmentation distributions, the historical trend has been to compare the values of the Peterson $\epsilon_Q$. This is an unfortunate tradition which shall be continued here. The main advantage of this approach is that $\epsilon_Q$ has an easily interpretable physical meaning, namely the parameter $\epsilon_Q$ is proportional to $m_{q\perp}^2/m_{Q\perp}^2$, where the transverse mass is defined as $m_{\perp}^2 \equiv m^2 + p_{\perp}^2$. The heavier the object is that combines with the charmed quark, the larger epsilon will be. Also, radiative effects will produce softer distributions which are again reflected in larger values of epsilon. CLEO has made a precise [27] measurement of the $D^{*+}$ fragmentation...
distribution and found $\epsilon_Q = 0.16 \pm 0.02$, and has also performed a fit to the fragmentation distribution of several $\Lambda_c$ modes [66] and determined $\epsilon_Q = 0.30 \pm 0.10$. The results of this analysis fit agreeably into this scheme. The $D$ mesons are softer than the $D^*$ mesons, as they should be since $D$’s are produced in the cascade of $D^*$’s, and the fragmentation distribution of charmed baryons is softer than charmed mesons. In conclusion, we note that mapping experimentally determined parameters of fragmentation functions to theoretical predictions for those parameters is an extremely perilous task. Data gathered at different ener-
gies by different experiments must be reconciled for choice of scaling variable
and more importantly QED and QCD radiative effects. Work by Bethke [67] has
attempted such a reconciliation, and determined \(< z_c > = 0.71 \pm 0.14\). This corre-
sponds to a Peterson \(\approx \epsilon = 0.04\), which is more in line with naive expectations.
Work has also been done by Galik [68] to evolve fragmentation distributions to
different energies (A detailed analysis of the CLEO D\(^+\) fragmentation distribu-
tion, with comparison to similar measurements at different energies, including
radiative effects, can be found elsewhere [27]).

Figure 5.12: The Fragmentation distributions for D\(^+\) → K\(_S^0\)\(\pi^+\) (solid squares) and
D\(^+\) → K\(_S^0\)\(\pi^+\pi^+\pi^-\) (solid circles). The two modes have been scaled by their relative
production ratios.
5.7.1 Estimates of Total Cross Sections

Table 5.13: Extrapolated Total Cross Sections (nb)

| Decay Mode            | $\sigma_{Total}$  |
|-----------------------|-------------------|
| $D^+ \rightarrow K^-\pi^+\pi^+$ | $0.63 \pm 0.06 \pm 0.10 \pm 0.09$ |
| $D^+ \rightarrow \overline{K}^0\pi^+$ | $0.74 \pm 0.11 \pm 0.12 \pm 0.11$ |
| $D^+ \rightarrow K^0\pi^+\pi^-\pi^+$ | $0.81 \pm 0.15 \pm 0.20 \pm 0.12$ |
| $D^0 \rightarrow K^0\pi^-\pi^+$ | $1.49 \pm 0.13 \pm 0.26 \pm 0.22$ |

Because of the large errors in the fragmentation distribution at low $x$, a measurement of the total cross section by summing the differential cross section will lose much of its statistical significance. Alternatively, we can attempt such a measurement by extrapolating from the part which is well measured. Both the Peterson and Andersson distributions predict about 63% of the fragmentation distribution resides above $x = 0.51$. We estimate the systematic error for the extrapolation procedure by varying the observed value of $\epsilon_Q$ by $\pm 1\sigma$ and determining the extrapolation factor. This, coupled with uncertainty in theoretical models and the inability to measure the end points of the fragmentation distribution well leads to an error of $\approx 15\%$. We base our extrapolation on the value of epsilon determined from the $\overline{K}^0\pi^+\pi^+$ and $K^-\pi^+\pi^+$ modes, which are determined with the most precision. Normalizing out the branching ratios (Table 5.1), we find the results presented in Table 5.13. The first error is the statistical and systematic combined in quadrature, the second is uncertainty in the branching ratio, and the third term is due to the extrapolation error. CLEO has recently measured [27] the total cross section for $D^0 \rightarrow K^-\pi^+$ to be $\sigma_{D^0} = 1.24 \pm 0.21$ (nb) which is consistent with this measurement of $\sigma_{D^0} = 1.49 \pm 0.36$ (nb). We also compare a previous CLEO [69] measurement for $D^+ \rightarrow K^-\pi^+\pi^+$ of $\sigma_{D^+} = 0.52 \pm 0.11$ (nb) agrees with our weighted average of the three decay modes $\sigma_{D^+} = 0.68 \pm 0.13$. The systematic error on the previous measurement may have been underesti-
mated.
CHAPTER 6

PENULTIMATE

In this chapter we examine several unique aspects of the production and decay of charmed mesons. We shall detail a search for a rare $D^0$ decay which may provide evidence for the role of hadronic final state interactions in charm decay. We measure the probability that a charmed meson will emerge from the $e^+e^-$ hadronization process in a state of non-zero angular momentum, and analyze the experimentally important transition rate for a charged $D^*$ meson to decay into a $D^0$ and a $\pi^+$. 

6.1 $D^0 \rightarrow K^0 \bar{K}^0$

6.1.1 Motivation

Decays of charmed hadrons have provided valuable insight into the dynamics of the weak interaction. The nonleptonic sector is believed to be the source of the radically different lifetimes and semileptonic branching ratios of the charged and neutral D mesons. The mechanisms which produce these effects are not fully quantified. A possible component of the solution is to shorten the $D^0$ lifetime by including a class of nonspectator decays known as W exchange, which are accessible to $D^0$'s at the Cabibbo allowed level. The contribution of these processes was initially anticipated to be small based on helicity arguments. The experimental observations $[38, 70]$ of the decay $D^0 \rightarrow \phi \bar{K}^0$ with a large branching fraction ($\sim 1\%$) was initially considered evidence for W exchange. An alternate theory $[40]$ proposed final state re-scattering of hadrons, not W exchange, as
the origin of this mode. The quark diagrams leading to $D^0 \to \phi \bar{K}^0$ through $W$ exchange and final state interactions are can be found in Figure 6.1. The extent to which either of these theories successfully explains $D^0 \to \phi \bar{K}^0$ remains to be resolved.

The decay $D^0 \to K^0 \bar{K}^0$ is uniquely suited to study the effect of final state interactions. At the quark level this decay proceeds through two classes of non-spectator diagrams (Figure 6.2). The unique feature of this decay is that both quarks present in the initial state are absent in the final state, which in this case is composed of an $s\bar{s}$ and a $d\bar{d}$ pair. There are two paths to the final state for each diagram, each of which contains one Cabbibo suppressed $W$ vertex. In the limit of exact $SU(3)_f$ symmetry, a $s\bar{s}$ can be popped from the vacuum on equal footing with a $d\bar{d}$ pair. We could then factorize the term $V_{ud} V_{cd}^* + V_{us} V_{cs}^*$, which then multiplies the matrix element for each of the two diagrams. This expression can be recognized as the product of the first two terms of the first two rows of the K-M Matrix. Unitarity of the K-M matrix demands

$$V_{ud} V_{cd}^* + V_{us} V_{cs}^* + V_{ub} V_{cb}^* \equiv 0$$

In the limit of vanishing $b \to u$ coupling, $V_{ud} V_{cd}^* + V_{us} V_{cs}^* \approx 0$, and the amplitude for this decay vanishes.

A calculation by Pham [71] based on the re-scattering of the modes $D^0 \to K^+ K^-$, and $D^0 \to \pi^+ \pi^-$, predicts $B(D^0 \to K^0 \bar{K}^0) \approx \frac{1}{2} B(D^0 \to K^+ K^-) \approx 0.3\%$. A branching ratio of 1% or greater for this decay could not be produced by final state interactions and would represent a violation of the standard model. If a substantial branching ratio were found for this mode, this would confirm the role of hadronic final state interactions. Conversely, if this decay was strictly ruled out, the case for nonspectator processes would be strengthened. Since
Figure 6.1: $D^0 \to \phi \bar{K}^0$ through a) W exchange and b)-c) final state inter-actions.

Figure 6.2: Quark diagrams contributing to the decay $D^0 \to K^0 \bar{K}^0$. a) W exchange, b) sideways “penguin.”
the general theoretical framework for describing charm decays does not fully implement either of these two processes, a better experimental understanding of this decay mode would lend valuable direction to charm decay theorists.

6.1.2 Signal Isolation

Experimentally this mode can be cleanly observed in $K_S^0 K_S^0$, given good $K_S^0$ mass resolution and reconstruction efficiency. In the previous chapter we have demonstrated the robustness of the CLEO detector in reconstructing D meson decay modes which contain a $K_S^0$. For quick comparison, we note that using the $\Upsilon(4S)$ data set, we have reconstructed $\sim 600$ high momentum $D^0 \rightarrow K_S^0 \pi^+ \pi^-$ decays. $D^0 \rightarrow K_S^0 K_S^0$ has a predicted branching ratio about twenty times smaller, and tacking on another factor of 3 to get $B(K^0 \rightarrow K_S^0) \cdot B(K_S^0 \rightarrow \pi^+ \pi^-)$ optimistically reduces our sample expectations to 10. Since we cannot hope to make an absolute measurement of this decay rate with these statistics, we choose to study the properties of this decay through normalization to a well known decay mode. The decay mode which is most similar to $K_S^0 K_S^0$ is clearly $D^0 \rightarrow K_S^0 \pi^+ \pi^-$. They are both final states containing four charged pions, and differ only by one secondary vertex. Normalization to this decay mode provides maximal cancelation of systematic errors.

To detect such a small signal we need to reduce the background to a minimum while maintaining reconstruction efficiency. We restrict our sample to those events in which the candidate $D^0$ has a momentum greater than 2.5 GeV. Despite the loss of the low momentum D’s, this cut has consistently produced invariant mass peaks with the highest signal to noise in all observed exclusive
charm decays. Since we are performing a normalization and not a cross section measurement, we are free to take advantage of the $D^0$'s which are copiously produced in $B$ decay. Unfortunately, $D^0$'s produced at low momentum have $K_S^0$ daughters in the momentum range where the $K_S^0$ reconstruction efficiency begins to fall off, producing as much as a 20% loss in efficiency. Coupled with the rising background at low momentum, this makes searching for this decay mode with low momentum $D^0$'s unprofitable.

The data for this analysis includes 113.7 pb$^{-1}$ taken on the $\Upsilon(4S)$ and 33 pb$^{-1}$ gathered on the $\Upsilon(3S)$. The event selection procedure and track corrections follow those outlined in section 5.1. Candidate $D^0$'s are then formed from two $K_S^0$ candidates. To further improve the purity of the sample we also require that the number of track pairs consistent with secondary vertices in an event be less than five, and that the $\chi^2_{\nu} \leq 3.0$ per $K_S^0$ candidate. Based on the measured $K_S^0$ FWHM as a function of momentum, we demand that each $K_S^0$ be within 2.5 standard deviations of the expected $K_S^0$ mass.

### 6.1.3 Background

We have analyzed the background from the decay $D^0 \rightarrow \bar{K}^0\pi^+\pi^-$. We simulated this mode using the resonant substructure measurement outlined in Table 5.2. For each Monte Carlo event we tagged the four final state $D^0$ daughters for this decay with matches to drift chamber tracks. We then ran the event through our $K_S^0K_S^0$ driver program, and analyzed $K_S^0K_S^0$ pairs only if one of the $K_S^0$'s was the true $D^0$ daughter and the vee finder has accidentally found one of the two $K_S^0$ tracks to be one of the other $D^0$ daughter pions. For the $K^-\pi^+$ mode, a plot of
these events in shown in Figure 6.3.

An anomalous enhancement occurs at the $D^{0}$ region. No such enhancement was found in either the nonresonant $K^{0}\pi^{+}\pi^{-}$ or $K^{0}\rho^{0}$ modes. To test that this was not a statistical fluctuation, we analyzed the properties of this distribution using the fact that the $K^{-}$ is polarized in the helicity zero state (section 4.2). The $K^{-}$ daughters are polarized with a $\cos^{2}{\Theta}_{P'P''}$ distribution in the $K^{-}$ rest frame with respect to the $K^{-}$ direction in the $D^{0}$ rest frame. We calculate this angle from the Monte Carlo for the $\pi^{-}$ daughter, and histogram the quantity for each “fake” $K_{S}^{0}K_{S}^{0}$ candidate. We show this distribution for all events and for events in the $D^{0}$ region in Figure 6.4. The events in the $D^{0}$ are clustered where $\cos{\Theta}_{P'P''} \leq 0.7$. This would correspond to the case where the $K^{-}$ daughter $\pi^{-}$ ends up almost at rest in the $D^{0}$ rest frame, and can end up being boosted in the same direction as the $D^{0}$ $\pi^{+}$ daughter. We note that while this effect seems to be caused by the unique kinematics of this particular decay, the efficiency for such processes is exceedingly small. The efficiency for a $D^{0} \rightarrow K^{-}\pi^{+}$ to cause a $K_{S}^{0}K_{S}^{0}$ candidate (where both $K_{S}^{0}$ candidates have been classified as good) with a mass in the $D^{0}$ region is of order 0.0006, while for a true $D^{0} \rightarrow K_{S}^{0}K_{S}^{0}$ decay the Monte Carlo efficiency is 0.194. We note, however that for a $D^{0} \rightarrow K^{0}\bar{K}^{0}$ branching ratio of 0.1%, and a $D^{0} \rightarrow K^{-}\pi^{+}$ rate of 5.6%, after scaling down the branching ratios to reach the four pion final state, we find the branching ratio times efficiency to be $2 \times 10^{-3}$ for $K^{0}\bar{K}^{0}$ and $7 \times 10^{-4}$ for $K^{-}\pi^{+}$, differing by only a factor of two! To avoid any such confusion, we calculate the invariant mass of of each $K_{S}^{0}$ candidate with the other two tracks which form the second $K_{S}^{0}$, and require that the mass not be consistent with a $K^{-}$. This eliminates this anomalous enhancement while only slightly reducing the $K_{S}^{0}K_{S}^{0}$ reconstruction efficiency. We also note that this faking caused by $K^{-}\pi^{+}$ does not seem to result from gross track mis-
Figure 6.3: Monte Carlo simulation of the decay $D^0 \rightarrow K^*-\pi^+$, $K^- \rightarrow K_S^0\pi^-$ which was passed through the $K_S^0K_S^0$ analysis program. One $K_S^0$ candidate was the correct $D^0$ daughter, while the other contained at least one other $D^0$ daughter. a) no cuts, b) with $K_S^0$ mass and $\chi^2_V$ cuts.
Figure 6.4: $\cos \Theta_{P'P''}$ distributions A) without and B) with a $D^0$ mass cut.
measurement errors, as the reconstructed $D^0$ has a momentum on average of 99% of the Monte Carlo generated momentum.

6.1.4 Calculation of Upper Limit

The mass spectrum for $K_S^0K_S^0$ candidates that have passed all cuts is displayed in Figure 6.5. We use a Monte Carlo procedure to determine the expected prop-

Figure 6.5: Invariant mass spectrum for $D^0 \rightarrow K_S^0K_S^0$. The $D^0$ candidate is required to have a momentum in excess of 2.5 GeV.
erties of the $D^0 \to K_S^0 K_S^0$ signal, where the $K_S^0$ is allowed to decay into the $\pi^+\pi^-$ mode. Subject to the cuts described above, we find the mean $1.8645 \pm 0.0002$ GeV, FWHM $0.029 \pm 0.003$ GeV and an overall reconstruction efficiency of $\epsilon_{K_S^0 K_S^0} = 0.170 \pm 0.007$. We fit the mass spectrum (Fig 3.) using a Maximum Likelihood method to a polynomial background and a Gaussian signal representing the $D^0 \to K_S^0 K_S^0$ decay. We observe a signal $N_{obs K_S^0 K_S^0}$ of $-2.7^{+2.7}_{-1.9}$ events at the $D^0$ mass, where the errors are statistical only. The stability of the observed signal area and errors has been analyzed subject to variations in $D^0$ mass, FWHM, and selection of background function. The results from the various fits are consistent, and we conservatively estimate a systematic error of 1.0 event. Figure 6.6 displays the $D^0 \to \overline{K}^0\pi^+\pi^-$ signal which is used to normalize the $K_S^0 K_S^0$ signal. In reconstructing this decay, we subject the $K_S^0$ to the same selection cuts as applied to the $K_S^0 K_S^0$ mode, and we require that the two $D^0$ daughter pions have a momentum greater than 200 MeV. The signal is fit to a Gaussian signal consistent with Monte Carlo simulation of this decay and a polynomial background. For this decay mode we find $N_{obs K_S^0 \pi^+\pi^-} = 811 \pm 64$, and $\epsilon_{K_S^0 \pi^+\pi^-} = 0.260 \pm 0.011$

The ratio of the branching fractions is obtained from the following prescription,

$$\frac{B(D^0 \to K_S^0 \overline{K}^0)}{B(D^0 \to \overline{K}^0\pi^+\pi^-)} = 2.91 \left( \frac{N_{obs K_S^0 K_S^0}}{\epsilon_{K_S^0 K_S^0}} \right) \left( \frac{\epsilon_{K_S^0 \pi^+\pi^-}}{N_{obs K_S^0 \pi^+\pi^-}} \right)$$

which yields $-0.015 \pm 0.016$. The factor 2.91 accounts for $B(K^0 \to K_S^0) \cdot B(K_S^0 \to \pi^+\pi^-)$. This ratio is converted to an upper limit on $B(D^0 \to \overline{K}^0 K^0)$ utilizing the most recent [60] Mark III value $B(D^0 \to \overline{K}^0\pi^+\pi^-) = 6.4 \pm 1.1\%$. From this we find a 90% confidence level upper limit of $B(D^0 \to \overline{K}^0 K^0) < 0.12\%$. The consistency of the normalization procedure has been checked by normalizing to the decay mode $D^0 \to K^-\pi^+$. This result is subject to systematic uncertainties in both $K_S^0$ reconstruction and $K^+$ identification efficiency. The upper limit determined in
Figure 6.6: Invariant mass spectrum for $D^0 \to K^0_S \pi^+ \pi^-$ which is used to normalize the $K^0_S K^0_S$ signal.

In this fashion gives $B(D^0 \to K^0 \bar{K}^0) < .17\%$.

Two other measurements of this decay have been publicized. The Mark III \cite{60} collaboration has determined the upper limit $B(D^0 \to K^0 \bar{K}^0) < .46\%$. A signal for this decay has been claimed by the E-400 \cite{72} experiment. Using the $D^*-D^0$ mass difference trick, they have observed a signal of $8.9 \pm 2.7$ events. They elect to normalize to the decay mode $D^0 \to K^+ K^-$, which would seem to maxi-
mize the prospects for systematic errors. They determine $\frac{B(D^0 \rightarrow K^0\bar{K}^0)}{B(D^0 \rightarrow K^+K^-)} = 0.4 \pm 0.3$. Normalizing away the denominator using the Mark III branching fraction $B(D^0 \rightarrow K^+K^-) = .51 \pm .11$, E-400 finds $B(D^0 \rightarrow K^0\bar{K}^0) = .20 \pm .16$. This is consistent with both the upper limit determined from this analysis and 0. It should also be noted that the E-400 group has not described any attempt to study the background to their signal, or examined the validity of the signal by exploring the $K_S^0$ side bands. The also have not normalized their $D^0 \rightarrow K^+K^-$ to any other $D^0$ decay mode.

In conclusion, more information is required to understand the decay $D^0 \rightarrow K^0\bar{K}^0$. CLEO is currently performing an analysis of this decay using a substantially larger ($\times 2$) data set taken with a new 52 layer drift chamber. CLEO has observed 5 $D^0$ decay modes with high statistical significance, and fails to observe a signal in this mode. CLEO has also determined that in their tracking system that a non negligible background for this decay mode can be produced from the decay $D^0 \rightarrow K^+\pi^+$, $K^- \rightarrow K_S^0\pi^-$. The E-400 group has observed a statistically weak effect at the $D^0$ mass in the $K_S^0K_S^0$ final state. The have not demonstrated the robustness of their normalization procedure, nor unambiguously attributed the signal to $D^0 \rightarrow K^0\bar{K}^0$.

6.2 $B(D^* \rightarrow D^0\pi^+)$ and Spinless D Meson Production.

In addition to extensive study of pseudoscalar charmed mesons, CLEO has made precise measurements of the charged and neutral vector D mesons. CLEO’s measurements [27] the total cross sections for the $D^{*+}$ and $D^{*0}$ mesons are summarized in Table [6.1]. It was first noticed by the author that CLEO’s
Table 6.1: Extrapolated Total Cross Sections for Vector D Mesons

| Decay Chain                  | $\sigma_{Total}$ (nb) |
|------------------------------|-----------------------|
| $D^{*+} \rightarrow D^0\pi^+$| $0.77 \pm 0.14$       |
| $D^0 \rightarrow K^-\pi^+$   |                       |
| $D^0 \rightarrow K^-\pi^+\pi^-$|                       |
| $D^{*0} \rightarrow D^0\pi^0$| $0.74 \pm 0.18$       |
| $D^{*0} \rightarrow D^0\gamma$|                       |
| $D^0 \rightarrow K^-\pi^+$   |                       |

unique situation of possessing a complete set of charged and neutral vector and pseudoscalar cross sections would allow for novel investigation of the hadronization process. Let us define $\sigma_p$ as the cross section for direct production of pseudoscalar charmed meson and $\sigma_v$ as the cross section for production of a vector charmed meson. We assume belief in isospin, making $\sigma_p$ and $\sigma_v$ the same for charged and neutral particles. Let us also define $\alpha$ as the branching ratio $B(D^{*+} \rightarrow D^0\pi^+)$ and $\beta$ as the branching ratio $B(D^{*+} \rightarrow D^+X)$. Where we have the constraint $\alpha + \beta = 1$. The total cross sections of $D^0$ and $D^+$ mesons are governed by

$$\sigma(D^0) = \sigma_p + \sigma_v + \alpha \sigma_v$$

$$\sigma(D^+) = \sigma_p + \beta \sigma_v$$

after applying the constraint this becomes

$$\sigma(D^0) = \sigma_p + (1 + \alpha) \sigma_v$$

$$\sigma(D^+) = \sigma_p + (1 - \alpha) \sigma_v$$

by adding and subtracting these two equations, we decompose the equations into one which contains $\sigma_p$ and one which contains $\alpha$. Solving these two equations for the aforementioned quantities we find the relations

$$\left(\frac{\sigma(D^0) + \sigma(D^+)}{2}\right) - \sigma_v = \sigma_p$$
\[
\left( \frac{\sigma(D^0) - \sigma(D^+)}{2\sigma_V} \right) = \alpha
\]

We note that in the derivation of the above equations we have excluded contribution from the $D^{*0}(2420)$. This is a candidate for the spin 2 charmed meson, its spin and parity have not been established. Its relative production rate to the $D^*$ is on order [73] 12 ± 5%. This state has only been observed in cascades to the $D^*$. Provided this state does not have a substantial decay rate directly into the ground states, these relations are unaffected. Using the results derived in the previous Chapter $\sigma_{D^*} = 0.68 \pm 0.13$, $\sigma_{D^{*0}} = 1.49 \pm 0.36$ (nb), and defining $\sigma_V = \sigma_{D^{*0}} = 0.74 \pm 0.18$ we derive

\[
\sigma_p = 0.35 \pm 0.26\text{(nb)}
\]

\[
\alpha = B(D^{*+} \to D^0\pi^+) = 0.53 \pm 0.29
\]

To evaluate our measurement of $B(D^{*+} \to D^0\pi^+)$, we first make the simple comparison of the $D^*$ to $D^{*0}$ cross section where both particles cascade into a $D^0$, which then decays into the $K^-\pi^+$ final state. This yields $B(D^{*+} \to D^0\pi^+) = 0.55 \pm 0.07 \pm 0.11$. Other measurements of this number include those of MARK I [74] (0.60 ± 0.15), MARK II [75] (0.44 ± 0.7), and MARK III [76] (0.55 ± 0.02 ± 0.06). The MARK II number does not assume conservation of isospin. We again find favorable agreement among the measurements. Lastly, using our measurement of the direct $D$ cross section to the amount of charmed particles which are observed in the vector state. We evaluate the quantity $\frac{\sigma_V}{\sigma_p + \sigma_V}$. Using our derived value for $\sigma_p = 0.35 \pm 0.26$ (nb) and a weighted average of the $D^*$ cross sections for $\sigma_V$ we derive

\[
\frac{\sigma_V}{\sigma_p + \sigma_V} = 0.68 \pm 0.18
\]

Traditionally the measurement was limited by the knowledge of $B(D^{*+} \to D^0\pi^+)$, which for some time suffered a 20% uncertainty. This method bypasses the need
for that number. Theoretical predictions for this ratio are not well established, however we note that simple spin statistics predicts 0.75. In retrospect, we note that the errors on the numbers derived here make them somewhat less competitive with previous measurements. We feel the originality of the method warrants its presentation, which may find application in the future. We also note that the means of these measurements are in fine agreement with expectations, attributing to the consistency of several distinct detector functions.
 CHAPTER 7

SUMMARY

We have made a broad study of the properties of charmed mesons produced in $e^+e^-$ annihilations. The hadronization properties of charmed quarks were studied in two ways. The fragmentation distributions of $D^0$ and $D^+$ mesons were compared using string inspired and independent fragmentation models. The Peterson function provided consistently poorer fits to the data. This was in part dominated by the highest $x$ bin which the function consistently undershot. It is noted that the Peterson function goes to 0 at $x = 1$ with a much higher rate $(1-x)^2$ than the other models $\sim (1-x)$. Radiative effects may soften the spectrum most severely at high $x$, however these effects have not been examined (this is in part due to the fact that there is significant feed down from the vector states in pseudoscalar fragmentation distributions, making vector mesons the laboratory of choice to study these effects). The Peterson function also predicts a much larger distribution at low $x$ than do the string functions, which sharply cutoff at $\approx x = 0.3$. This cutoff is more intuitively appealing, as catastrophic processes would be required in the hadronization process to produce mesons with very low $x$. We find the fragmentation distributions of the $D^+$ and $D^0$ mesons to be quite similar. The $D$ fragmentation distribution is softer than the $D^+$ and harder than the $\Lambda_c$, as expected. It was found that about 70% of the charmed mesons are produced with at least 1 unit of angular momentum. This is further supported by the large difference in the inclusive $D^0$ and $D^+$ inclusive cross sections.

The relative branching ratios were measured for three different $D^+$ decay modes. They were found to be in good agreement with previous measurements
done at SPEAR. No specific theoretical information can be gleaned from these relative rates. This is partly due to the fact that in $D^+$ decays there are often interfering diagrams (Figure 2.8) so rates cannot be purely determined. This is not the case in $D^0$ decays, where relative rates for two body decays often provide useful information. The second difficulty is that in general theoretical predictions tend to be limited to two body and quasi two body decays, and no such predictions exist for three and four body nonresonant decays. We have also searched for nonspectator decays of the $D_s^+$ as a contamination of the decay mode $D^+ \to \bar{K}^0\pi^+$. A novel technique was developed for this purpose, which would detect signs of $D_s^+$ reflections in two ways. Due to limited statistics, we were unable to make a definitive statement, although no real sign for this reflection was observed.

We also searched for the special decay mode $D^0 \to K^0\bar{K}^0$, which has been touted as evidence for hadronic final state interactions in charm decays. The E-400 group has measured a candidate signal for this decay and found $B(D^0 \to K^0\bar{K}^0) = 0.20 \pm 0.16 \%$. This analysis is unable to confirm the E-400 measurement, and places the 90 \% confidence interval upper limit of $B(D^0 \to K^0\bar{K}^0) < 0.12 \%$. In addition, we have determined that there is a non negligible background to this mode from the decay $D^0 \to K^+\pi^-$.

The theoretical understanding of charm decay has advanced greatly in recent years. This was fueled, in part, by experiments performed at $e^+e^-$ energies above the $\psi''$, and fixed target experiments. The ARGUS group at DORIS was the first to observe the decay $D^0 \to \phi\bar{K}^0$, and the tensor meson candidate $D^{*0}(2240)$. CLEO was the first experiment to perform a high statistics study of the $D^0$, $D^+$, and $D_s^+$ lifetimes. The E-691 group has precisely determined the
lifetimes of these three mesons, along with measuring several rare decay modes. One difficulty of these experiments has been the reliance on states reconstructed purely from charged tracks, neglecting final states with single or multiple neutral particles (this partially contributed to the undoing of the MARK III double tag method). The CLEO II detector, with exceptional charged and neutral particle reconstruction capabilities should make an excellent tool for studying charmed decays and spectroscopy. Further advances in charm decay theory will also require a better understanding of the charmed, strange, meson $D_s^+$, which has thus far only been observed in five decay modes. In conclusion, much has been learned, and much is to be learned about the charm sector. The study of charmed particles has deepened our knowledge of elementary particle physics. A fully quantified theory of charm decays will represent a great triumph for theorists and experimentalists alike.
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