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Studies of P-Wave Charm Meson Production at LEP with the OPAL Detector

by

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A thesis submitted to the
Faculty of Graduate Studies and Research
in partial fulfillment of the requirements
for the degree of
Doctor of Philosophy

Ottawa-Carleton Institute for Physics
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The undersigned hereby recommend to
the Faculty of Graduate Studies and Research
acceptance of the thesis,

STUDIES OF P-WAVE CHARM MESON PRODUCTION
AT LEP WITH THE OPAL DETECTOR

submitted by
T.M. Jones, B.Sc., M.Sc.

in partial fulfilment of the requirements
for the degree of Doctor of Philosophy

Chair, Department of Physics

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Date: April 22, 1997

Abstract

This thesis describes a study of P-wave charm mesons, particles composed of a heavy charm quark and a light up, down or strange anti-quark, which have one unit of orbital angular momentum between them. These mesons are produced in decays of Z^0 bosons and the results presented here are derived from over 4 million Z^0 decays recorded using the OPAL detector at CERN's Large Electron Positron collider. For each type of light quark, the symmetries introduced by the large mass of the charm quark imply the existence of four such charm mesons. In the limit where the charm quark mass is taken to infinity, the meson properties are determined only by the total angular momentum of the light quark. The production of three of these states, the $D_1^0(2420)$, the $D_2^{*0}(2460)$ and the $D_{s1}^+(2536)$, has been observed in hadronic Z^0 decays using data recorded with the OPAL detector at LEP.

The D_1^0 and D_2^{*0} are reconstructed through their decays to $D^{*+}\pi^-$ and their production rates in charm quark fragmentation and in decays of B hadrons are measured separately by applying a simultaneous fit to two flavour enriched samples of decays. The charm sample is obtained using candidates with scaled energy greater than 0.5 and small D^0 decay lengths, while the sample enhanced in $b\bar{b}$ events is selected by reconstructing the characteristic decay topology of the B hadron. In this analysis, the decay angles are used to measure the relative fractions of the D_1^0 and D_2^{*0} states and to distinguish their signals from the combinatorial background. The production of D_1^0 mesons in semileptonic b-decays is also observed. D_{s1}^+ production is observed by reconstructing both the $D^{*0}K^+$ and the $D^{*+}K_S^0$ final states and its production rate in charm fragmentation is measured for scaled energies larger than 0.6. This measurement constitutes the first observation of D_{s1}^+ production at LEP.

These studies resulted in the following measurements of the probabilities for charm quarks to produce P-wave mesons in the process of quark fragmentation:

$$f(c \rightarrow D_1^0) = 0.021 \pm 0.007 \pm 0.003$$

$$f(c \rightarrow D_2^{*0}) = 0.052 \pm 0.022 \pm 0.013$$

$$f(c \rightarrow D_{s1}^+) = 0.016 \pm 0.004 \pm 0.003.$$

In addition, the probabilities that the D_1^0 and D_2^{*0} are produced in decays of b-hadrons were determined:

$$f(b \rightarrow D_1^0) = 0.050 \pm 0.014 \pm 0.006$$

$$f(b \rightarrow D_2^{*0}) = 0.047 \pm 0.024 \pm 0.012.$$

and D_1^0 production in semileptonic b-decays was studied.

The measurements of the D_1^0 , D_2^{*0} , and D_{s1}^+ production rates in charm events are compared with predictions from recent models describing quark fragmentation in Z^0 decays. Their production properties are also compared with OPAL measurements of the production rates of the corresponding P-wave B mesons. Finally a comparison of D_1^0 production in semileptonic and inclusive decays of b-hadrons is presented.

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I must first thank all the members, past and present, of the OPAL collaboration, since without their dedicated effort to the stable operation of the detector, the reconstruction of the data, and the support of computing services at CERN, none of the work presented here would have been possible. The OPAL author list, as of January 1997, is included as an appendix, but there are many people who have made great contributions to the experiment who have left over the years and they should not be overlooked. Likewise, the machine physicists, engineers and technicians responsible for the construction and operation of LEP over the past decade deserve credit since their effort has yielded one of the most significant experimental programs in High Energy Physics.

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Chapter 1

Introduction

In 1913, Bohr described a model for the hydrogen atom in which the electron moves under the influence of the electric field of the positively charged and much heavier nucleus. Because the binding energy of the system is much smaller than the electron mass, the system can be described quite adequately using non-relativistic quantum mechanics which provided accurate predictions of the energies of spectral lines of hydrogen and other elements. Remarkably, 60 years later, it was found that the same non-relativistic model could be used to qualitatively describe the structure of mesons containing heavy quarks.

In 1964, Zweig[1, 2] and Gell-Mann[3] independently postulated the existence of quarks, to explain the properties of the large numbers of hadrons which were being discovered at that time. In their models, baryons such as protons and neutrons, consisted of three quarks while mesons such as pions and kaons, consisted of quark anti-quark pairs. Although this described the experimental data well, these quarks were generally regarded as mathematical conveniences rather than physical particles.

This view changed in light of measurements of the inelastic electron-proton scattering cross section carried out at SLAC in 1967 and 1968[4]. These results indicated that at high energies, electrons did not scatter from the proton as a whole, and it was Feynman, in 1969, who interpreted these results as a consequence of

scattering from point-like, spin $1/2$ constituents within the proton[5]. These had precisely the properties of the quarks postulated by Zweig and Gell-Mann, but at the time there was no dynamical theory which could explain the properties of hadrons in terms of their quark constituents. This changed in November of 1974, with the discovery of the J/ψ particle, which was quickly interpreted as a meson consisting of a new heavy quark-antiquark pair which was named “charm”. Soon, a rich spectrum of excited $c\bar{c}$ states was observed and since the charm quark mass was so large, the details of this spectrum were explained very well using a non-relativistic model in which a charm quark and charm anti-quark orbit each other.

While these developments in the understanding of hadronic structure were taking place, significant theoretical progress was also being made in the construction of a unified theory of the electromagnetic and weak interactions of matter. This developed into what is now referred to as the Standard Model of fundamental particles and its key aspects were confirmed when the $W^\pm$ and Z^0 particles, predicted by the theory, were observed in high energy $p\bar{p}$ collisions by the UA1 and UA2 experiments at CERN[6, 7]. Precision tests of the Standard Model could only be made by studying high statistics samples of these particles and it was for this reason that the Large Electron Positron (LEP) collider was constructed at CERN. LEP began operation in 1989, producing Z^0 particles by colliding 46 GeV beams of electrons and positrons. Since the Z^0 mass is large, it can decay directly to charm and bottom quark-antiquark pairs and thus, the LEP experiments are also an ideal place to study the properties of mesons containing heavy quarks.

All electroweak interactions studied at LEP are in excellent agreement with Standard Model predictions. However, there are some aspects of Z^0 decays which cannot be fully described within the context of the Standard Model. One of them, which is the primary focus of this thesis, is known as quark fragmentation, also called hadronization, which is the process whereby the initial quark-antiquark pairs, produced in Z^0 decays, become bound inside stable particles. To date, there

is no theory which can fully explain quark fragmentation and performing measurements to constrain the models which do exist is difficult. Some of the most useful studies are measurements of the production rates of particles containing the charm or bottom quarks originally produced in the Z^0 decay. Of these particles, the ones which can potentially provide the most information are highly excited states.

Just as hydrogen atoms can be excited from their ground state to states in which the electron and the nucleus have non-zero orbital angular momentum, so too can heavy mesons be produced with orbital angular momentum. This thesis focuses on studies of this type of particle, and in particular, presents measurements of production rates of orbitally excited “P-wave” charm mesons. These states have one unit of orbital angular momentum between a heavy charm quark and a light up, down or strange anti-quark and are observed via their decays to lighter charm mesons. They provide what is potentially an ideal way of studying quark fragmentation since, in addition to being excited particles carrying one of the original quarks from the Z^0 decay, their properties can be understood using simple symmetry arguments.

The following two chapters provide the background material necessary for the discussion of the studies carried out in this thesis. Chapter 2 describes the key concepts upon which the Standard Model of particle physics is based and how these ideas shape our view of the way particles interact at high energies. It ends with a discussion of the decays of hadrons containing bottom quarks which are expected to provide important tests of the Standard Model in the coming years. Although these are interesting in themselves, this discussion is relevant because bottom hadrons are also an important source of P-wave charm mesons in Z^0 decays.

Chapter 3 describes in more detail many aspects directly related to studying P-wave charm mesons at LEP. These include descriptions of Z^0 decays to quarks and a summary of several models which have been used to describe the way these quarks fragment to form hadrons. This is followed by a description of non-relativistic potential models which are used to describe the properties of mesons containing

a heavy quark. This illustrates the simple way in which the properties of P-wave charm mesons arise due to the large mass of the charm quark.

The chapters which follow the presentation of this theoretical framework focus on studies of the production of $D_1^0(2420)$ and $D_2^{*0}(2460)$ P-wave charm mesons in Z^0 decays. The final analysis chapter describes the first observation at LEP of the $D_{s1}^+(2536)$ meson, a P-wave charm meson containing a strange anti-quark. These measurements are then compared with predictions of a recent model in the final chapters.

Chapter 2

The Standard Model

The common properties of the fundamental particles can be used to classify them into groups. One such classification depends on the spin of the particle, and it is found that all the constituent particles of matter are fermions with spin $1/2$, while those which mediate the interactions between them are bosons with spin 1 . In the Standard Model, another important classification is made by considering the forces or interactions experienced by different types of particles. We know of only four types of forces, the weakest of which is gravity and the others being the electromagnetic force and the weak and strong forces. Forces between the fundamental fermions are mediated by the exchange of particles known as gauge bosons and the Standard Model attempts to describe all properties of matter in terms of these interactions. Gravity is so weak that it is irrelevant for the processes studied in high energy physics. Furthermore, a satisfactory quantum theory of gravity does not yet exist and thus, the Standard Model does not attempt to describe gravitational effects. The remaining three forces, however, are described extremely well.

An important classification can be made based on whether a particle experiences the strong interaction. Fermions which experience the strong force are called quarks and interact by the exchange of gluons. Bound systems of quarks such as protons, neutrons and pions, also undergo strong interactions and are collectively

referred to as hadrons. The remaining particles, such as the electron and neutrinos, do not interact strongly and are classified as leptons.

The lightest types, or flavours, of quarks are called the *up* and *down* quarks and these behave identically under the influence of the strong interaction. However, they have different electric charges of $+2/3$ and $-1/3$ relative to the electron charge of -1 and, being charged, all of these particles can experience electromagnetic interactions mediated by photons. Neutrinos carry no electric charge and only interact via the weak interaction, which is mediated by the exchange of $W^\pm$ and Z^0 bosons, which also couple to quarks and charged leptons.

The up and down quarks, together with the electron and the electron neutrino are grouped together to form one family of particles. For some unknown reason, nature has provided two additional copies of this family which apparently differ only by the mass scales of their members. Thus, there are twelve elementary fermions and all of their interactions are described by the Standard Model in terms of the exchange of vector bosons. Some important properties of these fundamental particles are listed in table 2.1.

Mathematically, the Standard Model is based on a description of relativistic quantum mechanics known as gauge field theory, in which the fields are mathematical objects used to describe the particles. Calculations in field theories of interacting particles cannot generally be performed exactly and perturbative expressions for observables are usually obtained instead. Although the mathematical expressions for the terms in the resulting perturbation series can be rather complicated, their structure is such that they can be uniquely represented by graphs referred to as Feynman diagrams. These are constructed using well defined “Feynman rules” which depend on the nature of the interactions involved.

An important characteristic of field theories, which severely limits the ones which can be used to describe the real world, is their ability to provide reliable predictions beyond first order in perturbation theory. The higher order Feynman diagrams in a perturbation series often contain infinities which, when considered

name	symbol	charge (e)	mass (MeV/c^2)
electron	e^-	-1	0.511
muon	μ^-	-1	106
tau	τ^-	-1	1777
electron neutrino	ν_e	0	$< 5.1 \times 10^{-6}$
muon neutrino	ν_μ	0	< 0.27
tau neutrino	ν_τ	0	< 24
down	d	-1/3	~ 10
up	u	+2/3	~ 5
strange	s	-1/3	~ 200
charm	c	+2/3	~ 1300
bottom	b	-1/3	~ 4300
top	t	+2/3	$175.0 \pm 3.9 \pm 4.5 \text{ GeV}/c^2$

Table 2.1: Quarks and leptons, their charges and masses. The corresponding anti-particles have opposite charges but equal masses.

alone, yield meaningless expressions. However, when all diagrams of a given order are considered, many of these divergencies cancel. The ones which remain can often be absorbed into the definition of a physical parameter of the theory such as a mass or coupling constant. Such theories are said to be renormalizable and those which are not would require an infinite number of arbitrary parameters to absorb all the divergencies which enter at each order in the perturbation series. This behaviour is deemed unphysical since nature appears to be characterized by a rather limited number of free parameters.

There is a certain degree of arbitrariness in the way the fermion fields in the Standard Model are defined. Since predictions of measured quantities must be unique, they cannot depend on arbitrary redefinitions of the fields involved and it is this property which essentially predicts the existence of force-mediating gauge bosons and fully specifies the ways in which they interact. These aspects of the Standard Model are described in the following sections.

2.1 Local Gauge Symmetries

Since all physical quantities depend only on the modulus-squared of a field's wavefunction, ψ , this field need only be specified up to an overall phase θ . In fact, this phase may be chosen differently at each point in space, requiring a description of the particle to be invariant under the unitary group of transformations

$$\psi(x) \rightarrow \psi'(x) = e^{i\theta(x)}\psi(x). \quad (2.1)$$

As in classical mechanics, the properties of the field can be described by its Lagrangian which is expressed in terms of the kinetic and potential energies involved. In quantum mechanics, the kinetic energy of a particle is represented by terms containing derivatives of the field, but these are not invariant under local phase transformations unless an additional vector field $A_\mu(x)$ is introduced which transforms as follows:

$$A_\mu(x) \rightarrow A'_\mu(x) = A_\mu(x) - \partial_\mu\theta(x). \quad (2.2)$$

This new field, $A_\mu(x)$, is introduced by the replacement of the ordinary partial derivatives of the field $\psi(x)$ by a covariant derivative:

$$\partial_\mu \rightarrow D_\mu \equiv \partial_\mu - iA_\mu. \quad (2.3)$$

In the case of electrodynamics, $\psi(x)$ describes an electron and the gauge field $A_\mu(x)$ is identified as the 4-potential from which physical electric and magnetic fields are obtained.

Other symmetries are observed in nature and the isospin invariance of the strong interaction led Yang and Mills to investigate the structure of Lagrangians describing nucleon fields invariant under isospin rotations [8]. In such a model, the nucleon field is a doublet consisting of the proton and neutron. It is transformed

under a local isospin rotation by

$$N \equiv \begin{pmatrix} p \\ n \end{pmatrix} \rightarrow N' = e^{i\tau \cdot \varepsilon(x)} N \quad (2.4)$$

where τ are the generators of the isospin group $SU(2)$ ¹.

In general, a representation of a multiplet of N fields, $\vec{\psi}(x)$ will transform under the group $SU(N)$ according to

$$\vec{\psi}(x) \rightarrow \vec{\psi}'(x) = e^{iT^a \varepsilon^a(x)} \vec{\psi}(x) \quad (2.5)$$

where $T^a, a = 1, \dots, N^2 - 1$ are $N \times N$ representations of the Hermitian generators of the group. For a Lagrangian to be invariant under such a transformation, one must introduce $N^2 - 1$ gauge fields $A_\mu^a(x)$ which transform like

$$A_\mu^a(x) \rightarrow A_\mu'^a(x) = A_\mu^a(x) - \partial_\mu \varepsilon^a(x) \quad (2.6)$$

and the corresponding covariant derivative

$$D_\mu \equiv \partial_\mu - iT^a A_\mu^a(x). \quad (2.7)$$

2.2 QCD

In the Standard Model, the strong interaction arises because of a three-fold gauge symmetry of quarks. A quark can be assigned one of three “colour” charges, and this description of quarks is referred to as Quantum Chromodynamics. The definition of which colour is which is arbitrary and hence, the physical properties of quarks must be invariant under a redefinition of these colours. This is insured mathematically by requiring that quark fields be represented by members of the

¹A common representation of the $SU(2)$ generators is related to the Pauli matrices by $\tau = \sigma/2$

$SU(3)$ symmetry group. Thus, as shown in the previous section, there are eight massless vector bosons associated with this symmetry which are called gluons.

A complete description of the strong interaction must also account for the kinematic properties of the gluons. Kinetic energy terms for gluons in the QCD Lagrangian contain the covariant derivative, shown in equation (2.7), and introduce terms containing products of three and four gluon fields. These terms generate a class of Feynman diagrams in which the gluons couple to themselves which is a consequence of the fact that gluons themselves carry colour charges. Because of this, quarks and gluons behave similarly in certain situations and when only the colour charge is relevant, quarks and gluons are often collectively referred to as “partons”, a name first used to describe the internal constituent parts of the proton. Although the group structure of QCD is similar to that of Quantum Electrodynamics, the difference between the electrically neutral photon and the self-coupling gluon fields produces completely different behaviours for these two theories.

A particularly important consequence results in different behaviours of the electromagnetic and strong coupling constants, α_{EM} and α_s , at different scales of momentum transfer in an interaction. The calculation of α_{EM} beyond first order in perturbation theory shows that it is not constant, but instead gets weaker at large distance scales or correspondingly, at small momentum transfers. However, the self-coupling of gluons produces the opposite behaviour whereby quarks experience only a weak force when they are close to each other which becomes stronger the farther they move apart. Conversely, quark interactions involving small momentum transfers experience large forces while the forces involved in high energy interactions are reduced.

This behaviour can be seen in the first order perturbative expression for α_s :

$$\alpha_s(\mu^2) = \frac{4\pi}{\beta_0 \log(\mu^2/\Lambda^2)} \quad (2.8)$$

in which $\beta_0 = 11 - \frac{2}{3}n_f$ depends on the number of quark flavours, n_f , accessible

at the scale, μ , of momentum transfer involved. Although α_s , when measured at a specific value of μ^2 , could be considered to be a fundamental constant, the lack of a unique scale at which to define it leads to the introduction of the scale independent parameter Λ . Since Λ is defined in terms of a perturbative expression for α_s , its value will depend on the order of the expansion and on the way in which the perturbative expression is renormalized. It will also take on different values depending on the number of quark flavours which may occur at particular energies. At the Z^0 energy scale of about 90 GeV, there are five active quark flavours and in the modified minimal subtraction ($\overline{\text{MS}}$) scheme, Λ is determined from experiment to be [9]

$$\Lambda_{\overline{\text{MS}}} = 220_{-60}^{+136} \text{ MeV}. \quad (2.9)$$

This corresponds to $\alpha_s(M_Z) = 0.122 \pm 0.007$. The fact that α_s is “small” at these energies is essential for reliable calculations of many properties of Z^0 decays using QCD.

Another consequence of the self-coupling nature of gluons is the fact that various configurations of quarks experience the strong force differently. It is found that attractive forces exist between quark-antiquark pairs and between groups of three quarks for cases where the combination has no net colour charge. This attractive force allows these two configurations to form bound states referred to as mesons and baryons, respectively. Other combinations of quarks do not form bound states since they experience repulsive forces. Although more complicated combinations of quarks, antiquarks and gluons might form bound states, such objects have not been positively identified in nature.

2.3 The Electroweak Model

It is found that the weak interaction affects the elementary fermions differently, depending on whether their spins are aligned with or opposed to their momentum vectors. In fact, the W^- boson will only couple to leptons in the latter case, that

is, to leptons in “left-handed” helicity states. Because the left-handed leptons and quarks in each family behave similarly under weak interactions it is natural to attempt to describe them as members of the $SU(2)$ symmetry group. However, the right-handed fermions do not behave in this way and are not grouped into doublets. Thus, the weak eigenstates of the fermion fields may be grouped as follows:

$$\begin{pmatrix} \nu_e \\ e^- \end{pmatrix}_L \quad \begin{pmatrix} \nu_\mu \\ \mu^- \end{pmatrix}_L \quad \begin{pmatrix} \nu_\tau \\ \tau^- \end{pmatrix}_L \quad (\epsilon^-)_R \quad (\mu^-)_R \quad (\tau^-)_R \quad (2.10)$$

$$\begin{pmatrix} u \\ d' \end{pmatrix}_L \quad \begin{pmatrix} c \\ s' \end{pmatrix}_L \quad \begin{pmatrix} t \\ b' \end{pmatrix}_L \quad (u)_R \quad (c)_R \quad (t)_R \quad (d)_R \quad (s)_R \quad (b)_R \quad . \quad (2.11)$$

The upper and lower fermions in the left-handed doublets are distinguished by their “weak hypercharge” quantum numbers of $T_3^f = +\frac{1}{2}$ and $T_3^f = -\frac{1}{2}$, respectively. The right-handed singlets are all assigned weak hypercharge quantum numbers of $T_3^f = 0$.

To preserve the $SU(2)$ symmetry of the left-handed fermions under the weak interaction, it is necessary to introduce three gauge bosons, but in the simplest form of the theory, these are required to be massless. This arises because explicit mass terms in the gauge boson Lagrangian render the theory non-renormalizable and thus, it loses its predictive power when calculations beyond first order are performed. The requirement that the gauge bosons be massless is in sharp contrast to experimental data since the $W^\pm$ and Z^0 gauge bosons are found to be some of the most massive particles yet discovered. Furthermore, if electromagnetic interactions are to be described simultaneously, it is necessary to relate the weak and electromagnetic coupling constants to prevent unphysical behaviour of the theory at high energies.

These issues led to the unification of the electromagnetic and weak forces experienced by the particles grouped into the left-handed weak doublets. In the unified Electroweak theory, these must be invariant under the combined symmetry

group $SU(2) \times U(1)$. When written in its simplest form, this theory still required that the gauge bosons be massless but in 1964, Higgs showed how the introduction of an additional doublet of spin-0 fields could solve this problem [10, 11]. It was shown that the gauge bosons would couple to this scalar doublet in a way which produced the required effective mass terms for the $W^\pm$ and Z^0 bosons, while leaving the photon massless. In this way, a unified theory of the electromagnetic and weak forces was constructed which describes all known properties of both interactions.

The fundamental parameters in the non-unified theories of electromagnetism and weak interactions could be taken to be the quantum of electric charge, e , the Fermi constant, G_F , and the intermediate vector boson masses, M_W and M_Z . Although unification of the theories replaced them by four new parameters, which can be taken as a weak coupling constant g , a gauge boson mass M_Z , the weak mixing angle θ_W and the unspecified mass of the Higgs boson, M_H , the combined theory is now renormalizable. The original parameters are also related to the new parameters by simple expressions:

$$e = g \sin \theta_W \quad (2.12)$$

$$\frac{G_F}{\sqrt{2}} = \frac{g^2}{8M_W^2} \quad (2.13)$$

$$M_W^2 = \rho M_Z^2 \cos^2 \theta_W. \quad (2.14)$$

The parameter ρ is unity to first order but can be calculated to higher orders using perturbation theory.

2.3.1 The CKM Matrix

The left-handed d' , s' and b' states shown in equation (2.11) are eigenstates of the weak interaction, but these turn out not to be the same as the quark eigenstates described by QCD. Instead, the two different sets of eigenstates are related by a unitary transformation which may be written in terms of the Cabibbo-Kobayashi-

Maskawa (CKM) matrix:

$$\begin{pmatrix} d' \\ s' \\ b' \end{pmatrix} = \begin{pmatrix} V_{ud} & V_{us} & V_{ub} \\ V_{cd} & V_{cs} & V_{cb} \\ V_{td} & V_{ts} & V_{tb} \end{pmatrix} \begin{pmatrix} d \\ s \\ b \end{pmatrix}. \quad (2.15)$$

For only two generations of quarks, the corresponding 2×2 matrix could be specified in terms of a single parameter, the Cabibbo mixing angle, θ_C , while for three generations, the full CKM matrix can be specified in terms of four parameters, three mixing angles and a single complex phase. Since the electroweak model treats left-handed and right-handed helicity states differently, it is not invariant under parity transformations. However, almost all processes are invariant under a combination of charge conjugation and parity reversal (CP). Nevertheless, small levels of CP violation have been observed in processes such as $K_L^0 \rightarrow \pi\pi$, and these might be accounted for in the Standard Model by the complex phase in the CKM matrix. This possibility has still not been tested to a satisfactory level and hence, there is great interest in accurate measurements of the CKM matrix elements.

The elements of the CKM matrix can be measured by studying processes involving weak transitions between hadronic states. Unfortunately, precision measurements require that the momentum distribution of the quark undergoing the weak interaction inside the hadron be well understood. Theoretical descriptions of this momentum distribution cannot be calculated since a perturbative treatment of QCD does not apply for momentum scales typical of quarks bound inside hadrons. Furthermore, when several hadrons are present in the initial or final states, additional non-perturbative strong interactions between them can complicate the theoretical prediction. Part of this problem can be avoided by studying semileptonic decay transitions, in which a hadron interacts with or decays into a lepton. This is because the leptonic states can be described perturbatively by the Electroweak theory, leaving only the interaction of the quark with the exchanged $W^\pm$ to be described, or parameterized by a phenomenological model.

Figure 2.1: Feynman diagram for a semi-leptonic process. The 4-momentum transfer resulting from the weak transition $q_i \rightarrow q_j$ of a quark in the initial hadron h is denoted by q , and $V_{q_j q_i}$ represents the relevant CKM matrix element.

2.3.2 Semi-leptonic Decays

The matrix element for a typical semileptonic interaction may be written

$$\mathcal{M} = \frac{g^2}{2\sqrt{2}} \bar{u}(\ell) \gamma_\mu (1 - \gamma_5) v(\bar{\nu}_\ell) \frac{1}{M_W^2 - q^2} \langle h' | J_h^\mu | h \rangle \quad (2.16)$$

where $g = e/\sin \theta_W$ is the weak coupling constant and the other factors involved are shown on the Feynman diagram for this process in figure 2.1. In this expression, J_h^μ is the weak hadronic current which may be written in terms of quark creation and annihilation operators as follows:

$$J_h^\mu = V_{q_j q_i} \bar{\psi}_{q_j} \gamma^\mu (1 - \gamma_5) \psi_{q_i} \quad (2.17)$$

where $V_{q_j q_i}$ is the relevant CKM matrix element. Although the leptonic vertex can be calculated precisely, the quarks at the hadronic vertex are bound inside hadrons

which introduces a dependence on the momentum transfer q^2 of the process. This less well known dependence is parameterized by introducing a function of q^2 called a form factor.

The CKM matrix element V_{q,q_i} is contained within the $\langle h' | J_h^\mu | h \rangle$ term and can only be determined from experiment using a theoretical calculation of the form factors involved. This description essentially requires knowledge of the hadronic wavefunctions and during the 1980's there was much attention devoted to the problem of calculating them. By the late 1980's, models had been developed to describe semileptonic decays, but based on these, it was expected, for example, that the semileptonic decay rate of B mesons should be dominated by decays to D and D* mesons. Considerable theoretical interest was generated when it was found experimentally that semileptonic decays to these final states accounted for only about 66% of the total semileptonic decay rate of B mesons[12].

At the same time, Isgur and Wise[13, 14] developed a powerful way of studying systems of heavy mesons in which they are viewed from the rest frame of the heavy quark. In such a reference frame, a weak $b \rightarrow c$ transition in which the charm quark receives no net momentum transfer leaves the remaining light quark in the B meson unaffected since the colour fields from the b and c quarks, acting on the light antiquark, are identical. These considerations allow the hadronic matrix elements between B and D or D* states to be expressed in terms of a single universal function $\xi(v \cdot v')$, where v and v' are the 4-velocity of the heavy quark in the initial and final states. With some modifications, it was possible to describe transitions to higher excited states, including P-wave charm mesons, in terms of the same Isgur-Wise function ξ and it was realised that these additional final states might account for the remaining fraction of semileptonic B decays not yet exclusively reconstructed.

Thus, there is good reason to study semileptonic B decays to final states involving P-wave charm mesons. Significant experimental progress has been made and it is well established that such transitions do occur at a rate which might well

account for the observed difference between the exclusive and inclusive semileptonic branching fractions in B-decays. Since these states are observed in semileptonic B decays one expects them to be present in hadronic decays as well, and it is for this reason that an inclusive measurement of their production rate becomes intriguing.

2.3.3 The Spectator Model and Hadronic B Decays

While accurate calculations of the overlap of initial- and final-state wavefunctions of heavy hadrons in semileptonic decays can be obtained using the techniques developed by Isgur and Wise, a much simpler way of viewing weak decays of heavy hadrons is the spectator model. This model considers the weak decay of a b-hadron to proceed as if the b-quark were essentially free, with no influence from the other quarks in the hadron. Although this description is greatly simplified, it can be used to account for many properties of these decays.

Figure 2.2 shows most of the relevant diagrams for understanding the decays of B-mesons. Figures 2.2(a) shows the spectator diagram for the semileptonic decay, while figures 2.2(b) and (c) show the corresponding hadronic decay diagrams. Figures 2.2(d-f) show other possible decays but these occur at rates much smaller than the principal spectator decay modes.

The hadronic spectator decays, are the dominant diagrams for the hadronic decays of heavy hadrons. Matrix elements which describe these processes are more complicated than those used to describe semileptonic decays since the production and interactions of the final state hadrons must be accounted for. However, it is found that to a good approximation the production and decay of the virtual W^- boson can be treated independently. This is referred to as factorization and is used in a model proposed by Bauer, Stech and Wirbel[15, 16] to quantitatively describe relations among the exclusive branching ratios of B-mesons. It assumes that all final state interactions can be contained within the initial and final state wavefunctions and weights the external and internal W-emission diagrams (b) and (c) by factors a_1 and a_2 in an effective Hamiltonian.

Figure 2.2: Diagrams for various B decay processes: (a) semileptonic decay, (b) external W emission, (c) internal W emission, (d) W exchange, (e) electromagnetic penguin decay and (f) hadronic penguin decay.

These parameters have been determined from fits to the branching ratios of exclusive B decays, measured at ARGUS[17] and CLEO[18], and favour values of $a_1 = 1.07 \pm 0.04 \pm 0.06$ and $a_2 = 0.23 \pm 0.01 \pm 0.01$. A simple interpretation can be provided by the idea of colour-matching, whereby diagram (c) is suppressed by a factor of $1/N_C \approx 0.3$ because of the requirement that the final state quarks form colourless hadrons. The success with which these ideas reproduce the ratios of many exclusive branching fractions indicates that they can be a very useful tool in understanding the properties of hadronic b-decays. More precise theoretical models for decays of b-hadrons include effects such as hard gluon bremsstrahlung and interference effects between final state quarks but it appears that much of b-decay phenomenology can be understood in terms of the naive spectator model.

2.4 Summary

This chapter has briefly described the key aspects of the theories out of which the Standard Model has been assembled. Although first motivated as a way to measure CKM matrix elements, the study of hadrons containing bottom quarks was seen to hold a wider interest because of the many properties which could be understood using the simple spectator model. The success of the spectator model is due largely to the fact that the b-quark mass is large, and as will be seen in the next chapter, large quark masses can simplify the description of hadronic structure in other ways as well.

The discussion of b-decays was not provided simply to give an example of a system containing a heavy quark. Charm quarks are almost always produced in the decays of b-hadrons and these result in sizeable branching fractions to D and D* final states. However, one also expects the decays of b-hadrons to lead to significant production rates of P-wave charm mesons. Although this thesis primarily focuses on the way in which P-wave charm mesons are produced in $Z^0 \rightarrow c\bar{c}$ events, since $Z^0 \rightarrow b\bar{b}$ events are also present, b-hadron decays produce an important second

source of these particles.

The following chapter describes our understanding of the properties of P-wave charm mesons and the way by which they are produced in the charm quark events at LEP. However, as will be seen in subsequent chapters, confronting these issues experimentally will require measuring their production rates separately in both $c\bar{c}$ and $b\bar{b}$ events. Both are interesting measurements and their implications are discussed separately in the concluding chapters.

Chapter 3

Theoretical Framework

3.1 Z^0 Production and Decays at LEP

Although the Z^0 boson was discovered using a $p\bar{p}$ collider, the most practical method for creating a “ Z^0 factory” is to collide electrons and positrons. The cross section for e^+e^- annihilation with the production of a hadronic final state is shown for a wide range of centre-of-mass energies in figure 3.1. There is a great resonant enhancement in the cross section at an energy corresponding to the Z^0 mass, M_Z .

The Z^0 resonance is characterized by its mass and width which are determined by measurements of the cross section at precisely determined centre-of-mass energies close to M_Z . Equally important are measurements of the strength with which the Z^0 couples to each type of fermion in the final state. The Z^0 couples with different strengths to left-handed and right-handed fermions which can be parameterized in terms of vector and axial coupling constants v_f and a_f . These are related to the electric charge Q_f and weak-hypercharge T_f^3 by $v_f = T_f^3 - 2Q_f \sin^2 \theta_W$ and $a_f = T_f^3$. To lowest order and neglecting the masses of the final state particles, the partial widths for $Z^0 \rightarrow f\bar{f}$ may be written in terms of these coupling constants as

Figure 3.1: Measured cross sections for $e^+e^- \rightarrow \text{hadrons}$ and $e^+e^- \rightarrow \mu^+\mu^-$ at various centre-of-mass energies.

Parameter	Measured	Predicted
Γ_{total} (MeV)	2494.6 ± 2.7	2497
Γ_{had} (MeV)	1743.6 ± 2.5	1743.8
$\Gamma_{\text{invisible}}$ (MeV)	499.5 ± 2.0	501.6
$\Gamma_{e^+e^-}$ (MeV)	83.96 ± 0.15	83.97
$\Gamma_{\mu^+\mu^-}$ (MeV)	83.79 ± 0.22	83.97
$\Gamma_{\tau^+\tau^-}$ (MeV)	83.72 ± 0.26	83.97
$R_c \equiv \Gamma_{c\bar{c}}/\Gamma_{\text{had}}$	0.1715 ± 0.0056	0.172
$R_b \equiv \Gamma_{b\bar{b}}/\Gamma_{\text{had}}$	0.2179 ± 0.0012	0.2156

Table 3.1: A comparison of some of the Standard Model parameters with recent measurements. Theoretical uncertainties in the predictions are small compared with the experimental uncertainties.

follows:

$$\Gamma_{f\bar{f}} = \frac{N_{C_f} G_F M_Z^3}{6\sqrt{2}\pi} (v_f^2 + a_f^2) \quad (3.1)$$

where N_{C_f} is a factor which is approximately 3 for quarks and exactly 1 for leptons. After finite-mass effects and higher-order corrections are applied, such an expression yields the predicted partial decay widths of the Z^0 shown in table 3.1[12]. These are all found to be in excellent agreement with world averages of experimental measurements [19]. From this table one sees that the Z^0 decays to hadrons approximately 70% of the time and of these hadronic decays, 17% produce a $c\bar{c}$ and 22% produce a $b\bar{b}$ final state.

The Z^0 has different coupling constants for fermions in left- and right-handed helicity states. This results in a polarization asymmetry which can be expressed in terms of a_f and v_f as follows:

$$A_{\text{pol}}^f \equiv -\frac{\sigma(h_f = +1/2) - \sigma(h_f = -1/2)}{\sigma(h_f = +1/2) + \sigma(h_f = -1/2)} = -\frac{2v_f a_f}{v_f^2 + a_f^2}. \quad (3.2)$$

Polarization asymmetries in Z^0 decays have only been measured for $Z^0 \rightarrow \tau^+\tau^-$ decays since the τ lifetime is short and the angular distribution of the τ decay

products is highly correlated to its polarization. While the lifetimes of c and b quarks are also short, they are very long compared with the time required for quarks to become bound inside stable hadrons and it is difficult to predict the way in which the helicity of the initial quark will be transferred to the hadron in this process. Although parity invariance of the strong interaction prevents asymmetries from being observed in strong decays of heavy mesons, it can produce a nonuniform population of their helicity states. This can sometimes complicate assumptions about the expected angular distributions of the decay products of hadrons produced in Z^0 decays.

3.2 Quark Fragmentation

Fragmentation refers to the process whereby quarks or gluons, produced in some high energy interaction, eventually evolve to form the stable hadrons which are observed in a detector. This process is presumably described by QCD, but the calculations required to make quantitative predictions are far too complex to be carried out. Hence, no theory for fragmentation exists which is based purely on first principles. Instead, a number of models have been developed which are inspired by our qualitative understanding of QCD. The entire process is depicted schematically in figure 3.2 and consists of three phases which, although conceptually distinct, are probably poorly separated in practice.

3.2.1 The Parton Shower

Since the Z^0 mass is much larger than that of even the b quark, all quarks from Z^0 decays are produced with high energies. At these energy scales the strong coupling constant has a value of approximately $\alpha_s \sim 0.12$ which is small enough that perturbative QCD calculations are believed to be reliable. These have been performed to calculate matrix elements describing the production of up to 4 partons in the final state, but it has been found that this alone is insufficient to describe

Figure 3.2: Schematic representation of the Z^0 fragmentation process. An e^+e^- pair annihilates to form the Z^0 (a) which subsequently decays to $q\bar{q}$ (b). The parton shower develops (c) until the nonperturbative fragmentation process takes over (d). Out of this emerge hadrons which can decay (e) before being observed in the detector (f). In reality, the boundaries between regions (c), (d) and (e) are not well defined.

some properties of hadronic events [20]. Better agreement is obtained by allowing successive $q \rightarrow qg$, $g \rightarrow gg$ and $g \rightarrow q\bar{q}$ parton processes to proceed until the available energy, Q , for subsequent parton branching falls below some tunable cutoff parameter $Q_{\min}$ which is typically of the order of 1 GeV.

The probabilities for each type of splitting can be calculated perturbatively to first order, with the effects of low energy gluons and gluon emission at small angles included analytically to all orders. These successive parton branchings give rise to a tree-like parton configuration. This is clearly only an approximation to the full perturbative expression since there are no interfering diagrams or loops. However, since loops are suppressed by factors of $1/N_c$ and the bulk of the interference effects are included in the parton splitting functions, the idea of a parton shower is considered to be a plausible description of the initial stages of fragmentation.

3.2.2 Fragmentation Models

The physics which describes the next stage of fragmentation is less well defined since it involves energy scales of the order of 1 GeV or lower. To describe such processes a number of fragmentation models have been proposed.

Independent Fragmentation

One simple model assumes each quark moving along the fragmentation axis with momentum p_L shares a fraction z of this momentum with $q\bar{q}$ pairs pulled from the vacuum. The $\bar{q}$ and the original quark form a meson and the remaining quark continues to fragment by pulling additional quarks out of the vacuum. The probability density $f(z)$ is known as the fragmentation function and a number of empirical parameterizations exist. These can be different for gluons or for the various quark flavours and generally contain adjustable parameters. These can be tuned to reproduce experimentally measured features such as the average scaled energy of the hadrons produced.

One fragmentation function is a formula proposed by Peterson *et al.* [21]:

$$f_Q(z) \propto \frac{1}{z \left(1 - \frac{1}{z} - \frac{\epsilon_Q}{1-z}\right)^2}. \quad (3.3)$$

It is commonly used to describe the resulting energy spectrum of hadrons containing heavy quarks which peaks towards $z = 1$. ϵ_Q is expected to scale like $1/m_Q^2$ and typical values used for the fragmentation of bottom and charm quarks are $\epsilon_c = 0.031$ and $\epsilon_b = 0.0038$.

Although baryon production can be incorporated in the model by allowing di-quarks to be pulled from the vacuum, there is no clear prescription for the treatment of gluons. Furthermore, in its simplest form, the independent fragmentation model is not Lorentz invariant and can violate total flavour, energy and momentum conservation to some degree. Its inability to reproduce particle correlations [22] or the “string effect” [23] cause this model to be disfavoured but nevertheless, independent fragmentation models are often used as benchmarks when comparing observed event properties with a variety of models.

Cluster Fragmentation

While independent fragmentation models form hadrons at each iteration, a more physical picture might be that the quarks and gluons produced by a parton shower rearrange themselves into loosely bound states which subsequently decay to lighter stable hadrons. These intermediate states are referred to as clusters and have masses typical of the $\sim 1 \text{ GeV}/c^2$ nonperturbative scale which is reached at the end of the parton shower. In this picture, gluons must first be split into $q\bar{q}$ pairs before the clusters are formed. Clusters may roughly correspond to wide, high-mass, strongly decaying resonances and their decays are often generated using phase space kinematics to produce lighter clusters, hadronic resonances or stable hadrons.

String Fragmentation

A third major class of fragmentation models is based on the fact that the strong coupling constant grows with increasing separation between partons. Thus, as quarks move apart, the energy stored in the gluon field which binds them increases. In this state the gluon field is believed to form a tube-like structure or string. Models for interquark potentials and lattice QCD studies indicate that the force experienced by the quarks at large distances is approximately linear and hence, the string is characterized by an effective spring constant, or energy per unit length of $\kappa \approx 1 \text{ GeV/fm}$. As the $q\bar{q}$ pair produced in a Z^0 decay move apart, their kinetic energy is transferred to the string and when this energy is sufficient to create a $q'\bar{q}'$ pair from the vacuum, the string breaks to produce two new strings connecting the $q\bar{q}'$ and $q'\bar{q}$. These have the energy of the original string distributed between them according to a fragmentation function such as the Peterson function shown in equation (3.3). This process continues until the energy stored in the string segments is comparable to hadronic masses whereupon they are transformed into pairs of hadrons. Gluon emission corresponds to transverse “kinks” on a string joining two quarks and hence, breaks in the string segment joining the quark and gluon will populate the region between them with hadrons. This is known as the “string effect” and is a feature of string fragmentation models which is in qualitative agreement with experimental observations [23].

3.3 QCD Potential Models

It was realized in 1975 that the masses of the excited $c\bar{c}$ states observed in e^+e^- collisions could be explained using non-relativistic potential models. These models used the fact that at short distances quarks couple weakly, while being confined strongly only at large distances. The short distance potential could therefore be calculated to first order in α_s , assuming a one-gluon exchange interaction while the confinement was introduced as a long range scalar potential. This resulted

in a spin-dependent potential which described the rich spectrum of $c\bar{c}$ states and predicted the observed partial decay widths.

It was later shown by Godfrey and Isgur [24] that the one-gluon exchange potential with a linear confining term could explain the mass spectra and decay widths of all mesons, provided a number of relativistic and strong interaction effects were accounted for. In particular, the meson masses were determined by solving the Schrödinger equation

$$H|\Psi\rangle = (H_0 + V(\mathbf{r}, \mathbf{p}))|\Psi\rangle = E|\Psi\rangle \quad (3.4)$$

with H_0 the relativistic energy of the quark-antiquark pair and $V(\mathbf{r}, \mathbf{p})$ the potential between them.

Although the relativistic effects are important, a non-relativistic approximation can describe many of the properties of a system involving heavy quarks. Therefore, we examine the features of the potential in a non-relativistic limit. In this limit, the kinetic and potential parts of the Hamiltonian may be reduced to

$$H_0 = \sum_{i=1}^2 \sqrt{p^2 + m_i} \rightarrow \sum_{i=1}^2 \left(m_i + \frac{p^2}{2m_i} \right) \quad (3.5)$$

and

$$V(\mathbf{r}, \mathbf{p}) \rightarrow V(\mathbf{r}) = H_{\text{conf}} + H_{\text{hyp}} + H_{\text{SO}} \quad (3.6)$$

where the terms in the potential may be written as a spin-independent confinement part, a hyperfine spin-spin interaction potential and terms describing spin-orbit interactions. These, in turn, may be written as follows:

$$H_{\text{conf}} = -\mathbf{F}_1 \cdot \mathbf{F}_2 \left(c + br - \frac{\alpha_s(r)}{r} \right) \quad (3.7)$$

$$\begin{aligned} H_{\text{hyp}} = & -\mathbf{F}_1 \cdot \mathbf{F}_2 \frac{\alpha_s(r)}{m_1 m_2} \left(\delta^3(\mathbf{r}) \frac{8\pi}{3} \mathbf{S}_1 \cdot \mathbf{S}_2 \right. \\ & \left. + \frac{1}{r^3} (3(\mathbf{S}_1 \cdot \hat{\mathbf{r}})(\mathbf{S}_2 \cdot \hat{\mathbf{r}}) - \mathbf{S}_1 \cdot \mathbf{S}_2) \right) \end{aligned} \quad (3.8)$$

$$\begin{aligned}
H_{\text{SO}} = & -\frac{1}{2r} \frac{\partial H_{\text{conf}}}{\partial r} \left(\frac{\mathbf{L} \cdot \mathbf{S}_1}{m_1^2} + \frac{\mathbf{L} \cdot \mathbf{S}_2}{m_2^2} \right) + \\
& -\mathbf{F}_1 \cdot \mathbf{F}_2 \frac{\alpha_s(r)}{r^3} \left(\frac{1}{m_1} + \frac{1}{m_2} \right) \left(\frac{\mathbf{L} \cdot \mathbf{S}_1}{m_1} + \frac{\mathbf{L} \cdot \mathbf{S}_2}{m_2} \right)
\end{aligned} \tag{3.9}$$

where $\langle \mathbf{F}_1 \cdot \mathbf{F}_2 \rangle = -4/3$ are colour factors for mesons arising from the $SU(3)$ generators of QCD. $\mathbf{S}_i$ are the spins of the quarks and $\mathbf{L}$ is their orbital angular momentum. The first term in H_{SO} describes the colour-magnetic interaction between the two quarks, while the second term accounts for the Thomas precession. This non-relativistic reduction is sufficient to describe the important features of the mass spectrum of P-wave charmed mesons.

3.4 Properties of Heavy P-wave Mesons

A meson composed of a heavy quark Q and a light $\bar{u}$, $\bar{d}$ or $\bar{s}$ anti-quark $\bar{q}$ corresponds to the case where $m_1 \gg m_2$. The properties of this configuration are identical to the case of a heavy anti-quark $\bar{Q}$ and a light quark q , but for the sake of definiteness, we consider the former situation. In this case, some terms in equation (3.9) are suppressed by factors of $1/m_Q$ and in particular, since the spin dependent terms are of the form $\mathbf{S}_i/m_i$, the effects of the heavy quark's spin are reduced. In this heavy quark limit the dominant spin-dependent terms arise only from the light-quark spin-orbit coupling and hence, the potential may be written¹

$$V(\mathbf{r}) = H_{\text{conf}} + c_q \mathbf{L} \cdot \mathbf{S}_q + \mathcal{O}(1/m_Q) \tag{3.10}$$

where

$$c_q = \frac{1}{2m_q^2} \left\langle \frac{4}{3} \frac{\alpha_s}{r^3} \right\rangle - \frac{1}{2m_q^2} \left\langle \frac{b}{r} \right\rangle. \tag{3.11}$$

¹The order in which the angular momentum is coupled is important. In this notation the order is $\mathbf{S}_q \times \mathbf{S}_Q \times \mathbf{L} = \mathbf{S}_Q \times \mathbf{L} \times \mathbf{S}_q$.

Thus, the properties of heavy-light mesons should depend only on the spin $\mathbf{S}_q$ of the light anti-quark and on its relative orbital angular momentum $\mathbf{L}$.

When $L = 0$, the meson mass depends only on the spin alignment of the heavy and light quarks. Anti-parallel spins correspond to the ground state pseudoscalar mesons, denoted D while aligned spins give rise to the D^* vector mesons. When the quarks are in a P-wave state, the spins and orbital angular momentum combine to form a more complicated multiplet of states which are collectively referred to as D^{**} mesons. Elements of the canonical basis with which these states can be represented are labeled using the notation $^{2S+1}L_J$ where S is the total spin and J is the total angular momentum. This basis consists of the 1P_1 singlet state and the triplet ($^3P_0, ^3P_1, ^3P_2$). Because $L = 1$, all these P-wave states have even parity and since the spin-1 singlet and triplet states have identical quantum numbers, the physical spin-1 states will generally be a mixture of the two.

The mass scale $\overline{M}$ of the P-wave system is determined by the spin-independent part of the Hamiltonian and the masses of individual states are obtained by diagonalizing the mass matrix. In the $^{2S+1}L_J$ basis, the unmixed states 3P_0 and 3P_2 have masses

$$M(^3P_2) = \overline{M} + c_2/2 \quad (3.12)$$

$$M(^3P_0) = \overline{M} - c_2 \quad (3.13)$$

while the two states with $J = 1$ mix as follows:

$$M \begin{pmatrix} ^3P_1 \\ ^1P_1 \end{pmatrix} = \begin{pmatrix} \overline{M} - c_2/2 & -c_2/\sqrt{2} \\ -c_2\sqrt{2} & \overline{M} \end{pmatrix} \begin{pmatrix} ^3P_1 \\ ^1P_1 \end{pmatrix}. \quad (3.14)$$

In almost all models c_2 is positive and in this case, the 3P_1 and 1P_1 states mix with a mixing angle given by $\tan \phi_0 = 1/\sqrt{2}$. The eigenvalues of this matrix are

$$M_\alpha = \overline{M} - c_2 \quad (3.15)$$

$$M_3 = \overline{M} + c_2/2. \quad (3.16)$$

corresponding to the two eigenstates

$$|\alpha\rangle = \sqrt{\frac{2}{3}}|{}^3P_1\rangle + \sqrt{\frac{1}{3}}|{}^1P_1\rangle = |P_1(\frac{1}{2})\rangle \quad (3.17)$$

$$|\beta\rangle = \sqrt{\frac{1}{3}}|{}^3P_1\rangle - \sqrt{\frac{2}{3}}|{}^1P_1\rangle = |P_1(\frac{3}{2})\rangle. \quad (3.18)$$

where $|L_J(j_q)\rangle$ are eigenstates in which the light quark has total angular momentum j_q with respect to the heavy quark. Thus, in the limit $m_Q \rightarrow \infty$ one obtains two degenerate doublets with properties determined only by the total angular momentum of the light quark. $j_q = 1 \pm s_q$.

3.5 Heavy Quark Spin Symmetry

The general consequence that the spin of a heavy quark is irrelevant in the limit $m_Q \rightarrow \infty$ is a result known as heavy quark spin symmetry. In this limit, the heavy quark spin commutes with the Hamiltonian and is therefore conserved. The conservation of total angular momentum then implies that both the heavy quark spin and the angular momentum of the light degrees of freedom are independently conserved. Several implications of this were described by Isgur and Wise [25]. In particular, this symmetry was used to predict the relative amplitudes for the decays of heavy mesons.

The decays of P-wave mesons are expected to be dominated by two-body final states such as $D^{**} \rightarrow D^*\pi$ or $D^{**} \rightarrow D\pi$. Of those decays allowed by kinematics, parity and angular momentum conservation of the strong interaction restrict the decay partial waves amplitudes to have $L = 0$ or $L = 2$. Clearly, the 3P_2 and 3P_0 states will decay as D- and S-waves, respectively, while the 3P_1 and 1P_1 states could decay with both amplitudes. However, the physical states are the mixtures shown in equation (3.18) and because of heavy quark spin symmetry, the $|P_1(\frac{1}{2})\rangle$ will decay as a pure S-wave, while the $|P_1(\frac{3}{2})\rangle$ state will decay through a D-wave.

In the decay $H \rightarrow (H'h)_L$ where the final state hadrons H' and h have relative orbital angular momentum L , the partial width is suppressed by a factor of p^{2L+1} , where p is the momentum of the light hadron h in the rest frame of the decaying meson. Thus, the amplitudes for states decaying as D-waves are suppressed which suggests that the $j_q = \frac{3}{2}$ doublet should form narrow resonances while the $j_q = \frac{1}{2}$ states should be very broad. The physical states corresponding to the members of the narrow doublet are denoted D_1 and D_2^* and, as will be described in section 3.6, have been experimentally observed.

3.5.1 Decay Properties

Because of the independent conservation of heavy and light quark degrees of freedom, the relative amplitudes for the decays of heavy mesons can be simply related to products of Clebsch-Gordon coefficients. This allows one to calculate the relative partial widths and the branching ratios for the various states and also to calculate angular distributions of decay products. Because the strong interaction is symmetric under charge conjugation, the following discussion, formulated in terms of the properties of D-mesons, applies equally well to the properties of anti-D mesons. This implied charge conjugation notation is used frequently throughout this thesis to simplify the discussion.

Partial Widths

In the heavy quark limit, both members of the $j_q = 3/2$ doublet decay via a D-wave into two body final states. The relative partial widths for decays to $D\pi$ and $D^*\pi$ are then as follows:

$$D_2^* \rightarrow \begin{cases} (D^*\pi)_{L=2} & : 3/5 \\ (D\pi)_{L=2} & : 2/5 \end{cases} \quad (3.19)$$

$$D_1 \rightarrow \begin{cases} (D^*\pi)_{L=2} & : 1 \\ (D\pi)_{L=2} & : 0 \end{cases} . \quad (3.20)$$

These decay widths are believed to saturate the total widths of the narrow resonances since decays involving higher mass ρ or η mesons in the final state are below threshold. However, these ratios are modified by factors of $|\mathbf{p}_\pi|^{2L+1}$ where $\mathbf{p}_\pi$ is the pion momentum in the rest frame of the decaying D_1^0 or D_2^{*0} . When $L = 2$ these can change the predicted branching ratios significantly.

Branching Ratios

Expected branching ratios for specific charge states can be obtained from the partial widths and isospin considerations. For example, the Clebsch-Gordon decomposition of the D_1^0 decay can be written

$$|D_1^0\rangle = -\sqrt{\frac{2}{3}}|D^{*+}\rangle|\pi^-\rangle + \sqrt{\frac{1}{3}}|D^{*0}\rangle|\pi^0\rangle. \quad (3.21)$$

After also including the kinematic factors of $|\mathbf{p}_\pi|^{2L+1}$, this leads to the following predictions for the D^{*0} branching ratios:

$$D_2^{*0} \rightarrow \begin{cases} D^{*+}\pi^- & : 0.19 \\ D^{*0}\pi^0 & : 0.10 \\ D^+\pi^- & : 0.47 \\ D^0\pi^0 & : 0.24 \end{cases} \quad (3.22)$$

$$D_1^0 \rightarrow \begin{cases} D^{*+}\pi^- & : 0.65 \\ D^{*0}\pi^0 & : 0.35 \end{cases} . \quad (3.23)$$

Similar arguments can be applied to the case when a P-wave charm meson contains a strange quark. Again, the $j_q = 3/2$ doublet is expected to form narrow resonances but in this case, decays to the $D^{*+}K^0$ and $D^{*0}K^+$ final states occur so close to threshold that their amplitudes are severely restricted by phase space. This suppression explains the very narrow width of the D_{s1}^+ for which measurements have only set upper limits. Because of isospin symmetry the D_{s1}^+ is expected to decay with equal probability to the $D^{*0}K^+$ and $D^{*+}K^0$ final states but again, this

is modified by kinematic factors of $|\mathbf{p}_K|^{2L+1}$. Taking these factors into account yields the expected branching ratios:

$$D_{s1}^+ \rightarrow \begin{cases} D^{*+}K^0 & : 0.35 \\ D^{*0}K^+ & : 0.65 \end{cases} \quad (3.24)$$

and predicts the ratio of these branching ratios to be about 1.8.

As was the case for the non-strange D_2^{*0} state, when phase space effects are neglected, the D_{s2}^{*+} should decay to D^*K and DK with branching ratios of 0.6 and 0.4, respectively. Again, suppression factors of the form $|\mathbf{p}_K|^{2L+1}$ severely restrict the D^*K amplitude and consequently D_{s2}^{*+} decays are expected to be dominated by the DK final state.

Angular Distributions

The decays of excited charm mesons proceed via the strong interaction so their decay amplitudes are invariant under parity transformations. This restricts the form of the D^{*+} and D_2^{*0} decay amplitude such that these decays can be expressed in terms of one complex parameter, while decays of the D_1^0 require two parameters. These two parameters can be rewritten in terms of the amplitudes for S- and D-wave decays. Since the narrow P-wave states will only undergo D-wave decays in the heavy quark limit, the angular distributions of the D_1^0 and D_2^{*0} decay products can be predicted unambiguously in this limit.

When studying the D^{*-0} decay chain

$$D^{*-0} \longrightarrow \begin{array}{c} D^{*+}\pi^- \\ \quad \searrow \\ \quad D^0\pi^+ \end{array} \quad (3.25)$$

the angle most useful for analysing the spin J of the D^{*-0} meson is referred to as the helicity angle α , shown in figure 3.3. This variable is useful because its distribution does not depend on the initial polarization of the decaying D^{*-0} . In

Figure 3.3: Definition of the angle α in the rest frame of the D^{*-+} . π_s^+ denotes the pion from the D^{*-+} decay, while the π^- is the pion emitted in the decay of the D^{*-0} .

the heavy quark limit, the angular distributions for the decays of the parent states with spin-parity J^P are as follows:

$$\frac{dN}{d(\cos \alpha)} \propto \begin{cases} 1 - \cos^2 \alpha & \text{for a } 2^+ \text{ decay,} \\ 1 + 3 \cos^2 \alpha & \text{for a D-wave } 1^+ \text{ decay.} \\ 1 & \text{for an S-wave } 1^+ \text{ decay.} \end{cases} \quad (3.26)$$

A derivation of these angular distributions can be carried out using the helicity formalism described by Chung[26] and Richman[27] and is shown explicitly in appendix A.

3.5.2 Production Properties

A naive description of heavy quark fragmentation can be expressed in terms of the probabilities for a heavy quark to acquire a light-quark with helicity h_ℓ such that a heavy hadron H with spin s and helicity h_s is formed. The probability for forming such a state is

$$P_{Q,h_Q \rightarrow s,h_s}^{(H)} = P_{Q \rightarrow s_\ell} \sum_{h_\ell} p_{h_\ell} |C_{s_\ell h_\ell; \frac{1}{2} h_Q}^{s h_s}|^2 \quad (3.27)$$

where $C_{s_\ell h_\ell; \frac{1}{2} h_Q}^{s h_s}$ are Clebsch-Gordon coefficients. $P_{Q \rightarrow s_\ell}$ is the probability that fragmentation will occur to a doublet in which the light quark degrees of freedom have

angular momentum s_ℓ , and p_{h_ℓ} are the conditional probabilities that they will have helicities h_ℓ . Parity invariance requires that $p_{h_\ell} = p_{-h_\ell}$ which fully specifies the relative populations of the helicity states for D^0 and D^{*+} production. For a charm quark with $h_Q = 1/2$, these are

$$P_{c,1/2 \rightarrow 0,0}^{(D)} = 1/4 \quad (3.28)$$

$$P_{c,1/2 \rightarrow 1,h_\ell}^{(D^*)} = \begin{cases} 1/2 & : h_\ell = 1 \\ 1/4 & : h_\ell = 0 \\ 0 & : h_\ell = -1 \end{cases} \quad (3.29)$$

For fragmentation to the $j_q = 3/2$ charm doublet, the probabilities h_ℓ can be written in terms of one parameter, $w_{3/2}$. Expressed in this way, the relative population of the D_1 and D_2^* helicity states are expected to be

$$P_{c,1/2 \rightarrow 1,h_\ell}^{(D_1)} = \begin{cases} 0 & : h_\ell = 3/2 \\ (1 - w_{3/2})/4 & : h_\ell = 1/2 \\ (1 - w_{3/2})/2 & : h_\ell = -1/2 \\ 3w_{3/2}/4 & : h_\ell = -3/2 \end{cases} \quad (3.30)$$

$$P_{c,1/2 \rightarrow 2,h_\ell}^{(D_2^*)} = \begin{cases} w_{3/2} & : h_\ell = 3/2 \\ 3(1 - w_{3/2})/4 & : h_\ell = 1/2 \\ (1 - w_{3/2})/2 & : h_\ell = -1/2 \\ w_{3/2}/4 & : h_\ell = -3/2 \end{cases} \quad (3.31)$$

The total production rates, obtained from summing over the final helicities, are then independent of $w_{3/2}$ and one expects the multiplicity of D_1^0 and D_2^{*0} states to occur with a ratio of $3/5$, just as expected from simple spin-counting arguments.

State	Mass (MeV/ c^2)	Width (MeV/ c^2)	Decay channels
D_1^0	2422.0 ± 2.1	$18.9^{+4.6}_{-3.5}$	$D^{*+}\pi^-$, $D^{*0}\pi^0$
D_2^{*0}	2458.9 ± 2.0	23 ± 5	$D^{*+}\pi^-$, $D^+\pi^-$, $D^{*0}\pi^0$, $D^0\pi^0$
D_1^+	2425 ± 3	26^{+9}_{-8}	$D^{*+}\pi^0$, $D^{*0}\pi^+$
D_2^{*+}	2459 ± 4	25^{+8}_{-7}	$D^{*+}\pi^0$, $D^{*0}\pi^+$, $D^+\pi^0$, $D^0\pi^+$
D_{s1}^+	2535.35 ± 0.34	< 2.3	$D^{*+}K^0$, $D^{*0}K^+$
D_{s2}^{*+}	$2573.2^{+1.7}_{-1.6} \pm 0.9$	$16^{+5}_{-4} \pm 3$	$D^{*+}K^0$, $D^{*0}K^+$, D^+K^0 , D^0K^+

Table 3.2: Observed masses, widths and allowed decay modes of the narrow P-wave charmed mesons.

3.6 Previous Experimental Measurements

Since their discovery in 1986, orbitally excited charmed mesons have been studied in detail at the DORIS and CESR e^+e^- colliders by the ARGUS and CLEO experiments, respectively. They have also been observed in the fixed target experiments E691 and E687 at Fermilab and some evidence for D_{s1}^+ production from neutrino interactions has been presented [28]. As expected, only the narrow states have been observed directly. The world averages of their measured masses and widths are shown in table 3.2.

A broad signal corresponding to the combined D_1^0 and D_2^{*0} states was first observed by the ARGUS experiment [29] using the $D^{*+}\pi^-$ decay mode. The composition of the signal was not determined at the time and a single Breit-Wigner was fit to the data yielding a broad width of 70 ± 21 MeV/ c^2 . The Fermilab fixed target experiment E691 [30] reconstructed the D_2^{*0} state through its decay to $D^+\pi^-$ and was able to unambiguously determine its mass and width since the D_1^0 cannot decay to this final state. With the D_2^{*0} mass and width fixed, they were then able to measure the signal properties of the D_1^0 in the $D^{*+}\pi^-$ channel. This analysis was later repeated by the ARGUS experiment [31]. The most precise determination of the masses and widths of the D_1^0 and D_2^{*0} have been made in this way by the E687 fixed target experiment [32] and by CLEO collaboration [33].

The charged isospin partner to the D_2^{*0} was observed in 1989 by ARGUS [34]

through its decay $D_2^{*+} \rightarrow D^0 \pi^+$. Again, the spin-one member of the charged $j_q = 3/2$ doublet cannot decay to this final state and the mass and width of the D_2^{*+} were unambiguously determined. The E691 experiment observed an enhancement in the $D^0 \pi^+$ spectrum which they attributed to feed-down of the $D_1^+ \rightarrow D^{*0} \pi^+$ decay where the neutral particles from the D^{*0} decay were not observed. It was not until 1994 that the D_1^+ state was fully reconstructed by CLEO, where the high quality of their electromagnetic calorimeter allowed them to obtain a high statistics sample of $D^{*0} \rightarrow D^0 \pi^0$ decays.

The $D_{s1}(2536)^+$ meson was also first seen by ARGUS in 1989 through its decay to $D^{*+} K_S^0$. Its mass was determined to be $2535.9 \pm 0.6 \pm 2.0 \text{ MeV}/c^2$ and since the mass difference between it and its decay products is only $27 \text{ MeV}/c^2$, the decay amplitude is severely restricted by phase space. Hence, its width is very narrow and only an upper limit of $2.3 \text{ MeV}/c^2$ has been placed on it. This enables one to also reconstruct the D_{s1}^+ in the $D^{*0} K^+$ final state, even when only the D^0 from the D^{*0} decay is observed. The spin-2 partner of the D_{s1}^+ was discovered by CLEO by reconstructing the $D^0 K^+$ final state. Its mass was found to be $38 \text{ MeV}/c^2$ higher than the D_{s1}^+ and its width was measured to be $16_{-4}^{+5} \pm 3 \text{ MeV}/c^2$.

The decays of the D_2^{*0} to $D^+ \pi^-$ and to $D^{*+} \pi^-$ have been observed by both ARGUS and CLEO which allows the predictions expressed in equation (3.22) to be tested. The ratio of these two decay modes is found to be [12]

$$\frac{\text{Br}(D_2^{*0} \rightarrow D^+ \pi^-)}{\text{Br}(D_2^{*0} \rightarrow D^{*+} \pi^-)} = 2.3 \pm 0.6 \quad (3.32)$$

which is seen to be in good agreement with the prediction of 2.5 obtained from equation (3.22). A more refined calculation [35] predicts a smaller ratio of 1.8 which is also compatible with the experimental ratio. The agreement is slightly worse for the corresponding ratio for D_2^{*+} decays, which has been found by CLEO to be

$$\frac{\text{Br}(D_2^{*+} \rightarrow D^0 \pi^+)}{\text{Br}(D_2^{*+} \rightarrow D^{*0} \pi^+)} = 1.9 \pm 1.1 \pm 0.3 \quad (3.33)$$

compared with the naive prediction of 3.7. Finally, the expected ratio of the two D_{s1}^+ branching ratios shown in equation (3.24) is seen to be 1.8 and can be compared with measurements made by CLEO and ARGUS yielding

$$\frac{\text{Br}(D_{s1}^+ \rightarrow D^{*0}K^+)}{\text{Br}(D_{s1}^+ \rightarrow D^{*+}K^0)} = \begin{cases} 1.1 \pm 0.3 & \text{CLEO} \\ 1.4 \pm 0.6 & \text{ARGUS} \end{cases} . \quad (3.34)$$

These measured branching ratios of excited P-wave charm meson do not have sufficient precision to provide rigorous tests of the theory presented here. These tests will be one of the goals of future experiments specifically designed to study charm mesons with high statistics [36].

3.7 Heavy Quark Symmetry Breaking

The properties of heavy mesons were seen to be described, at least qualitatively, using the heavy quark symmetry arguments outlined in the previous sections. However, it is also important to be aware of the extent to which the heavy quark limit is actually achieved in the charm system. The consequences of heavy quark symmetry described in section 3.5 were obtained by ignoring terms of $\mathcal{O}(1/m_Q)$ and higher, but these can be calculated order-by-order using an effective field theory of heavy quarks. In this Heavy Quark Effective Theory (HQET), observables are computed not as a perturbation series in a coupling constant or the number of loops, but instead as a power series in $1/m_Q$. The rules for evaluating Feynman diagrams in this theory can be obtained by expanding the QCD Lagrangian in powers of $1/m_Q$.

Violations of the heavy quark limit have already been seen in the predictions for the D_1 and D_2^* widths since these were modified by kinematic factors of $|\mathbf{p}_\pi|^5$ which arose due to the $D^* - D$ and $D_2^* - D_1$ mass differences. There is also the possibility that the D_1 may have a small S-wave component in its decay to $D^*\pi$ since the ratio of its width to that of the D_2^* is $0.82^{+0.27}_{-0.23}$, compared with the

expectation of 0.31, obtained by weighting the predicted partial widths shown in equations (3.19) and (3.20) by factors of $|\mathbf{p}_\pi|^5$. Thus, it may be appropriate to describe the D_1 decay in terms of S- and D-wave widths, Γ_S and Γ_D which have a relative phase φ between them. The resulting angular distribution of $\cos \alpha$, derived in appendix A, has the form

$$\frac{dN}{d(\cos \alpha)} = \frac{1}{2} \left(R + (1 - R) \cdot \frac{1}{2}(1 + 3 \cos^2 \alpha) + \sqrt{2R(1 - R)} \cos \varphi (1 - 3 \cos^2 \alpha) \right) \quad (3.35)$$

where $R = \Gamma_S/(\Gamma_S + \Gamma_D)$.

CLEO has set limits on the values of R and $\cos \varphi$ by fitting the measured $\cos \alpha$ distribution of their high statistics D_1^0 sample [33]. Their data prefers the values $R \approx 0$ and $\cos \varphi \approx 0$, but the 90% confidence limits on these parameters are not very restrictive. Figure 3.4 shows the 50, 90 and 99% exclusion contours in the R - $\cos \varphi$ plane and clearly shows that the D_1^0 could have a significant S-wave component to its decay, provided $\cos \varphi$ is large and negative. In any case, the angular distribution is completely consistent with the form shown in equation (3.26).

3.8 Summary

This chapter has presented the ideas of heavy quark spin symmetry to predict and explain the observed properties of P-wave charm mesons. The production of charm quarks in Z^0 decays provides an ideal means of studying the fragmentation process and the use of P-wave charm mesons to do so has many advantages. Because these P-wave states are more highly excited than, for instance, the D^* or ground state D mesons, their production rates are mainly due to fragmentation, with minimal additional contributions from strong decays of higher mass states. As will be seen in the subsequent chapters, interesting tests of fragmentation models can be performed by studying the production rates of P-wave states in $Z^0 \rightarrow c\bar{c}$ events.

Figure 3.4: 50% (lightest), 90% and 99% (darkest) exclusion contours of the parameters R and $\cos \varphi$ obtained from CLEO data [33]. The $\star$ denotes the position of the minimum value of χ^2 .

Chapter 4

LEP and the OPAL Detector

Electrons and positrons are accumulated and stored in the LEP synchrotron where they are accelerated to energies of roughly 46 GeV. The two beams are brought into collisions at four interaction regions at which the collaborations ALEPH, DELPHI, L3 and OPAL have detectors. These detectors record the events produced by the e^+e^- interactions. Since the operation of LEP directly influences the design and operation of the OPAL detector, it is worthwhile describing it in some detail.

4.1 The LEP Storage Ring

The LEP synchrotron is located underground in a tunnel approximately 27 km in circumference. Electrons and positrons are kept in counter-rotating circular orbits by dipole magnets and are accelerated from their injection energy by RF cavities which also replenish energy lost to synchrotron radiation.

The electrons and positrons which are injected into LEP are first produced and accelerated to moderate energies by a series of smaller synchrotrons and linear accelerators. Electrons are first accelerated to 200 MeV in a linear accelerator and passed through some converter material in which bremsstrahlung photons produce e^+e^- pairs. These are accelerated to 600 MeV by a second linear accelerator and are stored in an accumulator ring where they are cooled to reduce their energy

dispersion. Accumulated bunches of electrons and positrons are then injected into CERN's proton synchrotron (PS) which accelerates them to 3.5 GeV and transfers them to the super proton synchrotron (SPS) where they are further accelerated to 20 GeV prior to injection into LEP¹.

It is clearly desirable to operate LEP such that high luminosities are delivered to the four experiments. At the same time it is essential that the experiments experience tolerable backgrounds from synchrotron radiation and from particles scattered out of the beam. These competing factors make the stable operation of LEP a complex and challenging task.

Electrons and positrons circulate in the LEP ring in several "bunches", each containing of the order of 10^{11} particles. These have approximately Gaussian charge distributions which are characterized by horizontal and vertical widths for which typical dimensions at the interaction points are $\sigma_x^* \approx 150 \mu\text{m}$ and $\sigma_y^* \approx 12 \mu\text{m}$, respectively. Particles in each bunch do not follow a purely circular orbit around the ring but undergo transverse, or betatron oscillations with amplitudes denoted β_x and β_y in the horizontal and vertical planes. Also, energy lost due to synchrotron radiation and energy gained from the accelerating RF cavities cause longitudinal, or synchrotron oscillations. These parameters are related to the horizontal and vertical emittances, $\varepsilon_x = \sqrt{\sigma_x^* \beta_x}$ and $\varepsilon_y = \sqrt{\sigma_y^* \beta_y}$ which, being a measure of the area of phase space enclosed by the particles in the bunch, is conserved by Liouville's theorem of statistical mechanics. Thus, bunch dimensions can be made smaller, but only at the expense of larger betatron oscillations. This can lead to beam loss if the bunches scrape on accelerator components which limit the aperture of the machine.

The luminosity delivered to an experiment can be expressed in terms of the bunch current, I_b , the number of colliding bunches k_b , the orbit frequency f_0 and

¹Although operated this way in the past, LEP now performs injection at 22 GeV.

the bunch sizes at the interaction point, σ_x^* and σ_y^* , as follows:

$$\mathcal{L} = \frac{I_b^2 k_b}{4\pi\epsilon^2 f_0 \sigma_x^* \sigma_y^*} \quad (4.1)$$

where ϵ is the electron charge. It appears from this expression that the luminosity can be maximized by increasing the bunch current or the number of bunches or by minimizing the bunch sizes, σ_x^* and σ_y^* . However, there are several limiting factors which prevent the luminosity from being increased to arbitrarily high levels.

As bunch sizes are reduced or as the bunch current is raised, their charge density increases. Since bunches of electrons and positrons are brought into collisions at four points around LEP, the electrostatic forces between them can perturb the particles in each bunch from their stable orbits. The strength of these beam-beam interactions is quantified in terms of the beam-beam tune shift parameter[37], ξ , which will approach a limiting value $\bar{\xi}$ given by:

$$\bar{\xi} = \frac{I_b r_e}{2\pi\gamma e f_0} \frac{\beta_y^*}{\sigma_x^* \sigma_y^*}. \quad (4.2)$$

where γ is the Lorentz factor, and $r_e = e^2/mc^2$ is the classical electron radius. When LEP is operated with ξ near this limiting value, the luminosity is given by

$$\mathcal{L}_0 \approx \frac{k_b I_b \gamma \bar{\xi}}{2e r_e \beta_y^*} \quad (4.3)$$

and although this is still increased by raising bunch currents, the limiting value of $\bar{\xi}$ causes the bunch dimensions to increase as well. Excessive bunch currents then result in beam loss and high experimental backgrounds and are therefore to be avoided. Thus, bunch currents at LEP are limited to approximately 0.75 mA and it becomes increasingly difficult to raise the luminosity except by increasing the number of bunches, k_b .

In 1989-91, LEP operated as it was originally designed, colliding four bunches of electrons with four bunches of positrons. In this 4×4 scheme, peak luminosities

of the order of $11 \times 10^{30} \text{ cm}^{-2}\text{s}^{-1}$ were obtained [38]. To increase the luminosity further, LEP doubled the number of circulating bunches, running in an 8×8 configuration from the second half of 1992 to near the end of 1994. This resulted in peak luminosities of $23 \times 10^{30} \text{ cm}^{-2}\text{s}^{-1}$ but the equal spacing between the bunches required the introduction of electrostatic separators between the experiments to prevent parasitic collisions at these points.

For higher energies, LEP is limited not by $\bar{\xi}$, but instead by the RF power available to compensate for synchrotron losses. This limits the total current, $I_{\text{total}} = k_b I_b$, which can be stored in LEP at high energy. The expression for the luminosity shown in equation (4.1) can then be rewritten

$$\mathcal{L} = \frac{I_{\text{total}}^2}{k_b} \cdot \frac{1}{4\pi e^2 f_0 \sigma_x^* \sigma_y^*}. \quad (4.4)$$

Ideally, this would be maximized by placing the maximum possible total current in as few bunches as possible. However, since bunch currents are still limited to less than 1 mA, at least eight bunches are required to meet the luminosity goals of LEP's high-energy program. In 1995, LEP studied this mode of operation by colliding 4 "trains" of 4 closely spaced bunchlets to eliminate parasitic collisions. In this way, LEP achieved luminosities of over $25 \times 10^{30} \text{ cm}^{-2}\text{s}^{-1}$ at $\sqrt{s} = 136 \text{ GeV}$ with bunch currents half those used at Z^0 energies.

4.2 The OPAL Detector

The OPAL detector [39] was designed in the 1980's using technology which was well understood at the time. Since the physics topics studied at LEP are so varied, OPAL was designed to be a general purpose detector which would satisfy the requirements of most topics of interest. OPAL is built around the LEP beam pipe and has the cylindrical geometry characteristic of all detectors for colliding beam experiments. A spherical coordinate system is used, with the z -axis in the direction of the circulating electron beam. The angles θ and ϕ are then defined as the polar

Figure 4.1: The OPAL detector.

angle with respect to the z -axis and the azimuthal angle measured from the x -axis which points towards the centre of the LEP ring.

OPAL's subdetectors, shown in figure 4.1, start with a set of tracking detectors contained within a solenoidal magnetic field of approximately 0.4 Tesla. Beyond these lie time-of-flight counters, electromagnetic and hadronic calorimeters and muon detectors. In the region close to the beam pipe two luminosity monitors are installed which detect small angle e^+e^- scattering and complete the solid angle coverage of OPAL down to $\theta = 25$ mrad. The luminosity monitors will not be described further since all production rates measured in this thesis are normalized with respect to the number of observed hadronic Z^0 decays.

4.2.1 Tracking Detectors

The primary tracking detector in OPAL is a multi-wire drift chamber with a jet-cell geometry. It is approximately 4 m long with inner and outer radii of 25 cm and 185 cm, respectively. The cathode wire planes divide the detector into 24 identical azimuthal sectors, each of which contains 159 sense wires to detect the ionization produced by charged particles. These determine the r - ϕ coordinates of passing particles by the wire position and drift time, while the z -coordinate is determined by the division of the charge read out at each end of the wire. In addition, the total charge deposited on each wire allows up to 159 measurements of the energy loss of a charged particle which is essential for particle identification.

The transverse momentum, p_T , of a charged track is determined from the sagitta of the arc made as it curves in the magnetic field and momentum resolution is then determined by the precision with which the sagitta can be measured. The momentum resolution is also influenced by the magnitude of the magnetic field, the single hit-resolution in the r - ϕ plane and by multiple scattering. The intrinsic r - ϕ hit resolution is $\langle\sigma_{r\phi}\rangle \approx 120 \mu\text{m}$ while systematic uncertainties on hit positions arising from wire sag and field distortions are typically less than $30 \mu\text{m}$. Overall, the jet chamber measures the momentum of charged particles in the transverse plane with an accuracy given by [40]

$$\frac{\Delta p_T}{p_T} = \sqrt{(0.0015 \cdot p_T)^2 + (0.028)^2} \quad (4.5)$$

where the first term is due to the single hit resolutions and the second arises from multiple scattering.

Although accurate measurements of the z -coordinates of hits are important for the determination of a track's polar angle, the single-hit resolution in the z -direction obtained by charge division is only about 5 cm. To improve this measurement, a layer of wire chambers with drift direction parallel to the z -axis were installed outside the jet chamber. These z -chambers consist of 24 modules in ϕ .

each of which has 8 drift cells of 6 wires. These can provide up to six measurements with an accuracy of $\langle\sigma_z\rangle = 300\ \mu\text{m}$ at a radius of 1.85 m which greatly improves the θ resolution of tracks.

Although the z -chambers only cover polar angles in the range $43^\circ \leq \theta \leq 137^\circ$, improvements in the θ measurement are obtained for tracks which exit through the end-plate of the jet chamber. For these tracks, a loosened set of selection criteria for the outermost hits allows an efficient determination of the last wire to detect the passage of a particle. The knowledge of this wire's radius allows for an accurate determination of the polar angle in these cases.

4.2.2 Vertex Detectors

The study of short-lived particles such as B and D mesons is greatly aided by the ability to reconstruct their decay vertices which are displaced with respect to the interaction point. The secondary vertex reconstruction requires precise tracking capabilities in the region immediately surrounding the beam pipe. For this purpose, OPAL has a precision drift chamber inside the inner diameter of the jet chamber. This vertex detector has 36 drift cells, each containing 12 axial sense wires, parallel to the z -axis, and 6 stereo wires, inclined at 4° with respect to the z -axis. These allow measurements of the r - ϕ and z -coordinates of hits to be made with precisions of $\langle\sigma_{r\phi}\rangle \approx 50\ \mu\text{m}$ and $\langle\sigma_z\rangle \approx 700\ \mu\text{m}$, respectively.

In 1991 OPAL's beam pipe was replaced by one of reduced diameter around which a two-layer silicon microstrip detector was installed to provide high precision measurements of track trajectories close to the interaction point. This microvertex detector was constructed from several "ladders" of three $33 \times 60\ \text{mm}$ silicon wafers bonded end-to-end and arranged in two layers around the beam pipe. Originally the detector provided only r - ϕ information but was upgraded to measure the z -coordinates of hits and to extend the solid angle coverage. The parameters of the microvertex detectors at various stages are summarized in table 4.1.

	1991-92	1993-94	1995
Ladders (inner,outer)	11.14	11.14	12.15
Wafers/ladder (inner,outer)	3.3	3.3	3.5
intrinsic $\langle\sigma_{r\phi}\rangle$	10 μm	5 μm	5 μm
intrinsic $\langle\sigma_z\rangle$	-	13 μm	13 μm
$\langle d_0\rangle$	15 μm	17 μm	17 μm
$\langle z_0\rangle$	—	22 μm	22 μm
Single hit efficiency	$\sim 97\%$	$\sim 98\%$	$\sim 98\%$
Total thickness (X_0)	0.006	0.015	0.015

Table 4.1: Parameters of the silicon microvertex detector upgrades.

4.2.3 Electromagnetic Calorimeters

The OPAL electromagnetic calorimeters determine the energies of electrons, positrons and photons by measuring the amount of Čerenkov light produced in blocks of lead-glass by their electromagnetic shower. It provides discrimination between these and hadronic particles which generally do not shower. In addition, it is almost perfectly efficient for detecting minimum ionizing particles which is important for efficient triggering on Z^0 decays.

The calorimeters are assemblies of over 11,000 lead glass blocks fitted with either photomultiplier tubes or phototriodes. The blocks have typical transverse dimensions of approximately 10 cm $\times$ 10 cm and are more than 20 radiation lengths deep. This is sufficient to fully contain electromagnetic showers with energies up to those which will be encountered in collisions with centre-of-mass energies as high as 200 GeV.

The segmentation of the lead glass blocks could allow the positions of showers to be measured to within about 11 mm. However, the lead-glass blocks are located behind approximately 2 radiation lengths of material, mostly from the central detector pressure vessel and the solenoid coil. Since photons are likely to begin showering in this material before reaching the calorimeter itself, the spatial resolution is somewhat degraded. The resolution is improved by the addition of

	Barrel σ_E/E	Endcap σ_E/E
Intrinsic	$0.002 \oplus 0.063/\sqrt{E}$	$0.05/\sqrt{E}$
Measured	$0.015 \oplus 0.16/\sqrt{E}$	$0.018 \oplus 0.22/\sqrt{E}$
Corrected	$0.022 \oplus 0.13/\sqrt{E}$	—

Table 4.2: Intrinsic and measured energy resolutions for the OPAL's barrel and endcap lead-glass calorimeters. The corrected resolution for the barrel calorimeter is obtained using information from the presampler.

two layers of pre-sampling detectors, designed to measure the position of showers immediately in front of the lead-glass blocks. These pre-samplers are constructed from thin multi-wire chambers and allow positions of showers to be determined with accuracies of about 5 mm. The energy resolutions obtained using various components of the electromagnetic calorimeter is shown in table 4.2 [41].

4.2.4 Time of Flight Counters, Hadron Calorimeters and Muon Chambers

The remaining detector elements which participate in the reconstruction of hadronic Z^0 decays are a set of time-of-flight scintillation counters, the hadron calorimeters and muon chambers. These are all important for triggering and can provide useful information for particle identification. The time-of-flight system consists of 160 bars of scintillator attached to photomultipliers via light guides at each end. These are located only in the barrel region just outside the coil and measure the times of arrival of charged particles with an accuracy of ~ 250 ps. Although this is not useful for measuring the velocity of particles with momentum greater than a few GeV, it provides very useful trigger information.

The hadron calorimeter is incorporated into the iron flux return yoke of OPAL's solenoid. It consists of iron plates 10 cm thick between which are located low-cost ionization detectors operated as limited streamer tubes. Although the hadronic energy resolution of the calorimeter is poor ($120\%/\sqrt{E}$), its strip

readout granularity allows tracking of particles out to the muon drift chambers. The muon chambers form the outermost subdetector and hadrons encounter over 7 pion interaction lengths of material before reaching them. Thus, only muons with energies larger than about 3 GeV can penetrate and be detected. The muon detector itself consists of four layers of large area drift chambers with sense wires located between field shaping printed circuit boards. The efficiency for detecting isolated muons in the solid angle coverage is close to 100% while there is less than a 1% chance that isolated 5 GeV pions will produce signals in the muon chambers.

4.2.5 Trigger and Data Acquisition

When LEP circulates four bunches of electrons and positrons, beam collisions occur at a rate of 45 kHz. Since the time required to read, process and record all subdetector data is approximately 20 ms, this rate must be reduced to the order of 5 Hz to achieve an experimental dead time of a few percent. The trigger system[42] achieves this by correlating independent signals from geometrically segmented regions of the detector, and decides whether to record the detector information during the inter-bunch crossing time of $22 \mu\text{s}$.

For triggering on spatial coincidences of activity in different subdetectors, OPAL's geometry is segmented into overlapping regions of θ and ϕ . The granularity of the segmentation in ϕ varies from 10 to 24 bins, while all subdetectors have θ subdivided into 6 bins. Subdetectors provide total energy or particle multiplicities in each θ - ϕ bin which are made available to the central trigger logic. The trigger processor is configurable in software to initiate readout of the detector if the signal from a single subdetector exceeds some threshold, or if a spatial coincidence between different subdetectors is observed.

In 1993 and 1994 LEP ran in an 8×8 bunch configuration, doubling the bunch crossing rate. This mode of operation was not anticipated in the original trigger design which could not produce a trigger decision in this reduced time between bunch crossings. Therefore, the 90 kHz rate was reduced to one of a few kHz by

the implementation of a fast pre-trigger [43] which could generate a positive or negative response in the $11 \mu\text{s}$ bunch crossing interval. The design of the pre-trigger is similar to that of the trigger but uses only spatially segmented regions in ϕ to detect coincidences. LEP again operated in the 4×4 bunch configuration in 1995, and although no longer needed, the pretrigger was left operational since its presence does not degrade the detector performance.

4.2.6 Event Readout and Reconstruction

The front-end readout electronics of each subdetector is controlled by one or more local systems crates (LSC's) consisting of a local trigger unit and at least one single-board microcomputer running the OS-9 [44] real-time operating system. These are connected to the trigger processor and to the event builder by dedicated busses for the transfer of trigger information and detector readout data. As described in the previous section, trigger signals are provided by summed energies and particle multiplicities. Also, for the tracking chambers, trigger signals are formed by hardware which performs fast track finding[45, 46].

Upon receiving the trigger information, the trigger processor broadcasts either a reset signal or a trigger interrupt on the trigger bus. A positive trigger initiates an OS-9 event which schedules the readout of the front-end electronics. This data is then read into a buffer during which time the subdetector asserts its busy signal on the trigger bus preventing the generation of further triggers. Once the data is stored in the buffer, other processes running on the LSC format it and transfer it to the event-builder over the inter-crate data transfer bus.

The event-builder [47] is responsible for collecting the detector event information from the LSCs and merging them into complete events. This is essential since the differing time required by the various subdetectors to format the data means that sub-events may be transferred out of sequence. The event-builder can also inhibit triggers should its buffer memory become filled with partially transferred events. The fully merged events are then transferred over a dedicated bus to a

multiprocessor computer known as the filter for fast analysis and partial reconstruction.

The filter [48] is a dedicated multiprocessor computer which classifies, compresses, monitors and records the fully merged events. Originally the filter was implemented using a 4-processor Apollo DN10000 computer but was replaced in 1993 by a system of Hewlett-Packard 700-series microprocessors which perform the same functions. It performs online track finding using the data from the jet chamber allowing events to be classified according to the type of Z^0 decay or interaction recorded. Background events from cosmic rays or beam-gas interactions can then be rejected if desired. In addition, the filter fills a large number of histograms with data from the individual detectors which allows online monitoring of the data quality and detector performance.

Events which are not rejected by the filter are stored on a local hard disk in 20 Mbyte files which are then transferred over Ethernet to another microcomputer located on the surface where they are archived both on magneto-optical disks and magnetic tape. When online calibration constants become available from the LSC's, the files and the calibrations are distributed among a farm of HP workstations which fully reconstruct the events. This farm is also used to re-process archived data using new reconstruction code and calibration constants. The fully reconstructed files are then transferred over Ethernet to the workstation cluster on which most OPAL analysis is performed.

These data files are concatenated and archived on tape while a subset of the data contained in these files is stored permanently on a large pool of disks. These files are known as Data Summary Tapes² (DST's), and contain very little unprocessed data from individual subdetectors. Instead, all available information is combined to provide the best estimates of, for example, the momentum of charged tracks, energies of clusters in the electromagnetic calorimeters and positions of

²This term has a historical origin but is still used even though today the data is usually stored in files on disk.

reconstructed track segments in the muon chambers. This subset of the recorded data is sufficient for almost all types of analysis.

4.3 Monte Carlo Simulation

An essential tool for the analysis of OPAL data is the detector simulation program known as GOPAL [49]. This program provides a detailed description of the detector geometry to the GEANT [50] package which uses Monte Carlo methods to simulate the passage of particles through the detector material. This includes effects such as multiple scattering of charged particles, electromagnetic and hadronic interactions and the decays of long-lived weakly decaying particles such as K_S^0 's and Λ 's. The energy deposited in the "sensitive" regions of the detector, either by primary particles or by the products of their subsequent interactions can then be calculated. These results are then used to determine the expected response of the detector elements which is recorded in data structures equivalent to those read out by the front-end electronics. These data structures are then processed with the same software used to reconstruct real data, except that the information used to associate reconstructed quantities with the underlying particles in the event is retained.

There are many programs available for simulating Standard Model physics processes, such as Z^0 decay modes or small angle Bhabha scattering [51]. Generators also exist for simulating more exotic processes such as the postulated production and decay of the Higgs boson or supersymmetric particles. These programs all interface to the GOPAL program in the same way, by preparing a file containing the relevant event information. This includes the 4-momenta of each final state particle as well as codes identifying each particle type and data structures containing the history of particle decays in the event. Although some generators simulate the positions of particle decays, this information is not used and is instead re-generated by GOPAL.

The simulation of hadronic Z^0 decays used for the analyses presented here was performed using the Jetset 7.4 program [52]. This program simulates a parton shower which subsequently fragments into hadrons using the Lund string model [53]. Particle decays were simulated using tables of branching ratios obtained from the 1994 Review Of Particle Properties [54], where known, and from phenomenological arguments or additional string fragmentation where unknown [55]. In addition, numerous parameters in the program were tuned so that average event properties like particle momentum spectra and event shapes matched those observed in the OPAL data [56]. It is believed that this simulation was adequate for estimating backgrounds and efficiencies for the processes studied in subsequent chapters. Nevertheless, a number of cross checks were performed to either demonstrate that this is the case, or to estimate the extent to which inadequacies of the simulation influenced the measured results. These are described in detail in sections 6.5 and 7.2.4.

4.4 The OPAL Data Sample

OPAL and the other LEP experiments began taking data in 1989 and although LEP was not fully commissioned, recorded enough data to measure the total hadronic cross section to an unprecedented level of precision. However, it was not until 1990 that LEP was able to deliver the luminosity required for most analyses. Although LEP has been able to increase the total luminosity delivered to the experiments in most years, precision measurements of the Z^0 mass and width require operating at centre-of-mass energies between 88 and 95 GeV. As is seen from figure 3.1, this can result in a reduction in the rate for hadronic events by up to a factor of 3.

Figure 4.2 shows the integrated luminosity and the number of hadronic events recorded throughout the years 1990-96 and table 4.3 contains statistics for OPAL's data sets. In this analysis, only data recorded at $\sqrt{s} \sim M_{Z^0}$ from 1990 to 1995

Figure 4.2: Integrated luminosity delivered by LEP and the number of hadronic events recorded by OPAL shown as a function of time throughout the years 1990-96.

was used. Of this data, 25% was recorded with the silicon microvertex detector with only r - ϕ readout, while 60% was recorded with the upgraded silicon detector which had precise tracking measurements in both the r - ϕ and z coordinates.

4.4.1 Properties of Multihadronic Events

The environment in which D^{*-0} mesons are produced can have a significant impact on which measurements can be made and how they are carried out. To understand the benefits as well as the difficulties involved with studying D^{*-0} production at LEP, it is important to be aware of the general properties of hadronic events. This section describes many of the global event properties characteristic of hadronic events while the next section describes the way in which these are used to select a relatively pure sample of hadronic Z^0 decays.

Since both beams in LEP have the same energy of $\sim M_Z/2$, the e^+e^- collision

year	$\sqrt{s}$ (GeV)	$\int \mathcal{L}$ (pb ⁻¹)	hadronic events
1989	88.3–95.0	1.2	25801
1990	88.2–94.2	6.6	139893
1991	88.5–93.7	12.9	314378
1992	91.3	23.9	733059
1993	89.4–93.0	33.3	723500
1994	91.1	58.0	1710554
1995	89.3–92.9	33.6	729556
1995	130.0–140.0	5.8	5325
1996	161.0–172.0	22.0	41202

Table 4.3: A summary of OPAL’s data sets. Neither the data recorded in 1989 nor the high energy data recorded from November 1995 onward is used in this analysis.

will produce a Z^0 with no net momentum. Thus, when it decays to a $q\bar{q}$ final state, the quarks will be emitted back-to-back in OPAL. These quarks produce many charged and neutral particles in the fragmentation process but these are generally collimated into narrow “jets” of particles which are roughly aligned with the momentum of the initial parton. Hard gluon bremsstrahlung can occur early in the parton shower which can result in additional jets of hadrons. These features can be seen in figures 4.3 and 4.4 which show 2- and 3-jet events typical of hadronic Z^0 decays.

The number of charged particles in an event increases with the centre-of-mass energy and at $\sqrt{s} = 91$ GeV the average charged particle multiplicity is $\bar{n}_{\text{ch}} \approx 21$. This is shown in figure 4.5 where the mean charged particle multiplicity is plotted for e^+e^- collisions at various energies.

Simply identifying the production of hadrons in the final state cannot be used to select hadronic Z^0 decays due to the presence of additional production mechanisms. These include $Z^0 \rightarrow \tau^+\tau^-$, followed by the hadronic decays of one or both τ leptons, and the multiperipheral production process³ indicated in figure 4.6.

³These are also frequently referred to “two-photon” events, not to be confused with the production of two real photons via the electromagnetic process $e^+e^- \rightarrow \gamma\gamma$.

```

Run: event 4055: 42316 Date 930519 Time 75714 Ctrk(N= 32 Sump= 73.6) Ecal(N= 43 SumE= 30.4) Hcal(N=18 SumE= 9.0)
Ebeam 45.659 Evis 93.3 Emiss -2.0 Vtx ( -0.03, 0.02, 0.87) Muon(N= 3) Sec Vtx(N= 2) Fdet(N= 0 SumE= 0.0)
Bz=4.350 Thrust=0.9623 Aplan=0.0067 Oblat=0.0528 Spher=0.0358

```

Figure 4.3: A typical two-jet event in OPAL.

Run: event 4055: 37387 Date 930519 Time 72913 Ctrk(N= 40 Sump= 59.4) Ecal(N= 59 SumE= 45.7) Hcal(N=21 SumE= 15.1)
 Ebeam 45.659 Evis 88.6 Emiss 2.7 Vtx (-0.06, 0.11, 0.62) Muon(N= 1) Sec Vtx(N= 3) Fdet(N= 0 SumE= 0.0)
 Bz=4.350 Thrust=0.8815 Aplan=0.0564 Oblat=0.1056 Spher=0.1904

Figure 4.4: A typical three-jet event.

Figure 4.5: Mean charged track multiplicity observed in e^+e^- collisions at various centre-of-mass energies. The curve shows the expected behaviour: $\log\langle n_{ch} \rangle = a/\sqrt{\alpha_s(Q^2)} + b \log \alpha_s(Q^2) + \mathcal{O}(1)$, where a and b can be calculated using QCD [57].

Figure 4.6: The Feynman diagram for the multiperipheral production of a hadronic final state.

However, since τ decays necessarily produce non-interacting neutrinos and since the electrons in multiperipheral events are frequently scattered at only small angles, escaping undetected down the beam pipe, there is usually considerably less energy available for either of these processes to produce hadrons. There is also considerable momentum imbalance in these final states in contrast to the back-to-back jet topology characteristic of hadronic Z^0 decays. These features motivate the selection criteria used to reject $\tau^+\tau^-$ and multiperipheral events from the inclusive hadronic event sample.

4.4.2 Selection of Hadronic Z^0 Decays

Hadronic Z^0 decays are selected based on track and energy cluster multiplicities and on the total energy and energy imbalance observed in an event. A typical Z^0 decay yields about 30 charged and neutral hadrons which produce many tracks in the jet chamber and leave significant deposits of energy in the electromagnetic calorimeters. Since Z^0 decay occurs in the stationary lab frame, the energies observed in opposite hemispheres of the detector should be equal. Thus, OPAL uses

the following criteria to select a sample of hadronic Z^0 decays:

$$\begin{aligned} N_{\text{chg}} &\geq 5 \\ N_{\text{clus}} &\geq 7 \\ R_{\text{vis}} &> 0.1 \\ |R_{\text{bal}}| &< 0.65, \end{aligned}$$

where N_{chg} and N_{clus} are the numbers of charged tracks and electromagnetic energy clusters reconstructed. R_{vis} and R_{bal} are defined using the raw energy deposits in the calorimeter⁴ according to

$$R_{\text{vis}} = \frac{\sum E_{\text{raw}}}{2E_{\text{beam}}} \quad (4.6)$$

$$R_{\text{bal}} = \frac{\sum E_{\text{raw}} \cos \theta}{\sum E_{\text{raw}}} \quad (4.7)$$

where θ is the polar angle at which the clusters were recorded. The distributions of these quantities are presented in figure 4.7 and clearly show the effective rejection of $\tau^+\tau^-$ and multiperipheral events. The efficiency for selecting hadronic Z^0 decays using these criteria has been determined to be 98.4%[58] while remaining background from $\tau^+\tau^-$ and multiperipheral events is about 0.5%.

4.4.3 Jet Finding at LEP

The striking jet-like characteristics of the events shown in figures 4.3 and 4.4 are the result of hadrons forming along the initial directions of the primary quarks in the Z^0 decay. Hard gluon bremsstrahlung can produce three, four or in general n -jet events but these higher jet multiplicities are suppressed by additional powers of α_s per additional jet. Many analyses, including the ones described in this thesis, can benefit by only studying the set of particles assigned to the same jet since the

⁴In most analyses, these raw energies are corrected for material in front of the lead-glass blocks but this procedure is unnecessary when just identifying hadronic events.

Figure 4.7: Distributions of variables used to select hadronic Z^0 decays. The points with error bars are the measured distributions obtained from analysis of OPAL data while the open histograms show the corresponding distributions obtained from hadronic Monte Carlo events. The hatched distributions indicate the contributions from tau decays and multiperipheral events. Arrows indicate the positions of cuts used to select hadronic events. Each distribution is obtained with all other cuts applied.

tracks in different jets are usually completely uncorrelated. Thus, it is useful to describe the algorithms by which the set of all tracks in an event are partitioned into subsets corresponding to what we would visually describe as jets.

The JADE Algorithm

The input to most commonly used jet finding algorithms are the 4-vectors of the tracks and energy deposits observed in the detector. Although only the momenta of charged tracks and the energies of neutral particles are measured, the 4-vectors can be obtained by assuming tracks to be pions and electromagnetic clusters to be photons. One class of jet finding algorithms was developed by the JADE collaboration [59] and is commonly used either in its original form or in a number of slightly modified versions. This algorithm reduces the number of 4-vectors in an event by one by successively combining the pair which has the minimum scaled invariant mass.

$$y_{ij} = m_{ij}^2 / E_{\text{c.m.}}^2 = (p_i + p_j)^2 / E_{\text{c.m.}}^2, \quad (4.8)$$

replacing them by a new 4-vector $p_k = p_i + p_j$. The criteria used to compute the scaled invariant mass is not unique and many schemes exist [60, 61]. The measure shown in equation (4.8) is known as the “E0” scheme but there are several others which are in common usage, including the original JADE scheme which used $y_{ij} = 2E_i E_j (1 - \cos \theta_{ij}) / E_{\text{c.m.}}^2$. In general, these procedures all combine pairs of tracks which are kinematically “close” to each other and are iterated until all remaining pairs of 4-vectors have scaled invariant masses in excess of a quantity y_{cut} . This is known as the jet resolution parameter and the number of jets assigned to an event clearly depends on the value of y_{cut} used. This dependence is shown in figure 4.8 for the E0 scheme and has been used to measure the value of α_s at the Z^0 mass [60].

Figure 4.8: n -jet event fractions shown as functions of the jet resolution parameter, y_{cut} for the "E0" recombination scheme.

The Cone Algorithm

The JADE algorithm and its variants group tracks and energy clusters into jets based mainly on their kinematic properties. An alternative, developed for use in hadron collider experiments [62], is based on their geometric properties. In the “cone” scheme, jets are required to have a minimum energy ε and all tracks within a cone of half-angle R around the sum of their momentum vectors are assigned to the jet. This differs from the JADE algorithm in that not all 4-vectors are forced to be members of a jet and it is possible for some tracks to be left unassigned. The dependence of the n -jet event fractions on the parameters ε and R is also shown in figure 4.9. Monte Carlo studies show that better estimates of the primary parton momentum can be obtained using the cone algorithm[63] which motivate its use in the analysis described in this thesis.

4.4.4 Performance of the OPAL Detector

In this section, a discussion of the quality of OPAL’s data is presented. This evaluation focuses on the ability to measure quantities relevant for the analysis of P-wave charm mesons. Although of excellent quality in many other respects, properties of the data not relevant to this discussion have been omitted. The description concentrates mainly on the performance of the tracking detectors which, as will be seen in the subsequent chapters, are the most critical components for performing the analysis. However, before quantitative measures of detector performance can be presented, it is important to specify what quantities are actually measured.

OPAL Track Parameters

As described in sections 4.2.1 and 4.2.2, the paths of charged particles are recorded by the tracking and vertex detectors. The central jet chamber measures the momenta of tracks in the transverse plane, the z -chambers provide accurate measure-

Figure 4.9: n -jet event fractions shown as functions of the jet resolution parameter. R and ϵ for the “cone” scheme.

Figure 4.10: The geometry associated with the definition of OPAL's track parameters. The x - y projection, perpendicular to the beam axis, is shown in (a), while (b) shows the projection in the s - z plane. The point of closest approach to the origin is labeled "p.c.a.".

ments of their polar angles as they leave the jet chamber and the vertex detectors provide precise measurements along their trajectory close to the interaction point. A charged particle in a magnetic field follows a helical trajectory which requires five parameters to fully specify. Naturally, it is important that the track parameters be well defined for all reconstructed tracks but it is also very important that the derivatives of the equations defining the trajectories with respect to these parameters be non-singular. These requirements motivate OPAL's choice of track parameters, which is a convention also used by many other colliding beam experiments.

Figure 4.10 shows the way in which several quantities related to OPAL's set of track parameters are defined. Track parameters are specified at the point (x_0, y_0, z_0) on the helix which is closest to the OPAL origin but these can be

transformed to other reference points if required. The z_0 coordinate of the point of closest approach is itself used to define the track trajectory. However, rather than use x_0 and y_0 as fundamental parameters, it is more useful to use the track's azimuthal angle ϕ_0 and the transverse impact parameter $d_0 \equiv y_0 \cos \phi_0 - x_0 \sin \phi_0$ at this point. It is also suitable to define the transverse momentum in terms of the radius of curvature, ρ , using a parameter κ whose magnitude is given by $|\kappa| = 1/2\rho$. The sign of κ is positive if the azimuthal angle of the track's momentum vector increases as it moves along its trajectory and negative otherwise. Finally, rather than use the polar angle θ , it is more convenient to use $\tan \lambda$, where $\lambda = \pi/2 - \theta$ is known as the "dip angle".

The transverse momentum and momentum components at the point of closest approach are then defined using these parameters as follows:

$$p_T[\text{GeV}/c] = (1.5 \times 10^{-4}) \left| \frac{B_z[\text{kG}]}{\kappa[\text{cm}^{-1}]} \right| \quad (4.9)$$

$$p_x = p_T \cos \phi_0 \quad (4.10)$$

$$p_y = p_T \sin \phi_0 \quad (4.11)$$

$$p_z = p_T \tan \lambda. \quad (4.12)$$

Track trajectories can be approximated by a straight line at small distances from the origin or for small track curvature. This is the case for tracks produced at the primary vertex or in weak decays of heavy hadrons. In such cases, the equations for the trajectory can be specified in terms of the path length, s , measured from the point of closest approach:

$$x(s) = x_0 + s \cos \phi_0 = -d_0 \sin \phi_0 + s \cos \phi_0 \quad (4.13)$$

$$y(s) = y_0 + s \sin \phi_0 = d_0 \cos \phi_0 + s \sin \phi_0 \quad (4.14)$$

$$z(s) = z_0 + s \tan \lambda \quad (4.15)$$

The parametric equations for trajectories which cannot be approximated by

straight lines are more complicated. In these cases, the parameter s is defined as the arc-length along the trajectory, measured in the x - y plane from the point of closest approach. The azimuthal angle ϕ then varies as a function of s , having the value ϕ_0 when $s = 0$.

Track Parameter Resolution

The momentum resolution of charged tracks can be determined from a sample where the actual momentum is known. One such sample is non-radiative $Z^0 \rightarrow \mu^+\mu^-$ events for which each muon has a momentum of 45.6 GeV/ c , fixed by the beam energy. The measured momentum can then be compared with this expectation and it is through this type of analysis that the expression for the momentum resolution shown in equation (4.5) was obtained.

The impact parameter resolution is also of great importance since this limits the precision with which decay vertices from short-lived particles can be reconstructed. The d_0 resolution can also be studied using $\mu^+\mu^-$ events. Since the $\mu^\pm$ are known to be produced back-to-back, the sum of their transverse impact parameters should vanish. The distribution of this quantity is shown in figure 4.11 for tracks with and without hits in the silicon microvertex detector.

These estimates of the d_0 resolution apply only to this sample of tracks. In general, the resolution depends on many factors such as which subdetectors were used to reconstruct the track and the precision with which they were aligned and calibrated. Also, the effects of multiple scattering in the beam-pipe and in other detector material degrade the impact parameter resolution at low momentum.

The resolution of $\tan \lambda$, which determines a track's polar angle, is dominated by uncertainties in the longitudinal coordinates of its production point and the point at which it leaves the tracking chambers. Although the latter can be determined precisely using the z -chambers, the jet-chamber and vertex detector are optimized for precise measurements in the r - ϕ plane. However, an improved measurement of the z -coordinate of the primary vertex can be obtained from an average of the z_0

Figure 4.11: Impact parameter resolution determined from tracks in $Z^0 \rightarrow \mu^+ \mu^-$ events without (a) and with (b) hits in the silicon microvertex detector.

parameters of all tracks in an event. Its position is frequently used to constrain the production points of tracks to obtain a more precise determination of their polar angles. However, not all tracks can benefit from this procedure and it is sometimes necessary to constrain their production point by other means. This is the case for $\pi^+\pi^-$ pairs produced in the decays of K_S^0 mesons which can occur far from the interaction point. This case is discussed further in section 5.3.

Beam Spot Position

The precise measurement of track parameters is essential for reconstructing the displaced decay vertices from hadrons containing charm or bottom quarks. It is also important to have a good estimate of the position in OPAL at which the e^+e^- annihilation took place, since measurements of the separation between this primary vertex and a secondary vertex are limited by the precision with which both are measured. The average location and size of the LEP beam-spot can provide important information when fitting for the position of the primary vertex in an event.

As described in section 4.1, the beams at the interaction point have approximately Gaussian shapes with widths σ_x^* and σ_y^* in the horizontal and vertical planes. Although the operation of LEP requires a certain knowledge of these parameters, they can be measured directly by OPAL by correlating the observed locations of e^+e^- annihilations over short time periods during which the operation of LEP remains stable. In this way, the time-dependent average beam-spot positions and sizes are obtained. Displacements and sizes of the beam-spots averaged over each year are illustrated in figure 4.12 and all have typical dimensions of about $150\ \mu\text{m}$ by $12\ \mu\text{m}$ in the horizontal and vertical planes. Since typical decay lengths of heavy flavoured hadrons can be of the order of half a millimeter, the benefit of this knowledge prior to additional vertex finding is clear. Due to synchrotron oscillations along the beam axis, bunch lengths are much larger than these dimensions, giving average beam-spots which are about 1 cm in length.

Figure 4.12: Beam spot positions and sizes for the years 1990-95. The scales are in millimeters and are drawn at the OPAL origin.

Vertex Reconstruction

Charm and bottom hadrons decay weakly and hence, they have lifetimes long enough that the point at which they decay can often be separated by several millimeters from their production point. Thus, identifying the position of the e^+e^- annihilation and then identifying either displaced secondary vertices or tracks which do not pass through the primary vertex are powerful techniques often used to identify $Z^0 \rightarrow c\bar{c}$ and $Z^0 \rightarrow b\bar{b}$ events.

Although the average position of the LEP beam-spot provides a constraint on the e^+e^- interaction point, a more accurate determination can be obtained using the measured track parameters in a hadronic event. The position of the primary vertex is usually determined by finding the common point of intersection for these tracks in the event, a procedure which is somewhat technical in nature and is described in more detail in section 6.2.1 and 6.2.2. Since the strong decays

of hadrons created in the fragmentation process produce many tracks originating from the primary vertex, determining its position can generally be done reliably. The uncertainty in its position is then dominated by the precision with which the impact parameters of individual tracks are determined. Uncertainties in the primary vertex position typically range between $30\ \mu\text{m}$ and $150\ \mu\text{m}$ depending on the orientation of the jet axis. Since the tracking is less precise in the z -direction, the uncertainty in the z -coordinate of the primary vertex is typically $500\ \mu\text{m}$.

The positions of secondary vertices are usually reconstructed with somewhat less precision than the primary vertex since they are generally of lower track multiplicity and cannot benefit from the average beamspot constraint. Decay lengths can be calculated using simply the separation between two vertices, but improved decay length measurements can often be obtained by constraining the displacement to lie on along the jet axis or momentum vector of the decay vertex. Typical decay length resolutions for two-track vertices such as those obtained from $D^0 \rightarrow K^- \pi^+$ decays are of the order of $300\ \mu\text{m}$.

It is also common to determine the significance of the displacement between two vertices by dividing the measured decay length by its uncertainty. Vertices, resulting from misassigned tracks from the primary vertex would then have decay length significance distributions characterised by a Gaussian with unit width. Hence, requiring positive decay length significances greater than 2 or 3 can significantly enhance the purity of an event sample with secondary vertices. More complicated topologies, including the one shown in figure 4.13 can be reconstructed using specialized algorithms such as the one described in section 6.2.3.

Figure 4.13: An example of the secondary vertex topology reconstructed in a Monte Carlo event for the decay topology $Z^0 \rightarrow b\bar{b} \rightarrow D^{*+} X$ followed by $D^{*+} \rightarrow D^0 \pi^+$ and $D^0 \rightarrow K^- \pi^+$. The axes indicate the length scale in centimeters and the symbols show the true locations of the primary and secondary vertices.

4.5 Comparison with Other Facilities

The LEP accelerator and the OPAL detector complement many other high energy physics facilities around the world and it is useful to compare their merits and disadvantages, both in terms of their primary physics impact and for the study of P-wave charm mesons.

4.5.1 Lower Energy e^+e^- Colliders

Electron-positron colliders which operate near $\sqrt{s} \sim 10$ GeV are ideal for producing large samples of B decays because of the enhancement in the cross section due to the $\Upsilon(4S)$ resonance which decays to $B\bar{B}$ mesons 100% of the time. At these energies, pairs of u , d , s and c quarks are also produced via e^+e^- annihilation through an intermediate virtual photon. The ARGUS detector, which was located at the DORIS-II storage ring and the CLEO/CLEO-II detectors at Cornell, have dominated this area of research and both have accumulated large data samples both with energies on and off the $\Upsilon(4S)$ resonance.

B mesons are produced almost at rest and do not produce well-separated jets in the detector. Thus, analysis of B events is complicated by the overlapping decay topologies which can yield larger combinatorial backgrounds than are present at LEP in some analyses. However, the fact that the B-mesons decay at rest can be used to constrain the decay kinematics, making it easier to study some exclusive decay modes. The fact that particles are not collimated into narrow jets at these energies also makes the identification of π^0 's easier, which is essential for fully reconstructing final states which include D^{*0} mesons.

4.5.2 Fixed Target Experiments

Fixed target experiments such as the Fermilab E691 and E687 experiments are excellent laboratories for studying the properties of charm hadrons because of their very large data samples and the large boost given to the primary hadrons

in the lab. Pure samples of D^0 , D^+ and D_s^+ mesons were obtained by requiring a significant displacement between their points of production and decay. The D_1^0 , D_2^{*0} , D_2^{*+} and D_{s1}^+ states were all reconstructed with reasonably low backgrounds and provided precise measurements of their masses and widths. At the beam energies used in these experiments, the cross section for $b\bar{b}$ production is extremely low and consequently, B decays to P-wave mesons cannot be studied at present.

4.5.3 Hadron Colliders

The D0 and CDF detectors at the Fermilab $p\bar{p}$ collider currently dominate the field of hadron collider experiments. The total cross section for $p\bar{p}$ interactions is very high but since the cross section for the production of $b\bar{b}$ and $c\bar{c}$ pairs is small, very tight trigger requirements must be imposed. For the study of heavy flavour production this usually requires the presence of at least one high energy lepton which necessarily restricts the types of decays which can be studied. Although CDF has provided high quality measurements of b-hadron lifetimes and has reconstructed some of their exclusive decays, so far there have been no measurements involving P-wave charm mesons. This could change with proposed upgrades to the detector which would allow triggers to be initiated by tracks with large impact parameters. This would allow an almost inclusive sample of b-hadron decays to be recorded and will allow challenging measurements such as searches for CP violation to be performed.

Chapter 5

Reconstruction of Intermediate States

As described in section 3.5, P-wave charm mesons decay to lower mass hadrons and these, in turn, decay to the stable particles detected in OPAL. In order to observe the production of P-wave charm mesons it is necessary to reconstruct their chain of decays, but there are only a few modes for which this can be done effectively. The material in front of OPAL's electromagnetic calorimeter and its intrinsic granularity makes the full reconstruction of $D^{*0} \rightarrow D^0 \pi^0$ or $D^{*0} \rightarrow D^0 \gamma$ difficult in this class of events and the large combinatorial backgrounds associated with the direct 2-body D^0 , or 3-body D^+ decays further limit which P-wave charm meson decays can be observed in OPAL. However, D^{*+} and K_S^0 meson decays can be reconstructed cleanly in OPAL which allows three of the 6 states listed in table 3.2 to be observed. Furthermore, because of the kinematics of the $D_{s1}^+ \rightarrow D^{*0} K^+$ decay, this too can be reconstructed even when the π^0 or γ produced in the subsequent D^{*0} decay remains undetected.

5.1 D^{*+} Reconstruction

The $D^{*+}(2010)^+$ meson will decay to $D^0\pi^+$ and to $D^+\pi^0$ with relative branching fractions determined mainly by isospin¹:

$$\text{Br}(D^{*+} \rightarrow D^0\pi^+) = 68.1 \pm 1.3\% \quad (5.1)$$

$$\text{Br}(D^{*+} \rightarrow D^+\pi^0) = 30.8 \pm 0.8\% \quad (5.2)$$

with the small remaining fraction in the $D^+\gamma$ mode. The D^0 meson has many observed decay modes but most have neutral particles in the final state. Of the exclusive decays to charged particles, the most frequently used are

$$D^0 \rightarrow K^-\pi^+ \quad \text{Br} = 4.01 \pm 0.14\% \quad (5.3)$$

$$D^0 \rightarrow K^-\pi^+\pi^+\pi^- \quad \text{Br} = 8.1 \pm 0.5\% \quad (5.4)$$

$$D^0 \rightarrow K_S^0\pi^+\pi^-; K_S^0 \rightarrow \pi^+\pi^- \quad \text{Br} = 1.8 \pm 0.2\% \quad (5.5)$$

however the high combinatorial backgrounds associated with (5.4) and the smaller branching fraction of (5.5) make the first mode the preferred one for these studies.

5.1.1 Kinematics

The reconstruction of $D^{*+} \rightarrow D^0\pi^+, D^0 \rightarrow K^-\pi^+$ proceeds by first forming the invariant mass of all neutral two-particle combinations under the hypothesis that one is a kaon and the other a pion. The distribution of the $K^-\pi^+$ invariant mass, obtained after the complete D^{*+} selection criteria are applied, is shown in figure 5.1. This figure also shows the invariant mass distribution of $K^+\pi^+$ pairs which is expected to have the same shape as the component of the $K^-\pi^+$ distribution resulting from purely random track combinations. The excess above this combinatorial background near $M(K^-\pi^+) \sim 1.6 \text{ GeV}/c^2$ is due to $D^0 \rightarrow K^-\pi^+\pi^0$ decays

¹As first described in section 3.5.1, charge conjugation of all decays is implied.

where the π^0 is not reconstructed while the additional excess at lower masses arises from other partially reconstructed D^0 decays. The mass resolution in the region of the D^0 peak is about $30 \text{ MeV}/c^2$ and the subsequent selection of D^0 candidates is performed by requiring that the $K^-\pi^+$ invariant mass lie within about 2σ of the accepted D^0 mass:

$$1.79 < M(K^-\pi^+) < 1.94 \text{ GeV}/c^2. \quad (5.6)$$

After a $K^-\pi^+$ combination is found with invariant mass in this range, D^{*+} candidates are formed by adding additional pions with the appropriate charge. When reconstructing sequential decays such as this, it is common to consider the difference in mass between the $K^-\pi^+\pi^+$ combination and the previous $K^-\pi^+$ pair. By doing so, many systematic and statistical errors are cancelled which frequently results in a mass difference resolution superior to that of just the three-particle mass. This mass difference distribution, after the full D^0 selection has been applied, is shown in figure 5.2. The shape of the background component is well represented by random $K\pi\pi$ combinations obtained using both “wrong-sign” $K^+\pi^-\pi^+$ combinations and “rotated” candidates obtained with the additional π^+ selected from the hemisphere opposite the D^0 with its momentum reflected through the origin[64].

In the case of D^{*+} decay, there is only a small mass difference between the initial and final states, and the π^+ is emitted with a momentum of only $39 \text{ MeV}/c$ in the D^{*+} rest frame. Thus, the phase space available is severely restricted, both in the D^{*+} decay amplitude and in the combinatorial background. This results in the sharp peak superimposed on a rising background in the mass difference distribution. These features allow a relatively pure D^{*+} sample to be reconstructed making it an attractive intermediate channel through which to reconstruct D^{*-0} decays. Events in the peak of the mass difference distribution are selected by requiring

$$0.142 < M(D^{*+}\pi^+) - M(D^{*+}) < 0.149 \text{ GeV}/c^2. \quad (5.7)$$

Figure 5.1: Invariant mass distribution of $K^-\pi^+$ candidates (data points with error bars) compared with the wrong-sign $K^+\pi^+$ combinations (histogram) normalized above $2 \text{ GeV}/c^2$. These candidates are obtained after the full D^{*+} selection has been performed. The broad peak near $M(K^-\pi^+) \sim 1.6 \text{ GeV}/c^2$ is due to $D^0 \rightarrow K^-\pi^+\pi^0$ decays where the π^0 is not reconstructed. The excess at lower masses is due to other partially reconstructed D^0 decays.

Figure 5.2: Mass difference, $M(K^- \pi^+ \pi^+) - M(K^- \pi^+)$ distribution for $x_E(D^{*+}) > 0.2$. The signal is shown as points with error bars, while the histogram shows an estimate of the background shape, obtained using $K^+ \pi^- \pi^+$ combinations as well as “rotated” pion candidates, selected from the opposite hemisphere.

5.1.2 Energy Restrictions

Hadrons containing a heavy quark generally have higher energies than the lighter hadrons produced in the fragmentation process. Thus, a substantial reduction in the combinatorial background can result simply by requiring the energy of D^{*+} candidates to exceed some minimum value. It is convenient to make this selection based on the scaled energy, defined by $x_E(D^{*+}) = E(D^{*+})/E_{\text{beam}}$ which is in the range $0 < x_E < 1$. In this analysis, a minimum scaled energy of 0.2 is required for all D^{*+} candidates. The scaled energy is also used to define kinematic regions where different sets of additional cuts are applied.

5.1.3 Ionization Energy Loss

Although the kinematics of the D^{*+} decay allow one to obtain a relatively background-free sample, there is additional information which can be used to suppress the background even further. The combinatorial background in the D^0 reconstruction is dominated by $\pi^+\pi^-$ combinations and can be greatly reduced by requiring that one track be positively identified as a kaon. In this analysis, this was done by detecting the characteristic energy loss of kaons in the OPAL jet chamber.

A charged particle which travels a small distance Δx through a gas will deposit energy ΔE along its path due to ionization. In the limit as $\Delta x \rightarrow 0$, the rate of energy loss is described by the Bethe-Bloch equation [65].

$$\frac{dE}{dx} = -\frac{A_1 Q^2}{\beta^2} \left(A_2 + 2 \log(\gamma\beta) - \beta^2 - \delta/2 \right). \quad (5.8)$$

Here, Q and β are the particle's charge and velocity, $\gamma = (1 - \beta^2)^{-1/2}$, and A_1 , A_2 and δ are constants which are properties of the gas mixture. The distribution of the energy losses observed in hadronic events and $\mu^+\mu^-$ events is shown in figure 5.3.

Since the energy loss depends only on the velocity, $\beta = p/E$, and since the particle's momentum is also measured, the difference between the measured and expected energy loss may be used to test a particular mass hypothesis. Under a

Figure 5.3: dE/dx distribution observed in hadronic and $\mu^+\mu^-$ events.

correct mass assignment the probability of measuring a difference as large as the one observed will be distributed uniformly, while an incorrect assignment would yield probability distributions peaked towards zero. These probability measures are referred to as *weights* and can be given a sign according to whether the expected energy loss is greater or less than the observed value. The weight assigned to a particle y , under the hypothesis that it was of type x , is denoted $w_x(y)$.

Although more efficient selection of kaons can be achieved using an asymmetric cut on the signed dE/dx weight, this can introduce biases if the Monte Carlo does not precisely model any shifts in the mean value of dE/dx which could arise due to different gas constants from year to year. This potential source of bias can be eliminated by using symmetric cuts on the kaon weight at the expense of a slightly less efficient selection. Thus, kaons were selected in this analysis by requiring that $|w_K(K)| > 0.1$ giving an intrinsic efficiency of 90%. Since all energy loss curves converge at high momenta, this cut is only performed when $x_E(D^{*+}) < 0.5$.

5.1.4 Decay Angles

A further property of the D^0 meson which can be exploited to suppress the combinatorial background is the fact that it is a pseudoscalar meson which must decay isotropically in its rest frame. The angle formed between the kaon and the D^0 momentum vector in the D^0 rest frame is denoted θ^* and its definition is illustrated schematically in figure 5.4. For real D^0 decays, $\cos \theta^*$ is then distributed uniformly between -1 and 1 while for random $K^-\pi^+$ combinations, this distribution is found to peak towards ± 1 . This behaviour of the background is less pronounced at high energies but still motivates selecting $K^-\pi^+$ combinations only from the central region of the $\cos \theta^*$ distribution. Thus, different selection criteria are imposed for high and low values of scaled energy:

$$|\cos \theta^*| < 0.8 \quad \text{for } 0.2 < x_E < 0.5 \quad (5.9)$$

$$|\cos \theta^*| < 0.9 \quad \text{for } x_E > 0.5. \quad (5.10)$$

Figure 5.4: Definition of the angle θ^* in the rest frame of the decaying D^0 .

5.1.5 D^{*+} Efficiency and Purity

The efficiency with which D^{*+} mesons can be reconstructed using the selection described above is determined from Monte Carlo simulations of D^{*+} decays in the OPAL detector. The D^{*+} production in $Z^0 \rightarrow b\bar{b}$ and $Z^0 \rightarrow c\bar{c}$ events was simulated using the Jetset 7.4 Monte Carlo program, while the detector simulation and event reconstruction was performed using GOPAL, as described in section 4.3. Knowing how many D^{*+} mesons were originally generated and the number which were correctly reconstructed allows efficiencies to be calculated using $\varepsilon = N_{\text{recon}}/N_{\text{gen}}$ where the uncertainty is evaluated using $\sigma_\varepsilon = \sqrt{\varepsilon(1-\varepsilon)/N_{\text{gen}}}$. The efficiency is energy dependent, due in part to the behaviour of the momentum resolution shown in equation (4.5) but also because the selection criteria used to identify D^{*+} 's is different for $x_E > 0.5$ and $x_E < 0.5$. The reconstruction efficiency is shown as a function of the D^{*+} scaled energy, x_E , in figure 5.5.

5.2 D^0 Reconstruction

Although the decay $D^{*+} \rightarrow D^0\pi^+$ described in the previous section can be reconstructed with a high purity, much of the background suppression arises from the

Figure 5.5: D^{*+} reconstruction efficiency as a function of $x_E(D^{*+})$.

kinematics of the D^{*+} decay. When reconstructing D^0 's which are not produced in D^{*+} decay, one cannot apply the cut on the mass difference distribution and hence, there is significantly more background in the direct $K^-\pi^+$ invariant mass distribution. Furthermore, many D^0 mesons produced in Z^0 decays are secondary decay products of D^{*+} 's which constitute a background to direct D^0 production or production via a D^{*0} decay. These issues make the reconstruction of D^0 's in OPAL's hadronic environment a challenging task.

Since the D^0 selection is used for reconstructing D_{s1}^+ decays produced mainly in $c\bar{c}$ events, a significant gain in purity can be obtained by using only high-energy D^0 candidates, selected with $x_E(K^-\pi^+) > 0.45$. Again, the fact that the D^0 is a spin-0 particle can be exploited to improve the purity by requiring $|\cos\theta^*| < 0.9$ and tight cuts on the dE/dx weights for the kaon and pion can be applied to reduce the background further. In particular, this selection excludes the regions

Figure 5.6: Distributions of $p(K^-)$ and $p(\pi^-)$ for fake combinations (a) and for combinations from real $D^0 \rightarrow K^- \pi^+$ decays (b). The lines indicate the selection $p > 6$ GeV/c applied to both the K^- and the π^+ .

$$-0.05 < w_K(K^-) < 0.15 \quad (5.11)$$

$$-0.01 < w_\pi(\pi^+) < 0.01. \quad (5.12)$$

Much of the background which remains is due to random $K^- \pi^+$ combinations, many of which come from light-flavour $Z^0 \rightarrow u\bar{u}, d\bar{d}$ or $s\bar{s}$ events. Since real D^0 mesons are expected to produce displaced secondary vertices, some of this background can be reduced using a loose cut on the decay length significance of the $K^- \pi^+$ vertex of $L/\sigma_L > 0.5$.

Finally, some rejection of random $K^- \pi^+$ combinations can be achieved by tight cuts on the momentum of both particles. This is shown in figure 5.6 which motivates the requirement that $p(K^-) > 6$ GeV/c and $p(\pi^-) > 6$ GeV/c. The invariant mass distribution of Monte Carlo $K^- \pi^+$ candidates selected using these criteria is shown in figure 5.7 where the components of the background are each indicated separately.

Rough estimates of the purity of this signal can be obtained assuming the approximate production rates of D^0 , D^{*0} and D^{*+} mesons to be in the ratio 1 :

Figure 5.7: Invariant mass distribution of $K^-\pi^+$ combinations with background contributions indicated by hatched histograms. The signal after the π_s^+ veto was applied is shown in figure (b).

3 : 3, respectively. Since $\text{Br}(D^{*+} \rightarrow D^0\pi^+) \sim .68$ and $\text{Br}(D^{*0} \rightarrow D^0(\pi^0, \gamma)) \sim 1$, one then expects that only about 50% of the D^0 's reconstructed in the peak shown in figure 5.7(a) came from D^{*0} decays. Some of the D^0 mesons produced in D^{*+} decay can be rejected by vetoing any candidate which can be combined with an additional pion to obtain a mass difference consistent with a D^{*+} decay. Specifically, $K^-\pi^+$ combinations are rejected from the D^0 sample if another π^+ is found which gives $M(K^-\pi^+\pi^+) - M(K^-\pi^+) < 0.16 \text{ GeV}/c^2$. The sample of D^0 candidates obtained using this selection is used in the reconstruction of D_{s1}^+ mesons, described in chapter 7.

5.3 K_S^0 Reconstruction

K_S^0 mesons, used in the analysis of D_{s1}^+ production, are reconstructed using their decay mode $K_S^0 \rightarrow \pi^+\pi^-$ which has a branching fraction of 68.6%. The K_S^0 is relatively long lived and can produce secondary vertices displaced from the primary

vertex by several centimeters. This characteristic topology makes it possible to obtain a high purity K_S^0 sample since secondary vertices formed using random $\pi^+\pi^-$ combinations from the primary vertex will usually have very short decay lengths. Furthermore, although a K_S^0 can produce vertices significantly displaced from the primary vertex, its momentum vector will, in general, point from the primary vertex through the decay vertex. These two characteristics are used to identify K_S^0 decays while rejecting uncorrelated $\pi^+\pi^-$ combinations.

In the $r\phi$ plane, tracks have circular trajectories and the K_S^0 reconstruction begins by considering all pairs of oppositely charged tracks which intersect within the detector volume. At the points of intersection, the following quantities are calculated:

- The radial decay length, R_{int} , measured from the primary vertex to the point of intersection.
- The difference in the z -coordinates, Δz , of the tracks at the point of intersection.
- The angle ζ between the momentum vector, calculated at the point of intersection, and the flight vector from the nominal beamspot.
- The transverse impact parameter of the K_S^0 flight path with respect to the nominal beamspot.

Because some discrepancies exist between the shapes of the Δz distributions observed in data and simulated by GOPAL, it is unwise to make tight cuts on this parameter. Hence, only a loose cut of $|\Delta z| < 80$ cm is imposed, which is much larger than the 6 cm resolution in the z -direction provided by the jet chamber. K_S^0 candidates with momentum vectors pointing back to the primary vertex are selected by requiring $\zeta < 1^\circ$ and false track crossings are rejected by the requirement that the first hit on either track occur at a radius no more than 10 cm in front of their point of intersection. Finally, to insure that the K_S^0 candidate is contained

within the volume of the jet chamber and to reject a large fraction of random $\pi^+\pi^-$ combinations produced at the primary vertex. the radius of the point of intersection is required to satisfy $1 < R_{\text{int}} < 150$ cm.

Since the jet chamber's intrinsic hit resolution in the z -direction is poor, tracks which decay outside the more precise vertex detector have inferior measurements of their polar angle. As described in section 4.4.4, a track's polar angle resolution is usually optimized by constraining the z -coordinate of its production point to be that of the primary vertex. Although this is not possible for K_S^0 decays, an improvement in polar angle resolution can be obtained by constraining them to come from a common point in the s - z plane. This can improve the $\pi^+\pi^-$ invariant mass resolution from $9.3 \text{ MeV}/c^2$ to about $8 \text{ MeV}/c^2$, however this figure can depend on the kinematics of a particular K_S^0 sample. To reject spurious K_S^0 candidates which are completely inconsistent with the hypothesis of $\pi^+\pi^-$ production at a common point, the χ^2 from the s - z fit, which has only two degrees of freedom, is required to satisfy $\chi_{sz}^2 < 50$.

The invariant mass distribution of $\pi^+\pi^-$ pairs satisfying these criteria is shown in figure 5.8. Although much of the background under the K_S^0 peak could be reduced by requiring that one or both tracks have impact parameters with respect to the primary vertex inconsistent with zero, this introduces an undesirable loss of efficiency, especially at high momentum. The efficiency obtained using this selection applied to Monte Carlo events is shown in figure 5.3 and, apart from an overall track reconstruction efficiency, is found to be dominated by the geometric acceptance cut $1 < R_{\text{int}} < 150$ cm.

Figure 5.8: Invariant mass distribution of $\pi^+\pi^-$ pairs after K_S^0 selection cuts have been applied. The shaded regions indicate the background from misidentified photon conversions and $\Lambda \rightarrow p\pi^-$ decays.

Figure 5.9: Efficiency for the reconstruction of $K_S^0 \rightarrow \pi^+\pi^-$ as a function of momentum. The dashed curve shows the intrinsic geometric efficiency for the reconstruction of K_S^0 decays within the acceptance cut $1 < R_{\text{int}} < 150$ cm, normalized to the peak of the observed efficiency.

Chapter 6

D^{**0} Production in $b\bar{b}$ and $c\bar{c}$ Events

The measurements of D^{**0} production rates at low energy e^+e^- collider experiments indicate that these mesons are produced in charm fragmentation. There is also substantial evidence for D₁⁰ production in B decays and therefore, one would expect to be able to observe D^{**0} production at LEP energies from both of these sources. However, Z⁰ decays are quite different from events produced at lower energy e^+e^- colliders and a number of problems must be solved in order to make useful measurements of D^{**0} production at LEP. First, the reconstruction of D^{**0} mesons at LEP is complicated by the large combinatorial backgrounds which arise from the greater charged particle multiplicities in Z⁰ decays. It is also more difficult to separate D^{**0} mesons produced in charm quark fragmentation from those produced in decays of b-flavoured hadrons.

This chapter describes the selection of two event samples containing D^{**0} → D^{*+}π⁻ decays, one enriched in D^{**0} mesons produced in charm fragmentation and the other enriched from decays of b-hadrons. Although neither sample was completely pure, the D^{**0} production rate measured in one sample was used to correct for contamination in the other sample and *vice versa*. This was performed by fitting both samples simultaneously. The selection of each sample is described

in the following two sections. This is followed by a description of the method used to fit the resulting D^{*0} signals to simultaneously determine the production rates in $c\bar{c}$ and $b\bar{b}$ events. Finally, a discussion of the systematic uncertainties and an interpretation of the results is presented.

6.1 $Z^0 \rightarrow c\bar{c} \rightarrow D^{*0} X$ Selection

$D^{*0} \rightarrow D^{*+}\pi^-$ decays were observed by first reconstructing a D^{*+} candidate, as described in section 5.1, and then combining it with oppositely charged pions in the event. These pions were required to have momenta exceeding 2 GeV/ c to reduce the combinatorial background resulting from the numerous lower energy particles produced in the fragmentation process. Since hadrons containing one of the primary quarks from the Z^0 decay typically have high energies, the non-charm component of this sample was suppressed by requiring the scaled energy of $D^{*+}\pi^-$ combinations, $x_E = E(D^{*+}\pi^-)/E_{\text{beam}}$, exceed 0.5.

Although the average scaled energy of D^{*0} mesons produced in the decay of b-hadrons is somewhat lower than that of primary D^{*0} mesons, a significant fraction of $D^{*+}\pi^-$ combinations from $b\bar{b}$ events still have $x_E > 0.5$. However, since the average lifetime of b-hadrons is about 1.5 ps while the mean D^0 lifetime is only 0.4 ps, the $b\bar{b}$ component was further reduced by requiring the proper time at which the D^0 decayed to be small. The “apparent” proper time, t , is defined in terms of the D^0 mass, its reconstructed decay length and momentum, both measured in the x - y plane:

$$ct(D^0) = \frac{L_{xy} \cdot M_{D^0}}{P_{xy}}. \quad (6.1)$$

This is an accurate measure of the proper time for $c\bar{c}$ events, but it will be larger in $b\bar{b}$ events on account of the additional contribution from the b-hadron lifetime. The $b\bar{b}$ component was suppressed in this analysis by the requirement that $ct(D^0) < 0.03$ cm. This retained about 85% of all D^0 mesons produced in charm fragmentation while only 40% of all D^0 mesons produced as decay products of

b-hadrons passed this requirement.

After these criteria were applied, the efficiency for reconstructing D^{*0} mesons produced in charm fragmentation with $x_E > 0.5$ and decaying to $D^{*+}\pi^-$ was $18.7 \pm 0.8\%$. Although this selection also accepted about 9% of D^{*0} mesons produced in b-decays with $x_E > 0.5$, the softer energy spectrum of these mesons resulted in less than 2% of all D^{*0} mesons produced in b-decays being selected.

The behaviour of this selection was studied using Monte Carlo Z^0 decays with a full detector simulation in which the rate of D^{*0} production is known. Figure 6.1 shows the mass difference distribution, $M(D^{*+}\pi^-) - M(D^{*+})$, which resulted when the charm selection criteria were applied with the non-charm components of the background identified. It is clear that although light flavour and $b\bar{b}$ events do populate the distribution, the combinatorial background and especially the D^{*0} signal is dominated by charm events. The lower number of D_2^0 mesons reconstructed in this distribution reflects the lower rate at which they were generated in the Monte Carlo, an input parameter which had been tuned to reflect the relative rates observed at ARGUS and CLEO.

Applying this selection to OPAL's 1990-95 data sample yielded the mass difference distribution shown in figure 6.2. An enhancement is seen in this distribution at the position expected for D_1^0 and D_2^0 production. The statistics available in this analysis did not allow the D^{*0} masses and widths to be measured with a precision competitive with ARGUS, CLEO or the Fermilab experiment E687, but they were found to be consistent with the results from these experiments.

6.2 $Z^0 \rightarrow b\bar{b} \rightarrow D^{*0} X$ Selection

The average energy of D^{*0} mesons produced in the decays of b-hadrons is somewhat lower than that of D^{*0} mesons produced in charm fragmentation. Thus, to select a sample of $D^{*+}\pi^-$ candidates produced mostly in $b\bar{b}$ events, no explicit cut on x_E was made other than the preselection cut of $x_E(D^{*+}) > 0.2$. However, the

Figure 6.1: Mass difference distribution obtained from the $c\bar{c}$ selection criteria applied to Monte Carlo Z^0 decays. The double-hatched histogram shows the contribution from light-flavour events while the single-hatched distribution indicates the additional $b\bar{b}$ background. Arrows indicate the mass differences of the D_1^0 and D_2^0 states generated in the Monte Carlo.

Figure 6.2: Mass difference distribution obtained from the $c\bar{c}$ selection criteria. The hatched histogram shows mass difference distribution of the wrong-sign $D^{*+}\pi^+$ candidates.

large background level in the invariant mass distribution of all $D^{*+}\pi^-$ combinations with no x_E restriction then prevented the observation of a statistically significant D^{*-0} signal. Hence, alternative methods had to be used to purify the sample.

In this analysis, a $D^{*+} \rightarrow D^0\pi^+$ decay was first identified as described in section 5.1, and an attempt was then made to reconstruct a b-decay vertex, displaced with respect to the primary vertex, from which the D^0 could have been produced. Identifying displaced secondary vertices is commonly used to obtain event samples enriched in b-decays, but this analysis also used the assignment of tracks to such a vertex to increase the purity of $D^{*+}\pi^-$ candidates. This required a vertex reconstruction algorithm which could not only identify secondary vertices, but which could also assign to them the tracks from b-decays with high efficiency. The rejection of π^- candidates produced at the primary vertex provided the reduction in combinatorial background required to observe D^{*-0} decays in $b\bar{b}$ events.

Although several algorithms for identifying secondary vertices are in common use, none were suitable for making efficient use of the fact that a D^{*+} had already been reconstructed. It is appropriate to briefly describe these before discussing the algorithm used in this analysis. However, it is first necessary to discuss the way in which vertex fitting is performed in general.

6.2.1 Vertex Fitting

The fitting of functions to measured data is commonly performed using a χ^2 minimization method. This procedure can also be applied to determine the coordinates of a decay vertex from a set of tracks. One method for constructing the χ^2 function in terms of the vectors of measured track parameters $\mathbf{c}_i$ with covariance matrices $\mathbf{V}_i$ is to calculate the expected track parameters under the assumption that they pass through the vertex. In general, the parameters which define the vertex position can be denoted by $\mathbf{x}$, and the expected parameters for track i can be written

$\hat{\mathbf{c}}_i(\mathbf{x}, \mathbf{c}_i)$. The χ^2 function is then written

$$\chi^2(\mathbf{x}) = \sum_i (\hat{\mathbf{c}}_i(\mathbf{x}, \mathbf{c}_i) - \mathbf{c}_i)^T \mathbf{V}_i^{-1} (\hat{\mathbf{c}}_i(\mathbf{x}, \mathbf{c}_i) - \mathbf{c}_i). \quad (6.2)$$

This can be minimized with respect to the parameters $\mathbf{x}$ to obtain the best estimate of the vertex position for the given set of tracks. The minimal value of χ^2 may then be used to test the consistency of the assignment of this set of tracks to the vertex and the consistency of each track may be tested by examining their individual contribution to the total χ^2 .

This approach can involve complicated expressions relating $\hat{\mathbf{c}}_i$ to the measured track parameters and to the vertex parameters. An alternative scheme, used in this analysis, involves minimizing

$$\chi^2 = \sum_i (\hat{\mathbf{c}}_i - \mathbf{c}_i)^T \mathbf{V}_i^{-1} (\hat{\mathbf{c}}_i - \mathbf{c}_i) \quad (6.3)$$

where the functional dependence of $\hat{\mathbf{c}}_i$ on $\mathbf{c}_i$ and $\mathbf{x}$ is contained implicitly in a number of constraining equations which must all be satisfied simultaneously. Track parameters can be constrained through equations of the form $\mathbf{g} = \hat{\mathbf{c}} - \mathbf{f}(\hat{\mathbf{c}}, \mathbf{x}) = 0$ but additional constraints, such as energy and momentum conservation, can also be introduced in the same way. The problem of minimizing χ^2 subject to these constraints is solved by introducing a vector of Lagrange multipliers μ and minimizing the function

$$L = \sum_i (\hat{\mathbf{c}}_i - \mathbf{c}_i)^T \mathbf{V}_i^{-1} (\hat{\mathbf{c}}_i - \mathbf{c}_i) + 2\mu^T \mathbf{g}(\hat{\mathbf{c}}, \mathbf{x}) \quad (6.4)$$

with respect to $\hat{\mathbf{c}}_i$ and $\mathbf{x}$. This can be done by expanding the functions $\mathbf{g}$ to first order about an initial estimate of the parameters $\hat{\mathbf{c}}_i$ and $\mathbf{x}$ and determining corrections to them which will bring L closer to a minimum. This can be repeated until the change in L after additional iterations becomes smaller than some cutoff $\Delta\chi_{\min}^2$.

6.2.2 Vertex Reconstruction Algorithms

The algorithms used to reconstruct secondary vertices have historically fallen into two categories. The first, known as the “tear-down” algorithm begins by fitting a common vertex to all tracks in a particular jet or in the entire event in the case when it is the position of the primary vertex being determined. Each track will then contribute to the χ^2 of the fit and the track with the largest contribution is removed and the fit is repeated. This iterative process continues until no track contributes more than χ_{cut}^2 to the overall χ^2 of the fit. This algorithm is quite useful for determining the position of the primary vertex but in a multiple vertex topology, it generally converges to the vertex with the highest multiplicity. Hence, it may not find low multiplicity vertices and may not include as many tracks in those vertices which are reconstructed as might be possible with other algorithms.

An alternative, referred to as the “build up” algorithm starts by selecting a number of “seed tracks” which are believed to have been produced at a secondary vertex with some degree of confidence. Criteria often used to select seed tracks are that they be identified as leptons or that they have impact parameters inconsistent with production at the primary vertex. Generally, more than one seed track is required to start the algorithm since at least two are needed to obtain an initial estimate of the vertex position. The build-up algorithm continues by selecting tracks which are more significantly separated from the primary vertex than from the seed vertex. Each such track is used along with the current seed tracks to fit new vertex candidates and the track which yields the new vertex with the largest fit probability exceeding some cutoff, P_{min} , is retained. The process is iterated until no more tracks can be added in this way. Although this has the advantage that the selection of seed tracks naturally enriches the sample of tracks from secondary vertices, their selection is usually not highly efficient. Furthermore, the particular choice of seed tracks can sometimes bias the fit, yielding displaced vertices even in light flavour events where they should be absent.

In this analysis, the prerequisite that a D^{*+} be reconstructed before any vertex

finding is performed can provide significant additional information about the event topology. First, the two tracks from the D^0 decay form the only displaced vertex due to the decay of a charm hadron and since this decay is fully reconstructed, the D^0 momentum vector should point back to either the primary vertex in a $c\bar{c}$ event, or to a b-vertex in a $b\bar{b}$ event. Also, the identification of the pion from the $D^{*+} \rightarrow D^0\pi^+$ decay can provide some complementary information about the location of a b-vertex. Finally, the invariant mass of the vertex can be used to guide the selection of tracks, since combinations yielding masses exceeding those typical of b-hadrons are clearly unphysical.

The tear-down and build-up vertex finders are generally applied in a more inclusive environment and do not make explicit use of the information provided by a tagged D^{*+} in these events. Thus, a new algorithm for reconstructing the b-vertex in events containing a tagged D^{*+} decay was developed.

6.2.3 $B \rightarrow D^{*+}X$ Vertex Reconstruction

In this case, the decay vertices being reconstructed are somewhat more complicated than is assumed by the tear-down or the build-up algorithms. Figure 6.3 shows the important features of this vertex topology. It consists of two decay vertices in addition to the primary vertex, where the direction of the D^0 vertex with respect to the assumed b-vertex position should be constrained by the D^0 momentum vector. It is therefore desirable to use a fitting procedure which makes use of this information. It was for this reason that the method of constrained χ^2 minimization using Lagrange multipliers was employed — a technique which is sufficiently general to easily accommodate any assumed topology.

Once a D^{*+} candidate was reconstructed, the algorithm assigned tracks to the primary vertex, the reconstructed b-vertex, or to a third subset unassociated with either vertex in the following five steps:

1. Tracks in the event were assigned to jets using the “cone” algorithm described in section 4.4.3, with minimum jet energy $\varepsilon = 10$ GeV and cone

Figure 6.3: Decay vertex topology for the decay $B \rightarrow D^{*+} X$. The primary, b-decay and D^0 decay vertices are indicated, as are the decay products of the D^{*+} . Additional tracks, which may be present the primary and b-decay vertices, are also shown.

half-angle $R = 0.7$ rad. The D^{*+} 4-vector was used as a single “track” and, in practice, was uniquely assigned to a single jet due to the choice of ε and the requirement that $x_E(D^{*+}) > 0.2$.

2. The primary vertex of the event was redetermined using an implementation of the tear-down algorithm described above, where the average LEP beamspot position was used as a constraint. In this case, all high-quality tracks in the event were used, except those which form the D^{*+} candidate. If all tracks in the D^{*+} jet were assigned to the primary vertex then it is assumed that no B vertex exists and the procedure terminates.
3. Those tracks in the D^{*+} jet which were not assigned to the primary vertex were combined with the D^{*+} to form the $B \rightarrow D^{*+} X$ topology using a tear-down vertex finding approach. Tracks which were removed from this b-decay vertex at this stage were then inconsistent with either the primary or

the secondary vertex. These formed the set of “junk” tracks and did not participate in any further vertex reconstruction.

4. Next, the track assignment between primary and secondary vertices was optimized by fitting all vertex topologies in which one track in the D^{*+} jet was moved from the primary vertex to the b-vertex, or *vice versa*. The track assignment which yielded the greatest reduction in $\chi^2_{\text{total}} = \chi^2_{\text{prim}} + \chi^2_{\text{sec}}$ was retained and this step was iterated until no further reassignment of tracks would reduce the value of χ^2_{total} .
5. Finally, in an attempt to improve the efficiency for assigning the π^- from the D^{*-0} decay to the b-vertex, tracks which were initially assigned to the primary vertex but which did not increase χ^2_{total} by more than $\beta_{\chi^2} = 1.5$ were reassigned to the secondary vertex. Thus, tracks produced with low p_T with respect to the jet axis could be assigned to the secondary vertex, even if they were also consistent with production at the primary vertex.

After this procedure was performed, displaced vertex topologies were selected by requiring $S/\sigma_S > 2$ and $L/\sigma_L > -2$. Here, S and L are the decay lengths of the b-vertex with respect to the primary and the D^0 vertex with respect to the b-vertex, respectively, and σ_S and σ_L are the corresponding uncertainties. The comparison of these decay length significance distributions with Monte Carlo is shown in figures 6.4 and 6.5. The agreement indicates that the description provided by the Monte Carlo is adequate.

Further comparisons between data and Monte Carlo were made for the the association of leptons with the b-vertex and for the reconstructed b-vertex multiplicity, determined from the number of tracks with momentum greater than 2 GeV/c assigned to that vertex. These distributions are shown in figures 6.6(a) and 6.6(b). The vertex multiplicity comparison serves as a check that the charged track multiplicities of b-hadron decays in data and Monte Carlo were similar. The assignment of leptons is useful since semileptonic B decays, such as $\bar{B}^0 \rightarrow D^{*+}\ell^-X$ produce

Figure 6.4: Decay length significance distribution for reconstructed $B \rightarrow D^{*+}X$ event topologies with the cut $L/\sigma_L > -2$ imposed. The points with error bars indicate the distribution obtained from analysis of OPAL data, while the histograms are from Monte Carlo. Both data and Monte Carlo distributions have been normalized to the same number of hadronic Z^0 decays analysed. The arrow indicates the S/σ_S cut used in the selection of b-vertices.

Figure 6.5: Significance distribution for decay lengths of $D^0 \rightarrow K^- \pi^+$ vertices measured with respect to reconstructed b-vertices with the cut $S/\sigma_S > 2$ imposed. The points with error bars indicate the distribution obtained from analysis of OPAL data, while the histograms are from Monte Carlo. Both data and Monte Carlo distributions have been normalized to the same number of hadronic Z^0 decays analysed. The arrow indicates the L/σ_L cut used in the selection of b-vertices.

Figure 6.6: (a) Multiplicity of tracks with $p > 2$ GeV/c in b-vertices reconstructed in data (points with error bars) and in Monte Carlo (histogram). (b) The assignment of leptons to the primary vertex, the reconstructed b-vertex and to the class of "junk" tracks, incompatible with either vertex. In (b), the open histogram with solid points are from $D^{*+}\ell^{-}$ combinations, expected from $\bar{B}$ decays, while the hatched histogram with open circles shows $D^{*+}\ell^{+}$ combinations, characteristic of background.

leptons and D^{*+} mesons with opposite charge. The identification of electrons as tracks which shower in the electromagnetic calorimeter and muons as tracks which penetrate beyond the hadron calorimeter to the muon chambers will be described in more detail in section 6.7. This identification can be performed with high purity for electrons and muons with momentum greater than 2 GeV/c and one sees that the assignment of leptons between primary and secondary vertices is similar in data and Monte Carlo. It is difficult to draw quantitative conclusions from these distributions since the momentum spectrum of leptons is quite different from that of the π^{-} in $D^{*+0} \rightarrow D^{*+}\pi^{-}$ decays. Nevertheless, these distributions indicate that no major disagreement exists between the properties of b-vertices reconstructed in the data and Monte Carlo.

Figure 6.7: Mass difference distribution obtained from the $b\bar{b}$ selection criteria applied to Monte Carlo events. The hatched histogram indicates the non- $b\bar{b}$ background which is due almost entirely to charm events. Arrows indicate the mass differences of the D_1^0 and D_2^0 states generated in the Monte Carlo.

6.2.4 The D^{*-0} Signal in $b\bar{b}$ Events

The b-vertex reconstruction procedure provided the improvement in purity necessary to observe D^{*-0} production in $b\bar{b}$ events with a reasonably small fraction of the observed signal coming from charm events. Selecting π^- candidates with momentum greater than $2 \text{ GeV}/c$ only from the set of tracks assigned to the reconstructed b-vertex reduced the combinatorial background in the $D^{*-}\pi^-$ mass difference distribution such that a clear D^{*-0} signal could be seen when applied to Monte Carlo Z^0 decays. This is shown in figure 6.7 wherein the background components due to non- $b\bar{b}$ events are also identified. The distribution is clearly dominated by $b\bar{b}$ events with a small contribution from charm and an almost negligible light-flavour background. Hence, if the D^{*-0} rate in data is similar to that in the Monte Carlo, then D^{*-0} production from decays of b-hadrons should be observable.

The selection was applied to OPAL data but due to the more stringent vertexing requirements, this part of the analysis used only data in which a silicon

Figure 6.8: Mass difference distribution obtained from the $b\bar{b}$ selection criteria applied to OPAL data. The hatched entries indicate the mass difference distribution from the $D^{*+}\pi^+$ candidates and show no resonant structure. The expected mass differences of the D_1^0 and D_2^0 states are indicated by the arrows.

microvertex detector was present. The resulting mass difference distribution is shown in figure 6.8 which shows an enhancement in the region expected to be populated by D_1^0 decays. No such enhancement was present in the corresponding distribution obtained from wrong-sign $D^{*+}\pi^+$ combinations, which are expected to be uncorrelated. From this distribution alone, there is some evidence for D_1^0 production while evidence for the D_2^0 remains weak.

Again, a small fraction of the signal peak consisted of D^{*0} mesons produced in charm events which had to be accounted for to obtain an unbiased measurement of the production rate in $b\bar{b}$ events. However, since a relatively pure D^{*0} signal was obtained in both $b\bar{b}$ and $c\bar{c}$ enhanced samples, the rate measured in one could be used to determine the background in the other. This was accomplished in a self consistent way by determining both production rates using the simultaneous fitting procedure described in the following section.

6.3 Simultaneous Likelihood Fits

Frequently, mass- or mass-difference distributions containing a signal from a particular decay are fit using separate parameterizations for the signal and background shapes. In this analysis, the mass difference, x , resulting from decays of each state $D_J^{(*)0}$, is described by a Breit-Wigner distribution convoluted with the detector resolution, σ_x :

$$s_J(x) = (\text{BW}_J * G_{\sigma_x})(x) = \int \frac{\Gamma_J/2\pi}{((y - \Delta m_J)^2 + (\Gamma_J/2)^2)} \cdot \frac{1}{\sigma_x \sqrt{2\pi}} e^{-(x-y)^2/2\sigma_x^2} dy \quad (6.5)$$

where Γ_J is the width of the $D_J^{(*)0}$ state and $\Delta m_J = M(D_J^{(*)0}) - M(D^{*+})$. Also, Monte Carlo studies[66] indicated that the background was adequately described by the functional form

$$b(x) \propto e^{-B(x-m_\pi)}(x - m_\pi)^C \quad (6.6)$$

where m_π is the pion mass.

Because of the sparse statistics in some bins of these histograms, a binned minimum χ^2 fit would have been inappropriate. Thus, an unbinned maximum likelihood method was used to determine the D^{*+0} production rates. For each flavour-enhanced distribution, the likelihood function was constructed from the probability distributions for the D_1^0 and D_2^{*0} signals and for the background. In previous analyses, the mass difference alone was used to differentiate between D_1^0 and D_2^{*0} signals, but since it is now well established that these states have the distinct distributions of $\cos \alpha$ shown in equation (3.26), the angular information from their decay products was also used to distinguish the two states from each other and from the background.

The likelihood functions for the charm-enhanced and bottom-enhanced samples, were of the form

$$\mathcal{L}_c = \prod_i \left((f_c^{c\bar{c}} f_1^{c\bar{c}} + f_c^{b\bar{b}} f_1^{b\bar{b}}) s_1(x_i, \alpha_i) + (f_c^{c\bar{c}}(1 - f_1^{c\bar{c}}) + f_c^{b\bar{b}}(1 - f_1^{b\bar{b}})) s_2(x_i, \alpha_i) \right. \\ \left. + (1 - f_c^{c\bar{c}} - f_c^{b\bar{b}}) b_c(x_i, \alpha_i) \right) \quad (6.7)$$

$$\mathcal{L}_b = \prod_i \left((f_b^{c\bar{c}} f_1^{c\bar{c}} + f_b^{b\bar{b}} f_1^{b\bar{b}}) s_1(x_i, \alpha_i) + (f_b^{c\bar{c}}(1 - f_1^{c\bar{c}}) + f_b^{b\bar{b}}(1 - f_1^{b\bar{b}})) s_2(x_i, \alpha_i) \right. \\ \left. + (1 - f_b^{c\bar{c}} - f_b^{b\bar{b}}) b_b(x_i, \alpha_i) \right) \quad (6.8)$$

where x_i and α_i represent the mass-difference and the helicity angle for the i^{th} $D^{*+}\pi^-$ combination, respectively. The functions $s_1(x_i, \alpha_i)$ and $s_2(x_i, \alpha_i)$ represent the probability distributions for the D_1^0 and D_2^{*0} signals which are products of the convoluted Breit-Wigner functions, shown in equation (6.5), and the expected helicity angle distribution from equation (3.26):

$$s_1(x, \alpha) = (\text{BW}_1 * G_{\sigma_x})(x) \cdot \frac{1}{4}(1 + 3 \cos^2 \alpha) \quad (6.9)$$

$$s_2(x, \alpha) = (\text{BW}_2 * G_{\sigma_x})(x) \cdot \frac{3}{4}(1 - \cos^2 \alpha). \quad (6.10)$$

Since the signal in the charm-enhanced sample was expected to be a mixture of D^{*0} decays from $c\bar{c}$ and $b\bar{b}$ events, it was necessary to specify each of their fractions by the parameters $f_c^{c\bar{c}}$ and $f_c^{b\bar{b}}$ and likewise, the parameters $f_b^{b\bar{b}}$ and $f_b^{c\bar{c}}$ denote the fractions of the D^{*0} signal from $b\bar{b}$ and $c\bar{c}$ events in the bottom-enhanced sample. Of the total signal from $D^{*0} \rightarrow D^{*+}\pi^-$ decays produced in $c\bar{c}$ and $b\bar{b}$ events, $f_1^{c\bar{c}}$ and $f_1^{b\bar{b}}$ specify the corresponding fractions from $D_1^0 \rightarrow D^{*+}\pi^-$ decays, respectively.

The background in each flavour-enhanced distribution was parameterized separately by the functions $b_q(x, \alpha)$, each being of the form

$$b_q(x, \alpha) = \frac{B_q^{C_q+1} e^{-B_q(x-m_\pi)} (x - m_\pi)^{C_q}}{\Gamma(C_q + 1) P(C_q + 1, B_q(m_{\text{max}} - m_\pi))} \cdot \frac{1 + \beta_q \cos \alpha + \gamma_q \cos^2 \alpha}{2(1 + \gamma_q/3)} \quad (6.11)$$

where B_q , C_q , β_q and γ_q are free parameters and where $\Gamma(x)$ and $P(x, y)$ are

the gamma and incomplete gamma functions[67], respectively. This expression is normalized to unit area between the mass differences m_π and $m_{\max}$ and contains an explicit description of the background helicity. The need to accommodate an arbitrary helicity structure in the background was necessary since it was not known *a priori*, if it had a uniform $\cos \alpha$ distribution.

While the fractions $f_c^{c\bar{c}}$ and $f_b^{b\bar{b}}$ were free parameters in the fit, the parameters $f_b^{c\bar{c}}$ and $f_c^{b\bar{b}}$ were not. These were specified in terms of $f_c^{c\bar{c}}$ and $f_b^{b\bar{b}}$, respectively, as follows:

$$f_b^{c\bar{c}} = f_c^{c\bar{c}} \frac{N_c \cdot N_b^{\text{MH}} \varepsilon_b^{c\bar{c}}}{N_b \cdot N_c^{\text{MH}} \varepsilon_c^{c\bar{c}}} \quad (6.12)$$

$$f_c^{b\bar{b}} = f_b^{b\bar{b}} \frac{N_b \cdot N_c^{\text{MH}} \varepsilon_c^{b\bar{b}}}{N_c \cdot N_b^{\text{MH}} \varepsilon_b^{b\bar{b}}} \quad (6.13)$$

where N_q and N_q^{MH} were the number of $D^{*+}\pi^-$ combinations reconstructed and the number of hadronic Z^0 decays analysed for each flavour enhanced sample. The terms $\varepsilon_q^{q'\bar{q}'}$ represent the efficiencies for selecting events of flavour q using the $q'\bar{q}'$ -enhanced selection. Since the charm-enhanced and bottom-enhanced analyses were made for different ranges of scaled energy, these efficiencies were required to extrapolate the observed signal into unmeasured regions of scaled energy which introduced a dependence on the Monte Carlo fragmentation model used. This dependence is investigated as a systematic study in section 6.5.7.

With the masses and widths of the D_1^0 and D_2^{*0} fixed to their world average values, the likelihood fit for each sample enriched in events of flavour q was performed with 6 freely varying parameters: $f_q^{q'\bar{q}'}$ and $f_1^{q'\bar{q}'}$, which specified the signal properties in $q'\bar{q}'$ events, and B_q , C_q , β_q and γ_q , which defined the properties of the background. From $f_c^{c\bar{c}}$ and $f_b^{b\bar{b}}$, the mean multiplicity of $D^{*+0} \rightarrow D^{*+}\pi^-$ decays in $c\bar{c}$ and $b\bar{b}$ events is expressed:

Parameter	Value
N_c^{MH}	4395217
N_b^{MH}	3770108
$\varepsilon_b^{b\bar{b}}$	0.0545 ± 0.0036
$\varepsilon_c^{c\bar{c}}$	0.1134 ± 0.0049
$\varepsilon_c^{b\bar{b}}$	0.0147 ± 0.0019
$\varepsilon_b^{c\bar{c}}$	0.0070 ± 0.0013
$\varepsilon'_{cc}(x_E > 0.5)$	0.1865 ± 0.0078
$\varepsilon'_{bb}(x_E > 0.2)$	0.0628 ± 0.0042

Table 6.1: Parameters used to determine efficiency corrected rates from simultaneous fits to the $b\bar{b}$ and $c\bar{c}$ distributions.

$$\begin{aligned} \bar{n}_{Z^0 \rightarrow c\bar{c} \rightarrow D^{*+0}(x_E > 0.5)} \cdot \text{Br}(D^{*+0} \rightarrow D^{*+}\pi^-) = \\ \frac{N_c f_c^{c\bar{c}}}{N_c^{\text{MH}} \varepsilon'_{cc}} \cdot \frac{1}{\text{Br}(D^{*+} \rightarrow D^0\pi^+) \cdot \text{Br}(D^0 \rightarrow K^-\pi^+)} \end{aligned} \quad (6.14)$$

$$\begin{aligned} \bar{n}_{Z^0 \rightarrow b\bar{b} \rightarrow D^{*+0}(x_E > 0.2)} \cdot \text{Br}(D^{*+0} \rightarrow D^{*+}\pi^-) = \\ \frac{N_b f_b^{b\bar{b}}}{N_b^{\text{MH}} \varepsilon'_{bb}} \cdot \frac{1}{\text{Br}(D^{*+} \rightarrow D^0\pi^+) \cdot \text{Br}(D^0 \rightarrow K^-\pi^+)} \end{aligned} \quad (6.15)$$

where ε'_{qq} are the efficiencies for each selection, this time calculated only for the range of scaled energies for which the selection was made.

6.4 Rate Measurements

The numbers of hadronic events analysed and the efficiencies used in equations (6.14) and (6.15) are listed in table 6.1. From the mass-difference and helicity angle distributions obtained from the charm- and bottom-enhanced samples of $D^{*+}\pi^-$ combinations, the following production rates times branching ratios were

obtained from this analysis of OPAL data:

$$R_{x_E > 0.5}^{c\bar{c}} \equiv \bar{n}_{Z^0 \rightarrow c\bar{c} \rightarrow D^{*\ast 0}(x_E > 0.5)} \cdot \text{Br}(D^{*\ast 0} \rightarrow D^{*+}\pi^-) = 0.0054_{-0.0013}^{+0.0014} \quad (6.16)$$

$$R_{x_E > 0.2}^{b\bar{b}} \equiv \bar{n}_{Z^0 \rightarrow b\bar{b} \rightarrow D^{*\ast 0}(x_E > 0.2)} \cdot \text{Br}(D^{*\ast 0} \rightarrow D^{*+}\pi^-) = 0.0161_{-0.0036}^{+0.0037} \quad (6.17)$$

where, at this point, only the statistical uncertainty is shown. The fractions of these signals due to $D_1^0 \rightarrow D^{*+}\pi^-$ decays were determined to be

$$f_1^{c\bar{c}} = 0.56_{-0.15}^{+0.15} \quad (6.18)$$

$$f_1^{b\bar{b}} = 0.77_{-0.14}^{+0.16} \quad (6.19)$$

where again, only the statistical uncertainties are given. A detailed treatment of the systematic uncertainties associated with these measurements is described in the following section. The parameters $f_c^{c\bar{c}}$ and $f_1^{c\bar{c}}$, obtained from the likelihood fit, are correlated, as are $f_b^{b\bar{b}}$ and $f_1^{b\bar{b}}$. For the rate and $D_1^0 \rightarrow D^{*+}\pi^-$ fraction in charm events, the correlation coefficient is -0.089 , while in bottom events, it is -0.419 . Thus, it would be unwise to treat these parameters as uncorrelated quantities when propagating uncertainties in calculations derived from them. In the calculations which follow in section 6.6, these correlations have been taken into account.

The fitted mass difference distributions are shown in figures 6.9(a) and 6.10(a). The numbers of $D^{*\ast 0}$ decays observed in each signal peak were 147 ± 37 and 106 ± 24 , respectively. Clearly, the different shapes of the signal and background mass difference distributions result in only small correlations between the parameters which define them. Also, the fitted parameters β_q and γ_q indicated that the background helicity was consistent with an isotropic distribution in both flavour enhanced samples.

The effects of the helicity angle information are seen in figures 6.9(b) and 6.10(b) which show the corresponding mass-difference distributions, but this time only for $D^{*+}\pi^-$ combinations with $|\cos \alpha| > 0.7$. For this cut, a greater fraction of the D_1^0 component of the signal is retained relative to the D_2^{*0} compo-

ment and the background. These figures also show the fractions of the D^{*0} signals which were due to events of the opposite flavour. There is clearly a non-negligible fraction of D^{*0} decays from $b\bar{b}$ events populating the $c\bar{c}$ sample. Although there is a smaller fraction from $c\bar{c}$ events in the $b\bar{b}$ sample, it is still important to account for its presence.

Figure 6.9: Fits to the $c\bar{c}$ mass difference distribution. The top figure shows the inclusive fit while the bottom shows the behaviour for $|\cos \alpha| > 0.7$ which has a higher efficiency for the D_1^0 state than for the D_2^{*0} or for background. The hatched region shows the fraction of the signal which is due to D^{*+0} mesons produced in $b\bar{b}$ events.

Figure 6.10: Fits to the $b\bar{b}$ mass difference distribution. The top figure shows the inclusive fit while the bottom shows the behaviour for $|\cos \alpha| > 0.7$ which has a higher efficiency for the D_1^0 state than for the D_2^{*0} or for background. The hatched region shows the fraction of the signal which is due to D^{*0} mesons produced in $c\bar{c}$ events.

Source	$\Delta R_{x_F > 0.5}^{c\bar{c}}$	$\Delta R_{x_F > 0.2}^{b\bar{b}}$	$\Delta f_1^{c\bar{c}}$	$\Delta f_1^{b\bar{b}}$
Monte Carlo statistics	$+0.23$ -0.23×10^{-3}	$+1.0$ -1.0×10^{-3}	—	—
$\text{Br}(D^{*+} \rightarrow D^0 \pi^+) \cdot \text{Br}(D^0 \rightarrow K^- \pi^+)$	$+0.20$ -0.20	$+0.6$ -0.6	—	—
Resonance parameters	$+0.21$ -0.21	$+1.0$ -0.9	$+0.023$ -0.022	$+0.038$ -0.040
Resolution scale	$+0.15$ -0.23	$+0.6$ -0.6	$+0.021$ -0.024	$+0.007$ -0.000
B lifetime	$+0.00$ -0.04	$+0.3$ -0.3	$+0.001$ -0.001	$+0.000$ -0.000
Track resolution	$+0.04$ -0.00	$+1.2$ -0.2	$+0.003$ -0.007	$+0.002$ -0.005
Fragmentation parameters	$+0.15$ -0.08	$+0.1$ -0.0	$+0.006$ -0.004	$+0.000$ -0.001
$\cos \theta_\pi$ acceptance	$+0.34$ -0.65	$+0.6$ -0.7	$+0.004$ -0.006	$+0.001$ -0.001
D_1^0 polarization	$+0.06$ -0.23	$+0.3$ -0.2	$+0.008$ -0.014	$+0.009$ -0.004
Total	$+0.55$ -0.82×10^{-3}	$+2.2$ -1.8×10^{-3}	$+0.033$ -0.037	$+0.039$ -0.041

Table 6.2: Summary of systematic uncertainties for the D^{*0} production rates and D_1^0 signal fraction. The rates $R_{x_F > 0.5}^{c\bar{c}}$ and $R_{x_F > 0.2}^{b\bar{b}}$ correspond to those shown in equations (6.16) and (6.17) respectively.

6.5 Systematic Studies

So far, only statistical uncertainties have been quoted for the measured production rates and D_1^0 fractions for the signals in $c\bar{c}$ and $b\bar{b}$ events. There are several factors which can influence these results, some as simple as statistical uncertainties in branching ratios, but others which are more subtle such as the effects of selection criteria on the form of the helicity angle distributions used in this analysis. The contributions from each systematic effect are summarized in table 6.2 and are described individually in the following subsections.

6.5.1 Monte Carlo Statistics

Because the Monte Carlo samples used to determine the efficiencies are of finite size, there is some uncertainty in the rate measurements due to uncertainties in ε'_{bb} and ε'_{cc} , shown in table 6.1. These uncertainties could be reduced by generating more Monte Carlo events so they are, in a sense, statistical in nature. However, it

is useful to distinguish between the uncertainty due to limitations in the size of the data sample from those which, in principle, could be reduced if necessary. Thus, the uncertainty in the efficiencies resulting from limited Monte Carlo statistics, shown in the first row of table 6.2, are combined with the other sources of systematic uncertainty.

6.5.2 Branching Ratios

The branching ratios used in the calculation of the D^{*0} production rates are

$$\text{Br}(D^{*+} \rightarrow D^0 \pi^+) = (68.3 \pm 1.4)\% \quad (6.20)$$

$$\text{Br}(D^0 \rightarrow K^- \pi^+) = (3.83 \pm 0.12)\% \quad (6.21)$$

which, although obtained from world averages, are dominated by measurements made by ARGUS[68, 69] and CLEO[70, 71]. The uncertainties in the measured rates of D^{*0} production introduced through these branching ratios are shown in table 6.2 and are seen to be comparable in size to other sources of systematic uncertainty in this analysis.

6.5.3 Resonance Parameters

The Breit-Wigner parameters of the D_1^0 and D_2^{*0} resonances were fixed to world averages when performing fits to the mass difference distributions. Since these values are determined with limited precision, it is appropriate to study the dependence of the measured rates on the masses and widths of the D^{*0} resonances. Repeating the fits with the masses and widths of the D_1^0 and D_2^{*0} varied by one standard deviation resulted in the changes to the measurements shown in table 6.3. Adding the individual variations in quadrature resulted in the uncertainties listed in the third row of table 6.2.

Parameter	$R_{x_F > 0.5}^{c\bar{c}}$	$R_{x_F > 0.2}^{b\bar{b}}$	$f_1^{c\bar{c}}$	$f_1^{b\bar{b}}$
M_1 $^{+1.8}_{-1.8}$ MeV/ c^2	-0.15 $+0.11 \times 10^{-3}$	-0.38 $+0.34 \times 10^{-3}$	-0.000 $+0.001$	$+0.013$ -0.013
M_2 $^{+2.0}_{-2.0}$ MeV/ c^2	$+0.04$ -0.08	$+0.38$ -0.38	$+0.004$ -0.001	-0.015 $+0.016$
Γ_1 $^{+4.6}_{-3.5}$ MeV/ c^2	$+0.15$ -0.15	$+0.71$ -0.60	$+0.022$ -0.020	$+0.026$ -0.024
Γ_2 $^{+5.0}_{-5.0}$ MeV/ c^2	$+0.00$ -0.04	$+0.34$ -0.38	-0.004 $+0.008$	-0.020 $+0.024$
Total	$+0.21$ -0.21×10^{-3}	$+0.95$ -0.87×10^{-3}	$+0.023$ -0.022	$+0.038$ -0.040

Table 6.3: Systematic uncertainties associated with varying the $D^{*\pm 0}$ resonance parameters by ± 1 standard deviation of their measured values. The total systematic uncertainty is obtained by adding the individual contributions in quadrature.

6.5.4 Mass Difference Uncertainty Scale

The event-by-event mass difference uncertainty was obtained by propagation of errors using the track parameter covariance matrices. Although the reconstruction algorithms were carefully tuned to provide accurate estimates of these covariance matrices whenever possible, their properties in data and Monte Carlo are usually found to differ slightly. Thus, it is important to demonstrate that the measurements presented here do not depend too strongly on the precision with which they were determined.

This is most easily done by scaling the measured mass-difference uncertainties by a factor s_σ when convoluting the Breit-Wigner functions with the detector resolution. Based on Monte Carlo resolution studies, the discrepancy between the mass difference uncertainties determined in data and Monte Carlo is known to be less than 25%. The scale factor was varied between 0.75 and 1.25, resulting in the small variations listed in the fourth row of table 6.2.

6.5.5 B Lifetime

The Monte Carlo detector simulation, from which the efficiencies were determined, generated decays of D^0 and B mesons with lifetimes set to $\tau(D^0) = 0.415$ ps and $\tau(B) = 1.60$ ps, which are consistent with world averages[12]. Although the D^0

lifetime has been measured with a precision better than 1%, the uncertainty in the B lifetime is about 5%. Because the selection of D^{*-0} decays in $b\bar{b}$ events was based on the decay length of b-vertices, a discrepancy between the actual B lifetime and the value used in the simulation could result in biased efficiencies. To estimate the magnitude of any biases introduced in this way, the efficiencies were redetermined with distributions reweighted such that they corresponded to B lifetimes varied by $\pm 5\%$. The variations in the D^{*-0} production rates and $D_1^0 \rightarrow D^{*+}\pi^-$ fractions introduced by this uncertainty in the B lifetime are shown in the fifth row of table 6.2. The systematic uncertainties associated with the uncertainty in the D^0 lifetime were negligible.

6.5.6 Track Resolution

Although the Monte Carlo simulation was tuned to produce distributions which match those observed in data, it is important to estimate the sensitivity of measurements to any discrepancies in the modeling. In this analysis, the vertex reconstruction relied heavily on the correct determination of track parameters and their error matrices. In OPAL, it is common to study the dependence of measurements on the simulated track parameter resolutions by repeating the analysis of Monte Carlo events in which these resolutions have been artificially inflated or reduced. This is usually done by calculating the true track parameters and scaling the difference between these and the reconstructed track parameters by constant factors. In this analysis, the parameters d_0 , ϕ_0 , $\tan \lambda$ and z_0 were scaled by $\pm 10\%$, resulting in the variation in the measured D^{*-0} rates and $D_1^0 \rightarrow D^{*+}\pi^-$ fractions shown in the sixth row of table 6.2.

6.5.7 Fragmentation Parameters

D^{*-0} production rates were measured for restricted ranges of scaled energy but assumptions were still made about the shape of the D^{*-0} energy spectrum in $c\bar{c}$ and $b\bar{b}$ events simulated in the Monte Carlo. The Monte Carlo generated the energy

distribution of primary heavy hadrons using Peterson fragmentation functions of the form shown in equation (3.3) where the two Peterson parameters ϵ_c and ϵ_b were tuned to reproduce the mean scaled energies of D^0 , D^+ , D^{*+} and B mesons.

Peterson fragmentation of heavy quarks reproduces the observed behaviour of the high energy hadrons in the data, but it is very difficult to check that the low energy behaviour is also described well. In this analysis, the D^{*-0} production rate in $c\bar{c}$ events was measured for $x_E > 0.5$ but was then extrapolated to the full range of x_E to estimate the background from $c\bar{c}$ events in the $b\bar{b}$ -enhanced sample. The systematic uncertainties introduced in this way were then estimated by varying the Peterson fragmentation parameters over a range of values which reproduces the uncertainties in the mean scaled energies of primary quarks[72].

An additional source of model dependence arose from the simulation of inclusive decays of b-hadrons to D^{*-0} mesons. In the Jetset Monte Carlo program, this was done using string fragmentation at an energy scale corresponding to the b-hadron mass. The shape of the energy spectrum of D^{*-0} mesons produced in decays of b-hadrons was used to estimate the background from $b\bar{b}$ events in the $c\bar{c}$ -enhanced sample and it was important to verify that measurements made in this analysis did not depend too strongly on the details of the b-decay simulation. The systematic uncertainty associated with these effects were also estimated by varying the Peterson fragmentation parameters.

The Peterson fragmentation parameters were varied from their central values of $\epsilon_c = 0.031$ and $\epsilon_b = 0.0038$ over the range shown in table 6.4 from which the systematic uncertainty in the measured rates and D_1^0 fractions were determined. This range was determined from the uncertainties in OPAL measurements of the mean scaled energy of B and D mesons in Z^0 decays[73, 74]. It was found that these introduced a relatively small contribution to the overall systematic uncertainties of the measurements made in this analysis.

Parameter	$R_{\mathcal{E}\mathcal{E}}^{c\bar{c}} > 0.5$	$R_{\mathcal{E}\mathcal{E}}^{b\bar{b}} > 0.2$	$f_1^{c\bar{c}}$	$f_1^{b\bar{b}}$
$0.018 < \epsilon_c < 0.044$	$+0.04 \times 10^{-3}$ -0.04×10^{-3}	$+0.08 \times 10^{-3}$ -0.04×10^{-3}	$+0.002$ -0.002	$+0.000$ -0.001
$0.0018 < \epsilon_b < 0.0068$	$+0.15 \times 10^{-3}$ -0.08×10^{-3}	$+0.04 \times 10^{-3}$ -0.00×10^{-3}	$+0.006$ -0.004	$+0.000$ -0.001
Total	$+0.15 \times 10^{-3}$ -0.08×10^{-3}	$+0.09 \times 10^{-3}$ -0.04×10^{-3}	$+0.006$ -0.004	$+0.000$ -0.001

Table 6.4: Systematic uncertainties due to variation in the Peterson fragmentation parameters.

6.5.8 $\cos \theta_\pi$ Acceptance

Since each narrow $D^{*\pm 0}$ meson has non-zero spin, the angular distributions of its decay products depends on its helicity, or the alignment of its spin and momentum vectors. These distributions can be expressed in terms of the θ_π , which represents the angle at which the π^- is emitted with respect to the $D^{*\pm 0}$ boost vector, in the $D^{*\pm 0}$ rest frame, analogous to the angle θ^* , defined for the D^0 decay in section 5.1.4. Although an isotropic distribution of $\cos \theta_\pi$ would result if all helicity states were equally populated by the $D^{*\pm 0}$ production process, one cannot assume that this is the case. The minimum momentum requirement imposed on the π^- for $D^{*\pm} \pi^-$ combinations restricted $D^{*\pm 0}$ decays to the ranges of $\cos \theta_\pi$ shown in figure 6.11.

To estimate the systematic effects resulting from possible non-isotropic decay angle distribution in the data, the $D^{*\pm 0}$ production rates were determined in bins of $\cos \theta_\pi$ which were then fit to the functional form $1 + a \cos^2 \theta_\pi$. The parameters obtained from these fits were

$$a_{c\bar{c}} = -0.03 \pm 1.00 \quad (6.22)$$

$$a_{b\bar{b}} = 0.76 \pm 1.12 \quad (6.23)$$

The central values were used to reweight the isotropic distributions to estimate the relative efficiencies for non-isotropic decays. For the central values, the ratios of the isotropic to non-isotropic efficiencies, denoted ρ_π , were close to one and no correction to the central values of the rates has been applied. Nevertheless, the

Figure 6.11: Effective $\cos \theta_\pi$ acceptance introduced by restrictions on the π^- momentum for $c\bar{c}$ and $b\bar{b}$ selections. The hatched histograms show the accepted range of $\cos \theta_\pi$ after cuts on both $x_E(D^{*0})$ and $p(\pi^-) > 2 \text{ GeV}/c$ have been applied. The upper two histograms are for isotropic decay distributions while the lower two have been reweighted by the indicated functions.

Relative efficiency	$R_{x_F > 0.5}^{c\bar{c}}$	$R_{x_F > 0.2}^{b\bar{b}}$	$f_1^{c\bar{c}}$	$f_1^{b\bar{b}}$
$\rho_\epsilon^{c\bar{c}} = 1.00^{+0.10}_{-0.05}$	$-0.65^{+0.11}_{-0.08} \times 10^{-3}$	$+0.11^{+0.11}_{-0.08} \times 10^{-3}$	$-0.006^{+0.006}_{-0.003}$	$-0.001^{+0.001}_{-0.001}$
$\rho_\epsilon^{b\bar{b}} = 1.00^{+0.04}_{-0.03}$	$+0.04^{+0.04}_{-0.04} \times 10^{-3}$	$-0.65^{+0.54}_{-0.54} \times 10^{-3}$	$+0.002^{+0.002}_{-0.002}$	$+0.001^{+0.001}_{-0.000}$
Total	$+0.34^{+0.55}_{-0.65} \times 10^{-3}$	$+0.55^{+0.55}_{-0.65} \times 10^{-3}$	$+0.004^{+0.004}_{-0.006}$	$+0.001^{+0.001}_{-0.001}$

Table 6.5: Systematic uncertainties due to reweighting efficiencies to account for possible non-isotropic decay angle distributions.

changes in these ratios as $a_{c\bar{c}}$ and $a_{b\bar{b}}$ were varied within their statistical uncertainty were used to estimate the systematic uncertainty in the efficiency corrected $D^{*\ast 0}$ production rates. These contributions to the systematic uncertainty are also included in table 6.5.

6.5.9 D_1^0 Polarization

The $\cos \alpha$ distributions shown in equation (3.26) are used in the likelihood function to help distinguish between the D_1^0 and D_2^{*0} states. However, the form shown for the D_1^0 state in this equation is valid only when integrating over the full range of the decay angle, θ_π . As described in the previous section, only $D^{*\ast 0}$ mesons with a restricted range of decay angles are accepted due to the minimum momentum cut applied to the π^- candidates. While this does not affect the form of the distribution for the D_2^{*0} helicity angle, it can modify the D_1^0 helicity angle distribution and introduce a dependence on the average D_1^0 polarization. The modified form of the angular distribution for D_1^0 decays selected in a restricted range of $\cos \theta_\pi$ is derived in appendix A and is shown explicitly in equation (A.31). Because of parity invariance of the strong interaction, the D_1^0 polarization can be expressed in terms of one element of its polarization density matrix, denoted ρ_{00} .

To estimate the systematic effects which could result from a restricted $\cos \theta_\pi$ acceptance and a non-uniform D_1^0 polarization, the analysis was repeated using the distribution shown in equation (A.31) and with the ρ_{00} parameter varied between

Figure 6.12: $\cos \alpha$ distributions for D_1^0 decays in the heavy quark limit for restricted ranges of $\cos \theta_\pi$. The solid lines indicate the expected distribution when no restriction on $\cos \theta_\pi$ is imposed, while the dotted and dashed lines show the expected behaviour for extreme values of the polarization parameter ρ_{00} .

0 and 1. The limiting range of $\cos \theta_\pi$ acceptance were estimated from the distributions shown in figure 6.11 and roughly correspond to $\cos \theta_\pi > -0.6$ for the $c\bar{c}$ sample and $\cos \theta_\pi > -0.3$ for the $b\bar{b}$ sample. The resulting change in the shape of the $\cos \alpha$ distributions is shown in figure 6.12 and the resulting variation in the measured rates and $D_1^0 \rightarrow D^{*+}\pi^-$ fractions are summarized in table 6.6.

6.5.10 Cross Checks

The rates calculated using equations (6.14) and (6.15) rely on the accurate determination of the efficiencies using the Monte Carlo simulation of both the physics processes involved and the response of the OPAL detector to these events. Due to the presence of backgrounds, it is not possible to directly compare the Monte Carlo and data distributions on which selections are applied. However, agreement can be checked by varying the selection criteria since a disagreement between data and Monte Carlo would yield a systematic dependence of the measured rates on the specific values of the cuts used.

D_1^0 polarization	$R_{x_F > 0.5}^{c\bar{c}}$	$R_{x_F > 0.2}^{b\bar{b}}$	$f_1^{c\bar{c}}$	$f_1^{b\bar{b}}$
$0.0 < \rho_{00}^{c\bar{c}} < 1.0$	$+0.04$ -0.23×10^{-3}	$+0.27$ -0.08×10^{-3}	$+0.003$ -0.014	$+0.006$ -0.002
$0.0 < \rho_{00}^{b\bar{b}} < 1.0$	$+0.04$ -0.04×10^{-3}	$+0.04$ -0.15×10^{-3}	$+0.007$ -0.002	$+0.007$ -0.004
Total	$+0.06$ -0.23×10^{-3}	$+0.27$ -0.17×10^{-3}	$+0.008$ -0.014	$+0.009$ -0.004

Table 6.6: Systematic uncertainties associated with departures of the D_1^0 helicity angle distribution from the form shown in equation (3.28).

Interpretation of the resulting variations must be done with care due to the nature of statistical fluctuations, both in the data and the Monte Carlo. In the absence of evidence for systematic effects it is inappropriate to include variations in the measured parameters with the systematic uncertainty since this would amount to counting the statistical uncertainty twice[75]. However, it is important to demonstrate that the magnitude of such variations is consistent with purely statistical fluctuations.

To this end, the criteria used to select $D^{*\pm 0}$ decays in $c\bar{c}$ and $b\bar{b}$ events were varied over reasonable ranges of values. The resulting variations in the rates and D_1^0 fractions were generally consistent with statistical fluctuations. There was little evidence for large systematic effects in any of these studies and therefore, no systematic uncertainty due to variation of the $D^{*\pm 0}$ selection criteria is included in the measurements. The ranges over which the cuts were varied and the changes in the measured quantities are shown in table 6.7.

Parameter	$R_{x_E > 0.5}^{c\bar{c}}$	$R_{x_E > 0.2}^{b\bar{b}}$	$f_1^{c\bar{c}}$	$f_1^{b\bar{b}}$
$c\bar{c} p(\pi^-) > 2.0 \pm 0.5 \text{ GeV}/c$	$+0.08$ -0.80×10^{-3}	$+0.15$ -0.00×10^{-3}	$+0.039$ -0.020	$+0.001$ -0.007
$b\bar{b} p(\pi^-) > 2.0 \pm 0.5 \text{ GeV}/c$	$+0.19$ -0.11	$+1.68$ -0.34	$+0.016$ -0.037	$+0.025$ -0.037
$ct(D^0) < 0.03 \pm 0.02 \text{ cm}$	$+1.15$ -0.38	$+0.08$ -0.19	$+0.000$ -0.042	$+0.000$ -0.002
$S/\sigma_S > 2 \pm 2$	$+0.31$ -0.23	$+2.37$ -3.56	$+0.001$ -0.013	$+0.103$ -0.028
$L/\sigma_L > -2 \pm 2$	$+0.19$ -0.04	$+0.55$ -2.03	$+0.028$ -0.007	$+0.035$ -0.093
Range	$+1.15$ -0.38×10^{-3}	$+2.37$ -3.56×10^{-3}	$+0.042$ -0.037	$+0.103$ -0.028

Table 6.7: Variation in rates and $D_1^0 \rightarrow D^{*+}\pi^-$ fractions with varying the D^{*0} selection criteria.

6.6 Interpretation of Results

The measured D^{*0} production rates in Z^0 decays, including both statistical and systematic uncertainties are

$$\bar{n}_{Z^0 \rightarrow c\bar{c} \rightarrow D^{*0}(x_E > 0.5)} \cdot \text{Br}(D^{*0} \rightarrow D^{*+}\pi^-) = 0.0054_{-0.0013}^{+0.0014}(\text{stat.})_{-0.0008}^{+0.0006}(\text{syst.}) \quad (6.24)$$

$$\bar{n}_{Z^0 \rightarrow b\bar{b} \rightarrow D^{*0}(x_E > 0.2)} \cdot \text{Br}(D^{*0} \rightarrow D^{*+}\pi^-) = 0.0161_{-0.0036}^{+0.0037}(\text{stat.})_{-0.0018}^{+0.0022}(\text{syst.}) \quad (6.25)$$

while the D_1^0 fractions of the signals in $c\bar{c}$ and $b\bar{b}$ events are

$$f_1^{c\bar{c}} = 0.56_{-0.15}^{+0.15}(\text{stat.})_{-0.04}^{+0.03}(\text{syst.}) \quad (6.26)$$

$$f_1^{b\bar{b}} = 0.77_{-0.14}^{+0.16}(\text{stat.})_{-0.04}^{+0.04}(\text{syst.}). \quad (6.27)$$

While the observation of D^{*0} production in b-decays is important for understanding the decays of b-hadrons, it is the rate of D^{*0} production in charm events which can be compared with predictions from fragmentation models. However, such a comparison requires extrapolating the measurement made for $x_E > 0.5$ to the full range of scaled energies. This is performed using the Peterson fragmentation

model, implemented in Jetset.

Such a study indicated that approximately 64% of the D^{*0} mesons in charm events are produced with $x_E > 0.5$ while 88% are produced with $x_E > 0.2$ in $b\bar{b}$ events. Applying these corrections and varying the Peterson fragmentation parameters over the range of values previously indicated in table 6.4 resulted in the following production rates, corrected for the full range of x_E :

$$\bar{n}_{Z^0 \rightarrow c\bar{c} \rightarrow D^{*0}} \cdot \text{Br}(D^{*0} \rightarrow D^{*+}\pi^-) = 0.0084^{+0.0021}_{-0.0021}(\text{stat.})^{+0.0009}_{-0.0013}(\text{syst.})^{+0.0003}_{-0.0005}(\text{frag.}) \quad (6.28)$$

$$\bar{n}_{Z^0 \rightarrow b\bar{b} \rightarrow D^{*0}} \cdot \text{Br}(D^{*0} \rightarrow D^{*+}\pi^-) = 0.0183^{+0.0042}_{-0.0041}(\text{stat.})^{+0.0025}_{-0.0020}(\text{syst.})^{+0.0004}_{-0.0002}(\text{frag.}) \quad (6.29)$$

where the third error indicates the systematic uncertainty associated with the extrapolation in scaled energy. In many of the measurements which follow, this is added in quadrature with the systematic uncertainties previously discussed.

These measured quantities still include the branching ratios for $Z^0 \rightarrow b\bar{b}$ and $Z^0 \rightarrow c\bar{c}$. Also, since the branching ratios of the D_1^0 and D_2^{*0} to the $D^{*+}\pi^-$ final state are different, the term $\text{Br}(D^{*0} \rightarrow D^{*+}\pi^-)$ represents an average of these, weighted by the relative D_1^0 and D_2^{*0} production rates. It is more useful to calculate the quantities $f(q \rightarrow D_1^0)$ and $f(q \rightarrow D_2^{*0})$, which represent the probability that a particular flavour of quark will produce a $D_j^{(*)0}$ meson either when it fragments into hadrons, or in the subsequent decay chain. These can be computed from equations (6.28) and (6.29) using the relations

$$f_1^{q\bar{q}} \cdot \bar{n}_{Z^0 \rightarrow q\bar{q} \rightarrow D^{*0}} \cdot \text{Br}(D^{*0} \rightarrow D^{*+}\pi^-) = 2 \cdot \frac{\Gamma_{q\bar{q}}}{\Gamma_{\text{had}}} \cdot f(q \rightarrow D_1^0) \cdot \text{Br}(D_1^0 \rightarrow D^{*+}\pi^-) \quad (6.30)$$

$$(1 - f_1^{q\bar{q}}) \cdot \bar{n}_{Z^0 \rightarrow q\bar{q} \rightarrow D^{*0}} \cdot \text{Br}(D^{*0} \rightarrow D^{*+}\pi^-) = 2 \cdot \frac{\Gamma_{q\bar{q}}}{\Gamma_{\text{had}}} \cdot f(q \rightarrow D_2^{*0}) \cdot \text{Br}(D_2^{*0} \rightarrow D^{*+}\pi^-) \quad (6.31)$$

where $\Gamma_{q\bar{q}}/\Gamma_{\text{had}}$ is the fraction of hadronic Z^0 decays which produce quark-antiquark pairs of flavour q . As the identification of charge conjugate D^{*0} states is implied, the factor of 2 accounts for D^{*0} or $\bar{D}^{*0}$ production from either the primary quark or antiquark.

Evaluating $f(q \rightarrow D_1^0)$ and $f(q \rightarrow D_2^{*0})$ requires knowledge of the branching ratios $\text{Br}(D_1^0 \rightarrow D^{*+}\pi^-)$ and $\text{Br}(D_2^{*0} \rightarrow D^{*+}\pi^-)$. As described in section 3.5.1, one expects from isospin considerations that $\text{Br}(D_1^0 \rightarrow D^{*+}\pi^-) = 0.65$, which includes the effect of $|\mathbf{p}_\pi|^5$ kinematic factors. The D_2^{*0} branching ratio can be calculated assuming its width to be saturated by two-body decays to $D^*\pi$ and $D\pi$ final states. Using the world average of

$$\frac{\text{Br}(D_2^{*+} \rightarrow D^+\pi^-)}{\text{Br}(D_2^{*+} \rightarrow D^{*+}\pi^-)} = 2.3 \pm 0.6 \quad (6.32)$$

and isospin considerations, one expects that $\text{Br}(D_2^{*0} \rightarrow D^{*+}\pi^-) = 0.21 \pm 0.04$ which is consistent with the value of 0.19, derived using heavy quark spin symmetry arguments. Finally, the partial hadronic widths of the Z^0 can be fixed using the Standard Model expectations $\Gamma_{c\bar{c}} = 0.172$ and $\Gamma_{b\bar{b}} = 0.216$.

Using these quantities, the model-dependent probabilities that primary charm quarks form D_1^0 and D_2^{*0} mesons can be computed:

$$f(c \rightarrow D_1^0) = 0.021 \pm 0.007 \pm 0.003 \quad (6.33)$$

$$f(c \rightarrow D_2^{*0}) = 0.052 \pm 0.022 \pm 0.013. \quad (6.34)$$

Their sum represents the total probability that a charm quark will form a neutral member of the $j_q = 3/2$ doublet:

$$f(c \rightarrow D^{*0}) = 0.073 \pm 0.023 \pm 0.014. \quad (6.35)$$

The corresponding probabilities for the production of narrow D^{**0} states in decays of b-hadrons are

$$f(b \rightarrow D_1^0) = 0.050 \pm 0.014 \pm 0.006 \quad (6.36)$$

$$f(b \rightarrow D_2^{*0}) = 0.047 \pm 0.024 \pm 0.012 \quad (6.37)$$

from which the total probability is obtained:

$$f(b \rightarrow D^{**0}) = 0.097 \pm 0.035 \pm 0.017. \quad (6.38)$$

From these measurements, one sees that D^{**0} production is actually not a rare event — it is only the small product of branching ratios which makes them difficult to observe. The comparison of these measurements with predictions of a recent fragmentation model is presented in chapter 8.

6.7 D^{**0} Production in Semileptonic B Decays

The interest in studying P-wave charm meson production in semileptonic B-decays, introduced in section 2.3.2, arose from the observation that the sum of the measured exclusive semileptonic branching ratios did not account for their total inclusive semileptonic decay rate. These semileptonic decays to P-wave charm mesons have already been observed at ARGUS [76], LEP [77, 78, 79, 80], and CLEO [81]. The measured rates indicate that one should expect that the analysis described in section 6.2.4 would also be sensitive to D^{**0} production in semileptonic b-decays. In this section, an analysis is presented which is very similar to the one described in section 6.2.4, but which requires that a track identified as either an electron or muon be included in the b-decay vertex.

6.7.1 Identification of Leptons

Tracks are identified as being electrons or muons, depending on the way they deposit energy in the detector. Electrons induce electromagnetic showers which are absorbed in the lead-glass calorimeter while muons deposit very little energy in the detector material and penetrate the iron plates of the hadron calorimeter to be detected in the muon chambers. These characteristic features are used in their identification.

Neural Network Electron Identification

OPAL uses a neural network [82] technique, tuned using Monte Carlo events, to identify electrons. This made use of twelve input parameters which include.

- energy loss (dE/dx) measured in the jet chamber.
- the energy of an associated electromagnetic cluster, divided by the track momentum (E/p),
- the geometric distribution of the energy deposit in the electromagnetic calorimeter,
- the angular difference between the electromagnetic cluster and the path of the track, extrapolated beyond the tracking chambers.

This technique has been used to identify electrons from the decays of b- and c-quarks, allowing precision measurements of the electroweak parameters of b- and c-quarks to be performed[83]. The output of the neural-network is a real variable, NET_{el} which ranges from 0 to 1 and is a measure of the probability that a given track is an electron.

Electrons in the momentum range $p > 2 \text{ GeV}/c$ are selected by the cut $NET_{el} > 0.95$. The performance of this technique was studied using photon conversions, $\gamma \rightarrow e^+e^-$, which provides an almost pure sample of electrons. The probability of misidentification of pions was also studied using pions from $K_S^0 \rightarrow \pi^+\pi^-$.

identified using similar techniques to those described in section 5.3. As shown in figure 6.13, electrons are identified with an efficiency of about 70% for tracks with polar angles in the range $|\cos \theta| < 0.7$, with lower efficiencies in the end-cap region of the detector. Also shown is the hadronic background in the sample which is estimated to be less than 12%.

Muon Identification

The selection of muons used in this analysis was performed using the same method developed by members of OPAL to measure $\Gamma_{b\bar{b}}/\Gamma_{\text{had}}$ [84]. Tracks with momentum greater than 2 GeV/c and polar angles in the range $|\cos \theta| < 0.9$ were identified as muons if their trajectory, extrapolated through the calorimeter and hadronic calorimeter, matched an independently reconstructed track-segment in the muon chambers. The energy loss of the track, measured in the jet-chamber, was also required to be consistent with a muon hypothesis. The selection efficiency for muons produced in semileptonic b-decay was about 76% and had a purity close to 88%.

6.7.2 Selection of $(D^{*+}\pi^-)\ell^\pm$ Candidates

The D^{*+} reconstruction described in section 5.1 was used in this analysis. Once a D^{*+} candidate was identified, the tracks in the event were assigned to jets, just as they were in the initial step of the inclusive b-vertex reconstruction. However, before proceeding, a lepton with momentum greater than 2 GeV/c was required to be identified in the D^{*+} jet. Then, the b-vertex reconstruction proceeded as before except this lepton was always assigned to the b-vertex.

After the vertex assignment of tracks in the D^{*+} jet had been made, b-vertex candidates were selected by requiring a displaced b-vertex and D^0 -vertex topology. As before, this was performed using cuts on the significance of the decay lengths, S and L , of the b- and D^0 -vertices, respectively. These decay length significances were required to satisfy $S/\sigma_S > 2$ and $L/\sigma_L > -2$, where $\sigma_S < 0.1$ cm, identical to

Figure 6.13: Neural-network electron identification efficiency (a) and purity (b) as functions of $|\cos \theta|$ for tracks with $p > 2 \text{ GeV}/c$ determined using OPAL data. The shaded area in (b) shows the hadronic background contribution, determined from a Monte Carlo simulation.

those used for the selection of inclusive b-vertices without the lepton assignment. Because of the reduced background associated with semileptonic b-decays, the momentum cut imposed on π^- candidates was relaxed to $p(\pi^-) > 1 \text{ GeV}/c$.

The mass difference distribution of $D^{*+}\pi^-$ combinations in this $D^{*+}\pi^-\ell^-X$ sample, for π^- hypotheses selected from the set of tracks assigned to the b-vertex is shown in figure 6.14. Figures 6.14(a) and 6.14(b) show this mass difference distribution for opposite sign, $D^{*+}\ell^-$, and same-sign $D^{*+}\ell^+$ combinations, respectively. The opposite sign correlation is expected from the $b \rightarrow c\ell^-\bar{\nu}$ transitions in weak decays of b-hadrons, while $D^{*+}\ell^+$ combinations result from random $D^{*+}\ell$ correlations or from misidentified leptons. A clear peak is observed in figure 6.14(a) at the mass difference corresponding to D_1^0 production, while little evidence for D_2^0 production in semileptonic b-decays is present.

6.7.3 Signal Fitting

The signal shown in figure 6.14 was fit using the methods described in section 6.3 but in this case, the D^{*+0} background from charm events was completely negligible. Also, due to the sparse statistics, the helicity distribution of the background was assumed to be flat and its parameterization was not allowed to vary in the fit. For convenience in setting a limit on D_2^0 production, the fitting procedure was modified so that the independent rates of D_1^0 and D_2^0 production were the independent parameters in the fit. In this way, the number of D_1^0 decays observed in this distribution was determined to be $17.3_{-5.6}^{+6.3}$ while the number of reconstructed D_2^0 decays was found to be consistent with zero.

The efficiency for the selection of $(D^{*+}\pi^-\ell^-)$ combinations was determined to be $\varepsilon = 0.0706 \pm 0.0085$, and from the fit to the mass difference distribution shown in figure 6.15, the mean multiplicity of semileptonic b-decays producing the $D^{*+0} \rightarrow D^{*+}\pi^-$ final state was obtained:

$$\bar{n}_{Z^0 \rightarrow b\bar{b} \rightarrow D_1^0 \ell^-} \cdot \text{Br}(D_1^0 \rightarrow D^{*+}\pi^-) = (1.24_{-0.40-0.13}^{+0.45+0.25}) \times 10^{-3}. \quad (6.39)$$

Figure 6.14: Mass difference distribution obtained from $D^{*+}\pi^-$ combinations from displaced vertices containing $D^{*+}\ell^-$ pairs (a) and $D^{*+}\ell^+$ pairs (b).

Figure 6.15: A fit to the $D^{*+}\pi^-$ mass difference distribution from vertices containing $D^{*+}\ell^-$ candidates.

which represents the average over $\ell = e$ and $\ell = \mu$ final states. This is quoted only for the D_1^0 state since no evidence for D_2^0 decays was observed. Therefore, only the corresponding limit was determined for D_2^0 production in semileptonic b-decays. At the 95% confidence level, the mean multiplicity of semileptonic b-decays producing $D_2^0 \rightarrow D^{*+}\pi^-$ can be excluded above

$$\bar{n}_{Z^0 \rightarrow b\bar{b} \rightarrow D_2^{*0} \ell^-} \cdot \text{Br}(D_2^{*0} \rightarrow D^{*+}\pi^-) < 1.00 \times 10^{-3}. \quad (6.40)$$

The systematic uncertainties associated with these measurements were determined using the same procedure described in section 6.5 and have been included in the 95% confidence limit on D_2^0 production. The systematic uncertainty due to the lepton selection was negligible.

Experiment	$\text{Br}_1 \times 10^{-3}$	$\text{Br}_2 \times 10^{-3}$
This analysis	$2.9^{+1.1+0.6}_{-1.0-0.3}$	< 2.4
ALEPH[79]	$1.68^{+0.40+0.36}_{-0.28-0.29}$	< 1.29
DELPHI[80]	$4.1 \pm 1.0 \pm 0.5$	
OPAL[77]	$3.1 \pm 0.7 \pm 0.7$	

Table 6.8: Comparison of LEP measurements of $\text{Br}_J \equiv f(b \rightarrow \bar{B}) \cdot \text{Br}(\bar{B} \rightarrow D_J^{(*)0} \ell^- \bar{\nu} X) \cdot \text{Br}(D_J^{(*)0} \rightarrow D^{*+} \pi^-)$.

The measured rates can be used to compute the following product of branching ratios for D_1^0 and D_2^{*0} production in b-decays:

$$\text{Br}_1 \equiv f(b \rightarrow \bar{B}) \cdot \text{Br}(\bar{B} \rightarrow D_1^0 \ell^- \bar{\nu} X) \cdot \text{Br}(D_1^0 \rightarrow D^{*+} \pi^-) = (2.9^{+1.1+0.6}_{-1.0-0.3}) \times 10^{-3} \quad (6.41)$$

$$\text{Br}_2 \equiv f(b \rightarrow \bar{B}) \cdot \text{Br}(\bar{B} \rightarrow D_2^{*0} \ell^- \bar{\nu} X) \cdot \text{Br}(D_2^{*0} \rightarrow D^{*+} \pi^-) < 2.4 \times 10^{-3}. \quad (6.42)$$

A comparison of these results with previous measurements is shown in table 6.8.

It can be seen that the results of this study are not as precise as those obtained from other analyses which were specifically designed to measure these branching ratios. Nevertheless, the results obtained here are in agreement with these measurements, providing additional support for the consistency of the analysis of inclusive D^{*+0} production in b-decays.

Chapter 7

D_{s1}^+ Production

A search for $D_{s1}^+(2536)$ production at LEP is another attractive analysis since this state has a very narrow width and decays near threshold. Although it can be fully reconstructed using its $D^{*+}K_S^0$ decay mode:

$$\begin{array}{lcl}
 D_{s1}^+ & \longrightarrow & D^{*+}(2010)K^0 \\
 & & \swarrow \quad \searrow \\
 & & K_S^0 \longrightarrow \pi^+\pi^- \\
 & & D^0\pi^+ \\
 & & \swarrow \\
 & & K^-\pi^+
 \end{array} \tag{7.1}$$

this is not necessarily the most attractive final state, since it is suppressed by a number of small branching ratios. Alternatively, the D_{s1}^+ can be observed using its $D^{*0}K^+$ decay mode by partially reconstructing the decay chain:

$$\begin{array}{lcl}
 D_{s1}^+ & \longrightarrow & D^{*0}(2007)K^+ \\
 & & \swarrow \\
 & & D^0(\pi^0, \gamma) \\
 & & \swarrow \\
 & & K^-\pi^+.
 \end{array} \tag{7.2}$$

In this analysis, the π^0 or γ from the D^{*0} decay is unobserved but because of the small D^{*0} - D^0 mass difference, a peak can still be observed near threshold in the $M(D^0 K^+) - M(D^0)$ mass difference distribution which is broadened only slightly by the incomplete reconstruction. This chapter describes an analysis in which D_{s1}^+ production was observed in hadronic Z^0 decays by reconstructing both decay modes (7.1) and (7.2).

7.1 Partial Reconstruction of $D_{s1}^+ \rightarrow D^{*0} K^+$

The D^{*0} decays to either $D^0 \pi^0$ or $D^0 \gamma$ and hence, the product of branching ratios for the decay shown in (7.2) is:

$$\text{Br}(D^{*0} \rightarrow D^0 \pi^0, \gamma) \cdot \text{Br}(D^0 \rightarrow K^- \pi^+) = \text{Br}(D^0 \rightarrow K^- \pi^+) = 0.0383 \pm 0.0012. \quad (7.3)$$

As described in section 5.2, the D^0 cannot be tagged by the presence of a transition pion, as was the case for $D^{*+} \rightarrow D^0 \pi^+$ decays, so a pure sample of D^0 's was obtained using tight cuts on dE/dx and on the momenta of the K^- and π^+ . In addition, a significant gain in purity was obtained by requiring $x_E(D^0) > 0.45$. The D^0 reconstruction efficiency for $x_E(D^0) > 0.45$ was about 18% but the efficiency for reconstructing the entire D_{s1}^+ decay chain was lower, due to the additional requirements imposed when selecting K^+ track candidates.

7.1.1 $D^0 K^+$ Selection

A large amount of combinatorial background can arise when pions are accepted as K^+ candidates. It is essential to reduce this background and in OPAL, the most effective way to do this is to use dE/dx information. Thus, both the K^- from the D^0 decay and the K^+ candidates were required to satisfy

$$-0.05 < w_K(K^\pm) < 0.15. \quad (7.4)$$

The remaining background from fake K^+ candidates peaked at low momenta and was suppressed by the requirement that $p(K^+) > 5.5 \text{ GeV}/c$. Finally, to suppress backgrounds further, and also to reduce the number of D_{s1}^+ mesons reconstructed in B decay, the remaining $D^0 K^+$ candidates were required to have scaled energies in excess of 0.57. Because of the missing neutral particle, this corresponds to $x_E(D_{s1}^+) > 0.6$ which was also the requirement imposed on the fully reconstructed D_{s1}^+ decays described later in section 7.2.1.

The efficiency for this selection was determined from a sample of 8000 Monte Carlo $D_{s1}^+ \rightarrow D^0 K^+$ decays generated using Jetset 7.4 and reconstructed using the GOPAL detector simulation. After D^0 reconstruction and the selection of an appropriate K^+ candidate, the overall D_{s1}^+ reconstruction efficiency in the accepted range of scaled energy was $\varepsilon = 9.0 \pm 0.7\%$.

7.1.2 Signal Properties

The D^{*0} decays to either $D^0 \pi^0$ or $D^0 \gamma$ and because neither the π^0 nor the γ were reconstructed, the mass difference obtained from the D^0 and the K^+ was broadened slightly. Ignoring detector resolution effects for the moment, the shape of mass difference distribution, $\Delta M(D^0 K^+) = M(D^0 K^+) - M(D^0)$, can be calculated in terms of the masses of the particles involved, assuming that the π^0 or γ is emitted isotropically in the D^{*0} rest frame. This gives the following kinematic limits for the $\Delta M(D^0 K^+)$ distribution:

$$524.39 < \Delta M < 532.06 \text{ MeV}/c^2 \quad D^{*0} \rightarrow D^0 \pi^0 \quad (7.5)$$

$$517.03 < \Delta M < 541.41 \text{ MeV}/c^2 \quad D^{*0} \rightarrow D^0 \gamma. \quad (7.6)$$

and to a good approximation, the mass difference distribution has the form of two uniform distributions lying between these kinematic limits, with areas weighted by the D^{*0} branching ratios. This is shown in figure 7.1, along with the expected shape after smearing with a nominal detector resolution of approximately $4 \text{ MeV}/c^2$.

Figure 7.1: Mass difference distribution for partially reconstructed D_{s1}^+ decays (a) without detector resolution effects and (b) convoluted with the expected detector resolution of $4 \text{ MeV}/c^2$.

The D_{s2}^{*+} can decay directly to $D^0 K^+$ but since it is farther from threshold, it is a broader resonance with a width of $\Gamma = 15_{-4}^{+5} \text{ MeV}/c^2$. Detector resolution will further degrade its peak in the mass difference distribution and one does not expect to be able to measure a significant D_{s2}^{*+} production rate with the current OPAL data sample.

7.1.3 Monte Carlo Analysis

After the selection described in section 7.1.1 was applied, only a few hundred $D^0 K^+$ combinations with $\Delta M(D^0 K^+) < 0.8 \text{ GeV}/c^2$ remained. Maximal use was made of the available statistics by fitting the distribution using an unbinned maximum likelihood method, as was performed in the analysis of D^{*+0} production. In this case, the probability density distribution for the D_{s1}^+ signal was obtained by convoluting the raw mass difference distribution shown in figure 7.1(a) with the expected detector resolution. Ignoring the possible contribution from the D_{s2}^{*+} state, the

likelihood function was of the form

$$\mathcal{L} = \prod_i (f_1 s_1(x_i) + (1 - f_1) b(x_i)) \quad (7.7)$$

where the background parameterization $b(x_i)$ was of the form

$$b(x) \propto e^{-B(x-m_K)}(x - m_K)^C \quad (7.8)$$

in which m_K is the charged kaon mass. As described in section 7.1.2, the signal function, $s_1(x)$ was obtained by convoluting the uniform distributions allowed by the D^{*0} decay kinematics with the Gaussian resolution. This was expressed as follows:

$$s_1(x) = \frac{\text{Br}(D^{*0} \rightarrow D^0 \pi^0)}{2(m_1^{\pi^0} - m_0^{\pi^0})} \left(\text{erf} \left(\frac{x - m_1^{\pi^0}}{\sigma \sqrt{2}} \right) - \text{erf} \left(\frac{x - m_0^{\pi^0}}{\sigma \sqrt{2}} \right) \right) + \frac{\text{Br}(D^{*0} \rightarrow D^0 \gamma)}{2(m_1^{\gamma} - m_0^{\gamma})} \left(\text{erf} \left(\frac{x - m_1^{\gamma}}{\sigma \sqrt{2}} \right) - \text{erf} \left(\frac{x - m_0^{\gamma}}{\sigma \sqrt{2}} \right) \right) \quad (7.9)$$

where $m_0^{\pi^0}$, $m_1^{\pi^0}$, m_0^{γ} and m_1^{γ} are the lower and upper kinematic limits for the two decay modes, shown in equations 7.5 and 7.6, and σ is the mass difference resolution.

To demonstrate that this procedure yielded unbiased estimates of the number of D_{s1}^+ events in the distribution, the analysis was performed on a Monte Carlo sample of 5.8 million hadronic Z^0 decays. After the $D^0 K^+$ selection was applied, 27 D_{s1}^+ decays remained. The fit to the mass difference distribution for this sample is shown in figure 7.2 from which the fitted number of D_{s1}^+ decays was estimated to be 28.5 ± 8.1 , consistent with the number actually present. The expected contribution from D_{s2}^{*+} decays is also shown in this figure and one can see that neglecting its contribution in the likelihood fit cannot influence the measured D_{s1}^+ rate.

It was also important to demonstrate that the reconstructed signal had the expected properties. Thus, the fit to the Monte Carlo distribution was also per-

Figure 7.2: A fit to the $D^0 K^+$ mass difference distribution from 5.8 million Monte Carlo hadronic Z^0 decays. The dashed line indicates the shape of the fitted background function under the signal. Hatched entries under the signal indicate correctly reconstructed D_{s1}^+ decays, while those at higher mass differences indicate the contribution from $D_{s2}^{*+} \rightarrow D^0 K^+$ decays.

Figure 7.3: $D^0 K^+$ mass difference distribution from OPAL's full 1990-95 data set.

formed with the signal described by a Gaussian function. Allowing the position and width to vary freely yielded a mass difference of $525.7 \pm 2.1 \text{ MeV}/c^2$, consistent with the generated value of $528.3 \text{ MeV}/c^2$ and a width of 6.0 ± 2.2 consistent with the expectation of $7 \text{ MeV}/c^2$ derived from the high statistics Monte Carlo sample of D_{s1}^+ decays. Thus, this Monte Carlo analysis indicates that the reconstruction and fitting procedure provided reliable estimates of the D_{s1}^+ production rate.

7.1.4 Analysis of OPAL Data

Figure 7.3 shows the mass difference distribution obtained when the selection described in section 7.1.1 was applied to the full OPAL data sample. From the maximum likelihood fit, the number of D_{s1}^+ decays was determined to be $N(D_{s1}^+) = 28.7 \pm 8.3$. Although this was obtained with the masses fixed to world averages and the resolution estimated from Monte Carlo, a separate fit using a Gaussian signal converged to a mass difference of $527.3 \pm 2.2 \text{ MeV}/c^2$ and a width of $5.6 \pm 2.2 \text{ MeV}/c^2$. Again, these values were consistent with those expected for

$D_{s1}^+ \rightarrow D^{*0}K^+$ decays and, as expected from the Monte Carlo study, no significant resonant structure from D_{s2}^{*+} decays was observed.

The mean multiplicity of D_{s1}^+ mesons produced per hadronic Z^0 decay was determined from:

$$\bar{n}_{Z^0 \rightarrow D_{s1}^+} \cdot \text{Br}(D_{s1}^+ \rightarrow D^{*0}K^+)_{x_E > 0.6} = \frac{N(D_{s1}^+)}{N_{\text{MH}} \cdot \varepsilon \cdot \text{Br}(D^0 \rightarrow K^-\pi^+)} \quad (7.10)$$

where N_{MH} is the number of hadronic events analysed and ε is the D_{s1}^+ reconstruction efficiency. Thus, the 28.7 ± 8.3 events obtained from the fit correspond to

$$\bar{n}_{Z^0 \rightarrow D_{s1}^+} \cdot \text{Br}(D_{s1}^+ \rightarrow D^{*0}K^+)_{x_E > 0.6} = 0.0019 \pm 0.0005. \quad (7.11)$$

In this expression, the quoted uncertainty is statistical only. The determination of the corresponding systematic uncertainty is discussed in the following section.

7.1.5 Systematic Studies

The systematic effects which could influence the measurement shown in equation (7.11) were evaluated using techniques similar to those described for the analysis of $D^{*0} \rightarrow D^{*+}\pi^-$. A summary of the individual contributions to the systematic uncertainty is shown in table 7.1 and a description of each component follows.

Monte Carlo Statistics

For this measurement, the statistical uncertainty in the reconstruction efficiency was included with the sources of systematic uncertainty, as was also done for the $D^{*0} \rightarrow D^{*+}\pi^-$ measurements.

Source	Uncertainty
Monte Carlo statistics	± 0.00015
$\text{Br}(D^0 \rightarrow K^- \pi^+)$	± 0.00006
Signal parameterization	$+0.00019$ -0.00021
Background parameterization	$+0.00010$ -0.00004
Tracking resolution	$+0.00005$ -0.00005
Monte Carlo b/c efficiency	$+0.00003$ -0.00004
Monte Carlo $D^{*0} \varepsilon \cdot \text{Br}$	± 0.00002
Total	± 0.00028

Table 7.1: Summary of systematic uncertainties for the rate determined from the $D_{s1}^+ \rightarrow D^{*0} K^+$ channel.

Branching Ratios

The uncertainty in the world average of $\text{Br}(D^0 \rightarrow K^- \pi^+)$ has a small contribution to the uncertainty in the D_{s1}^+ production rate in this channel. However, this branching ratio is determined with a precision of approximately 3% and does not limit this measurement.

Signal Parametrization

With limited statistics, it is difficult to determine whether the treatment of the D_{s1}^+ signal in the mass difference distribution is appropriate. Although motivated by the kinematic properties of $D_{s1}^+ \rightarrow D^{*0} K^+$ decays, it is important to demonstrate that the measured rate does not depend critically on these assumptions. The effects of a possible misdetermination of the mass difference resolution were investigated by repeating the fits with the resolution scaled by $\pm 50\%$. Also, the signal was fit using a pure Gaussian with mass and width allowed to vary freely. This formed the largest contribution to the overall systematic uncertainty.

Background Parametrization

The mass difference distribution was fitted using a background function of the form $e^{-B(x-m_K)}(x-m_K)^C$. Although this is seen to describe the background shape well, especially near threshold, it is appropriate to study different parameterizations. Two additional parameterizations which contain the expected threshold behaviour were studied. First, the Weibull function was used which has the following expression when normalized to unit area between m_K and $m_{\max}$:

$$W(x) = \frac{\left(\frac{x-m_K}{B}\right)^{C-1} \exp\left(-\left(\frac{x-m_K}{B}\right)^C\right)}{B/C(1 - \exp(-((m_{\max} - m_K)/B)^C))}. \quad (7.12)$$

Also, a second order polynomial in $(x - m_K)$ multiplied by a square root threshold factor, as well as a constant background function were studied. The variation in the measured rate due to these different treatments is shown in table 7.1 and was found to be relatively small.

Tracking Resolution

Systematic uncertainties due to the imperfect modelling of the track parameters were estimated by redetermining the efficiencies from Monte Carlo in which the track parameter resolution was scaled by $\pm 10\%$. This had a small effect on the measured D_{s1}^+ production rate in this channel which is reflected in the systematic uncertainty shown in table 7.1.

Monte Carlo b/c Efficiency

Because of the requirement that $x_E(D^0 K^+) > 0.57$, the fraction of D_{s1}^+ decays reconstructed in $b\bar{b}$ events was expected to be small. However, unlike the study of D^{*0} production, the D_{s1}^+ signal is not observed separately in a b-enhanced sample so in this case, the fraction of D_{s1}^+ decays from b-events was estimated from Monte Carlo. Since the efficiency for reconstructing $D_{s1}^+ \rightarrow D^{*0} K^+$ decays in $c\bar{c}$ and in $b\bar{b}$

events may differ, the measured D_{s1}^+ production rate depends on the fraction of $b\bar{b}$ events present.

To determine the dependence of the measured $D_{s1}^+ \rightarrow D^{*0}K^+$ rate on the fraction of decays in $b\bar{b}$ events, the efficiencies were redetermined with the b-fraction varied by $\pm 50\%$. Since the difference in efficiencies was less than 2.5% and the fraction of b-events is small, this effect was found to give only a very small contribution to the overall systematic uncertainty.

D^{*0} Decay Modes

It can be seen from figure 7.1 that the $D^{*0} \rightarrow D^0\pi^0$ and $D^{*0} \rightarrow D^0\gamma$ decay modes produce different shaped signals in the D^0K^+ mass difference distribution. Differences in the kinematics of these two decay modes can result in different efficiencies for D_{s1}^+ reconstruction depending on the D^{*0} decay mode. To estimate the size of this contribution to the systematic uncertainty, the efficiencies were calculated for each D^{*0} decay mode separately and the weighted mean was redetermined with the D^{*0} branching ratios, determined from world averages, varied by their uncertainty. This was also found to introduce an almost negligible systematic variation in the measured rate.

Cross Checks

The same discussion of the variation of selection criteria presented in section 6.5.10 applies to this measurement. To determine whether the result shown in equation (7.11) was stable under changes in the selection criteria, the following cuts were modified:

- The minimum x_{D^0} was varied from 0.40 to 0.55.
- The minimum $p(K^+)$ was varied from 4.0 to 6.5 GeV/c.
- The minimum $p(K^-)$ and $p(\pi^+)$ were varied from 5.0 to 8.0 GeV/c.

- dE/dx weights for kaon identification were varied from ± 0.05 to ± 0.15 .
- The minimum D^0 decay length significance was varied from 0 to 1.

In all cases, the changes in the measured rate were consistent with statistical fluctuations. Thus, no contribution to the systematic uncertainty was attributed to the selection criteria.

7.1.6 Results

Including the systematic uncertainty, the mean D_{s1}^+ multiplicity per hadronic Z^0 decay was determined to be

$$\bar{n}_{Z^0 \rightarrow D_{s1}^+} \cdot \text{Br}(D_{s1}^+ \rightarrow D^{*0}K^+)_{x_E > 0.6} = 0.0019 \pm 0.0005 \text{ (stat)} \pm 0.0003 \text{ (syst)}. \quad (7.13)$$

The fraction of charm quarks which fragment to form D_{s1}^+ mesons could be determined from this quantity but this would require assuming a value for the $D_{s1}^+ \rightarrow D^{*0}K^+$ branching ratio. Because the D_{s1}^+ width is expected to be saturated by decays to $D^{*0}K^+$ and $D^{*+}K^0$, one would expect this branching ratio to be fixed by isospin and kinematic factors. However, a measurement of D_{s1}^+ production in both decay channels would allow a model independent measurement of the D_{s1}^+ production rate to be made. The reconstruction of D_{s1}^+ mesons which decay to the $D^{*+}K_S^0$ final state is described in the next section.

7.2 Full Reconstruction of $D_{s1}^+ \rightarrow D^{*+}K_S^0$

Although the partial reconstruction technique benefited from increased statistics, the full reconstruction of the D_{s1}^+ decay chain has other advantages. The fact that all particles in the final state can be reconstructed results in an improved signal resolution which can compensate, to some degree, for the smaller branching ratios involved.

The product of branching fractions for decay sequence shown in equation (7.1) is

$$\frac{1}{2}\text{Br}(D^{*+} \rightarrow D^0\pi^+)\cdot\text{Br}(D^0 \rightarrow K^+\pi^-)\cdot\text{Br}(K_S^0 \rightarrow \pi^+\pi^-) = 0.00897 \pm 0.00034 \quad (7.14)$$

which is less than 1/4 that of the corresponding product of branching fractions for the $D^{*0}K^+$ decay mode. Furthermore, as described in section 3.5.1, it is expected that kinematic factors will reduce this D_{s1}^+ branching fraction by a factor which can be as small as 1/2 due to the smaller mass difference between the initial and final states. Therefore, one only expects to reconstruct a small number of D_{s1}^+ mesons decaying to the $D^{*+}K_S^0$ final state.

7.2.1 $D^{*+}K_S^0$ Selection

Because of the small product of branching ratios in this final state, the selection was motivated by the need to retain a high reconstruction efficiency. The reconstruction of D^{*+} and K_S^0 mesons was described in sections 5.1 and 5.3, and the only additional requirements imposed were that the D^{*+} mass difference lie within $0.143 < \Delta M(D^0\pi^+) < 0.148 \text{ GeV}/c^2$ and that K_S^0 candidates have their $\pi^+\pi^-$ invariant mass in the range $0.45 < M(\pi^+\pi^-) < 0.54 \text{ GeV}/c^2$. D^{*+} and K_S^0 mesons which satisfy these criteria were selected if $x_E(D^{*+}K_S^0) > 0.6$ and for this region of scaled energy, the reconstruction efficiency was $10.9 \pm 1.0\%$. This efficiency included the $K_S^0 \rightarrow \pi^+\pi^-$ branching ratio since K_S^0 mesons sometimes interact hadronically in the detector material before they decay.

7.2.2 Signal Properties

Only upper limits have been set on the natural width of the D_{s1}^+ meson and the shape of the reconstructed signal should be described by a Gaussian distribution with a width determined by the resolution of the detector. Figure 7.4(a) shows this mass difference resolution calculated using $\Delta M(D^{*+}K_S^0) = M(D^{*+}K_S^0) - M(D^{*+})$

Figure 7.4: Mass difference resolutions obtained using (a) $\Delta M(D^{*+}K_S^0) = M(D^{*+}K_S^0) - M(D^{*+})$ and (b) $M(D^{*+}K_S^0) - M(D^{*+}) - M(K_S^0) + M_{K_S^0}^{P.D.G.}$. The distributions are fitted using double Gaussians and widths of the narrow component are indicated.

but a significant improvement was obtained by also subtracting the reconstructed K_S^0 mass. This measure of the $D^{*+}K_S^0$ mass difference is shown in 7.4(b) which is calculated using

$$\Delta M(D^{*+}K_S^0) = M(D^{*+}K_S^0) - M(D^{*+}) - M(K_S^0) + M_{K_S^0}^{P.D.G.} \quad (7.15)$$

where the world average K_S^0 mass was added to restore the familiar mass difference scale. The mass difference determined in this way had a resolution of $3.2 \pm 0.4 \text{ MeV}/c^2$. No contribution from D_{s2}^{*+} decays was expected in this analysis because of its larger width and since its branching ratio to $D^{*+}K_S^0$ is expected to be only $\sim 5\%$.

To demonstrate that this method is actually sensitive to the expected levels of D_{s1}^+ production in multihadronic Z^0 decays, the selection was applied to 4 million Monte Carlo events. The resulting mass difference distribution is shown in figure 7.5 and was fit using the likelihood function shown in equation (7.7) with

Figure 7.5: The $D^{*+}K_S^0$ mass difference distribution resulting from the analysis of 4 million Monte Carlo hadronic Z^0 decays. The hatched entries under the fitted Gaussian indicate correctly reconstructed $D_{s1}^+ \rightarrow D^{*+}K_S^0$ decays. The background level under the signal is indicated with the dashed curve.

the same background parameterization but with the signal described by a single Gaussian with its width fixed to the expected resolution. The fitted number of D_{s1}^+ decays was 6.0 ± 2.7 , consistent with the 6 $D_{s1}^+ \rightarrow D^{*+}K_S^0$ decays which were actually present in the region consistent with smearing due to detector resolution.

7.2.3 Rate Measurement

The same selection was applied to the full OPAL data sample resulting in the mass difference distributions shown in figure 7.6. Approximately 6 entries fell in a mass difference region consistent with D_{s1}^+ production. To fit this distribution a number of background parameterizations were studied but the fact that no $D^{*+}K_S^0$ candidates had mass differences below the D_{s1}^+ signal caused all to have unexpected threshold behaviours. Without entries near the kinematic limit of $\Delta M = m_{K^0}$.

Figure 7.6: Fits to the $D^{*+}K_S^0$ mass difference distribution from the full OPAL data set of 4.45 million hadronic Z^0 decays. The dashed curve indicates the background level under the signal.

these functions could potentially underestimate the background level under the signal. Therefore, the background shape was constrained to that determined from the fit to the Monte Carlo distribution which did have the expected threshold behaviour. The fitted curve shown in figure 7.6 is the result of this procedure.

The fitted number of $D_{s1}^+ \rightarrow D^{*+}K_S^0$ decays was determined to be $N(D_{s1}^+ \rightarrow D^{*+}K_S^0) = 5.9^{+2.8}_{-2.3}$. This corresponds to a mean multiplicity per hadronic Z^0 decay of

$$\bar{n}_{Z^0 \rightarrow D_{s1}^+} \cdot \text{Br}(D_{s1}^+ \rightarrow D^{*+}K^0)_{x_E > 0.6} = 0.0010^{+0.0005}_{-0.0004} \quad (7.16)$$

where only the statistical uncertainty is shown.

7.2.4 Systematic Uncertainties

The systematic error associated with this measurement is estimated in much the same way as was performed for the rate measured in the $D^{*0}K^+$ final state. The individual contributions are listed in table 7.2 and discussed further in this section.

Source	Uncertainty
Signal parameterization	+0.00001 -0.00000
Background parameterization	+0.00005 -0.00002
Tracking resolution	+0.00003 -0.00002
Monte Carlo statistics	± 0.00009
$\text{Br}(D^{*+} \rightarrow D^0 \pi^+) \cdot \text{Br}(D^0 \rightarrow K^- \pi^+)$	± 0.00004
Total	+0.00011 -0.00009

Table 7.2: Summary of systematic uncertainties for the rate determined from the $D_{s1}^+ \rightarrow D^{*+} K_S^0$ channel.

Including the systematic uncertainty, the result is

$$\bar{n}_{Z^0 \rightarrow D_{s1}^+} \cdot \text{Br}(D_{s1}^+ \rightarrow D^{*+} K^0)_{x_E > 0.6} = 0.0010_{-0.0004}^{+0.0005} (\text{stat}) \pm 0.0001 (\text{syst}). \quad (7.17)$$

with the error clearly dominated by the statistical uncertainty.

Systematic uncertainties arising from intermediate branching ratios and from the statistical uncertainty in the reconstruction efficiency have been described in sections 6.5 and 7.1.5 for the $D^{*-0} \rightarrow D^{*+} \pi^-$ and $D_{s1}^+ \rightarrow D^{*0} K^+$ analyses and do not warrant further discussion for this measurement. However, the variation of the $D^{*+} K_S^0$ selection criteria and the fitting of the mass difference distribution are sufficiently different that some discussion is worthwhile.

Signal Parametrization

Although the $D_{s1}^+ \rightarrow D^{*+} K_S^0$ signal was described well by a Gaussian, its width could not be allowed to vary as a free parameter due to the sparse statistics in the mass difference distribution. Instead, the width was fixed to a value expected from Monte Carlo studies but the accuracy of this determination was difficult to estimate. In this analysis, the D_{s1}^+ signal was fit to a Gaussian distribution with a width of $3.2 \text{ MeV}/c^2$, but this was varied between 2.8 and $6.3 \text{ MeV}/c^2$. This range was determined by reducing the width of the narrow component of the signal in

figure 7.4(b) by its uncertainty and by increasing it to the weighted mean of the widths of the narrow and wide components. It is believed that this range is large enough to accommodate the actual resolution obtained in the analysis. In any case, because of the low level of background in the mass difference distribution, the variation of the width has an almost negligible effect on the fitted number of $D_{s1}^+ \rightarrow D^{*+} K_S^0$ decays.

Background Parametrization

The sparse statistics in this analysis made it difficult to determine the correct shape with which to describe the background. As described above, one expects there to be some threshold behaviour present at the kinematic limit of the mass difference distribution, but the lack of any $D^{*+} K_S^0$ combinations with mass differences less than that obtained for the D_{s1}^+ signal results in any freely varying background distributions underestimating the expected background in this region. Thus, the shape of the background was constrained to that expected from the fit to the Monte Carlo distribution which did exhibit the expected threshold behaviour.

It is important to check that the measured rate does not depend too critically on this somewhat subjective treatment of the background. This was done by repeating the fits to the mass difference distribution using the standard background parameterization of $e^{-B(x-m_{K^0})(x-m_{K^0})^C}$, a Weibull function and a second order polynomial with a square-root threshold behaviour. In these studies, all background parameters were allowed to vary freely and consequently tended to increase the measured rate. Nevertheless, the details of the background parameterization did not introduce a dominant systematic error.

Tracking Resolution

Again, the effect of imperfect simulation of the track parameter resolution was studied by redetermining the reconstruction efficiencies from Monte Carlo which had track parameter resolutions scaled by $\pm 10\%$. The resulting contribution to

the systematic uncertainty is shown in table 7.2.

Variation of Selection Criteria

Because so few D_{s1}^+ decays were reconstructed in this channel, it was important to determine whether the signal was merely a statistical fluctuation. The stability of the measured rate was studied by varying the following criteria:

- The width of the mass difference window used to select D^{*+} candidates was varied between 4 and 6 MeV/c^2 .
- The width of the $\pi^+\pi^-$ mass window used to select K_S^0 candidates was varied between 50 and 130 MeV/c^2 .
- The requirement that the K_S^0 tracks have hits in the z -chambers or the jet-chamber end-plate was removed.

Again, the sizes of fluctuations in the measured rate as these cuts were modified were found to be small compared to the statistical uncertainty and no systematic error was assigned to the selection criteria.

7.3 Interpretation of Results

Since the width of the D_{s1}^+ meson is expected to be saturated by the $D^{*0}K^+$ and $D^{*+}K^0$ final states, the mean D_{s1}^+ multiplicity per hadronic Z^0 decay for $x_E > 0.6$ is the sum of the two measurements shown in equations (7.13) and (7.17):

$$\bar{n}_{Z^0 \rightarrow D_{s1}^+ (x_E > 0.6)} = 0.0029_{-0.0006}^{+0.0007} (\text{stat}) \pm 0.0003 (\text{syst}). \quad (7.18)$$

Because of the $x_E > 0.6$ requirement, this measurement is dominated by D_{s1}^+ production in charm events, but since no explicit separation of $b\bar{b}$ and $c\bar{c}$ components was performed, this rate also includes a small contribution from $b\bar{b}$ events. In the absence of a separate measurement of D_{s1}^+ production in a $b\bar{b}$ enhanced sample, this

contribution was determined from the Monte Carlo model. This study indicated that D_{s1}^+ production in $b\bar{b}$ events should contribute about four events in the $D^{*0}K^+$ channel and one in the $D^{*+}K_S^0$ channel. However, because of the inherent model dependence, these numbers were varied by $\pm 50\%$ to estimate the corresponding systematic uncertainty. Subtracting this estimated contribution from $b\bar{b}$ events gave the following production rate in charm fragmentation:

$$\bar{n}_{Z^0 \rightarrow c\bar{c} \rightarrow D_{s1}^+ (x_E > 0.6)} = 0.0025_{-0.0006}^{+0.0007} (\text{stat}) \pm 0.0002 (\text{syst}) \pm 0.0002 (\text{b-frac}) \quad (7.19)$$

where the last error indicates the uncertainty in the $b\bar{b}$ contribution, estimated by varying the D_{s1}^+ rate in $b\bar{b}$ events by $\pm 50\%$.

To obtain the probability that a charm quark will produce a D_{s1}^+ meson in the fragmentation process, equation 7.19 can be scaled by the fraction of D_{s1}^+ mesons expected to be produced in the observed range of scaled energy. Again, this extrapolation of the production rate to lower scaled energies required input from the Monte Carlo fragmentation model and introduced a further systematic uncertainty. This extrapolation gives

$$\bar{n}_{Z^0 \rightarrow c\bar{c} \rightarrow D_{s1}^+} = 0.0056_{-0.0013}^{+0.0015} (\text{stat}) \pm 0.0006 (\text{syst}) \pm 0.0008 (\text{frag}) \quad (7.20)$$

where the second systematic error contains the uncertainty from both the fraction of $b\bar{b}$ events and from the extrapolation to lower scaled energies, estimated using the same procedure described in section 6.6. From this quantity, the probability that a charm quark will produce a D_{s1}^+ meson as it fragments can be calculated using

$$\bar{n}_{Z^0 \rightarrow c\bar{c} \rightarrow D_{s1}^+} = 2 \cdot \frac{\Gamma_{c\bar{c}}}{\Gamma_{\text{had}}} \cdot f(c \rightarrow D_{s1}^+). \quad (7.21)$$

Thus, using $\Gamma_{c\bar{c}}/\Gamma_{\text{had}} = 0.172$, the following probability was obtained:

$$f(c \rightarrow D_{s1}^+(2536)) = 0.016 \pm 0.004 \pm 0.003 \quad (7.22)$$

where all systematic uncertainties have been combined to give the final error.

Chapter 8

Discussion

The production of D^{*-0} and D_{s1}^+ mesons has now been observed using the OPAL data sample. It is now appropriate to place the measurements of their production rates in context by comparing them with other measurements and with recent predictions from fragmentation models.

8.1 Comparison with Fragmentation Models

The quark fragmentation simulated by Jetset is an essential tool for many analyses at LEP. Nevertheless, it is still only an *ad-hoc* model of fragmentation containing a very large number of parameters which can be adjusted to modify its behaviour. While these can be tuned to reproduce experimentally observed distributions their physical significance can be difficult to establish. But it is not necessarily appropriate to compare measured distributions with the predictions of Jetset, except in this context of tuning parameters, since a primary goal of Jetset is to accommodate, rather than to predict, the properties of hadronic events.

An alternative approach to constructing fragmentation models is to simplify the description such that only the most important ideas are included. This has been the motivation for a number of models which attempt to describe hadron production rates in e^+e^- annihilation using as few parameters as possible. One

model in particular [85] expresses the number of hadrons of a specific type per Z^0 decay as

$$\langle N \rangle = C \cdot \frac{2J+1}{C_B} \cdot (\gamma_s)^{N_s} \cdot e^{-E_{\text{bind}}/T} \quad (8.1)$$

where J and E_{bind} are the spin and binding energy of the hadron, γ_s is a suppression factor for the N_s strange-quarks in the hadron, C_B is a separate normalization for baryons and T is an effective temperature in the hadronization process.

The suppression of strange quark production over that of light u and d quarks is expected because of its larger quark mass. In the string fragmentation model it can be related to the effective string tension κ by

$$\gamma_s = \exp(-\pi(m_s^2 - m_u^2)/\kappa). \quad (8.2)$$

assuming u and d quarks to have the same mass. Many estimates of the value of γ_s have been obtained at LEP and other accelerators by comparing the ratio of production rates of hadrons which differ only by the flavour of one quark. All yield values close to $\gamma_s \approx 0.3$.

The different normalization for baryon production rates, introduced via C_B which is defined to be unity for mesons, is expected due to the different way the quark-antiquark pairs produced in the fragmentation process must be rearranged to form baryons or mesons. The effective temperature T arises from the idea that fragmentation represents a thermodynamic process in equilibrium. This last assumption is rather suspect since the densities and energies involved indicate that fragmentation occurs far from equilibrium, but nevertheless, it motivates a method by which the hadron binding energy can be introduced.

The binding energy, $E_{\text{bind}} = M_h - \sum_i m_{q_i}$, of a hadron, h , introduces a sensitivity to the mass difference between strange and light quarks. Hence, the parameters γ_s , $\Delta m = m_s - m_u$, T , C and C_B can be varied so as to fit the production rates of light mesons and baryons measured at LEP. This procedure was carried out in [85] and the fitted parameters determined in this way were used to predict the

production rates of hadrons containing heavy quarks. The predictions which can be compared with the measurements described in this thesis are as follows:

$$\langle (D_1^0 + D_2^{*0}) \rangle_{c\text{-jet}} = 0.085 \quad (8.3)$$

$$\langle D_{s1}^+ + D_{s2}^{*+} \rangle_{c\text{-jet}} = 0.028. \quad (8.4)$$

In principle, these predictions have uncertainties associated with them which arise from the fitted parameters of the model. These uncertainties were not quoted. Although it is desirable to compare experimental results with predictions from a range of models, such a treatment is difficult in this case since the content of other models is already included in equation 8.1. Nevertheless, the predicted rate of D^{*0} production can be compared directly with $f(c \rightarrow D^{*0}) = 0.073 \pm 0.023 \pm 0.013$ from equation (6.35), and is seen to be in reasonable agreement.

It is more difficult to make a comparison of the measured rate of D_{s1}^+ production since only the $J = 1$ state has been observed in this analysis. However, equation (8.1) can be used to predict the ratio of the D_{s2}^{*+} and D_{s1}^+ production rates as follows:

$$\frac{\langle N_{D_{s2}^{*+}} \rangle}{\langle N_{D_{s1}^+} \rangle} = \frac{5}{3} \exp \left(- (M_{D_{s2}^{*+}} - M_{D_{s1}^+}) / T \right) = 1.46. \quad (8.5)$$

This was evaluated using the parameter $T = 0.29$ GeV, determined from the fits to the production rates of light hadrons. Thus, the rate of D_{s1}^+ production is predicted to be

$$\langle D_{s1} \rangle_{c\text{-jet}} = 0.011. \quad (8.6)$$

The measured rate from equation (7.22) of $f(c \rightarrow D_{s1}^+) = 0.016 \pm 0.004 \pm 0.003$ is seen to be somewhat larger than this prediction, but is still statistically consistent.

A determination of the strange-quark suppression factor for spin-1 P-wave charm mesons can be made from the ratio of equations (7.22) and (6.33). This ratio is

$$\frac{f(c \rightarrow D_{s1}^+)}{f(c \rightarrow D_1^0)} = 0.76_{-0.32}^{+0.35} \pm 0.20 \quad (8.7)$$

which, again, is seen to be somewhat higher than, but still consistent with, the expected value of $\gamma_s \approx 0.3$.

8.2 Comparisons with P-wave B meson Production

OPAL[86], ALEPH[87] and DELPHI[88] have all observed the production of P-wave B mesons in hadronic Z^0 decays. These mesons are composed of a heavy b-quark combined with a light $\bar{u}$, $\bar{d}$ or $\bar{s}$ antiquark with one unit of orbital angular momentum. These analyses involved the reconstruction of an inclusive sample of displaced vertices from which the charge of the decaying b-hadron was determined on a statistical basis. Using this technique, OPAL, ALEPH and DELPHI reconstructed samples containing ~ 2000 b-decay vertices out of their data collected between 1991 and 1994. Since P-wave B hadrons decay strongly, these vertices are paired with tracks originating from the primary vertex and a resonant enhancement is observed near threshold in $B^+\pi^-$ and B^+K^- combinations.

OPAL measured

$$\frac{\text{Br}(Z^0 \rightarrow \bar{b} \rightarrow B^{*-0} \rightarrow B^{(*)+}\pi^-)}{\text{Br}(Z^0 \rightarrow \bar{b} \rightarrow B^+)} = 0.18 \pm 0.04 \quad (8.8)$$

while ALEPH and DELPHI presented their measurements with the branching ratio $\text{Br}(B_{u,d}^{*-} \rightarrow B_{d,u}\pi^\pm) = 2/3$ factored out:

$$\frac{\text{Br}(Z^0 \rightarrow \bar{b} \rightarrow B_{u,d}^{*-})}{\text{Br}(Z^0 \rightarrow \bar{b} \rightarrow B_{u,d})} = \begin{cases} 0.279 \pm 0.016^{+0.071}_{-0.081} & \text{ALEPH} \\ 0.27 \pm 0.02 \pm 0.06 & \text{DELPHI} \end{cases} \quad (8.9)$$

These are both quite consistent with OPAL's measurement. In addition, OPAL studied BK flavour correlations from which the measurement

$$\frac{\text{Br}(Z^0 \rightarrow \bar{b} \rightarrow B_s^{*-0} \rightarrow B^{(*)+}K^-)}{\text{Br}(Z^0 \rightarrow \bar{b} \rightarrow B^+)} = 0.026 \pm 0.008 \quad (8.10)$$

was obtained. These rates are to be understood as a superposition of both members of the $j_q = 3/2$ doublet of P-wave B mesons since the widths and experimental resolution are such that individual states are not resolved.

It is interesting to compare these results with the production rates of P-wave charm mesons, described in this thesis. Using $\text{Br}(B^{*-0} \rightarrow B^{(*)+}\pi^-) = 2/3$, $\text{Br}(B_s^{*-0} \rightarrow B^{(*)+}K^-) = 1/2$ and $f(\bar{b} \rightarrow B) = 0.378 \pm 0.022$ [12], these OPAL measurements imply

$$f(\bar{b} \rightarrow B^{*-0}) = 0.10 \pm 0.02 \quad (8.11)$$

$$f(\bar{b} \rightarrow B_s^{*-0}) = 0.020 \pm 0.007. \quad (8.12)$$

Since only the rate of D_{s1}^+ production was determined, it is appropriate to compare it with $f(\bar{b} \rightarrow B_{s1}^0)$ which should differ from equation (8.12) by a factor of $3/8$, determined from the usual spin-counting arguments. In this way, one obtains

$$f(\bar{b} \rightarrow B_{s1}^0) = 0.007 \pm 0.002. \quad (8.13)$$

Equations (8.11) and (8.13) can then be compared directly to the analogous rates for P-wave charm meson production, shown in equations (6.35) and (7.22):

$$f(c \rightarrow D^{*-0}) = 0.073_{-0.022}^{+0.023} {}_{-0.016}^{+0.013}. \quad (8.14)$$

$$f(c \rightarrow D_{s1}^+) = 0.016 \pm 0.004 \pm 0.003. \quad (8.15)$$

Qualitatively, these rates are seen to be similar. They are compatible, given the uncertainties involved. One might expect them to be identical given that the fragmentation process should become insensitive to the quark flavour in the heavy quark limit. However, since the uncertainties are still large, it is not possible to draw strong conclusions from this comparison.

8.3 Properties of Inclusive B Decays

One can also compare D^{*+0} production rates in semileptonic and hadronic b-decays. Although no evidence for D_2^0 production was observed in semileptonic b-decays, the D_1^0 production rates can be compared and from equations (6.39), (6.25) and (6.27), one obtains

$$\frac{\bar{n}_{Z^0 \rightarrow b\bar{b} \rightarrow D_1^0 \ell \bar{\nu} X}}{\bar{n}_{Z^0 \rightarrow b\bar{b} \rightarrow D_1^0 X}} = \frac{\bar{n}_{Z^0 \rightarrow b\bar{b} \rightarrow D_1^0 \ell \bar{\nu} X} \cdot \text{Br}(D_1^0 \rightarrow D^{*+} \pi^-)}{f_1^{b\bar{b}} \cdot \bar{n}_{Z^0 \rightarrow b\bar{b} \rightarrow D^{*+0} X} \cdot \text{Br}(D_1^0 \rightarrow D^{*+} \pi^-)} = 0.10 \pm 0.05 \pm 0.02. \quad (8.16)$$

where ℓ represents either an electron or muon. This ratio is consistent with the trend set by similar comparisons made from charm meson production rates in $Z^0 \rightarrow b\bar{b}$ events:

$$\frac{\bar{n}_{Z^0 \rightarrow b\bar{b} \rightarrow D^{*+} \ell \bar{\nu} X}}{\bar{n}_{Z^0 \rightarrow b\bar{b} \rightarrow D^{*+} X}} = 0.12 \pm 0.02 \quad (8.17)$$

$$\frac{\bar{n}_{Z^0 \rightarrow b\bar{b} \rightarrow D^0 \ell \bar{\nu} X}}{\bar{n}_{Z^0 \rightarrow b\bar{b} \rightarrow D^0 X}} = 0.12 \pm 0.01 \quad (8.18)$$

$$\frac{\bar{n}_{Z^0 \rightarrow b\bar{b} \rightarrow D^+ \ell \bar{\nu} X}}{\bar{n}_{Z^0 \rightarrow b\bar{b} \rightarrow D^+ X}} = 0.10 \pm 0.02 \quad (8.19)$$

Performing these comparisons is useful since many quantities which are not precisely known, such as $\text{Br}(D_1^0 \rightarrow D^{*+} \pi^-)$ and $f(b \rightarrow B)$, are cancelled in the ratio. Although some fraction of the inclusive rates to D^0 and D^+ mesons arises from inclusive decays to D^{*+} mesons, the branching ratios for these processes also cancel in the ratio. Thus, one might expect these ratios to approach a universal value since the hadronization of light degrees of freedom in the B decay should become uncorrelated with the behaviour of the heavy quark. The level of agreement observed in these ratios is interesting.

Chapter 9

Conclusions

9.1 Summary of Results

The research described in this thesis has resulted in measurements of the production rates of some of the P-wave charm mesons produced in hadronic Z^0 decays. These measurements were used to obtain the following probabilities with which charm quarks produce the D_1^0 , D_2^{*0} and D_{s1}^+ states in the fragmentation process:

$$f(c \rightarrow D_1^0) = 0.021 \pm 0.007 \pm 0.003$$

$$f(c \rightarrow D_2^{*0}) = 0.052 \pm 0.022 \pm 0.013$$

$$f(c \rightarrow D_{s1}^+) = 0.016 \pm 0.004 \pm 0.003.$$

These measurements demonstrate that the production of P-wave charm mesons in Z^0 fragmentation is not a rare process. The fragmentation of charm quarks is seen to yield one of the mesons studied roughly 10% of the time and one can infer similar production rates for the remaining D_1^+ , D_2^{*+} and D_{s2}^{*+} P-wave states which were not studied in this analysis. Thus, the production of these excited states must be accounted for in any description of the process of quark fragmentation.

The situation was seen to be similar in the analysis of the decays of b-hadrons to the D_1^0 and D_2^{*0} states. The frequencies with which these P-wave states are

produced in b-decays were found to be

$$\begin{aligned} f(b \rightarrow D_1^0) &= 0.050 \pm 0.014 \pm 0.006 \\ f(b \rightarrow D_2^{*0}) &= 0.047 \pm 0.024 \pm 0.012. \end{aligned}$$

which implies that they constitute a sizeable fraction of all possible decay modes. As in the study of charm fragmentation, these measurements demonstrate that models which do not account for P-wave charm meson production in the decays of b-hadrons are inadequate.

9.2 Discussion of Results

The discussion presented in chapter 3 motivated many aspects of the study of P-wave charm mesons. While the LEP environment is not suitable for precision measurements of their masses, widths or branching ratios, it was argued that these states could provide an ideal laboratory in which to study quark fragmentation. It was demonstrated in the subsequent chapters that the technical challenges associated with their reconstruction in hadronic Z^0 decays and the separation of their production into $c\bar{c}$ and $b\bar{b}$ components could be overcome. The resulting measurements of D^{**0} production in $Z^0 \rightarrow c\bar{c}$ events show that these states are commonly produced at rates compatible with recent phenomenological models used to describe quark fragmentation. Although the determination of these rates relies on several theoretical assumptions, these can all be understood using the simple heavy quark symmetry arguments presented in chapter 3.

Similarly, the production of D_{s1}^+ mesons in charm fragmentation was observed, providing the first evidence at LEP for their presence in hadronic Z^0 decays and to date, this measurement remains the only LEP result on D_{s1}^+ production in charm fragmentation. The production rate was also found to be in reasonable agreement with the quark fragmentation model studied.

Finally, the study of D^{**0} production in semileptonic b-decays allowed an in-

interesting comparison with the inclusive production rate to be made. The ratio of semileptonic to inclusive D^{*-0} production was seen to be in agreement with similar ratios for the production of other charm mesons in the final state which can be explained using the simple spectator model of b-decays. However, it is unfortunate that none of these measurements have provided stringent tests of either fragmentation models or b-decay phenomenology. The qualitative agreement seen in all cases is interesting but ultimately it is important to more rigorously challenge the limits of these simple views of quark fragmentation and the decays of b-hadrons.

9.3 Future Outlook

In October of 1995, LEP ended its physics program devoted to the study of the Z^0 boson and the properties of its decays. The decision to operate at significantly higher energies is clearly justified by the entirely new class of physics processes which can now be studied but it means that there will be no additional statistics available for studying quark fragmentation. Although there are tentative proposals to operate LEP at Z^0 energies in the future, it is unlikely that the statistics gained would be sufficient to significantly improve the studies of P-wave charm mesons presented here. The only other facility at which Z^0 decays are studied is the SLD detector at the Stanford Linear Collider, but the size of their currently available data sample are much less than those of LEP experiments. Nevertheless, proposed upgrades to this collider could allow the accumulation of 1.4 million hadronic Z^0 decays and the high quality of the SLD vertex detectors might allow competitive measurements to be performed there.

While our ability to perform more detailed studies of quark fragmentation remains limited, it is likely that significant improvements can be made in measurements of the properties of P-wave charm mesons and their production in b-decays. Perhaps the most significant facility may be a new interaction point under construction at Fermilab where the CHARM2000 collaboration intends to conduct a

fixed target experiment. The number of fully reconstructed $D^0 \rightarrow K^- \pi^+$ decays which could be recorded by this detector facility in one year is of the order of 10^8 , which would allow measurements similar to those conducted at the Fermilab experiments E687 and E791 to be performed with statistical significance improved by a factor of 50. The SLAC B-factory, also now under construction, should allow the measurements previously performed by ARGUS and CLEO to be repeated with improved statistical precision. However, unlike these experiments, the SLAC B-factory is an asymmetric collider, producing $B\bar{B}$ pairs with a forward boost in the detector. It will then be possible to reconstruct B-decays individually, providing a cleaner environment in which to study both inclusive and exclusive B-decays to P-wave charm mesons.

Thus, although the remaining analyses involving P-wave charm mesons at LEP are now severely limited, the future study of these mesons at new experiments remains bright and there is good reason to believe that they will play a useful role in understanding hadronic structure and b-decays in the future.

Appendix A

Angular Distributions

The $\cos \alpha$ distributions described in section 3.5.1 may be calculated from the matrix elements which describe the sequence of two-body decays. Since these processes must conserve angular momentum, the distributions can be computed using only angular momentum and symmetry considerations. These are the principles behind the helicity formalism, described in [26, 27], which can be used to calculate the angular distributions without knowledge of the full matrix elements describing the strong decays. This formalism and its application to the calculation of the angular distributions are described in the following sections.

A.1 The Helicity Formalism

The flight direction of an unstable meson with spin J may be used to define an initial quantization axis in its rest frame. This may be used to define a basis, $|JM\rangle$, for states of angular momentum J with projection M along this axis. If this meson undergoes a two-body decay,

$$J \rightarrow 1 + 2 \tag{A.1}$$

then the flight path of particle 1 may be used to define another axis, the orientation of which may be specified by polar angles θ and ϕ with respect to the flight direction of J . The two-particle final state may be expressed in terms of the helicities, λ_1 and λ_2 which correspond to the projections of the spins of particles 1 and 2 aligned with and opposed to this axis, respectively. These two bases are then related by the following transformations:

$$|JM, \lambda_1 \lambda_2\rangle = N_J \int d\Omega D_{M\lambda}^{J*}(\phi, \theta, -\phi) |\phi \theta, \lambda_1 \lambda_2\rangle \quad (\text{A.2})$$

and

$$|\phi \theta, \lambda_1 \lambda_2\rangle = N_J \sum_M D_{M\lambda}^J(\phi, \theta, -\phi) |JM, \lambda_1 \lambda_2\rangle \quad (\text{A.3})$$

where $N_J \equiv \sqrt{J(J+1)/4\pi}$, $\lambda \equiv \lambda_1 - \lambda_2$ and the particular choice of the third Euler angle in the rotation matrices has been chosen for convenience.

With the appropriate normalization of two-particle momentum states, the amplitude for the decay of J may be written:

$$\begin{aligned} \mathcal{A}(J \rightarrow 1 + 2) &= \langle \mathbf{p} \lambda_1; -\mathbf{p} \lambda_2 | \mathcal{M} | J, M \rangle \\ &= 4\pi \left(\frac{w}{p} \right)^{1/2} \langle \phi \theta, \lambda_1 \lambda_2 | JM \lambda_1 \lambda_2 \rangle \langle JM \lambda_1 \lambda_2 | \mathcal{M} | JM \rangle \\ &= N_J \mathcal{A}_{\lambda_1 \lambda_2}^J D_{M\lambda}^{J*}(\phi, \theta, -\phi) \end{aligned} \quad (\text{A.4})$$

where $\mathcal{A}_{\lambda_1 \lambda_2}^J$ are known as the helicity amplitudes of the decay and since $\mathcal{M}$ is rotationally invariant, they cannot depend on any of M , θ or ϕ .

In the case of D_1^0 decay, it will be useful to express these helicity amplitudes in terms of the partial wave amplitudes for the decay. This may be done by expressing

the partial wave basis $|JM\ell s\rangle$ in terms of $|JM\lambda_1\lambda_2\rangle$ as follows¹:

$$|JM\ell s\rangle = \sum_{\lambda_1\lambda_2} |JM\lambda_1\lambda_2\rangle \langle JM\lambda_1\lambda_2|JM\ell s\rangle \quad (\text{A.5})$$

$$= \sum_{\lambda_1\lambda_2} \sqrt{\frac{2\ell+1}{2J+1}} C_{\ell 0; s \lambda}^{J \lambda} C_{s_1 \lambda_1; s_2 \lambda_2}^{s \lambda} |JM\lambda_1\lambda_2\rangle. \quad (\text{A.6})$$

The helicity amplitudes are then related to the partial wave amplitudes by

$$A_{\lambda_1\lambda_2}^J = 4\pi \left(\frac{w}{p}\right)^{1/2} \langle JM\lambda_1\lambda_2|\mathcal{M}|JM\rangle \quad (\text{A.7})$$

$$= 4\pi \left(\frac{w}{p}\right)^{1/2} \sum_{\ell s} \langle JM\lambda_1\lambda_2|JM\ell s\rangle \langle JM\ell s|\mathcal{M}|JM\rangle \quad (\text{A.8})$$

$$= \frac{1}{N_J} \sum_{\ell s} \sqrt{2\ell+1} f_{\ell s} C_{\ell 0; s \lambda}^{J \lambda} C_{s_1 \lambda_1; s_2 -\lambda_2}^{s \lambda} \quad (\text{A.9})$$

where the partial wave amplitudes, $f_{\ell s}$, are defined by

$$f_{\ell s} = \sqrt{4\pi} \left(\frac{w}{p}\right)^{1/2} \langle JM\ell s|\mathcal{M}|JM\rangle. \quad (\text{A.10})$$

A.2 Sequential Two Body Decays

In this analysis we are concerned with angular distributions in sequential two-body decays of the form

$$J \longrightarrow \begin{array}{c} 1+2 \\ \downarrow \\ 3+4. \end{array} \quad (\text{A.11})$$

¹Clebsch-Gordon coefficients are denoted $C_{j_1 m_1 j_2 m_2}^{J M}$ which is equivalent to the frequently used $\langle j_1 j_2 m_1 m_2 | j_1 j_2 J M \rangle$ notation.

The helicity of particle 1 is not observed and must be summed over when calculating the full matrix element. Thus, the amplitude may be expressed

$$A(J \rightarrow 1 + 2: 1 \rightarrow 3 + 4) \propto \sum_{\lambda_1} D_{\lambda_1, \lambda_3 - \lambda_4}^{s_1 \pi}(\phi', \theta', -\phi') B_{\lambda_3 \lambda_4} D_{M, \lambda_1 - \lambda_2}^{J \pi}(\phi, \theta, -\phi) A_{\lambda_1 \lambda_2} \quad (\text{A.12})$$

where θ', ϕ' are the polar angles between the momentum vectors of particles 1 and 3 in the rest frame of particle 1 and where $B_{\lambda_3 \lambda_4}$ are the helicity amplitudes for the second decay. Hence, for the decay chain in equation 3.25, θ' corresponds to the angle α defined in equation 3.26.

The angular distribution of the decay products may be written in terms of this amplitude and a density matrix, $\rho_{MM'}$ as follows:

$$\begin{aligned} \frac{dN}{d\Omega d\Omega'} \propto \sum_{M, M', \lambda_1, \lambda_1', \lambda_2, \lambda_3, \lambda_4} & \rho_{MM'} D_{\lambda_1, \lambda_3 - \lambda_4}^{s_1 \pi}(\Omega') D_{\lambda_1', \lambda_3 - \lambda_4}^{s_1 \pi}(\Omega') |B_{\lambda_3 \lambda_4}|^2 \\ & \times D_{M, \lambda_1 - \lambda_2}^J(\Omega) D_{M', \lambda_1 - \lambda_2}^{J \pi}(\Omega) A_{\lambda_1 \lambda_2}^* A_{\lambda_1' \lambda_2} \quad (\text{A.13}) \end{aligned}$$

where Ω and Ω' represent the angles $(\phi, \theta, -\phi)$ and $(\phi', \theta', -\phi')$, respectively.

The helicity amplitudes $A_{\lambda_1 \lambda_2}$ and $B_{\lambda_3 \lambda_4}$ must reflect the symmetries of the strong interaction and since parity is conserved its components are not independent. If η_i represents the intrinsic parity of particle i , the parity of the final state is $\eta_1 \eta_2 (-1)^{L_{12}}$ where $L_{12} = J - s_1 - s_2$. Hence, parity conservation requires that

$$A_{\lambda_1 \lambda_2} = \eta_J \eta_1 \eta_2 (-1)^{J - s_1 - s_2} A_{-\lambda_1 -\lambda_2} \quad (\text{A.14})$$

and constrains $B_{\lambda_3 \lambda_4}$ similarly.

A.3 The Decays $D^{**0} \rightarrow D^{*+}\pi^-$

The decay chains relevant in this thesis are

$$\begin{array}{ccc} D_J^{(*)0} & \longrightarrow & D^*(2010)^+\pi^- \\ & & \downarrow \\ & & D^0\pi^+ \end{array} \quad (\text{A.15})$$

in which case $\lambda_2 = \lambda_3 = \lambda_4 = 0$ and $s_1 = 1$. Parity conservation requires that the D_1^0 and D_2^{*0} states decay with an even partial wave and while both S- and D- wave decays are allowed for the D_1^0 , angular momentum conservation requires that the D_2^{*0} decay only via a D-wave. The angular distributions for D_2^{*0} decay are then simpler to evaluate and will be treated first.

A.3.1 $D_2^{*0} \rightarrow D^{*+}\pi^-$

In this case, equation A.13 may be written

$$\frac{dN}{d\Omega d\Omega'} \propto \sum_{M,M',\lambda_1,\lambda_1'} \rho_{M,M'} D_{\lambda_1,0}^1(\Omega') D_{\lambda_1',0}^{1*}(\Omega') |B_{00}|^2 \times D_{M,\lambda_1}^2(\Omega) D_{M',\lambda_1'}^{2*}(\Omega) A_{\lambda_1,0}^* A_{\lambda_1',0}. \quad (\text{A.16})$$

Integrating over $d\Omega$ introduces a delta function:

$$\int d\Omega \left(D_{M,\lambda}^J(\Omega) D_{M',\lambda'}^{J*}(\Omega) \right) = \frac{4\pi}{2J+1} \delta_{M,M'} \delta_{\lambda,\lambda'} \quad (\text{A.17})$$

which gives

$$\frac{dN}{d\Omega'} \propto |A_{-1,0}|^2 |D_{-1,0}^1(\Omega')|^2 + |A_{0,0}|^2 |D_{0,0}^1(\Omega')|^2 + |A_{1,0}|^2 |D_{1,0}^1(\Omega')|^2 \quad (\text{A.18})$$

Parity conservation imposed by equation A.14 implies that $A_{10} = -A_{-10}$ and $A_{00} = 0$ and the trace of the density matrix is unity which yields

$$\frac{dN}{d\Omega'} \propto |A_{10}|^2 (|D_{10}^1|^2 + |D_{-10}^1|^2). \quad (\text{A.19})$$

and since $d_{-mm'}^j = (-1)^{m-m'} d_{mm'}^j$, this gives the expression in equation 3.26:

$$\frac{dN}{d(\cos \alpha)} \propto |d_{10}^1(\cos \alpha)|^2 \propto 1 - \cos^2 \alpha. \quad (\text{A.20})$$

A.3.2 $D_1^0 \rightarrow D^{*+} \pi^-$

This decay is more complicated since parity conservation implies that $A_{\lambda_2 \lambda_3} = A_{-\lambda_2 -\lambda_3}$ in which case there are now two independent, nonvanishing helicity amplitudes, A_{10} and A_{00} . Thus, the angular distribution may be written in terms of these as follows:

$$\begin{aligned} \frac{dN}{d(\cos \alpha)} &\propto |A_{10}|^2 (|D_{10}^1|^2 + |D_{-10}^1|^2) + |A_{00}|^2 |D_{00}^1|^2 \\ &= 2|A_{10}|^2 |d_{10}^1|^2 + |A_{00}|^2 |D_{00}^1|^2 \\ &= |A_{10}|^2 \sin^2 \alpha + |A_{00}|^2 \cos^2 \alpha. \end{aligned} \quad (\text{A.21})$$

It is now convenient to express these in terms of partial wave amplitudes using equation A.10. This gives

$$A_{10} = \frac{1}{3} \left(S + \frac{D}{\sqrt{2}} \right) \quad (\text{A.22})$$

$$A_{00} = \frac{1}{3} (S - \sqrt{2} D) \quad (\text{A.23})$$

where S and D are the partial wave amplitudes f_{01} and f_{21} in the notation used in equation A.10. Hence, the angular distribution can be expressed

$$\frac{dN}{d(\cos \alpha)} = \frac{1}{2} \left(|S|^2 + |D|^2 \left(\frac{1 + 3 \cos^2 \alpha}{2} \right) + \frac{S^* D + D^* S}{\sqrt{2}} (1 - 3 \cos^2 \alpha) \right). \quad (\text{A.24})$$

This can be expressed in terms of the S- and D-wave widths $\Gamma_S = |S|^2$ and $\Gamma_D = |D|^2$ which differ by a relative phase φ as follows:

$$\frac{dN}{d(\cos \alpha)} \propto \frac{1}{2} \left(\Gamma_S + \Gamma_D \left(\frac{1 + 3 \cos^2 \alpha}{2} \right) + \sqrt{2\Gamma_S \Gamma_D} \cos \varphi (1 - 3 \cos^2 \alpha) \right) \quad (\text{A.25})$$

and when the relative widths are denoted $R = \Gamma_S/(\Gamma_S + \Gamma_D)$, one obtains the expression shown in equation 3.35:

$$\frac{dN}{d(\cos \alpha)} = \frac{1}{2} \left(R + (1 - R) \cdot \frac{1}{2}(1 + 3 \cos^2 \alpha) + \sqrt{2R(1 - R)} \cos \varphi (1 - 3 \cos^2 \alpha) \right). \quad (\text{A.26})$$

In the heavy quark limit, the $\ell = 2$ amplitude is expected to dominate the decay of the D_1^0 which corresponds to $R = 0$ and leads to the expression in equation 3.26:

$$\frac{dN}{d(\cos \alpha)} = \frac{1}{4} (1 + 3 \cos^2 \alpha). \quad (\text{A.27})$$

A.4 Decay Angle Dependence

The result shown in equations A.25 and A.20 were obtained by integrating over the decay angles Ω of the D^{**0} . When only a restricted range of decay angles is accepted, the angular distribution of $\cos \alpha$ is modified and can introduce a dependence on the polarization of the D^{**0} . To derive this dependence, one integrates equation A.13 over the angles ϕ and ϕ' using

$$\int d\phi D_{M,\lambda}^{J*}(\phi, \theta, -\phi) D_{M',\lambda'}^J(\phi, \theta, -\phi) = 2\pi d_{M,\lambda}^J(\theta) d_{M',\lambda'}^J(\theta) \delta_{M,M'} \delta_{\lambda,\lambda'} \quad (\text{A.28})$$

to obtain

$$\frac{dN}{d(\cos \theta_\pi) d(\cos \alpha)} \propto \sum_{M,\lambda} \rho_{M,M} |d_{\lambda,0}^1(\alpha)|^2 |A_{\lambda 0}|^2 |d_{M,\lambda}^J(\theta_\pi)|^2. \quad (\text{A.29})$$

The angular distribution of the D_2^0 state remains of the form $1 - \cos^2 \alpha$, but that of the D_1^0 state is modified as follows:

$$\begin{aligned} \frac{dN}{d(\cos \theta_\pi) d(\cos \alpha)} &\propto |A_{10}|^2 \frac{\sin^2 \alpha}{2} \sum_M \rho_{MM} (|d_{M,1}^1(\theta_\pi)|^2 + |d_{M,-1}^1(\theta_\pi)|^2) + \\ &|A_{00}|^2 \cos^2 \alpha \sum_M \rho_{MM} |d_{M,0}^1(\theta_\pi)|^2. \end{aligned} \quad (\text{A.30})$$

Parity invariance implies that $\rho_{-1-1} = \rho_{11} = \frac{1}{2}(1 - \rho_{00})$ and hence,

$$\begin{aligned} \frac{dN}{d(\cos \theta_\pi) d(\cos \alpha)} &\propto \\ &\frac{\sin^2 \alpha}{4} \left[(1 + \cos^2 \theta_\pi) + \rho_{00}(1 - 3 \cos^2 \theta_\pi) \right] |A_{10}|^2 + \\ &\frac{\cos^2 \alpha}{2} \left[(1 - \cos^2 \theta_\pi) - \rho_{00}(1 - 3 \cos^2 \theta_\pi) \right] |A_{00}|^2. \end{aligned}$$

When the helicity amplitudes are written in terms of the S - and D -wave amplitudes using equations A.22 and A.23 one obtains

$$\begin{aligned} \frac{dN}{d(\cos \theta_\pi) d(\cos \alpha)} &\propto \\ &\frac{\sin^2 \alpha}{8} \left[(1 + \cos^2 \theta_\pi) + \rho_{00}(1 - 3 \cos^2 \theta_\pi) \right] \left\{ 2\Gamma_S + \Gamma_D + 2\sqrt{2\Gamma_S\Gamma_D} \cos \varphi \right\} + \\ &\frac{\cos^2 \alpha}{2} \left[(1 - \cos^2 \theta_\pi) - \rho_{00}(1 - 3 \cos^2 \theta_\pi) \right] \left\{ \Gamma_S + 2\Gamma_D - 2\sqrt{2\Gamma_S\Gamma_D} \cos \varphi \right\}. \end{aligned}$$

In the heavy quark limit where $\Gamma_S = 0$, one can compute the integral of this distribution over a restricted range of $\cos \theta_\pi \equiv y$ to obtain

$$\begin{aligned} \frac{dN}{d(\cos \alpha)} &= \int_{y_{\min}}^1 \frac{dN}{dy d(\cos \alpha)} dy \\ &\propto \frac{\sin^2 \alpha}{8} \left[\frac{4}{3} - y_{\min} - \frac{y_{\min}^3}{3} + \rho_{00}(y_{\min}^3 - y_{\min}) \right] + \\ &\cos^2 \alpha \left[\frac{2}{3} - y_{\min} + \frac{y_{\min}^3}{3} - \rho_{00}(y_{\min}^3 - y_{\min}) \right]. \end{aligned} \quad (\text{A.31})$$

In this analysis it was necessary to impose a minimum momentum cut on the π^- which effectively limits the range of $\cos \theta_\pi$ accepted. This changes the shape of the $\cos \alpha$ distribution slightly and introduces a dependence on ρ_{00} which is shown in figure 6.12.

Appendix B

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