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**Inclusive Hadron Production and Two Particle Correlations
in e^+e^- Annihilation at 29 GeV Center-of-Mass Energy**

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Ph.D. Thesis
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Abstract

We have studied hadron production in e^+e^- annihilation at 29 GeV center-of-mass energy using the PEP-4 Time Projection Chamber Detector. The inclusive cross sections and mean multiplicities for $\pi^\pm$, $K^\pm$ and $(p + \bar{p})$ production have been measured using ionization energy loss to separate particle species. We find on average $10.7 \pm 0.6 \pi^\pm$, $1.35 \pm .13 K^\pm$ and $0.60 \pm 0.08 (p + \bar{p})$ per multihadron event. The differential cross section is well described by a number of Monte Carlo hadronization models. In addition, we have observed correlations in rapidity space for identified pions and kaons. Short-range KK correlations provide evidence for local flavor compensation during hadronization. Long-range $\pi\pi$ and KK correlations indicate that the initial partons carry flavor. We also observe significant long-range πK correlations as a result of heavy quark decays.

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Chapter 1

Introduction

One of the successes of the last 10 years in particle physics has been the development of the QCD theory of the strong interactions. Quantum Chromodynamics, a renormalizable theory based on the principle of local gauge invariance, describes the interaction of quark and gluon fields that carry a “color” quantum number. One of the peculiarities of the model, however, is that these fundamental, colored particles are never observed as physical states. Instead, all known strongly interacting particles (hadrons) are described as color-neutral quark-antiquark (meson) or three quark (baryon) bound states. The fact that colored objects never appear as free particles is known as the principle of confinement.

Because the strength of the coupling constant, α_s , depends on the four momentum transfer, Q^2 , strong interaction theory can be divided into two regimes. For large momentum transfers, the coupling is small and the theory describes the interaction of quasifree quarks and gluons. At low Q^2 , however, α_s is large; here the theory must describe bound state hadrons. The high energy (short distance) behavior of QCD can be explored using perturbation theory. The long distance behavior is more problematic. Recently, advances in describing the static behavior of QCD have been made using lattice gauge theories¹. In addition, non-relativistic potential models have proved to be important tools for studying the spectroscopy of heavy quark bound states². The process by which quarks are transformed into hadrons has been explored using a number of phenomenological models. Because an adequate treatment of hadronization must include the details of hadron spectroscopy and decays, these models are often implemented in computer simulations (Monte Carlos).

The reaction $e^+e^- \rightarrow \text{hadrons}$ provides an important laboratory for studying hadronization because the initial state is simple and well understood. Electro-weak theory fully describes e^+e^- annihilation into a virtual photon or Z^0 and the subsequent creation of a quark-antiquark pair. The strong interactions contribute only to the final state transformation of this $q\bar{q}$ pair into observable particles. In this thesis we investigate hadronization by studying e^+e^- annihilation at a center-of-mass energy of 29 GeV using the Time Projection Chamber (TPC) detector at the PEP storage ring. We present a measurement of the charged pion, kaon and proton differential cross sections and an analysis of two particle correlations between identified hadrons. Correlation studies provide a useful method for separating the perturbative and confinement regimes. Since the initial quark-antiquark pair is created with large relative momenta, the hadrons containing these quarks will be widely separated in phase space. Thus, this process gives rise to long-range quantum number correlations. In contrast, confinement is dominated by low Q^2 phenomena and therefore generates local, short-range correlations.

The paper is organized as follows. In Chapter 2 we review strong interaction theory and describe current models of hadronization and confinement. We discuss the experimental evidence supporting these models and outline our experimental program for further testing them. Chapter 3 is a description of PEP and the PEP-4 detector facility. Since the analysis presented in this thesis relies exclusively on charged particles tracked in the TPC, we discuss the Time Projection Chamber in greatest detail. The remaining detector elements are described briefly. The PEP-4

event selection and reconstruction software are outlined in Chapter 4. This chapter also includes a description of the detector simulation Monte Carlo. Chapter 5 reviews the theory of ionization energy loss (dE/dx) in gases and discusses the dE/dx method of particle identification. Particular emphasis is placed on the design constraints imposed on the TPC by this method. In Chapter 6 we present our measurement of the differential cross sections for charged pion, kaon and proton production and compare these measurements to the predictions of several Monte Carlo models. A study of rapidity distributions and two particle correlations in rapidity space is presented in Chapter 7. Finally, in Chapter 8 we summarize our results and discuss the issues that can be clarified by future experiments.

Chapter 2

The Phenomenology of Hadronization

The reaction $e^+e^- \rightarrow \text{hadrons}$ can be divided into processes that occur on vastly different length scales. The initial e^+e^- annihilation into a virtual photon and the subsequent production of a parton system occur at large momentum transfers and hence small distances. For these short length scales, strong interaction theory can be described as the interaction of quasifree quarks and gluons. In contrast, the confining forces that transform quarks and gluons into hadrons are generally thought to be soft. Thus, the hadronization process is a long wavelength phenomenon. Because the momentum transfer involved in hadronization is small, quarks and gluons fragment into particles with small momenta transverse to the initial parton direction. Thus, each initial parton is transformed into a "jet" of hadrons.

In this chapter we discuss the phenomenology of hadronization. We begin with a review of the quark-parton model and the QCD theory of strong interactions. We compare the experimental data on jet production in e^+e^- annihilations to the predictions of the QCD improved parton model and conclude that this model adequately describes the hard scattering processes $e^+e^- \rightarrow q\bar{q}$ and $e^+e^- \rightarrow q\bar{q}g$. We then discuss the long wavelength behavior of the strong interactions. We describe the current phenomenological picture of color confinement and discuss a number of Monte Carlo models of hadronization. Finally, we outline our experimental program for studying hadronization.

2.1 The Quark-Parton Model

In 1961 Gell-Mann³ and Ne'eman⁴ observed that all known mesons and baryons could be classified as members of representations of the gauge group SU(3). In their scheme, mesons were found in the singlet and octet representations while baryons were classified as singlets, octets and decuplets. At the time, the absence of other representations of SU(3) was a puzzle. In particular, why was the **3** representation missing? In 1964, Gell-Mann⁵ and Zweig⁶ demonstrated that the observed multiplet structure could be explained if the fundamental building blocks of hadrons were in fact members of an SU(3) triplet. Mesons then consisted of a triplet and an anti-triplet:

$$\mathbf{3} \otimes \bar{\mathbf{3}} = \mathbf{1} \oplus \mathbf{8}$$

Baryons were composed of three triplets:

$$\mathbf{3} \otimes \mathbf{3} \otimes \mathbf{3} = \mathbf{1} \oplus \mathbf{8} \oplus \mathbf{8} \oplus \mathbf{10}$$

In order to match the known hadron quantum numbers, the triplets, called quarks, were required to have spin $\frac{1}{2}$ and baryon number $\frac{1}{3}$. One quark "flavor", u , had charge $\frac{2}{3}e$ and the two others, d and s , had charge $-\frac{1}{3}e$.

The quark formalism did not depend on the SU(3) triplets being physical particles. The belief that quarks were "purely mathematical entities"⁷ was for a time encouraged by the negative results of fractional charge searches. A change in the status of the elusive quark occurred in

the late 1960's with the measurement of deep-inelastic electron-proton scattering in high energy collisions⁷. Just as the existence of hard scattering led Rutherford to postulate the presence of a massive core within the atom, the observation of large angle electron scattering in e - p collisions provided evidence for the existence of constituents, dubbed partons, within the proton. The possibility that these partons were quarks was recognized almost immediately⁸. Since that time, data from μ - p and ν - p scattering have supported this viewpoint⁹.

The basic tenet of the quark-parton model is that the interaction of a hadron at very high energies can be described as the quasifree scattering of one of its constituents. This model is essentially a formulation of the impulse approximation; it is based on the assumption that the high momentum transfer interaction occurs on a time scale that is short relative to the time for interactions between partons within the hadron. Under such conditions, the cross section for e - p scattering can be factorized into three terms. The initial distribution of parton states within the proton is described by a probability distribution:

$$f_i(x) = \begin{array}{l} \text{the probability that the proton contains a parton } i \text{ that} \\ \text{carries a fraction } x \text{ of the proton's momentum} \end{array}$$

The hard scattering is considered an electron-parton interaction, summed incoherently over possible initial states. The result of this scattering is to knock a parton out of the proton. The strong interactions then transform this parton into hadrons. The decay or fragmentation of the parton is parameterized by another probability distribution. Because hadronization is a low Q^2 process, it cannot produce particles with large momentum transverse to the parton direction of flight. Therefore, it is natural to describe fragmentation in terms of the longitudinal momentum fraction z . The quark fragmentation function thus is defined:

$$D_i^h(z) = \begin{array}{l} \text{the probability that parton } i \text{ fragments so as to produce a} \\ \text{hadron } h \text{ that carries a fraction } z \text{ of the parton's momentum} \end{array}$$

(see Fig. 2.1a).

The cross section for $e^+e^- \rightarrow \text{hadrons}$ can also be described as an initial hard parton scattering ($e^+e^- \rightarrow q\bar{q}$) and a transformation of the final state partons into observable particles (see Fig. 2.1b). Because the initial e^+ and e^- are pointlike, no summation over initial states is necessary. The assumption that the hard scattering and hadronization processes factorize implies that the probability that a final state parton fragments into a hadron with a given momentum fraction z should not depend on the hard scattering process. Thus, the same fragmentation functions should describe e^+e^- and e - p scattering.

2.2 Color and Quantum Chromodynamics

Since the formulation of Gell-Mann and Zweig's model of flavor SU(3), two new quarks, c and b , have been discovered. The extension of the model to include these quarks is straightforward. All known hadrons continue to be found in $q\bar{q}$ and qqq states. This pattern can be explained if we introduce an addition symmetry, SU(3)_{color}. In this scheme, quarks are color triplets while all hadrons are color singlet states.

A number of theoretical and experimental results support the existence of SU(3)_{color}. The Δ^{++} is composed of three identical u quarks in a totally symmetric configuration. If quarks are spin $\frac{1}{2}$ fermions, however, the Pauli principle requires that their wave function be antisymmetric under interchange. This contradiction can be overcome if we postulate that the quarks carry another quantum number, i.e. color, and are in an antisymmetric state with respect to color interchange. In order to create an antisymmetric combination of 3 particles, there must be at least 3 color states. Evidence supporting the hypothesis that there are exactly 3 colors comes from the measurement of R , the ratio of the total hadronic cross section to the cross section for μ pair production in e^+e^- annihilation. The QED result for this ratio is:

$$R = \frac{\sigma(e^+e^- \rightarrow \text{hadrons})}{\sigma(e^+e^- \rightarrow \mu^+\mu^-)} = \sum_{i=1}^N e_q^2 \times (\text{number of colors}) \quad (2.1)$$

Figure 2.1: The parton model picture of a) deep inelastic e - p scattering and b) e^+e^- scattering.

Figure 2.2: Contributions to the QED photon propagator a) to lowest order and b) to second order in α

where N is the number of quark flavors that can be produced with the available energy and e_q is the charge of quark i . At PEP and PETRA energies, where there are 5 quark flavors (u, d, s, c and b) and $\sum e_q^2 = \frac{11}{9}$, the measured value of R is 3.95 ± 0.08^{10} . Three colors are also needed to make the experimental value of the π^0 lifetime agree with theory¹¹.

The incorporation of $SU(3)_{\text{color}}$ into a true theory of the strong interactions occurred in the early 1970's. Quantum Chromodynamics (QCD) is a renormalizable field theory based on the principle of local gauge invariance under $SU(3)$ color transformations. The theory describes the interaction of color triplet quarks ($\psi_{i=1,3}$) through the interchange of color octet gauge bosons ($A_{a=1,8}^\mu$), called gluons. The strong interaction Lagrangian has the form:

$$\mathcal{L} = -\frac{1}{4}F_{a\mu\nu}F_{a\mu\nu} + i\bar{\psi}_j\gamma_\mu D_{jk}^\mu\psi_k - \bar{\psi}_j M_{jk}\psi_k \quad (2.2)$$

where

$$F_{a\mu\nu}^\mu = \partial^\mu A_a^\nu - \partial^\nu A_a^\mu + g f_{abc} A_b^\mu A_c^\nu \quad (2.3)$$

and D_{jk}^μ is the covariant derivative

$$D_{jk}^\mu = \delta_{jk}\partial^\mu - ig\{T_a\}_{jk}A_a^\mu \quad (2.4)$$

The T matrices are defined in terms of f_{abc} , the fundamental structure constants of $SU(3)$, by the commutation relation:

$$[T_a, T_b] = if_{abc}T_c \quad (2.5)$$

The first term in Eq. 2.2 is the Lagrangian for the gluon gauge field in the absence of quarks. The non-abelian nature of the field (i.e. the fact that the A_a^μ do not commute) indicates that the gluons interact with each other. The second term in the Lagrangian represents the interaction of quarks and gluons. The final term gives the quarks mass. This term is usually considered negligible for u, d and s quarks. Note that all quark and gluon interactions are governed by a single coupling constant g and that the strong interactions are flavor independent.

The presence of a gluon self coupling has profound effects on the structure of QCD. This effect is most clearly seen in the Q^2 dependence of the effective coupling constant α_s . The concept of an effective coupling constant is familiar from QED and can be understood in terms of vacuum polarization. In QED, the electromagnetic coupling $\alpha \cong 1/137$ is defined by the behavior of the theory at low energy. The Coulomb interaction is reproduced in lowest order QED because the photon propagator, $D_{\mu\nu}(q)$, is proportional to $1/q^2$. When the order α corrections to the

propagator are calculated (see Fig. 2.2), the result contains a logarithmically divergent term and a finite, q^2 dependent result¹²:

$$D_{\mu\nu}(q) = -\frac{ig_{\mu\nu}}{q^2} \left[1 - \frac{\alpha}{3\pi} \log \frac{M^2}{m^2} + \frac{2\alpha}{\pi} \int_0^1 dz z(1-z) \log \left(1 - \frac{q^2 z(1-z)}{m^2 - i\epsilon} \right) \right] \quad (2.6)$$

where m is the electron mass and M is a large mass cutoff. The divergent term is absorbed into the definition of the electromagnetic charge:

$$e^2 \rightarrow e^2 \left(1 - \frac{\alpha}{3\pi} \log \frac{M^2}{m^2} \right) \quad (2.7)$$

The argument associated with this "renormalization" is that the measured charge e corresponds to the value of the coupling *after* all higher order corrections have been made. The bare coupling in the Lagrangian is not a measurable quantity; hence, the fact that it diverges has no significance. The finite term in Eq. 2.6 is physically measurable and contributes to the observed atomic Lamb shift. The inclusion of this finite, q^2 dependent term in the electromagnetic coupling means that Coulomb's law is no longer correct. We can at least formally restore the $1/q^2$ dependence by absorbing the finite term in our definition of α . This restoration is made at a price, however. The electromagnetic coupling constant now has a q^2 dependence; in the field theorist's jargon it has become a "running" coupling constant. To first order in α ¹²:

$$\alpha(Q^2) = \alpha(Q^2 = 0) \left(1 + \frac{\alpha}{3\pi} \ln \frac{|q|^2}{m^2} \right) \quad (2.8)$$

where $Q^2 \equiv -q^2$. As the momentum transfer increases, the electromagnetic coupling becomes larger. The increase of the effective electric charge can be interpreted as the result of vacuum polarization. If a positive charge is placed in space, it will polarize the vacuum, creating a cloud of negative charges around it that shields part of its bare charge. A high energy probe can penetrate part of this cloud and therefore observe a higher value of e .

In QCD, as in QED, higher order corrections force a renormalization of the coupling constant and the introduction of a Q^2 dependence. The presence of gluon-gluon interactions, however, changes not only the magnitude but also the sign of this Q^2 dependence. The strong coupling constant α_s runs much faster than α , and it runs in the opposite direction¹³.

$$\alpha_s(Q^2) = \frac{\alpha_s(Q_0^2)}{1 + (33 - 2N_f) \ln(Q^2)/Q_0^2} \equiv \frac{12\pi}{(33 - 2N_f) \ln(Q^2/\Lambda^2)} \quad (2.9)$$

where N_f is the number of quark flavors that contribute in Fig. 2.3b and Λ is the one free parameter of the theory and must be set by experiment.

The form of Eq. 2.9 holds an interesting promise for understanding the strong interactions. At large Q^2 , α_s becomes small (the theory is "asymptotically free,"¹⁴) while at low energies α_s becomes large. The fact that the coupling constant is small for large values of Q^2 means that perturbative calculations should converge and also provides some theoretical justification for the parton model. The behavior of QCD at low energies is not well understood, largely because the tool of perturbation theory is not applicable. It is hoped that the theory can explain the absence of free quarks (confinement) and describe the transformation of these quarks into hadrons.

Since the observed participants of any experiment are hadrons rather than quarks and gluons, one cannot ignore the non-perturbative regime. The success of the quark-parton model provides a key to the merging of perturbative and non-perturbative processes. The parton model approach of factoring the initial and final state interactions into probability distributions is retained in perturbative QCD. Only the hard scattering process is treated using perturbation theory. One modification of parton model formalism is necessary, however. While structure functions and

Figure 2.3: Contributions to the QCD gluon propagator a) to lowest order and b) to next order in α_s .

fragmentation functions are scale independent in the parton model (i.e. depend only on x and z and not on the momentum transfer), QCD introduces logarithmic variations with Q^2 :

$$f_i(z) \rightarrow f_i(z, Q^2), \quad D_i^h(z) \rightarrow D_i^h(z, Q^2)$$

These variations, called scaling violations, result from the fact that the ability to resolve the difference between a quark and a nearly collinear quark plus gluon is a function of the energy of the probe.

The question of when the use of perturbation theory is valid is often hotly argued among theorists. The most common viewpoint, however, seems to be the one expressed by Politzer¹⁶: "Use perturbation theory where it is not obviously wrong." In other words, if the perturbation theory results are well behaved (e.g., if the first order term is small relative to zeroth order and higher order terms are likely to be smaller still) and if there are no obvious non-perturbative effects that are larger than the perturbative result, it is likely that the perturbative answer is meaningful. One problem with this approach, unfortunately, is that when experimental results disagree with QCD calculations, the disagreement can usually be ascribed to the non-validity of the perturbative approach rather than to a fundamental fault in the theory.

As an example of the use of perturbative QCD, we present results for the QCD modified cross section for $e^+e^- \rightarrow \text{hadrons}$. The zeroth order cross section is defined in Eq. 2.1 and corresponds to the process shown in Fig. 2.4a. The QCD corrections to the process are shown in Fig. 2.4b and Fig. 2.4c. Fig. 2.4b shows the corrections to the QED vertex and to the quark propagators due to virtual gluon emission. The interference of these diagrams with the Born graph of Fig. 2.4a leads to a first order in α_s correction. Fig. 2.4c represents the case of real gluon emission from the quark or antiquark. Together, these diagrams modify the total cross section

$$\sigma_{\text{Tot}} = 3 \sum_q e_q^2 \left[1 + \frac{\alpha_s(q^2)}{\pi} \right] \sigma_{\mu\mu} \quad (2.10)$$

Thus, the prediction for R is now

$$R = \frac{\sigma(e^+e^- \rightarrow \text{hadrons})}{\sigma(e^+e^- \rightarrow \mu^+\mu^-)} = R_{\text{QED}} \left[1 + \frac{\alpha_s(q^2)}{\pi} \right]$$

Figure 2.4: Diagrams contributing to the first order α_s cross section for $e^+e^- \rightarrow \text{hadrons}$. a) the zeroth order cross section b) corrections due to virtual gluon exchange and c) contribution due to the radiation of real gluons. The interference between diagrams a) and b) gives a first order α_s correction

where R_{QED} is defined by Eq. 2.1.

For center-of-mass energies of roughly 30 GeV (the range of PEP and PETRA experiments), the result meets the criteria of reasonability expressed above. For $\alpha_s \sim .2$, the QCD correction to the total cross section is less than 10%. There are no apparent difficulties in factorizing out the non-perturbative effects. The values of R measured at PEP and PETRA are consistent with QCD if α_s is between .16 and .27¹⁰.

Fig. 2.4c has an obvious physical interpretation. It represents the process $e^+e^- \rightarrow q\bar{q}g$, i.e., gluon bremsstrahlung from the quark or antiquark. Gluons, like quarks, carry color. We therefore do not expect to observe them as physical particles. Instead, each gluon is transformed into hadrons. Because the hadronization process is soft, these hadrons should have limited momentum with respect to the gluon direction. It is thus possible to hunt for such "jets" experimentally (see the next section for details). A calculation of the cross section for gluon emission can therefore be used to test perturbative QCD.

The total rate of gluon bremsstrahlung can be calculated in first order perturbative QCD and it diverges for low gluon energies (infrared divergence) or small quark-gluon opening angles (collinear divergence). The infrared problem is identical to that of photon bremsstrahlung and is solved in the same manner. The difficulty arises only because we are asking a question that is not physically meaningful. There is no way to distinguish between bremsstrahlung of a real, zero energy photon (or gluon) and the case where no particle at all was emitted. The divergences in Fig. 2.4c are exactly cancelled by divergences in the cross section for virtual gluon emission (Fig. 2.4b). A similar situation exists for the collinear divergence. It is impossible to distinguish between a quark jet and collinear quark plus gluon jets. The physically sensible region for calculating the 3 parton cross section is defined by the (non-perturbative) angular spread of the jets. Within these boundaries, the cross section (integrated over all event orientations) is

$$\frac{d^2\sigma}{dx_1 dx_2} = \sigma_{Tot} \left\{ \frac{2\alpha_s}{3\pi} \frac{(x_1^2 + x_2^2)}{(1-x_1)(1-x_2)} \right\} \quad (2.11)$$

where x_1 and x_2 are the quark and antiquark energies, respectively, scaled to E_{beam} . This

cross section is of course symmetric in the quark and antiquark variables and demonstrates the previously discussed divergence when the quark or antiquark energy approaches the beam energy.

2.3 Evidence for Jets in e^+e^- Annihilation

The QCD improved parton model is based on the belief that the forces governing confinement are soft. If this is the case, the confinement process cannot produce particles with large momenta transverse to the parton direction. The quark and gluon fragmentation products will form collimated "jets" that follow the initial parton directions. The first experimental evidence for jet production in e^+e^- collisions came from the MARK I detector at SPEAR. Each event's jetiness was measured using the sphericity algorithm¹⁰. In this method a tensor is constructed from the particle momenta

$$T^{\alpha\beta} = \frac{\sum p_i^\alpha p_i^\beta}{\sum |\vec{p}_i|^2} \quad \alpha, \beta = x, y, z \quad (2.12)$$

where the sums run over all observed particles. If T is diagonalized, and its eigenvalues are ordered

$$\lambda_1 \geq \lambda_2 \geq \lambda_3$$

(note $\lambda_1 + \lambda_2 + \lambda_3 = 1$) the corresponding eigenvectors $\hat{n}$ obey the equation

$$\lambda_\alpha = \frac{\sum_i (\vec{p}_i \cdot \hat{n}_\alpha)^2}{\sum_i |\vec{p}_i|^2} \quad (2.13)$$

The eigenvector $\hat{n}_1$ is in the direction that minimizes $p_\perp^2$ and is called the sphericity axis. The sphericity is then defined

$$S = \frac{3}{2} (\lambda_2 + \lambda_3) \quad (2.14)$$

S lies between 0 and 1, with 1 implying a completely spherical event and 0 corresponding to extreme jetiness. The MARK I compared the measured sphericity distribution to two Monte Carlo models, a phase space generator and a jet model where particles were produced according to longitudinal phase space and with a gaussian distribution in $p_\perp$. At energies above 5 GeV, the jet model gave a significantly better representation of the data. In addition, the distribution of the direction of the jet axis was consistent with the $(1 + \cos^2\theta)$ form expected for spin 1/2 partons.

Because the transverse momentum introduced by confinement depends on the hadronization length scale and not on the e^+e^- center-of-mass energy, the average value of $p_\perp$ should be independent of $\sqrt{s}$. Thus, the angular dispersion of jets should decrease as the center-of-mass energy is increased. As a result, the mean sphericity should decrease with $\sqrt{s}$. This phenomenon has been observed¹⁷; at PEP and PETRA energies jet structure is so pronounced, it can be seen on an event by event basis (see Fig. 2.5). One consequence of this jet narrowing is the increased ability to recognize the wide angle gluon bremsstrahlung predicted in Eq. 2.11. Three jet events were first observed at PETRA in 1979¹⁸. Evidence supporting the existence of gluon bremsstrahlung can be obtained by examining the eigenvalues of the sphericity tensor (Eq. 2.12). We have already seen that a small value of S is evidence for two-jet structure. Since the q , $\bar{q}$ and g in a three-jet event are constrained to lie in a plane, evidence that events not consistent with two-jet structure are planar supports the view that they are the result of gluon emission. The aplanarity is defined by

$$A = \frac{3}{2} \lambda_3 = \frac{\sum p_{out}^2}{\sum |\vec{p}_i|^2} \quad (2.15)$$

where p_{out} is the momentum out of the plane defined by n_1 and n_2 . Planar events have small values of A . A totally spherical event has $\lambda_1 = \lambda_2 = \lambda_3 = \frac{1}{3}$ and thus has $A = \frac{1}{2}$. Fig. 2.6 shows the sphericity and aplanarity distributions for the PEP-4 TPC experiment. The majority

View 1

View 2

View 3

Figure 2.5: A typical two jet event as seen by this experiment. The three views show the event projected along the three axes of the sphericity tensor. The solid lines represent the observed charged particles, with the length of the lines proportional to the particles' projected momenta. The dashed lines show the two jet axes, as found by a clustering algorithm, and the dotted lines indicate the beam direction.

Figure 2.6: The distribution of eigenvalues of the sphericity tensor, as measured by this experiment. The majority of events have low sphericity, demonstrating that 2 jet structure dominates. However, approximately 10% of the events are planar. These planar events provide evidence for hard gluon bremsstrahlung.

of events show a two jet structure, while about 10% are planar. There are very few spherical events. Studies of the angular distribution of the 3 jets by the TASSO¹⁰ and PLUTO²⁰ groups indicate that the gluon has spin 1.

The study of inclusive hadron spectra also provides evidence that the QCD improved quark-parton model is a good description of $e^+e^- \rightarrow \text{hadrons}$. As discussed in the previous section, the quark fragmentation function is expected to scale with the momentum fraction z in the parton model. QCD introduces logarithmic deviation from scaling. By studying the dependence of the fragmentation function on center-of-mass energy, the MARK II and TASSO collaborations have observed approximate scaling and small scaling violations^{21,22}.

Because most of the events at PEP and PETRA have a two-jet structure, it is reasonable to study inclusive hadron production relative to the jet axis. A useful variable for such studies is rapidity, defined by

$$y = \frac{1}{2} \ln \left(\frac{E + p_{\parallel}}{E - p_{\parallel}} \right) \quad (2.16)$$

where $p_{\parallel}$ is the component of momentum parallel to the jet axis and E is the particle's energy. A particle's rapidity can be related to its velocity along the jet axis by the expression

$$y = \text{arctanh}(\beta_{\parallel}) \quad (2.17)$$

If the particles produced during quark and antiquark fragmentation are created with limited transverse momentum and according to longitudinal phase space, the resulting distribution of hadrons is expected to be flat in rapidity:

$$\int \frac{d^3p}{E} \rightarrow \int d^2p_{\perp} e^{-\frac{p_{\perp}^2}{2\sigma_{\perp}}} \underbrace{\frac{dp_{\parallel}}{E}}_{dy} \quad (2.18)$$

The maximum value of y increases with $\ln(E_{\text{beam}})$ since

$$y_{\text{max}} = \frac{1}{2} \ln \left(\frac{2E_{\text{beam}}}{m_{\pi}} \right) \quad (2.19)$$

Since most experiments cannot separate pions, kaons and protons over the whole momentum range, it is common to calculate y using the pion mass to determine E . Measurements of $d\sigma/dy_{\pi}$ have been done for energies between 5 and 32 GeV^{23,24}. These distributions exhibit the expected rapidity plateau (aside from a small dip at $y = 0$). We present results on inclusive rapidity distributions for identified pions, kaons and protons in Chapter 7.

2.4 Hadronization and Confinement

While the high energy behavior of QCD can be explored using perturbation theory, the infrared behavior of the theory is not well understood. It is generally hoped that the rise of the coupling constant for long distances results in a confinement mechanism, explaining the fact that all observed particles are color singlets. We can ask, however, whether such a speculation is reasonable within the field theoretical framework. Fortunately, there exists a solvable gauge theory exhibiting the properties of confinement. This theory is 1+1 dimensional QED (the Schwinger model)²⁵. An analysis of this theory can perhaps provide a clue to the mechanisms responsible for confinement in QCD. The QED Lagrangian in one space and one time dimension is

$$\mathcal{L} = -\frac{1}{4}F_{\mu\nu}F^{\mu\nu} + \bar{\psi}i\gamma^{\mu}\partial_{\mu}\psi - g\bar{\psi}\gamma^{\mu}\psi A_{\mu} \quad (2.20)$$

where g is the coupling constant between the vector field A_{μ} and the massless fermion field ψ . If we impose the boundary condition that at time $t = 0$ the state consists of a fermion-antifermion pair receding from each other at close to the speed of light (the analog of high

Figure 2.7: The mechanism for charge screening in the Schwinger model. a) As the initial charged fermions separate, a constant electric field develops between them. b) fermion-antifermion pairs are produced in the field and align themselves to cancel the electric field from the primary charges.

energy $q\bar{q}$ production), the solution to these equations has the surprising property that the asymptotically observed particles are massive dipole states. This total screening of charge is the result of the extreme polarizability of the vacuum; the mechanism is displayed in Fig. 2.7. In one spatial dimension, the electric field from a charge is independent of distance. Thus, the field configuration for two oppositely charged particles moving in opposite directions is a uniform field of strength g (and thus uniform energy per unit length $g^2/4\pi$) between the particles and no field elsewhere. The solution to the Schwinger model shows that the vacuum will polarize in the region of uniform field, producing dipole pairs that exactly cancel that field. The physical states are bosons of mass $g/\sqrt{\pi}$.

The Schwinger mechanism provides a useful paradigm for hadron production in jets. Experimental data indicate that long distance behavior of QCD is transversely confined (i.e. the transverse dimension of jets is roughly 350 MeV/c). It is therefore reasonable to approximate the long-range strong interactions between a quark and an antiquark by a longitudinal flux tube (or "string") joining the q and $\bar{q}$. Although the electric field of QED is replaced by a color field, it is possible to make polarization arguments similar to those above²⁶. In this case, the polarization phenomenon can be interpreted as the creation of quark-antiquark pairs from the vacuum. The analog of charge-screening in the QED case is color-screening; only color neutral objects emerge asymptotically.

Let us examine this mechanism using a semiclassical picture of hadronization²⁷. In this model, quarks are approximated as the classical trajectories of massless particles in $1 \oplus 1$ dimension interacting via a constant force κ . Because the quarks are massless, they always move at the speed of light and the equation of motion is

$$\frac{dp}{dt} = \pm \kappa \quad (2.21)$$

where the sign of the force depends on whether the quark is to the left or the right of the other interacting particle. In analogy to the Schwinger model, we set the magnitude of κ equal to $g^2/4\pi$. The Hamiltonian for a bound state of a quark (with momentum p_1 and position x_1) and an antiquark (with momentum p_2 and position x_2) can then be written

$$H = |p_1| + |p_2| + \frac{g^2}{4\pi} |x_2 - x_1| \quad (2.22)$$

In the bound state center-of-mass frame, the motion of the pair is a "yo-yo" mode (see Fig. 2.8a). The energy of the bound state shifts between the kinetic energies of the quark and antiquark ($|p_1|$

Figure 2.8: The motion of a $q\bar{q}$ bound state in the semiclassical string model of Andersson *et al.*²⁷ a) in the $q\bar{q}$ center-of-mass and b) in a frame moving with respect to the $q\bar{q}$ center-of-mass. The cross-hatched areas show where the color field is non-zero.

and $\{p_2\}$ and the potential energy stored in the color field ($g^2/4\pi|x_2 - x_1|$). In the center-of-mass frame, the bound state has no overall momentum; the quark and antiquark turning points occur at the same time. The motion of the $q\bar{q}$ bound state in a frame moving with velocity $v = \beta c$ with respect to the center-of-mass is shown in Fig. 2.8b. Since the quark and antiquark equations of motion in the center-of-mass frame are

$$x = \pm t, \quad E = \pm p \quad (2.23)$$

the Lorentz boosted equations are

$$t' = \gamma(t - \beta x) = \gamma(1 \mp \beta)t, \quad x' = \gamma(x - \beta t) = \gamma(1 \mp \beta)p \quad (2.24)$$

In this frame, the quark and antiquark turning points occur at different times.

The mechanism for hadron production in e^+e^- annihilation is represented in Fig. 2.9. The process begins with the creation of an initial $q\bar{q}$ pair at $x = 0$, $t = 0$. The quark and antiquark travel in opposite directions at the speed of light, creating a color field between them. At a later time, $q\bar{q}$ pairs are created in this field. If pairs $q_1\bar{q}_1$ and $q_2\bar{q}_2$ appear at space-time positions (x_1, t_1) and (x_2, t_2) respectively, and if q_1 and $\bar{q}_2$ combine to form a meson, that meson will have energy and momentum

$$E = \frac{g^2}{4\pi} |x_2 - x_1| \quad (2.25)$$

$$p = \pm \frac{g^2}{4\pi} (t_2 - t_1)$$

where the plus sign corresponds to the case where $x_2 > x_1$. To insure that this $q_1\bar{q}_2$ state is a physical particle, Andersson *et al.*²⁷ introduce a mass quantization condition requiring that the

Figure 2.9: The space-time development of a $q\bar{q}$ jet. As the primary quark and antiquark separate, the field between them is neutralized by $q\bar{q}$ pairs pulled from the vacuum. Note that the mesons produced at earliest times are the slowest.

pair production point (x_2, t_2) occur along the hyperbola

$$\left(\frac{g^2}{4\pi}\right) |(x_2 - x_1)^2 - (t_2 - t_1)^2| = m^2 \quad (2.26)$$

where m is the mass of a known meson.

It is worthwhile to make some comments about the space-time structure of the configuration described above. Eq. 2.26 implies that the $q\bar{q}$ creation points are not causally connected and hence are not time-ordered in any frame-independent sense. However, the fact that pair production can occur only in the regions of non-zero color field (the hatched regions in Fig. 2.9) means that in any Lorentz frame the final hadrons are on average ordered in velocity (or rapidity), with the slowest particles being produced first. This structure is known as an "inside-out" cascade. One important result of this approximate ordering is that the q and $\bar{q}$ pulled from the vacuum tend to appear in hadrons close together in phase space. As a result, the model predicts that a particle's quantum numbers should be compensated by other hadrons nearby in rapidity. An exception to this rule applies, of course, to hadrons containing the initial partons. Their quantum numbers should be balanced by particles in the opposite jet.

The fact that each $q\bar{q}$ pair pulled from the vacuum in Fig. 2.9 appears at a single point in space-time is possible because the q and $\bar{q}$ are massless and are produced with zero relative momentum. If we extend the above model to give quarks mass or transverse momentum, energy conservation requires that a finite portion of the flux tube disappear when the partons materialize. This process corresponds to quantum mechanical tunneling. In the Schwinger model, the probability for a pair of particles with transverse momentum $p_\perp$ and energy $E_T = \sqrt{m^2 + p_\perp^2}$ to appear as a physical state is^{28,29}

$$P(p_\perp) = \exp\left(-\frac{4\pi^2 E_T^2}{g^2}\right) \quad (2.27)$$

Figure 2.10: The mechanisms for producing a) mesons and b) baryons from the color field in the model of Casher *et al.*²⁸. The labels 1,2 and 3 refer to the color states of the quarks.

If we assume that this result holds for QCD as well, the tunneling mechanism produces a gaussian suppression of transverse momentum and a dynamic mass suppression for heavy quarks. Using quark masses of $m_u = m_d \sim 0.35$ and $m_s \sim 0.5$, we find a suppression of strange quark production of about 0.3. The rate of $c\bar{c}$ production from the vacuum in this model is $\sim 10^{-11}$.

While soft production of heavy quarks is highly suppressed, hard production of charm and bottom quarks from the virtual photon is copious. For quarks with masses that are non-negligible relative to Q^2 , the flux tube picture must be modified slightly. The assumption that the initial partons separate at close to the speed of light is no longer valid; the initial partons slow down as they separate. Since particles are only produced in regions of non-vanishing chromodynamic field, heavy quark events exhibit a suppression of the production of high z particles from the vacuum. Thus, the fraction of energy carried by mesons containing primary heavy quarks is on average larger than for the mesons containing the primary quark in light quark events. This kinematic argument was first made by Suzuki³⁰ and by Bjorken³¹, who predicted for large masses M that

$$\langle z \rangle \simeq 1 - \frac{1 \text{ GeV}}{M} \quad (2.28)$$

We shall see in the next section that the Lund group predicts an even harder heavy quark fragmentation spectrum. Studies of charm and bottom fragmentation show that³²

$$\begin{aligned} \langle z \rangle_{\text{charm}} &= 0.56 \pm 0.03 \\ \langle z \rangle_{\text{bottom}} &= 0.78 \pm 0.03 \end{aligned}$$

where $z = E_{\text{Hadron}}/E_{\text{Beam}}$. These data are consistent with the expectation that heavier quarks carry more of the jet energy.

Experimental results indicate that the production rate for baryons in high energy jets is reasonably large³³. The inclusion of baryons into the above framework has proceeded along two paths. In one treatment^{34,35} baryons are viewed as quark-diquark bound states, where a diquark is a qq color anti-triplet. A diquark-antidiquark pair is pulled from the vacuum in the same

manner as a $q\bar{q}$ pair. The suppression of baryon production relative to mesons then is due to the fact that diquarks are heavier than quarks and are thus suppressed by the mass term in Eq. 2.27. Although such a model reproduces the known data quite well (see Chapter 6 for details), it has the peculiar feature of treating composite objects (diquarks) as primary. This difficulty is avoided in the work of Casher *et al.*²⁶, who demonstrate that baryon production is a natural consequence of the color structure of the flux tube and the fact that gluons carry color (*i.e.* the fact that QCD is non-abelian). They argue that because the initial $q\bar{q}$ pair is created from a photon, it is in a color singlet state. The color flux tube joining the quark and antiquark is thus in a well defined color configuration, corresponding to the exchange of diagonal chromostatic field lines (*i.e.* only gluons corresponding to the SU(3) matrices T_3 and T_8 are exchanged). It is possible therefore to calculate the probability that $q\bar{q}$ pairs of each of the 3 colors will be produced from the field. The most likely case is that the vacuum produced pair have the same color as the initial pair. Then, the $q\bar{q}$ align with the color field, the flux tube breaks and mesons are produced. If, however, the vacuum pair has a different color than the primary quarks, the $q\bar{q}$ pair cannot neutralize the color field. Instead, the pair reverse align, producing a region of inverse direction field between them (see Fig. 2.10). The production of a third pair corresponding to the remaining color is necessary for the flux tube to break. This mechanism produces qqq and $\bar{q}\bar{q}\bar{q}$ states, *i.e.* baryons.

2.5 Monte Carlo Models of Fragmentation

The above description of hadronization is essentially probabilistic. The production of hadrons is a stochastic process with a space-time ordering that is only defined on average. In order to test these ideas, the experimentalist needs specific, quantitative predictions for the distribution of final state particles. These predictions must include the effects of energy-momentum conservation and of resonance production and decay. A number of Monte Carlo models have been developed to meet this need. Although these models differ significantly in the detailed implementation of hadronization, they all reflect several basic beliefs:

- QCD is the correct theory of the strong interactions. Perturbative QCD adequately describes the short distance behavior of quarks and gluons. Low Q^2 processes should not mask these hard scatterings.
- The long range behavior of the strong interactions is dominated by confinement. When colored partons separate, a color field develops between them. This field is neutralized locally, leading to short range compensation of flavor and color quantum numbers.

In this section we describe how these concepts are implemented in several popular Monte Carlo models. Understanding these models and comparing their predictions is important for two reasons. First, the interpretation of many measurements is done by comparing the data with Monte Carlo predictions. The experimentalist must determine whether such interpretations are sensitive to the details of the model. For example, studies of 3-jet events as a method for measuring α_s show that the value obtained for the strong coupling constant depends on the gluon fragmentation scheme used³⁶. On the other hand, the fact that gluon bremsstrahlung is necessary to explain the data is *not* model dependent. In addition, because the models differ significantly in their treatment of hadronization, a comparison of their predictions can perhaps be used as a tool for studying the confinement mechanism.

One of the earliest jet Monte Carlos was the Feynman-Field (FF) model³⁷. This model was never intended to be a realistic picture of hadronization. Its purpose was to parameterize the available data on jet production and thus allow a comparison of the jets produced in different processes. In FF, quark fragmentation is described recursively. An initial quark q_a is considered a source for the color field and a quark-antiquark pair ($q\bar{q}_b$) is produced in that field. The antiquark $\bar{q}_b$ then combines with q_a to form a meson, leaving q_b to continue the process. The cascade continues until some cutoff momentum is reached. The momentum carried by meson

$M(q_a, \bar{q}_b)$ is chosen from a probability distribution

$$f(\eta) d\eta = \begin{array}{l} \text{the probability that Meson } M(q_a, \bar{q}_b) \text{ leaves a fraction} \\ \text{of energy-momentum between } \eta \text{ and } \eta + d\eta \text{ to the} \\ \text{remaining jet} \end{array} \quad (2.29)$$

If we define z to be the light cone variable

$$z = \frac{(E + p_z)_{\text{Meson}}}{(E + p_z)_{\text{quark}}} \quad (2.30)$$

then η can be expressed as $\eta = 1 - z$. Feynman and Field parameterized the function f as

$$f(\eta) = 1 - a_F + 3a_F(\eta)^2 \quad (2.31)$$

In combining the $q_a \bar{q}_b$ pair, FF consider the possibility of vector and pseudoscalar spin states. Although spin counting would predict a production ratio $V/P = 3:1$, the relative masses of the states can introduce a suppression of the vector production rate. The FF model treats $a_v = V/P$ as a free parameter. Fitting to the available data gives a value $a_v \sim 1$. Similarly, FF set the rate of strange to nonstrange quark production experimentally to $u:d:s = 1:1:0.4$. (Results published since that time indicate that 1:1:0.3 gives a better representation of the data³².) In the FF model, transverse momentum relative to the jet axis is introduced by giving each produced $q\bar{q}$ pair equal and opposite $p_\perp$. The $p_\perp$ is drawn from a gaussian distribution with a $\sigma=320$ MeV/c.

The original FF model has a number of defects, which were first pointed out by the authors themselves:

- The model does not conserve energy or momentum since a massless quark is transformed into a number of massive mesons.
- Because the original FF Monte Carlo generates single jets, it always ends with an unused quark or antiquark. Thus, quantum numbers are not conserved.
- The model does not include baryons or higher spin resonances.
- Hard gluon bremsstrahlung is not included.
- The model is not Lorentz invariant.

In spite of these difficulties, FF was quite successful in reproducing the data. In particular, single particle inclusive distributions, rapidity distributions and transverse momentum spectra could all be adequately described with a single set of parameters. To correct some of the above problems, several extended FF Monte Carlos have been written. Most notable among these are the Hoyer³³ and Ali³⁹ models. In these Monte Carlos, energy and momentum conservation is approximately restored either by boosting and rescaling the energies in the event or by rejecting events that do not conserve energy within tolerance. In $e^+e^- \rightarrow q\bar{q}$ events, quantum number conservation is imposed by combining the remaining end quark and antiquark from the two jets into a soft meson. Baryon production is included using a diquark model³⁴ and a hard fragmentation function for heavy quark initial states is introduced. Finally, radiation of hard gluons is added.

The extended Feynman-Field Monte Carlos are called independent fragmentation models because the quark, antiquark and gluon each fragment as if the others are not present (see Fig. 2.11a). In these models the cross section for hard gluon bremsstrahlung is calculated using first or second order QCD (see Eq. 2.11), with energy and angular cutoffs set by requiring a minimum parton-parton invariant mass of a few GeV/c². Quarks are then allowed to decay, using the FF formalism described above. The fragmentation scheme for gluons varies from model to model. Ali splits the color octet gluon into a $q\bar{q}$ pair, with equal probabilities for u , d and s production, and splits the energy between the q and $\bar{q}$. Hoyer gives all of the energy either to

the q or $\bar{q}$. Both models then fragment the quark and/or antiquark using the FF fragmentation functions. After the quark, antiquark and gluon fragment, energy-momentum conservation is imposed either by rescaling the energies of the final state particles or by varying the angles between the parton jets axes. An alternative approach to this independent fragmentation scheme is the Lund string model⁴⁰, which we describe below.

In the Lund Monte Carlo⁴¹, two-jet events are generated in a similar manner to the extended FF models. The fragmentation starts at the initial quark or antiquark and proceeds via $q\bar{q}$ production from the vacuum. Meson energies are again selected using a fragmentation function and mesons are given transverse momentum by generating $q\bar{q}$ pairs with equal and opposite $p_{\perp}$. There is an important philosophical difference between Lund and FF, however. The Lund group views jet production as resulting from the evolution of a color flux tube ("string"), as described in the previous section. Hadrons are produced from the flux tube rather than from decays of the initial partons. The Lund model, therefore, does not allow single jet production. Jets can only be defined in terms of a source and a sink of the color field. The fact that the Monte Carlo treats the entire color system rather than creating independent jets allows an approach that both conserves energy-momentum and is Lorentz invariant. In addition, the Lund group has shown that the fragmentation function formalism used in both the FF and Lund models is not inconsistent with the inside-out cascade structure we expect from simple space-time arguments. Since in any Lorentz frame, the slowest mesons are the first to be produced, it is always possible to boost to a frame where the meson containing the primary quark is created first. Once that meson has been created, the remaining system is a jet that looks identical to the initial jet, except that its total energy is lower. The fragmentation function approach is equivalent to performing a series of Lorentz transformations to insure that at each stage of jet fragmentation the quark initiating the jet appears in the next meson created.

While almost any fragmentation function is consistent with the inside-out structure, only a limited range of functions fulfill the requirement of "left-right" symmetry. The final particle state should not depend on whether we choose to fragment from quark to antiquark or from antiquark to quark. The FF fragmentation function and the initial versions of Lund did not satisfy this consistency condition. Recently⁴², the Lund group has shown that a symmetric fragmentation function must have the form

$$f_{\alpha\beta}(z) = N_{\alpha\beta} \left(\frac{1}{z}\right) z^{a_{\alpha}} \left(\frac{1-z}{z}\right)^{a_{\beta}} e^{-bm_{\perp}/z} \quad (2.32)$$

where α and β correspond to the flavors of the initial quark and the quark member of the vacuum produced $q\bar{q}$ pair respectively. In practice, the Lund group has used the same value of a for all quark species.

The Lund belief that hadronization can be understood using string dynamics leads to their interpretation of a gluon as a momentum-carrying bend or "kink" in a color string. Let us take a $q\bar{q}g$ event as an example (see Fig. 2.11b). Since QCD contains a quark-antiquark-gluon color conserving vertex, a gluon can be interpreted as having a color and an anticolor index. Thus a $q\bar{q}g$ event consists of color 1 quark, a color 2 antiquark and a color 12 gluon. Color field lines stretch from the quark to the gluon and from the gluon to the antiquark. The flux tube therefore consists of two straight segments and a kink. In the Lund model, pair production occurs along the straight segments of the string. The kink is incorporated in a meson created from q and $\bar{q}$ breaks on either side of the bend. Once breaks have occurred on either side of the gluon kink, the $q\bar{q}g$ event looks like two $q\bar{q}$ strings, plus a leading meson at the gluon corner. Neither of these $q\bar{q}$ strings is at rest in the overall center-of-mass frame. Since these strings fragment in their own rest frames, the fragmentation products appear boosted in the center-of-mass, preferentially populating quark-gluon and antiquark-gluon regions relative to the region between quark and antiquark. The particles tend to lie on hyperbolas in momentum space whose asymptotes are the initial parton directions. Measurements of the distribution of particles in 3 jet events first made by the JADE group at PETRA⁴³ and later confirmed by the TPC⁴⁴ indicate that this string mechanism provides a better description of the data than the independent fragmentation models.

Figure 2.11: The development of a $q\bar{q}g$ event in a) an independent fragmentation model and b) the Lund string model. The arrows represent the momentum space distribution of particles. In independent fragmentation models, fragmentation products follow the initial parton directions. In Lund, the color field lines stretch from the quark to the gluon to the antiquark. Particles are produced along hyperbolas whose asymptotes are the $q\bar{q}g$ directions.

The string picture of gluon dynamics also provides an elegant solution to the problem of merging jets. In an independent fragmentation scheme, the distribution of fragmentation products for a quark jet and for roughly collinear quark and gluon jets with the same total energy are different. Therefore, the predicted particle distributions are dependent on the angular and energy cutoffs for perturbative gluon bremsstrahlung. In the string model, the distribution of final state hadrons can be made independent of which configuration is present by an appropriate treatment of the gluon kink⁴⁵.

All of the models discussed above treat hadronization as the vacuum production of $q\bar{q}$ pairs and the combination of neighboring partons into mesons and baryons. The role of the gluon in these Monte Carlos is reduced to a limited number of hard emissions (plus of course the implicit role it plays in the confinement process). An alternate approach is offered by the class of QCD-cluster models⁴⁶. Proponents of the QCD Monte Carlos argue that the creation of a direct link between partons and mesons is dangerous and unnecessary. For example, the imposition of the mass constraint in Eq. 2.26 is solely a result of requiring that the outgoing states correspond to pions, kaons or other physical particles. An alternative approach would be to think of these states as clusters that then decay into mesons and baryons. The QCD models go even further, however. They also remove the $q\bar{q}$ vertex as the primary mode for color screening. Instead, in these models the initial quarks radiate gluons and these gluons then radiate more quarks and gluons according to perturbative QCD. The process continues until the invariant mass of the remaining parton-parton pairs reaches some predefined mass cutoff. The results of this parton shower are then combined to form low mass colorless clusters. (see Fig. 2.12). The clusters then decay into hadrons, either according to two-body phase space or through a parameterization of low energy particle production data. Since the rate for this branching process can be calculated using perturbative QCD, the cluster formalism does not rely on fragmentation functions. The cluster decay products are chosen by spin counting and phase space factors; there are no $u:s$ or diquark:quark suppression factors. The relative $\pi:K:p$ production ratio is the result of kinematic suppression. These ratios are tuned by varying the mass cutoff for cluster production. The observed $p_{\perp}$ distribution is also generated solely as a byproduct of cluster decays.

A number of QCD based cluster models now exist. They differ in their treatment of where to cut off the perturbative shower, how to handle soft gluons and how to parameterize the cluster decays. One of the earliest cluster models is the Field-Wolfram Monte Carlo⁴⁷. Here, the perturbative shower is extended to very low energies (typical cutoffs are 1-2 GeV) and clusters

Figure 2.12: The evolution of a $q\bar{q}g$ event in a QCD cluster model. The dotted ellipses represent color singlet clusters.

are transformed into hadrons via two body decays. The Webber model⁴⁸ also uses a small cutoff and two-body cluster decays. However, the treatment of soft gluons in this model is quite different. QCD calculations show that the cross section for multiple soft gluon radiation shows significant interference effects. The origin of the interference is the fact that a soft (long wavelength) gluon cannot resolve multiple collinear gluons and therefore cannot separately couple to them. Therefore, the total rate of multiple soft gluon emission is lower than a naive calculation would predict. It can be shown that this interference effect is equivalent to an angle ordering for successive gluon emission⁴⁹. The Webber model introduces angle ordering into the Monte Carlo. This ordering can produce measurable changes in the predicted particle spectra. For example, the Webber model predicts a slight dip in the central rapidity region for e^+e^- annihilation.

A different approach to QCD clusters has been taken by Gottschalk⁵⁰. While Field-Wolfram and Webber push the perturbative showering to low masses where the QCD calculations are no longer reliable, the cutoff for perturbative radiation in the Gottschalk (CIT) model is large (about 5 GeV). The high mass clusters that result are then allowed to decay into lower mass clusters using a string formalism similar in many ways to Lund, without the mass constraint of Eq. 2.26. These low mass clusters then decay into hadrons. The parameterization of the cluster decay spectrum is made by fitting low energy data⁵¹.

2.6 Overview and Experimental Program

Measurements of R and the observation of jet structure in multihadronic annihilation events provide support for the belief that the process $e^+e^- \rightarrow \text{hadrons}$ occurs through the production of two back-to-back partons. The angular distribution of jetlike events indicates that these partons have spin $\frac{1}{2}$. The observation of planar 3-jet events at PEP and PETRA provides evidence for emission of vector bosons, presumably gluons. Studies of heavy quark fragmentation functions show that charm and bottom mesons carry a significant fraction of the jet energy and that the heavier quark (bottom) has a harder fragmentation. It appears, therefore, that the quark-parton model augmented by perturbative QCD successfully describes the hard scattering process in e^+e^- annihilation.

In contrast, there is surprisingly little experimental data that tests our basic assumptions about hadronization. The picture of local color neutralization, which is central to all current models of hadronization, remains essentially untested. We might ask whether the success of the previously described Monte Carlo models in reproducing the observed inclusive hadron spectra provides evidence that the basic beliefs of the models are correct. Unfortunately, it does not. The shape of the inclusive spectra is a reflection of longitudinal phase space. It indeed results from the jetlike nature of hadronization, but tells us little about the mechanisms involved. The overall normalization of the Monte Carlo predictions contains even less information. The models all contain adjustable parameters that are tuned to reproduce the observed multiplicities.

The demonstration that color is locally neutralized must therefore be based on the analysis of particle correlations. Studies of long-range (opposite jet) charge correlations by previous experiments have already provided evidence that the primary partons produced in e^+e^- annihilation are charged^{52,53}. Attempts to use short-range charge correlations to examine local charge neutralization are complicated, however, by the large number of resonances produced during hadronization. Resonance decays produce trivial correlations in the data, masking the correlation signal expected from hadronization.

The problem of resonance decays can be partially overcome by studying particles other than pions. In this thesis, we present a measurement of the strength of the correlations between kaons. Because these particles are less common than pions and because few resonances decay into K^+K^- pairs (the ϕ being a notable exception), strangeness is a much cleaner way to label specific $q\bar{q}$ pair than charge. Our experimental program is

1. To measure the inclusive π , K and p cross sections in $e^+e^- \rightarrow \text{hadrons}$. We expect the shape of these distributions to be consistent with the predictions of longitudinal phase space models. It is important, however, to verify this point experimentally for each particle species. In addition, measurements of the total production rates and momentum dependence of the cross sections can be used to tune the Monte Carlo models.
2. To study correlations between particles, in particular to study correlations between kaons. The primary goal of this analysis is to provide evidence that the picture of hadronization via local color screening is reasonable. Since this mechanism is incorporated in all the Monte Carlo models described above, we expect them all to reproduce the data, at least qualitatively. Providing quantitative measurements that can perhaps distinguish between the models is only a secondary goal.

Chapter 3

PEP and the PEP-4 Detector

The study of multihadronic final states in e^+e^- annihilation requires a facility capable of detecting particles over close to 4π steradian solid angle with high efficiency, even in large multiplicity or extremely jetlike events. The goal of the PEP-4 experiment is to provide such detection both for charged particles and photons and in addition to provide mass identification for all stable charged particles over close to the whole momentum range available at PEP. The central detector of the PEP-4 system is the Time Projection Chamber (TPC). A large volume drift chamber, the TPC provides simultaneous charged particle tracking and mass identification with a high level of segmentation. In this chapter we describe the e^+e^- storage ring PEP and the PEP-4 detector. Since the analysis in this thesis relies exclusively on data from the TPC, we describe the other detector elements only briefly.

3.1 PEP

The PEP storage ring is located at the Stanford Linear Accelerator Center (SLAC) and has a circumference of about 2.2 km. The ring was completed in 1980 and has run at a single center-of-mass energy of 29 GeV since that time. PEP consists of six bending arcs and six straight sections. Electrons and positrons are accelerated in the 2 mile long SLAC linac and are injected into the northwest and southwest straight sections. During normal operation, the ring contains 3 bunches each of e^+ and e^- , providing collisions at an intersection point in each of the straight sections. Crossings occur every 2.45 μ sec. Each bunch has a transverse size of about 500 μ m in the plane of the accelerator and 50 μ m out of the plane. The bunch length is approximately 1.5 cm. Typical peak luminosities are $2-4 \times 10^{31} \text{ cm}^{-2} \text{ sec}^{-1}$. This corresponds to a total rate for $e^+e^- \rightarrow \text{hadrons}$ of approximately 10^{-2} events per sec.

3.2 The PEP-4 Facility

The PEP-4 detector is located in the northeast straight section of PEP (Interaction Region 3). A schematic view of the facility is shown in Fig. 3.1 and its elements are summarized in Table 3.1. The detector has a cylindrical geometry, with the beamline located along the axis of the cylinder. During the 1983 running cycle, a 4 kG axial magnetic field was provided by an aluminum coil located 107 cm from the beamline. Future data will be taken using a 14 kG superconducting coil. The region inside the magnet coil is filled with pressurized gas. The operating mixture is normally 80% argon-20% methane at 8.5 Atm. The gas is contained at the inner radius by an aluminum pressure wall, at the outer radius by the magnet coil and at the endcaps by the magnetic pole-tips. Three detector elements are contained within the pressure volume: a cylindrical drift chamber at small radius (IDC), the TPC itself (which we describe in the next section) and electromagnetic calorimeter modules mounted on the pole-tips (PTC). Outside the gas volume and the coil are the outer drift chamber (ODC), six electromagnetic calorimeter

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Figure 3.1: The PEP-4 Facility

Detector	Fraction of 4π	Location	Comments
Inner Drift Chamber (IDC)	95%	13–19 cm radius	4 layers
TPC	97% [†]	20–100 cm radius	up to 185 wire hits, 15 pad hits
Magnet Coil	–	1.075 m radius	4 kG, Al coil
Outer Drift Chamber (ODC)	77%	1.19–1.24 m radius	3 layers
Barrel Calorimeter (HEX)	75%	1.2 m radius	limited Geiger mode, 10 X_0 deep
Pole Tip Calorimeter (PTC)	18%	endcaps	proportional mode
Muon detection	98%	endcap and barrel	extruded Al proportional tubes

[†] ≥ 2 pad rows and 15 wires

Table 3.1: TPC Detector Elements

modules (HEX) and muon detectors. The ODC is mounted directly on the magnet coil and the HEX is located around it. Muon chambers are mounted both in the barrel and forward regions.

The inner drift chamber (IDC) is 1.2 m long and contains 4 axial layers. Each layer consists of 60 sense wires, equally spaced in azimuth. Neighboring layers are offset by half a cell spacing (3°). The IDC is a main element of the charged particle trigger. It provides a fast timing signal, used to determine if an event is synchronous with beam crossing and if the tracks point back to the beamline. During the 1983 running, portions of two IDC layers experienced voltage breakdown. The neighboring layers were therefore ganged together in the trigger; the IDC was essentially a two layer chamber for triggering purposes. Although IDC data can in principle improve the momentum resolution of tracks found in the TPC, this technique has not been used with the present data.

The pole-tip calorimeter is a gas sampling calorimeter consisting of 51 layers of Pb-laminate and proportional wires. The calorimeter is 13.5 radiation lengths deep and contains 60° stereo views, with projective readout. The angular resolution is 8 mrad and the detector covers 18% of 4π . In addition to providing information on low angle photons and electrons, the PTC serves as the primary luminosity monitor for the PEP-4 detector.

The ODC is 3 m long and covers 77% of the solid angle. It contains 3 axial layers and is operated in 1 Atm 80% Argon–20% methane gas. Each layer contains 60 sense wires. The ODC, like the IDC, is primarily a triggering device. It is also used to tag photons that convert in the magnetic coil package.

The HEX is a gas, Pb-laminate sampling calorimeter run in limited Geiger mode. It consists of 6 modules, each 10 radiation lengths deep. The 40 gas gaps per module contain three 60° stereo views, one anode wire and two cathode strips. Each module is divided into two submodules in depth (27 and 13 layers, respectively). The detector covers 75% of the solid angle. During the 1983 running, the HEX used a 1 atm 96% argon–4% ethyl bromide gas mixture. This gas was discovered to be highly corrosive and eventually disabled 2 of the 6 modules. The modules have since been repaired and a different gas will be used in the future.

The PEP-4 facility has muon identification over 98% of 4π . The muon chambers consist of extruded aluminum proportional tubes operating in 1 Atm 80% argon–20% methane gas. Muons passing through the barrel region cross 4 layers of detector and at least 90 cm of iron. The

forward regions contain 3 muon layers. Tracks here pass through at least 50 cm of iron.

3.3 The Time Projection Chamber

The Time Projection Chamber (TPC) is a large volume drift chamber with a length of 2 m and a radius of 1 m (see Fig. 3.2). The axis of the TPC coincides with the beamline and the e^+e^- interaction point is at the center of the chamber. Charged particles produced at the interaction point pass through the TPC, curve in the solenoidal magnetic field and leave a trail of ionization in the gas. The ionization electrons drift to the TPC endcaps under the influence of an axial electric field. Here, they are detected using multiwire proportional chambers. Up to 15 three-dimensional space points are measured for each charged particle, with a typical accuracy of $185\ \mu\text{m}$ in the bending plane (x - y) and $350\ \mu\text{m}$ in the drift (z) direction. In addition, up to 183 measurements of the track's ionization energy loss (dE/dx) are made. These dE/dx measurements provide a means of separating charged particle species (see Chapter 5). In the following subsections, we describe the main elements of the TPC.

3.3.1 Electric Field Shaping

The TPC drift field is established by a central conducting membrane held at negative high voltage, normally $-75\ \text{kV}$. Since the the TPC endcaps are held at ground potential, an axial electric field is produced in each half of the detector. To insure a uniform electric field throughout the TPC volume, a series of field shaping electrodes are placed at the inner and outer radii of the TPC. These structures are known as the inner and outer field cages. The inner field cage is a hexagon made of fiberglass-epoxy laminate (Nema-G10); the outer field cage is a cylinder of the same material. The surfaces of both field cages are plated with a series of equipotential guard rings. The voltage drop between rings is established by a precision resistor-divider network. The first and last rings are connected (via resistors) to the central membrane and TPC endcaps, respectively.

3.3.2 Sectors

Each endcap of the TPC is divided into six multiwire proportional chambers, called sectors (see Fig. 3.3). The sectors are wedge shaped and occupy the region between $r = 20\ \text{cm}$ and $r = 100\ \text{cm}$. Each chamber contains 183 sense wires, with a wire spacing of $4\ \text{mm}$. Under 15 of the sense wires, the cathode plane is segmented into $7.5\ \text{mm}$ square pads. These pad rows are almost equally spaced, with the first and last rows located 23.6 and $95.2\ \text{cm}$ from the beamline, respectively. The standard sector coordinate system is defined so that the ξ axis points radially outward along the sector centerline and the η axis runs in the wire direction. The z axis points in the beam direction, perpendicular to the sector plane.

The sector wire geometry is shown in Fig. 3.4a. The sense wires lie in a plane $4\ \text{mm}$ above the grounded cathode. Field shaping wires are placed in the same plane, half way between the sense wires. In addition to improving azimuthal gain uniformity, these field wires substantially reduce cross talk between neighboring sense wires. A wire plane with $1\ \text{mm}$ pitch lies $4\ \text{mm}$ above the sense plane, approximately balancing the field of the cathode below. This grid separates the avalanche and drift regions. When a track drifts to the TPC endcaps, the ionization follows the electric field lines through the grid and is collected on the sense wires, where amplification occurs. Since no field lines originating in the drift region terminate on the grid plane, the grid is transparent to the ionization electrons. This argument does not hold for positive ions produced at the sense wire, however, due to azimuthal spreading of the avalanche. Approximately 40% of the positive ions liberated during the amplification process are collected on the grid, while about 20% reach the field wires. The remaining positive ions pass through the grid into the TPC drift region and slowly move to the high voltage midplane. This feedback of positive ions into the drift volume is a major source of electrostatic distortions (see Section 4.4).

TIME PROJECTION CHAMBER

Figure 3.2: The TPC detector

Figure 3.3: A TPC Sector.

(a)

(b)

Figure 3.4: a) The sector wire geometry and b) the relative wire and pad positions. This plot shows, as well, the coupling of the cathode pads to ionization signals on the wires above.

Plated feedthrough holes drilled in each cathode pad are used to transfer the pad signals to the sector backplane. Wire signals are transferred via printed circuitry along the sector edge. Both wire and pad signals are then fed to motherboards that hold the preamplifiers. In order to fit in the requisite space, each preamp is only 50 mm x 18 mm x 6 mm in size. Because the gain in a proportional chamber is a strong function of density and hence (for a fixed pressure) of temperature, power consumption on the preamplifiers must kept to a minimum. (For the TPC sectors, a 1°C change in temperature results in a 3% change in gain.) The typical dissipation per preamplifier is 100 mW. This heat is removed by a water cooling system anchored to the sectors and in thermal contact with the electronics. During the design of the electronics, special care was taken to avoid ferromagnetic substances (which could distort the magnetic field shape near the sectors) and materials that outgas electronegative substances (which would poison the active gas volume).

To minimize the problem of ground loops, the preamplifiers were designed with current drive output. The input charge from the chamber is integrated and an output signal with risetime below 50 ns is produced. The decay time of the signal is 5 μ sec; the dynamic range of the preamplifier allows more than 20-40 minimum ionizing tracks before saturation.

The output of the 16000 preamplifier channels is transmitted via twisted pair cable to high pressure feedthroughs. The signals then travel on coaxial ribbon cable to the remainder of the electronic chain. Shaper amplifier circuitry transforms the preamplifier signals to roughly gaussian output with a peaking time of about 250 nsec. Each TPC wire amplifier also contains a computer programmable discriminator. The discriminator outputs are used by the PEP-4 trigger (see Section 3.4).

The shaped wire and pad signals are the input to charge coupled devices (CCD's). A CCD is actually an analog shift register. The input to the CCD is sampled at a fixed frequency and the resulting voltage is stored in a cell or bucket of the device. As the CCD is clocked, the content of each bucket is shifted to the next cell. The devices used by the PEP-4 detector (Fairchild 321A) contain 455 buckets per channel. The TPC data is read into the CCD's at a rate of 10 MHz. Since the electron drift velocity in the TPC is about 5 cm/ μ sec and the maximum drift distance is 1 m, the history of an event fits into 200 CCD buckets.

If an event pretrigger (described below) is present, the electronics system goes into readout mode. The CCD clocking frequency is reduced to 20 kHz, allowing time for each bucket to be digitized (9 bit accuracy). Signals above a preset threshold are then fed to a large data buffer and transferred to the online computer. The threshold ("lower limit RAM") for each digitizer channel is individually set. The appropriate threshold value is determined from the pedestal and rms noise levels measured for the channel. Typically, 5 to 7 CCD buckets read above threshold per channel per track.

Both the dE/dx and tracking functions of the TPC require precise electronic calibration. Since the intrinsic dE/dx resolution is of the TPC is of order 3%, the response of the wire electronics must be known to 1%. The spatial resolution of the chamber depends on our ability to determine the center of gravity of each ionization cluster. A 5% error in the relative calibration of neighboring pads results in an 80 μ m error in reconstructed position. The electronic calibration of the TPC is performed in two stages. First, pedestal levels are determined as a function of bucket number. Second, the shape of the gain curve is measured. Electronic pedestal and noise levels for each channel are found by setting the digitizer threshold to zero and initiating readout of the complete CCD. A least-squares straight line fit to the output gives the pedestal level, the variation of pedestal with bucket number (over the full CCD readout, dark current gives a pedestal slope of typically 1-2% of the pulse height for a minimum ionizing track) and the rms noise. The shape of the gain curve is measured by pulsing each preamplifier input using a precision pulser. This calibration is used to correct for nonlinearities in the electronic gain (in particular in the CCD's). The pulser has a differential linearity of better than 1% and an integral linearity of .2%. Coupling of the calibration signal to both the wire and pad electronics is done by pulsing the grid above the sense wires. This method insures that the coupling is the same for all pads within a row. (The strategy fails for the pads at the edge of the sector, since they have

a different capacitive coupling to the grid wires. A satisfactory scheme for calibrating edge pads has never been implemented.) The wire coupling, however, varies with wire length. Although we correct for this effect in the calibration, the absolute wire gain (chamber plus electronics) is in fact measured using radioactive sources. The procedure is described in Chapter 5. For both wires and pads, the shape of the gain curve (input vs. output) is characterized by an 11 point multilinear fit⁶⁶. Deviations of the measured response from the fit are typically less than 1% for most of the dynamic range. Stability between calibrations is approximately 1% as well.

3.4 Trigger

The PEP-4 trigger electronics consists of two stages. A "pretrigger" with a rate of about 1 kHz selects event candidates within 2 μ sec of beam crossing. This decision time allows about 500 nsec before the next crossing for clearing analog registers if no pretrigger is present. The time needed for ionization electrons to drift a full meter in the TPC is approximately 20 μ sec. A second level trigger uses this 20 μ sec of TPC data to reduce the trigger rate to 1-3 Hz. In the rest of this section, we describe the criteria used to define track candidates in the pretrigger and final trigger.

Both the pretrigger and trigger electronics use the TPC wire discriminator data after the output of many channels have been ganged together to form larger blocks called "majority units." The logical **OR** of signals from neighboring sectors are joined to form six overlapping "supersectors" and the 183 supersector signals are divided into 23 groups of eight wires each. Each group of eight is connected to a majority logic circuit. This circuit fires, giving a majority unit hit, if it detects signals on a predetermined fraction (usually set to 3 out of 8) of the channels within a 1-2 μ sec time window.

The purpose of the charged particle pretrigger is to find track candidates by matching hits in the inner drift chamber with signals from either the TPC or outer drift chamber. Since most ionization in the TPC does not reach the endcap within the 2 μ sec pretrigger decision time, TPC data are only used to select low angle tracks that penetrate the endcaps. The ODC is used to define tracks at angles of greater than 45° with the beamline. A pretrigger track in the IDC is defined by requiring matching hits in two azimuthally aligned cells within a ± 28 nsec time window. These hits must be coincide with beam crossover to within 0.5 μ sec. The resulting signals are then combined (logical **OR**) to form overlapping 30° segments. Similar procedures are used to produce 30° swathes in the ODC. A majority unit hit in a given TPC supersector is combined with two 30° sections of the IDC and a coincidence is defined as a track. Similarly, azimuthally consistent IDC and ODC sections are combined. A valid charged pretrigger consists of two such "tracks." From the trigger standpoint, two particles in the same supersector or the same 30° ODC section are counted as one track.

The second stage charged particle trigger relies exclusively on TPC wire data. For events with a valid pretrigger, the output of each TPC majority unit is divided into 64 time slices, corresponding to the full meter of drift in the TPC. The trigger is designed to search for patterns of majority hits as a function of time. It requires that these patterns be consistent with beam associated tracks. A valid charged trigger requires two independent "tracks", where a track is found by either of two sets of trigger logic. The "Ripple trigger" is used to find inclined tracks. The "Majority trigger" recognizes tracks roughly parallel to the midplane.

The Ripple trigger searches for series of majority units firing successively with decreasing radius. When a unit fires, it activates a search for signals in any of the next 3 lower radius majority units in the supersector. (The 3 unit search allows for gaps due to inefficiencies at sector boundaries) A ripple consists of a set of majority units hits in one supersector, terminating with one of the two inner radius groups. The time of arrival of the innermost majority unit must coincide with a timing signal set to select events coming from the beam crossing point.

Because the Ripple search is initiated by TPC signals themselves, this trigger does not recognize tracks where all of the ionization arrives at the same time. A second trigger, the TPC

Majority trigger, is used to select tracks making an angle of more than 84° with the beamline. This trigger simply asks how many majority units in each supersector fire during a time window about beam crossing. If more than a preset fraction fire, a Majority trigger is satisfied in that supersector. Both the Ripple and Majority trigger can localize tracks to within ± 25 cm from the beam interaction point.

In addition to the charged particle triggers, the PEP-4 detector has several neutral energy and combined charged plus neutral triggers. We shall not discuss these triggers here.

Chapter 4

Event Reconstruction and Event Selection

The physics analysis presented in this thesis is based on the study of data collected during Spring 1983. In this chapter we describe the procedure by which data taken during that running cycle were reduced to 23,000 reconstructed multihadron events. We only discuss those aspects of the analysis relevant for the TPC detector.

The PEP-4 analysis system is shown schematically in Fig. 4.1. The data are reduced in 5 stages or passes. The online acquisition system removes triggers not consistent with beam-related events and writes the remaining data to magnetic tape. Electronic calibration and pattern recognition are performed in the first offline pass, along with additional event filtering. At this stage, the data are processed through track fitting and dE/dx software, using nominal reconstruction constants. The results of this pass are used to refine the values of such run dependent constants as the average gain, the electron drift velocity and the average beam-beam interaction point. The data are then reprocessed using the final constants. Finally, physics samples are selected.

The system described above is a response to the size and complexity of the analysis software. Because the reconstruction programs are so large, running the complete analysis chain in one step is inefficient. Since only a small fraction of the code can be contained in the physical memory of the computer at one time, the computer continually shifts instructions to and from disk. Running too large a fraction of the analysis in one pass degrades system performance. In addition, the multi-pass structure allows us to run the data through the early stages of analysis while studying detector systematics, such as gain variations and distortions. Corrections for such systematic problems are not done until later in the analysis chain.

4.1 Pass 1: Data Acquisition and Online Event Filtering

Both the acquisition and the analysis of PEP-4 data are done using VAX 11/780 computers. As a result, the same software is used online and offline. The advantage of this approach is twofold. First, software developed for offline use can be transferred to online processing quickly and easily. For example, the online event filtering program was run as an offline process during early data taking. Second, a subset of events (randomly sampled) can be run through the complete analysis online. This sample provides a continual monitor of data quality.

The goal of the online filtering program, called PREANALYSIS, is to verify that a good trigger is present. Deadtime requirements place an upper bound on the processing time for online filtering of about 100 msec per event. (The total event rate is 1-3 Hz.) The charged particle PREANALYSIS searches for the tracks that triggered the TPC and tests whether these tracks point back to the beam crossing position. The distance of closest approach to the interaction point is calculated by fitting the output of the TPC majority units and raw data from selected pad rows in the TPC. These fits allow us to make tighter cuts on the event vertex position (thereby eliminating beam-gas and cosmic ray backgrounds) than are possible using the hardware trigger. The cut values vary with the quality of the fit, but typically require the track to extrapolate back

Figure 4.1: The PEP-4 analysis chain.

Figure 4.2: The CCD output as a function of bucket number for 3 contiguous pads. The z position for each pad hit is determined using a parabolic fit to the top 3 buckets. The data from neighboring pads is then combined to form a three-dimensional cluster.

to the interaction point to within ± 14 cm in the beam direction (z axis) and ± 10 cm in the plane perpendicular to the beam (xy plane). Events too complicated to be unraveled from the trigger latches are kept. Neutral particle triggers are subjected to PREANALYSIS cuts requiring that the energy deposited reproduces a reasonable shower shape. The online rejection rate for Spring 1983 was approximately 30%.

4.2 Pass 2: Pattern recognition

The first step in analysing TPC data is to translate electronic hits from the sense wires and cathode pads into space points, called clusters. We define these clusters in the coordinate system shown in Fig. 3.3. The η axis points in the wire direction, the ξ axis points in the radial direction at the centerline of the sector and the z axis points along the beamline. A wire cluster is a peak in z for a given wire number (ξ), while a pad cluster is a peak in the η - z plane for a given pad row (ξ).

To find wire clusters, the output signal for each channel is examined as a function of CCD bucket number. Contiguous buckets are combined to form z hits. The peak pulse height and z position of each cluster is determined by fitting the top 3 CCD buckets to a parabola and then performing an electronic calibration. The first stage of pad clustering is performed in the same manner as for wires. If hits with roughly the same z are found on contiguous pads, however, these hits are combined (see Fig. 4.2). In PASS 2, the η position of a pad cluster is calculated as the average of the η positions of the pads in the cluster, weighted by the pulse height seen on each pad. The ξ cluster position is set to be the centerline of the pad row and the z position is

the weighted average of the z positions of the individual pad hits. Later in the analysis chain, these η and ξ positions are refined using track information and wire pulse height data.

The TPC pattern recognition software at present works exclusively on pad data. This choice was made because the three dimension character of this data decreases the combinatorial background significantly, speeding up the software and decreasing the number of false tracks. The price paid for these advantages, however, is a decrease in efficiency for short tracks (a track must cross at least 2 pad rows to have any chance of being found). Future improvements of the pattern recognition software may use TPC wire data to augment the present algorithms. A discussion of the pattern recognition algorithms can be found in Hadley (1983).⁶⁰

The results of the pattern recognition programs are used for a second stage data filter. As in PREANALYSIS, the PASS 2 filter rejects events where the charged tracks do not pass sufficiently close to the nominal interaction point. Because a proper pattern recognition has already been done, the PASS 2 filter recognizes complicated junk events, such as electromagnetic showers from cosmic rays and beam-gas events, which the online filter kept as being too difficult to untangle. In addition, the distance of closest approach cuts are reduced to ± 10 cm in the z direction and ± 5 cm in the η - ξ plane. Because the momentum fitting at this stage of the analysis is still rather crude, no cuts requiring momentum balance are applied in PASS 2.

4.3 Pass 3: Monitoring of Constants

The track reconstruction and particle identification software for the TPC requires a large number of time dependent calibration constants. Many of these constants can best be determined from the data itself. Thus, reconstruction is an iterative procedure. The data is first analyzed using nominal constants and the results of this analysis are used to produce new constants. The 1983 PASS 3 processing was performed exclusively for measuring analysis parameters. No output tapes were produced.

The constants measured in this pass require TPC tracking information. Many require dE/dx data as well:

- The beam-beam interaction point is determined for each run by averaging the fitted event vertex for all multihadron events in the run.
- The variations of drift velocity with time are monitored by recording the time of arrival of clusters associated with tracks that cross the TPC midplane.
- The variations of sector gain with time and with drift distance are found by monitoring the value of dE/dx for minimum ionizing pions. This procedure is described in the Chapter 5.
- The size and shape of electrostatic distortions are determined by measuring the deviation of the wire cluster position from the fitted track as a function of radius. The correction of electrostatic distortions is described in the next section.

During future data taking these constants will be measured by online data sampling. Such an approach was not possible during the 1983 running first because the final software was not yet ready when the data were taken and second because the online computer could not handle the load. Time has solved the first problem and an upgrade of the computer (to a VAX 11/782) should solve the second.

4.4 Pass 4: Track Reconstruction

The process of track fitting can be divided into three steps. First, wire clusters are associated with the tracks previously found using pads; the wire pulse heights and the track data are then used to refine the pad space points. Second, the refined space points are refit. Finally, the fitted tracks are constrained to come from an event vertex.

A wire cluster is associated with a track if its z position is within ± 1 cm of the trajectory determined from the pad data. If more than one track fulfills this requirement, the cluster is called ambiguous. Because correct particle identification requires that the average energy loss be known to a few per cent, a stricter set of requirements is placed on the wire data used to calculate dE/dx . Because the recovery time of the analog electronics is long, clusters are eliminated for dE/dx calculations if there is another hit on the wire within ± 3 cm. To reduce biases due to wire cross-talk (which is about 5%), a hit is removed if one of its neighbors is saturated. With these cuts, the fraction of tracks with more than 80 dE/dx samples is 65%.

During the initial clustering of pad hits, the position of the avalanche along the pad row is approximated by a simple weighted average of the positions of the cathode pads. The spatial resolution of the sectors can be significantly improved if the average is replaced by a gaussian fit to the central position. For those space points where three cathode pads read above threshold (40% of the points), the pulse height, width and position of the gaussian can be simultaneously measured. If the cluster consists of only two pads, the width must be used as input to the gaussian fit. This width depends critically on track geometry. The intrinsic pad response is broadened if the track crosses the wires at an oblique angle (see Fig. 3.4). The width is also modified by diffusion (which is proportional to the square root of the distance the ionization cluster has drifted) and $E \times B$ effects (which depend on the angle the track makes with the sense wire). The width of the pad response as a function of these variables has been measured using three pad clusters and has been parameterized. We refer the reader to previously published results for the detailed form of these corrections⁵⁴.

Before a final fit to the refined spaced points is made, the clusters are corrected for electrostatic distortions. Inhomogeneities in the bulk properties of the material used to build the TPC field cages are one source of such distortions. In addition, positive ions liberated during the avalanche process drift into the main volume of the TPC, producing a radial electric field. These space charge effects produce the largest distortions seen in the detector. Because background radiation is preferentially emitted in the beam direction, the density of positive ions in the TPC decreases with radius. This ion distribution, coupled with the induced negative charge buildup on the inner radius field cage, produces a field that decreases with distance from the beamline. (In fact, close to the inner field cage the distortion field is not truly radial but instead is everywhere perpendicular to the hexagonal field cage.)

The result of a radial field on ionization tracks in the TPC is presented schematically in Fig. 4.3. For tracks that make an angle λ with the TPC midplane, we observe an apparent displacement Δz :

$$\Delta z = \Delta \xi \tan \lambda \quad (4.1)$$

Similarly, an apparent azimuthal displacement ($\Delta \eta = \Delta z \tan \phi$) is observed for tracks crossing the wires with an angle ϕ . A second azimuthal distortion is present due to $\mathbf{v} \times \mathbf{B}$ forces. Together these effects produce a displacement

$$\Delta \eta = \Delta \xi [\tan \phi + \mu B] \quad (4.2)$$

where μ is the mobility of electrons in the TPC gas. Combining Eq. 4.1 and Eq. 4.2 by eliminating $\Delta \xi$ gives

$$\Delta \eta = \frac{\Delta z}{\tan \lambda} [\tan \phi + \mu B] \quad (4.3)$$

Given a model of the field cage inhomogeneities and of the distribution of space charge within the detector, it is possible to calculate the electric field everywhere and hence the size of the distortions. However, the amount of background in the detector varies significantly with time and with the tune of the accelerator. Evidence indicates that time variations in the amount of charge buildup on the field cage cause time dependent distortions as well. Our approach therefore has been to measure the distortions in z on a track by track basis and to then relate these z distortions to η distortions (which bias the momentum measurement). For purely radial electric field, this relationship is defined by Eq. 4.3. A more complicated form is needed if the field

Figure 4.3: The mechanism responsible for electrostatic distortions a) in the η - z plane and b) in the η - ξ plane. A radial electric field produces apparent distortions both in z and azimuth. Further azimuthal distortions are caused by $\mathbf{v} \times \mathbf{B}$ forces.

Figure 4.4: Deviation of the measured z cluster position from the predicted value as a function of wire number. Predictions are obtained by fitting the region between wires 80 and 150. The above plot shows the mean deviations per wire, averaged over one run.

disturbance has azimuthal components. In practice, the analog of Eq. 4.3 is determined using a computer program that models the field produced by a sheath of positive charge close to the conducting field cage. The parameters of the electrostatic model are then tuned to reproduce the observed *shape* of the distortions as a function of radius and azimuth. Tracks are used to measure the overall size of the effect. For tracks with nonzero $\tan \lambda$, the z distortion is measured using wire clusters. A fit is made to the central portion of the track and the deviations from that fit are measured as a function of radius (see Fig. 4.4). For tracks with small $\tan \lambda$ and for those tracks where we cannot use wire data to determine the distortion size (for example, if an event contains two tracks in one sector with approximately the same value of λ , the wire hits for the two tracks cannot be distinguished), a nominal correction is made based on the average distortion size measured for that run in PASS3.

There is one class of distortions that cannot be measured using the above technique. An overall rotation of the track can always be accommodated in the track fit. We have studied these rotations using cosmic ray data, where the same track can be measured in back to back sectors and have tuned the electrostatic model to minimize such rotations. With the present analysis we can require that all tracks come from a common vertex without introducing an overall momentum bias. However, the measured momentum resolution is degraded by close to a factor of two by limitations in the distortion corrections.

Since the electrostatic distortions cause radial displacement of the track ionization, they influence dE/dx as well as spatial resolution. The distortion correction software is used to remove this effect as well. At this stage of the analysis, the dE/dx for each track is calculated and a preliminary mass identification is made. We defer discussion of the particle identification properties of the TPC to the next chapter.

After the space points have been corrected for electrostatic distortions and preliminary particle identification has been performed, the final track fitting software is called. The momentum of each track is fitted in the TPC and the fit is extrapolated to the beamline. The extrapolation includes multiple scattering and energy loss in the material before the TPC. These corrections

are made using the dE/dx mass assignment. The final stage of the fitting process is to include an event vertex in the fit. The mean beam position for each run is measured in PASS 3 to an accuracy of about 200μ in the x and y directions and about 500μ in z . Tracks consistent with being from the primary vertex are simultaneously refit under the constraint that they come from a common vertex. The location of this vertex is constrained to be at the nominal vertex position for the run, within the errors set by beamsize and our resolution for measuring the beam position. For high momentum tracks, this vertex constraint improves the resolution by close to a factor of two. The final momentum resolution achieved with the 1983 data is $(dp/p)^2 = (0.06)^2 + (0.035p)^2$, where p is the momentum in GeV/c. The first term is due to multiple scattering in the TPC gas. The second represents our fitting resolution.

4.5 Pass 5: Event Selection

Once the event reconstruction is complete, a final sample of events is selected. Separate data summary files are created for multihadron final states, $\tau\bar{\tau}$ events and two photon candidates. In this section we describe the selection criteria for e^+e^- annihilation into multihadron final states. These criteria have been developed using a Monte Carlo simulation of the PEP-4 detector. The simulation is described in the next section.

The main backgrounds to our multihadron sample are beam-gas scattering, two photon interactions ($e^+e^- \rightarrow e^+e^- X$) and multiprong τ decays. The goal of the selection software is to select single photon multihadron events with high efficiency and well defined acceptance while rejecting these backgrounds. During event selection, we wish to consider only tracks that have a good probability for coming from the beam-beam interaction point, that are adequately measured and that are well within the TPC acceptance. We therefore impose the following criteria:

- The track must pass the nominal interaction point within 10 cm in the beam direction and within 6 cm in the plane perpendicular to the beam.
- The track must make an angle of at least 30° with the beamline.
- The fitted track must meet either the condition $dc < 0.3$ or $dc/c < 0.3$, where c is the track curvature and dc is the estimated error on the measured curvature.
- The measured momentum must exceed 120 MeV/c. (Primary tracks with lower momentum will stop in the beampipe, pressure wall and inner field cage.)

Because there is approximately 20% of a radiation length in front of the TPC, a significant fraction of the photons convert to e^+e^- pairs. It is necessary to identify such pairs since γ conversions can make radiative bhabha or $\tau\bar{\tau}$ events look like multihadron candidates. Electrons are identified either by dE/dx information or by geometrical reconstruction of the pair. The efficiency for finding pairs is about 70% and the fraction of hadrons misidentified as electrons is a few percent. The criteria for event selection are

- Each event candidate is required to have at least 5 non-electron tracks meeting the above track cuts.
- To eliminate two photon events and beam-gas background, we require that the energy E_{ch} seen in charged particles (including electrons) exceed half the beam energy and that the sum of the momentum components in the beam direction be less than 40% of E_{ch}/c .
- The sphericity axis is calculated and the event is divided into two hemispheres. To eliminate tau candidates, we require events to have at least one hemisphere with either an invariant mass above 2 Gev or more than 3 non-electron tracks.

We estimate (via Monte Carlo) that our backgrounds are $0.3 \pm 0.1\%$ from $\tau\bar{\tau}$ events, $0.8 \pm 0.6\%$ from two photon events, and less than 0.1% from beam-gas and Bhabha scattering events⁵⁷. Before radiative corrections, the event acceptance is 77%.

4.6 Event Simulation

Although event simulation is not strictly speaking part of the analysis chain, it is an essential ingredient for interpreting the output. The PEP-4 analysis system includes two Monte Carlo simulations. One generates raw data for all of the detectors. The other produces space points and fits them to reproduce the output of the tracking software. Although generating raw data is useful for evaluating the analysis software, it is a slow task (a typical multihadron event contains 10–30,000 words of data). For calculating the acceptance, where high statistics are necessary, we therefore use the TPC Fast Monte Carlo (FMC). Typically, 10,000 multihadron events can be generated in 12 hours on a VAX 11/780 computer.

The FMC begins by generating an event using any of the physics Monte Carlos described in Chapter 2. Standard acceptance corrections have been calculated using the "symmetric" Lund generator. The effects of initial state radiation are simulated using the program of Berends and Kleiss⁵⁶. The event vertex is chosen by smearing the position of the interaction point to account for the PEP beam size. All charged and neutral particles are tracked through the detector and long lived particles (π , K, Λ) are decayed. Pad hits are generated for each charged track and the particle's truncated mean is calculated. For tracks with enough hits to be found by the pattern recognition software, the measured momentum is determined by fitting the pad points. At this stage the response of the other PEP-4 detectors is simulated. A data summary file with the same format as for the data is written, with an additional block giving information about the generated tracks.

The Monte Carlo models the PEP-4 detector as a series of slabs, each of which is assumed to have uniform density. Slabs are positioned to represent the beampipe and pressure walls. A single slab simulates the full distance between pad rows in the TPC; individual wire hits are therefore not generated. Charged particles are tracked from slab to slab including the effects of energy loss, multiple scattering, bremsstrahlung, decay in flight and nuclear interaction. Neutral hadrons can interact in the material and photons can convert to e^+e^- pairs. For each charged track, the η , ξ and z positions where the particle crosses the pad rows are translated into pad hits using a resolution function that includes the effects of diffusion, track angle and saturation. Pad hits are randomly removed to simulate the fraction of electronic channels that have readout problems. Inefficiencies due to sector boundaries and track overlap are also included. The Monte Carlo does not simulate the electrostatic distortions in the TPC. The resolution and the quoted measurement errors on the inner pad rows are increased, however, to reproduce the degradation of momentum resolution.

The program calculates N , the number of wires associated with each track, including the effects of geometry, sector boundaries, track overlap and electronic readout problems. The particle's dE/dx is calculated using the track's average $\beta\gamma$ in the TPC (see Section 5.4). This value is smeared according to a gaussian distribution with a width that varies with the number of wires: $\sigma \sim 3\% \sqrt{180/N}$. The effects of miscalibration and gain variations are mimicked by smearing a fraction (about 10%) of the tracks with a wider gaussian ($\sigma \approx 6\%$).

The TPC pattern recognition software is replaced by a simple algorithm. All pad hits that are within 3 cm of only one track are uniquely assigned to that track. Ambiguous hits are assigned only to one track. All tracks with 3 or more hits are considered found. The pad hits assigned to each track are fit to a helix to determine the particle's momentum in the TPC. The tracks are extrapolated to the beamline using a multiple scattering error determined from the track's momentum and a mass assignment based on the measured value of dE/dx . A vertex constrained fit is then done. Fig. 4.5 shows a comparison of the data and Monte Carlo for a number of variables used to define track quality. The positions of the track cuts listed in Section 4.5 are shown as well. Fig. 4.6 compares the data and Monte Carlo results for several distributions used to select multihadron events and Fig. 4.7 shows quantities related to dE/dx . In all cases, the simulation reproduces the experimental distributions adequately.

The fact that the Monte Carlo does not generate raw data does, however, introduce limitations on its accuracy. One serious problem is that the program does not necessarily reproduce pattern

Figure 4.5: A comparison of the data (the points with error bars) and the predictions of the Monte Carlo simulation (the histograms) for quantities used to select beam associated tracks. a) the curvature error dc/c , b) the distance of closest approach to the beamline in the bending plane, c) the distance of closest approach to the event vertex in z and d) the momentum spectrum. The arrows show the position of the track cuts described in Section 4.5.

Figure 4.6: A comparison of the data (the points with error bars) and the predictions of the Monte Carlo simulation (the histograms) for the distribution of tracks and energy in multihadron events. a) the number of charged hadrons, b) the energy found in charged hadrons, E_{ch} , c) the momentum balance along the beam direction, p_x/E_{ch} and d) the invariant mass per hemisphere, where every event is divided in half using the sphericity axis.

Figure 4.7: A comparison of the data (the points with error bars) and the predictions of the Monte Carlo simulation (the histograms) for dE/dx quantities. a) the number of dE/dx samples per track (numbers below 0 mean that pads rather than wires were used, b) the dE/dx distribution for minimum ionizing pions, c) the measured dE/dx for tracks with momenta between 410 and 500 MeV/c and d) the measured dE/dx for tracks with momenta between 3.7 and 6.7 GeV/c.

recognition inefficiencies. As a result, the treatment of conversion electrons (which are always produced as pairs close together in space) and of $\pi^\pm$ and $K^\pm$ decays (where the simulation must decide whether the pattern recognition software recognizes that there are two tracks) can be problematic. In addition, the fact that electronic saturation sets in at different pulse heights for different channels means that the calculation of the truncated mean for highly ionizing tracks involves larger than normal uncertainties. While these difficulties involve only a small fraction of the tracks and therefore do not significantly affect the distributions shown in Fig. 4.5 through Fig. 4.7, they can cause systematic problems for some physics analysis. In general, we have chosen to cut our data to avoid regions where the simulation is questionable. For example, by imposing a cut on the minimum number of wires found on the track, we can reject a large fraction of the conversion pairs while only losing a small fraction of primary tracks. Such a cut is used in Chapter 7, where a reliable estimate of the mass misidentification background is essential. Similarly, because the Monte Carlo has difficulty in simulating electronic saturation, we choose not to study tracks whose mean dE/dx is greater than 40 keV/cm.

Chapter 5

Particle Identification via dE/dx

In recent years, particle identification through the measurement of ionization energy loss (dE/dx) has been advocated by a number of physicists^{59,60,61}. The appeal of this method is largely a result of the failure of alternate approaches for relativistic particles. Time of flight identification becomes difficult for tracks with momenta above a GeV/c . Čerenkov systems often have difficulty in separating tracks that are close together in space. Ring imaging Čerenkov detectors provide an exciting possibility for the future⁶². However, no large systems of this kind are presently functioning.

The dE/dx method relies on the fact that the most probable value of the ionization energy loss in a sample of gas is a function of the incident particle's velocity rather than its momentum. Thus, a simultaneous measurement of momentum and dE/dx is equivalent to a measurement of momentum and velocity and therefore determines the particle's mass. In this chapter we describe how dE/dx information is used to identify particles. We begin with a discussion of ionization energy loss in gases. We show that the properties of the dE/dx distribution place important constraints on detector design and strict tolerances on detector calibration and discuss how these tolerances are met in the TPC. Finally, we discuss algorithms for particle identification.

5.1 The Theory of Ionization Energy Loss

When a particle passes through a gas it loses energy through collisions with electrons and nuclei. Because the energy transfer in such collisions is inversely proportional to the mass of the target particle⁶³, dE/dx losses are dominated by collisions with electrons. (In contrast, the angular deflection of the incident particle is dominated by collisions with the higher Z nuclei.)

It is useful to distinguish two regimes of energy transfer. For energy transfers large compared to atomic binding energies, electrons act essentially free. In this case the cross section per electron per unit energy follows the Rutherford scattering formula⁶⁴,

$$\left(\frac{d\sigma}{dE}\right)_{\text{quasi-free}} = \frac{2\pi e^4}{m_e c^2} \frac{1}{\beta^2 E^2} \int_0^E f(E') dE' \quad (5.1)$$

where $f(E')$ is the atomic oscillator strength. The integral of $f(E')$ gives the probability that the electron is in an atomic state whose energy is less than E , i.e., a state where the Rutherford formula is valid. For energy transfers comparable to the atomic binding energy, electrons act like harmonically bound charges. The cross section peaks sharply at the ionization potentials of the atom. This behavior can be described for a low density gas⁶⁵

$$\left(\frac{d\sigma}{dE}\right)_{\text{resonance}} = \frac{2\pi e^4}{m_e c^2} \frac{f(E)}{\beta^2 E} \left[\ln \left(\frac{2m_e c^2 \beta^2 \gamma^2}{E} \right) - \beta^2 \right] \quad (5.2)$$

It was first shown by Fermi⁶⁶ that Eq. 5.2 must be modified for large $\beta\gamma$. The expression was derived under the assumption that the electric field seen by an electron far from the moving

charge is unaffected by intervening matter. This is in fact not the case. Large transverse electric fields can polarize the medium. Electrons at large distances from the moving charge see a smaller electric field than the predicted value in vacuum. This phenomenon, known as the density effect, compensates for the logarithmic rise of the cross section with $\beta\gamma$. For a gas with complex dielectric constant ϵ the correction to the cross section has the form⁶⁴

$$\left(\frac{d\sigma}{dE}\right)_{\text{density effect}} = \frac{2\pi e^4}{m_r c^2} \frac{f(E)}{\beta^2 E} \left[\ln \left(\frac{1}{|1 + \beta^2 \gamma^2 (1 - \epsilon)|} \right) + \beta^2 (1 - \Re(\epsilon)) \right] \quad (5.3)$$

$$+ \frac{2\pi e^4}{m_r c^2} \frac{1}{\beta^2 (\hbar\omega_p)^2} \left(\beta^2 - \frac{\Re(\epsilon)}{|\epsilon|^2} \right)$$

where ω_p is the plasma frequency of the gas and $\Re(\epsilon)$ is the real part of the dielectric constant. The term on the second line represents the emission of Čerenkov radiation and is small in our case. Omitting this term and combining Eq. 5.1 through Eq. 5.3, we find

$$\frac{d\sigma}{dE} = \frac{2\pi e^4}{m_r c^2} \frac{f(E)}{\beta^2 E} \left[\ln \left(\frac{2m_r c^2 \beta^2 \gamma^2}{E |1 + \beta^2 \gamma^2 (1 - \epsilon)|} \right) + \beta^2 \Re(\epsilon) \right] \quad (5.4)$$

$$+ \frac{2\pi e^4}{m_r c^2} \frac{1}{\beta^2 E^2} \int_0^E f(E') dE'$$

This expression defines the energy dependence of the cross section for a charged particle interacting with a single electron. What is actually measured in a particle detector, however, is the total energy deposited within a sample of finite thickness. In a 4 mm sample of Argon-Methane gas at 8.5 atm, a minimum ionizing particle undergoes about 100 low energy collisions (M level), about 5 moderate energy collisions (K and L level) and about 10 Rutherford type collisions. The sum of these contributions gives a rather complicated spectrum of observed dE/dx per sample. A quantitative treatment of this energy loss is given in Appendix A; we merely outline the method here.

The majority of collisions in a gas correspond to low energy transfer interactions with far away electrons. The number of such collisions per unit path length is distributed according to a Poisson distribution with a mean value N

$$\left\langle \frac{dN}{dx} \right\rangle = \int \rho \frac{d\sigma}{dE} dE = N \quad (5.5)$$

where ρ is the number of electrons per unit volume. Since the oscillator strength $f(E)$ is highly peaked near the ionization potential I , the distribution of energy loss has a mean value of about $N I$. If the peaking is sharp enough, the width of the distribution is $\sqrt{N I}$. The close collisions corresponding to the Rutherford scattering limit occur much less often. Since $d\sigma/dE$ is inversely proportional to E^2 , large energy transfer collisions are especially rare. The chance of such a collision occurring in a given sample is much smaller than one. The effect of these collisions is to produce a $1/E^2$ falloff on the high end of the dE/dx distribution. This falloff is called the "Landau tail"⁶⁶ and the distribution of energy loss per sample is often referred to as the "Landau distribution." Combining the Poisson distribution of soft collisions and the hard "Landau tail" produces a very asymmetric distribution. On the low side, the distribution rises rapidly (the width being determined by the distribution of the soft collisions). The shape of the falling edge is determined by the probability distribution for K and L level collisions. At high dE/dx , the falloff is dominated by Rutherford collisions and is essentially $1/E^2$. (The true $1/E^2$ regime is beyond the dynamic range of the TPC. The high dE/dx tail we measure exhibits a $1/E^3$ falloff.)

The shape of the observed dE/dx distribution in the TPC is shown in Fig. 5.1. The width of the distribution is quite large, about 55% fwhm, and the data show the expected tail at high dE/dx . From the point of view of particle identification, this Rutherford scattering contribution to the distribution contains little useful information. For relativistic incident particles, it is the

Figure 5.1: Distribution of energy loss in a 4 mm sample of 8.5 Atm Argon-Methane gas. The histogram shows the measured dE/dx distribution in the TPC. The solid line gives the prediction of the model describe in Appendix A, with the normalization of the model chosen to give the same truncated mean as the data.

number of low energy transfer collisions that is sensitive to the speed. This number can be estimated by taking the mean of a fixed fraction of the dE/dx samples with the lowest pulse height. The dE/dx resolution is not very sensitive to the choice of fraction within the range of .4 to .75. We choose to estimate the most probable energy loss per sample by the mean of the lowest 65% of the samples. We refer to this quantity as the particle's measured dE/dx . For cases where track overlap prevents reliable determination of the dE/dx from the wire data, the truncated mean is calculated using pulse heights measured by the cathode pads. Since each pad signal integrates over several wires, a pad hit has more statistical significance than a wire hit. We switch to pad dE/dx if $N_{\text{pads}} \geq 0.4 N_{\text{wires}}$.

The ionization path length collected on each sense wire in the TPC is a function of the angle the track makes with the TPC endcap. Thus, the effective dE/dx sample thickness varies from track to track. The dependence of the truncated mean on sample thickness is not trivial. Because the shape of the dE/dx distribution changes with sample thickness (for greater thicknesses rises the contribution of hard scattering processes becomes more important), the truncated mean does not merely scale with path length. Instead, there is a correction term that is logarithmic in the sample thickness (see Appendix A). In the TPC analysis we correct for this logarithmic variation before quoting the dE/dx value.

Fig. 5.2 shows the predicted value of dE/dx in the TPC for the 5 stable charged particle species as a function of momentum. At low momentum each curve exhibits a $1/\beta^2$ falloff. This falloff resulted from the fact that a rapidly moving charge spends less time near a given electron and hence transfers less energy to it. For electrons, the $1/\beta^2$ region is at momenta below the lower edge of the plot. After passing through a minimum, the dE/dx for each species rises logarithmically. This "relativistic rise," which we can attribute to the $\ln(\beta^2 \gamma^2)$ term in Eq. 5.4, is a result of the relativistic increase in the incident particle's transverse electric field. The eventual flattening of the dE/dx curve (the "plateau region") is due to the onset of the density effect described previously.

5.2 Implications for Detector Design and Calibration

The shape of the Landau distribution and the dependence of the measured dE/dx on $\beta\gamma$ provided important constraints on the design of the TPC. Because the intrinsic width of the distribution is 55% and the typical π -K separation in the relativistic region is 12%, good particle identification requires many independent measurements of the ionization. One goal in the design of the TPC was to reach a final dE/dx resolution of approximately 3%.

If ionization energy loss is measured many times along a track, the measurements are independent only if the sample thickness is larger than the mean free path of the electrons released in atomic collisions. Since the transverse diffusion distance for an electron that has drifted a meter in the TPC is $1.20 \pm .01$ mm at 8.5 atm and $1.84 \pm .06$ mm at 1.5 atm in a 4 kG magnetic field⁶⁰, the 4 mm wire spacing in the sectors is well matched to the choice of gas. By running at high pressure, we increase the number of collisions per sample. This decreases the statistical fluctuations within each cell and improves the overall resolution. Allison and Cobb quote the following formula (derived using an atomic physics model) for the dE/dx resolution of a proportional chamber filled with argon⁶⁰:

$$R(\% \text{ fwhm}) = 96n^{-0.40} (xP)^{-0.32} \quad (5.6)$$

The variation with number of samples (n) is essentially an inverse square root, as expected from statistical arguments. The variation with pressure (P) and sample thickness (x) is slower than square root because the effect of higher energy collisions becomes more important as the mean number of collisions per sample increases. Thus, pressurizing the TPC to 8.5 atm improves the dE/dx resolution by a about a factor of two.

Good resolution is not the only requirement for doing particle identification via dE/dx . In addition, the systematic variations in gain must be small compared to the resolution. In the TPC,

Figure 5.2: The predicted value of dE/dx as a function of Momentum for the 5 stable charged particle species.

systematic gain variations must be controlled at the 1% level. Random fluctuations, for example gain variations due to changes in the sense wire thickness, do not bias the measured dE/dx if they are small relative to the 55% width of the dE/dx distribution. To reduce systematic variations in chamber gain, tight tolerances on the positioning of the wire planes were maintained during sector construction. To maintain 1% uniformity, it is necessary to control the size of the gap between the cathode and sense wire planes to $20\mu\text{m}$ and between the sense and grid plane to $26\mu\text{m}$ ⁶⁷. In addition, since a 1°C change in gas temperature causes a 3% change in wire gain⁶⁸, temperature variations must be minimized. Since all the wire preamplifiers are located on one edge of the sector, such gradients are almost inevitable. As part of the initial testing of each sector, a detailed map of the gain as a function of wire number and position along the wire was created. These maps are used to remove the remaining systematic gain variations. The rms variation of gain in a sector is roughly 3-4%⁶⁹.

In principle, the sector maps should correct for all systematic gain variations. However, the orientation of the chambers in the TPC and the flow pattern of the gas are different in the final experiment and in the test setup used during mapping. Thus, the temperature profiles of the sectors are different. In order to correct for this effect, *in situ* calibrations are necessary. Three rods with radioactive sources (Fe^{60}) plated on them are mounted in each sector. These rods are located at 0°, 30° and -16° with respect to the centerline of the sector. Separate calibration runs are taken with each rod positioned so that its sources irradiate the sense wires. (During normal data taking, all the sources are moved to a position where the sources are shielded from the sense wires.) Measurements of the pulse heights obtained at the three source positions are used to correct the detailed gain maps made during the initial sector testing. These corrections are typically of order 1-2%. The Fe^{60} source calibration also provides an absolute energy calibration for each wire. This data is used in conjunction with the electronic calibration (see Section 3.3.3) to correct the measured wire pulse height.

5.3 Corrections to the Data

The size of the systematic variations in gain across each sector remains stable for long periods of time. Source calibrations taken over a period of several months show that the relative gain as a function of position changes by less than 0.3%. It is therefore adequate to perform source calibrations about once per month. These calibrations, however, cannot correct for short term variations in overall gain due to changes in gas temperature, pressure and composition. Many such changes are diurnal in character and thus must be monitored on a run by run basis. Although attempts were made to correct the gain using the temperatures and pressures measured in the detector, these efforts were not adequate. In practice, the most effective technique is to measure the gain for each run using the data itself.

We measure the gain by looking at tracks with momenta between 400 and 650 MeV/c. In this region pions are minimum ionizing and electrons are close to plateau. Kaons in the $1/\beta^2$ region produce a tail on the high side of the electrons. Since the pion dE/dx is well separated from the dE/dx of other particles of the same momentum, we can cut to remove electron and kaon contamination and then fit the remaining distribution of truncated means to a gaussian. During the PASS 3 data processing, we obtain a gain constant for each run by making such a fit. In addition to the momentum cut, we require in PASS 3 that each track be completely contained in a sector and have at least 100 good dE/dx samples. The truncated mean for these tracks is defined to be 12.1 keV/cm because early theoretical calculations predicted this value for $(dE/dx)_{\text{min}}$.

Since electron capture by electronegative contaminants in the gas can cause biases in dE/dx as a function of drift distance, the rate of electron capture is also monitored. Each track is divided in half and the difference in dE/dx between the two halves is plotted as a function of the difference in the average drift distance. The slope of this curve is the electron capture and is typically 5% over one meter. The electron capture exhibits diurnal variations, suggesting that

Figure 5.3: The observed variation of the truncated mean for minimum ionizing pions with sample thickness. The data have already been corrected for the logarithmic dependence predicted in Appendix A.

the rate of outgassing in the detector is temperature dependent

After run to run gain corrections to the data are made, it is possible to study dE/dx systematics as a function of geometric variables. Our approach is to plot the value of the dE/dx obtained for minimum ionizing pions as function of a number of variables. If the data exhibit any dependence on these variables, the systematic variations are removed with *ad hoc* corrections.

Fig. 5.3 shows the measured value of dE/dx for minimum ionizing pions as a function of the log of the track length collected on each sense wire (λ approximated by $L = \ln(0.4 \tan \lambda)$ where λ is the angle the track makes with the TPC endcap) for data processed through the standard analysis chain. The data exhibits an overall slope, which we attribute to inadequate correction of the log length dependence of the truncated mean (the dependence used in the analysis was that derived in Appendix A). In addition, there is a noticeable increase in dE/dx for tracks that are parallel to the endcap ($\lambda = 0$). The sample length dependence can be parameterized

$$(dE/dx) = (dE/dx)_{\text{mcua}} - 0.379 - 0.539L \quad (5.7)$$

After this correction is applied, the rise for $\lambda = 0$ tracks can be clearly seen in Fig. 5.4. A likely explanation of the effect is cross talk in the wire electronics which adds coherently when all pulses are close together in time. This explanation is supported by the fact that no $\lambda = 0$ effect is seen when dE/dx is calculated using pads instead of wires. (The signals from separate pad rows are

Figure 5.4: Dependence of the truncated mean for minimum ionizing pions on the sine of the angle the track makes with the endcap. The sample length dependence in Fig. 5.3 has already been corrected using Eq. 5.7.

well separated on the sector, while all wire signals travel along nearby paths.) We correct the wire data using the formula:

$$\langle dE/dx \rangle = \langle dE/dx \rangle_{\text{wire}} - 0.65e^{-7.26 \sin(\lambda)}. \quad (5.8)$$

During the preliminary analysis of the TPC data, a final systematic effect was discovered. For events with several tracks in a single sector, the average dE/dx for the later tracks differed by about 1% from the average dE/dx of the first track to reach the chamber. This difference was attributed to a failure of the TPC electronics to return quickly to pedestal level after being pulsed. To check this hypothesis, we studied the full electronic pulse shape by turning the digitizer thresholds to zero on selected channels and taking Fe^{56} spectra. Fig. 5.5 shows the results of the study. After being pulsed, the electronics exhibits an undershoot of approximately 0.5% of the peak pulse height. This undershoot lasts for the rest of the CCD readout and therefore biases the measured peak height on subsequent hits. The data were reanalysed correcting all wire pulse heights for the presence of earlier clusters. This reanalysis removed the observed dependence of the dE/dx on the number of tracks per sector.

Fig. 5.6 shows the dE/dx data as a function of momentum after all the corrections and calibrations described previously were applied. Only tracks with at least 80 dE/dx samples are included in the plot. This corresponds to about 65% of the tracks. In the $1/\beta^2$ region, the bands

Figure 5.5: Measured pedestal recovery of the TPC electronics as a function of bucket number. The pulse peak occurs in bucket 0 and the peak height (offscale on this plot) is normalized to 1. For this study, the digitizer thresholds have been set to zero.

for all particle types are clearly separated. Above the dE/dx crossovers, π , e and K - p bands are seen. It is not possible to separate Kaons and protons in the relativistic region by eye. A similar histogram is shown in Fig. 5.7 for tracks where dE/dx is calculated using 8 or more pad rows. Here the resolution is considerably worse, but it is still possible to distinguish pions from electrons.

Fig. 5.8a shows the distribution of the truncated mean for those tracks in Fig. 5.6 with momenta between 400 and 650 MeV/c. The solid lines show a fit to the data, where the dE/dx resolution is assumed to be gaussian. The fitted resolution is $3.65 \pm .2\%$. In the case of pad dE/dx (Fig. 5.8b), a resolution of $\sim 9\%$ gives a reasonable representation of the data.

Fig. 5.9 shows the dE/dx resolution (measured using minimum ionizing pions) as a function of the square root of the number of dE/dx samples (N). The line shows the predicted dependence if the resolution scales as $1/\sqrt{N}$. This form is a reasonable approximation to the data.

XBL 842-10061

Figure 5.6: Measured value of dE/dx for all tracks with at least 80 usable dE/dx samples.

XBL 846-10652

Figure 5.7: Measured value of dE/dx for all tracks where 8 or more pads are used to calculate the truncated mean.

Figure 5.8: The distribution of dE/dx for tracks with momenta between 400 and 650 MeV/c and a) with at least 80 usable dE/dx samples, or b) where at least 8 pads have been used to calculate the truncated mean.

Figure 5.9: Dependence of the resolution on number of samples

5.4 Measured $\beta\gamma$ Dependence of the Truncated Mean

Although the theoretical calculations reproduced in Appendix A represent the data fairly well, they are not adequate for doing particle identification in the relativistic region. To determine the dependence of the truncated mean on $\beta\gamma$, we have measured the shape of the dE/dx curve using cosmic ray muons at low $\beta\gamma$ ($2 < \beta\gamma < 100$), electrons from the multihadron event sample at the edge of the plateau region and Bhabha candidates for the plateau. The momentum resolution for cosmic ray events is worse than for real data because the track fitting software cannot constrain the particles to a predetermined vertex. To reduce biases in the measurement of the muon momentum due to the limited resolution of the TPC and the hardness of the cosmic ray spectrum, we therefore combine for each event the momentum measured in back to back sectors of the TPC. The resulting dE/dx plot is shown in Fig. 5.10. The solid line shows the predictions of our atomic physics calculations. The dashed line gives an empirically corrected form used for our physics analysis. This empirical shape has been parameterized as a polynomial because polynomials can be rapidly calculated in the software. The requirement that local structure caused by the polynomial form of the approximation remain less than 0.3% throughout the $\beta\gamma$ range $1 < \beta\gamma < 10000$ means that a ninth order polynomial is necessary.

5.5 Particle Identification Algorithms

The question of particle identification can be approached in two different ways. For studies of inclusive and semi-inclusive production, a statistical separation of particle species is adequate. In this case we can take momentum slices and simultaneously fit for the number of π , K and p. This technique allows us to use all of the available information, including the number of tracks found in ambiguous dE/dx regions. This method is used to determine the charged hadron cross sections and particle fractions and is described in Chapter 6. For many studies, however, it is necessary to determine on a track by track basis the probability that a given particle is a π , K,

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Figure 5.10: The dependence of the truncated mean on $\beta\gamma$. The plot was made by combining data from cosmic ray muons, conversion electrons and wide angle Bhabha events. The curves show the prediction of the atomic physics model described in Appendix A and of a polynomial fit to the data.

Figure 5.11: A representation of the method used to find the dE/dx χ^2 for a track.

p or ϵ . For this reason we introduce the concept of a dE/dx χ^2 .

For a track with momentum p and truncated mean E , we define the χ^2 for a given mass hypothesis to be the square of the distance between the point (p, E) and the theoretical dE/dx curve, measured in units of standard deviations (see Fig. 5.11). The χ^2 therefore depends on both the dE/dx and momentum uncertainties. When calculating χ^2 , we estimate the dE/dx uncertainty using the value obtained for minimum ionizing pions, scaled according to the assumption that $\sigma_{dE/dx}$ is proportional to the square root of the truncated mean and inversely proportional to the square root of the number of dE/dx samples. The momentum uncertainty is taken from the output of the track fitting routines, but is corrected for the mass dependence of the multiple scattering component of the uncertainty.

Because the number of standard deviations between particle species is a strong function of momentum (see Fig. 5.12), a cut on χ^2 itself will produce a sample whose purity has a strong momentum dependence. Fig. 5.13 shows the kaon purity (fraction of tracks called kaons that are true kaons), as estimated by Monte Carlo, for a cut that calls everything with $\chi_K^2 < 4$ a kaon. Such a method is reasonable for studies that can tolerate high background rates but need good acceptance.

For many applications, however, high purity is necessary. The flavor correlation studies in Chapter 7 fall into this category. Such studies require an algorithm that uses our knowledge of the expected separation between particle species and the known population of each species. We present here our choice of such an algorithm. For a given momentum, the population of particle species a should be distributed in χ^2 according to the function:

$$\frac{d\sigma}{dp} d\chi_a = N_a(\chi_a) = \frac{d\sigma}{dp} \frac{1}{\sqrt{2\pi}} e^{-\chi_a^2/2} d\chi_a \quad (5.9)$$

Figure 5.12: Separation between species as a function of Momentum

Therefore, we can describe the fraction of all particles with a given E and p that are of species a as

$$W_a(\chi_a^2) = \frac{N_a(\chi_a)}{\sum N_i(\chi_i)} = \frac{f_a(p)e^{-\chi_a^2/2}}{\sum f_i(p)e^{-\chi_i^2/2}} \quad (5.10)$$

where $f_a(p)$ is the fraction of hadrons with momentum p that are of species a . Cutting on $W_a(\chi_a^2)$ therefore produces a sample with uniform purity. The values of $f(p)$ used in eq. Eq. 5.10 are calculated by fitting the measured particle fractions vs. momentum (see Chapter 6). We find

$$\begin{aligned} f_\pi(p) &= (0.2/p)^2 \\ f_\pi(p) &= \text{Max}(0.84936 - 0.13499 \log(p), 0.0) \\ f_K(p) &= \text{Max}(0.10934 + 0.09163 \log(p), 0.0) \\ f_p(p) &= \text{Max}(0.04130 + 0.04337 \log(p), 0.0) \end{aligned} \quad (5.11)$$

In the low momentum region, where the species are well separated, a cut on W does not seriously effect the efficiency. In the high momentum region, however, this is not the case. Since the dE/dx resolution is inversely proportional to the square root of the number of dE/dx samples, tracks with large momenta and few dE/dx samples will have reasonable χ^2 for all mass hypotheses. The weighting scheme above means that these tracks are not likely to pass a tight W

Figure 5.13: The purity of the kaon sample obtained with the requirement that each track have a χ^2 of ≤ 4 for the kaon mass hypothesis. These results were obtained using the Lund Monte Carlo and the FMC detector simulation.

cut. Such a cut therefore implicitly introduces a cut on the number of dE/dx samples. Although this difficulty affects all particle species, it is especially relevant in the case of kaons, since the kaon dE/dx band overlaps that of both pions and protons. We defer detailed discussion of acceptances to Chapters 6 and 7.

Chapter 6

Measurement of Inclusive Cross Sections

The charged pion, kaon and proton production cross sections and the charged hadron fractions were first measured with the TPC using data collected in Spring 1982⁶⁶. These data were insufficient for separating kaons and protons above 1 GeV/c. Improvements both in the quality and quantity of data obtained since then have made such a separation possible. In this chapter we describe the method by which the hadronic fractions and cross sections are obtained.

6.1 Track Selection

The criteria for selecting multihadron events have been discussed previously (see Section 4.5). Within these events, we impose the following cuts on the tracks used for the inclusive cross section studies:

- To remove tracks not associated with the primary interaction, we require that the extrapolated orbit miss the nominal beam interaction point by less than 6 cm in the plane perpendicular to the beamline and 10 cm in the beam direction.
- To insure that the track's momentum is well measured, we require that the curvature c and its uncertainty dc satisfy either $dc/c < 0.3$ or $dc < 0.3 \text{ (GeV/c)}^{-1}$.
- To remove tracks with poorly determined dE/dx , we require at least 80 dE/dx samples per track.
- Below 2 GeV/c, we estimate the number of protons plus antiprotons by twice the number of antiprotons. This is because the number of low momentum protons created during nuclear interactions in the beampipe, pressure wall and TPC field cage is comparable to the number of primary protons.
- For this measurement we require that each track make an angle λ with respect to the TPC midplane of between 5° and 45° . The 5° cut is motivated by the presence of coherent crosstalk on $\lambda = 0$ tracks (see Section 5.3). The 45° cut serves two purposes. First, it removes tracks that penetrate the TPC endcap and hence do not deposit ionization on all the sense wires. Rejecting these tracks with a geometric cut reduces the systematic uncertainty on the acceptance, since the fraction of such tracks that meet the 80 sample dE/dx cut described above is sensitive to the Monte Carlo simulation of the distribution of the number of wire hits found per track. Second, the cut removes tracks that pass through the beampipe and pressure wall obliquely and have a larger chance of interacting in the material before the TPC. The choice of angle cuts is quite conservative and is designed to minimize the systematic error associated with acceptance corrections. If we loosen the cuts to accept all tracks with angles between 0° and 60° , the results do not change within the quoted systematic errors.

Figure 6.1: The measured value of dE/dx as a function of momentum for all tracks meeting the cuts described in Section 6.1. The dE/dx band bounded by the solid lines results from conversion electrons.

Figure 6.2: The momentum spectrum for tracks with measured dE/dx within $\pm 7\%$ of the electron band. Clear peaks are seen at the π , K and p dE/dx crossovers. The solid curve shows the results of a fit to data. The dashed lines show the e , π , K and p contributions to that fit.

The result of these combined cuts is to reduce the track acceptance to about 45%. The efficiencies are dominated by geometrical acceptance (85%) and the requirement of at least 80 usable dE/dx samples (65%). We defer discussion of the momentum and species dependence of the acceptance to Section 6.3.

6.2 Fitting the Raw Numbers

Fig. 6.1 shows the distribution of measured dE/dx as a function of momentum for all tracks meeting the cuts described above. The large number of electrons in the plot is due to the conversion of photons in the 0.2 radiation lengths of material before the TPC. We have decided to remove these electrons statistically before measuring the hadron cross sections. Fig. 6.2 shows the momentum spectrum for all tracks that have a measured dE/dx within $\pm 7\%$ of the predicted dE/dx for electrons. The electron spectrum rises at low momentum, reaches a maximum at about 100 MeV/c and then falls off slowly. Peaks are clearly seen corresponding to the π -e, K-e and p-e crossovers. The solid line superimposed on the histogram is a fit to the data, representing the combined contributions of electrons, pions, kaons and protons. The individual contributions of each particle species are shown by the dashed curves. During the fitting, the positions of the hadron crossovers are fixed, but the widths of the peaks and the number of particles of each type are allowed to vary. The electron spectrum is approximated by the form:

$$N_{\text{electron}} = C_1 p^{-\lambda_1} - C_2 p^{-\lambda_2} \quad (6.1)$$

The results of this fit are used to remove conversion pairs from the standard dE/dx distributions.

Figure 6.3: The distribution of truncated means as a function of momentum after electrons have been subtracted using the statistical method described in the text. The vertical lines on the plot indicate the binning used to determine the pion, kaon and proton cross sections.

Figure 6.4: The ratio of measured dE/dx to expected dE/dx for pions of the same momentum for all tracks in Fig. 6.3 with momenta between 3.7 and 6 GeV/c.

The number of electrons expected in each dE/dx and momentum bin is calculated and that number of tracks is subtracted from the data.

The dE/dx distribution that remains after electrons are removed is shown in Fig. 6.3. The plot indicates as well the boundaries of the bins used for the cross section measurements. These bins were chosen to meet two criteria. First, the separation in average dE/dx between the measured particle species and its nearest neighbor must be at least 1 standard deviation (see Fig. 5.12). Second, the slices in momentum must be wide enough to insure that resolution smearing between neighboring bins does not significantly bias the mean momentum within a bin or introduce large correlations between the cross sections measured in neighboring bins.

Two different techniques are used to measure the raw number of pions, kaons and protons as a function of momentum. In the low momentum region, the dE/dx curves are well separated and the number of particles of each species is determined by counting the number of tracks in each dE/dx band. This method is used for kaons with momenta below 600 MeV/c and for antiprotons below 1.35 GeV/c. At high momentum, the dE/dx separation between species is not large enough to use this technique. Instead, for each momentum slice we perform a simultaneous fit to the number of pions, kaons and protons. Since the dE/dx bands for the three species are roughly parallel at high momentum, we have chosen to project the measured dE/dx distribution along the predicted dE/dx curve for pions. This projection is equivalent to normalizing the measured dE/dx for each track to the predicted value of the energy loss for a pion with the same momentum. Fig. 6.4 shows such a projection for tracks in the momentum range $3.7 < p < 6$ GeV/c. The distribution shows a clear peak at $R \equiv (dE/dx)/(dE/dx)_\pi = 1$, corresponding to pions. A second peak at lower R contains kaons and protons. To determine the number of particles of each species, we perform a maximum likelihood fit to this normalized dE/dx distribution. The

Figure 6.5: The distribution of dE/dx , normalized to the expected dE/dx for pions, for minimum ionizing tracks. The dashed line shows the result of a gaussian fit. The solid line is a fit to the form of Eq. 6.2.

method requires an accurate knowledge both of the mean dE/dx (i.e. the position of each peak) and the distribution of dE/dx about that mean. The method used to parameterize the mean energy loss as a function of $\beta\gamma$ has already been described in Section 5.4. We discuss here our studies of the shape of the distribution of truncated means.

We have examined this shape using minimum ionizing tracks. Fig. 6.5 shows $\mathcal{R}$, the ratio of measured dE/dx to the expected dE/dx for pions, for tracks with momenta between 400 and 650 MeV/c. The distribution has been cut so that electrons are offscale. The dashed line shows a gaussian fit to the data. The χ^2 per degree of freedom for this fit is 5.3 and $\sigma = (3.64 \pm 0.02)\%$. The fact that a gaussian shape does not reproduce the distribution well is not surprising. If we assume that the dE/dx distribution has a width that depends on the number of dE/dx samples on the track, we do not expect the distribution summed over tracks with different numbers of dE/dx samples to remain gaussian. We can approximate the sum of gaussians with the same center, x_0 , but different widths by the form:

$$N(x) = \begin{cases} \frac{N_0}{2\pi} \left(\frac{\sigma_1\sigma_2}{\sigma_1 - \sigma_2} \frac{1}{(x - x_0)^2} \right) \left[e^{-(x-x_0)^2/2\sigma_1^2} - e^{-(x-x_0)^2/2\sigma_2^2} \right] & \text{for } x \neq x_0 \\ \frac{N_0}{2\pi} \left(\frac{\sigma_1 + \sigma_2}{2\sigma_1\sigma_2} \right) & \text{for } x = x_0 \end{cases} \quad (6.2)$$

Fitting Fig. 6.5 to a distribution of this form reduces the χ^2 per degrees of freedom to 1.1 for the

p (GeV/c)	$\pi^\pm$	K $^\pm$	p, $\bar{p}$
0.20 - 0.25	4864 $\pm$ 70	-	-
0.25 - 0.30	6274 $\pm$ 79	-	-
0.30 - 0.41	12094 $\pm$ 110	367 $\pm$ 19	-
0.41 - 0.50	9513 $\pm$ 98	406 $\pm$ 20	-
0.50 - 0.61	10302 $\pm$ 102	607 $\pm$ 25	178 $\pm$ 19
0.61 - 0.74	10256 $\pm$ 102	844 $\pm$ 30	262 $\pm$ 23
0.74 - 0.82	5035 $\pm$ 72	416 $\pm$ 22	144 $\pm$ 17
0.82 - 0.90	4352 $\pm$ 72	492 $\pm$ 28	210 $\pm$ 20
0.90 - 1.00	4794 $\pm$ 77	509 $\pm$ 40	240 $\pm$ 22
1.00 - 1.35	-	-	744 $\pm$ 70
1.35 - 1.65	7231 $\pm$ 101	1277 $\pm$ 59	480 $\pm$ 31
1.65 - 2.01	-	1177 $\pm$ 52	-
2.01 - 2.23	2692 $\pm$ 70	666 $\pm$ 50	181 $\pm$ 81
2.23 - 2.72	4422 $\pm$ 74	-	-
2.72 - 3.67	4807 $\pm$ 72	-	-
3.67 - 4.95	2927 $\pm$ 56	1094 $\pm$ 53	420 $\pm$ 35
4.95 - 6.69	1501 $\pm$ 41	715 $\pm$ 37	298 $\pm$ 26

Table 6. The raw number of pion, kaons and protons found in each momentum bin. For those bins where the numbers are obtained by fitting the dE/dx distributions to several peaks, the quoted statistical errors include correlations between the parameters of the fit.

values $\sigma_1 = 5.39\%$ and $\sigma_2 = 2.60\%$. The results of this fit are shown as a solid line on the plot.

The ratio σ_1/σ_2 determines the amount of nongaussian tail in Eq. 6.2. Since the distribution of the number of dE/dx samples per track is not dependent on the particle's momentum, we expect this ratio to be constant. The value of σ_1 , however, is momentum and species dependent, for several reasons. First, the intrinsic dE/dx resolution varies inversely with the square root of the mean energy loss per sample. Second, because we are fitting the dE/dx distribution for finite size momentum bins and the π , K and p dE/dx curves are not exactly parallel, the observed width for kaons and protons is broadened due to variations in the value of $\mathcal{R}$ within each bin. Finally, smearing effects due to finite momentum resolution further broaden the measured distribution by introducing tracks that rightly belong in neighboring momentum slices. Given a momentum bin of width Δp that is centered at momentum p , we can parameterize σ_1 , the variation in $\mathcal{R}$ for particles of species i by the form:

$$\sigma_i^2 = \frac{\sigma_0^2}{(dE/dx)_n} \left(\frac{dE/dx}{(dE/dx)_n} \right) + \frac{(\Delta \mathcal{R}_i)^2}{12} + \left(\frac{\partial \mathcal{R}_i}{\partial p} \right)^2 (\Delta p)^2 \quad (6.3)$$

where σ_0 and $(dE/dx)_n$ are the values of the resolution and the average truncated mean for minimum ionizing particles, $(dE/dx)_n$ is the predicted value of the truncated mean for pions at momentum p , $R_i \equiv (dE/dx)_i / (dE/dx)_n$ and $\Delta \mathcal{R}_i$ is the variation in the value of $\mathcal{R}_i$ over the momentum bin. The first term accounts for the variation of resolution with energy loss. The second corrects for the variation of $\mathcal{R}_i$ within the momentum bin (assuming a uniform population in dE/dx within the bin). The final term corrects for momentum smearing from neighboring bins.

The $\mathcal{R}$ distributions for several momentum slices are shown in Fig. 6.6, along with the results of fits to the number of pions, kaons and protons. The numbers obtained in these fits are also given in Table 6.1. All fitting is done using the CERN program MINUIT⁷⁰. The quoted statistical errors include correlations between the fitted parameters. We have investigated the sensitivity

Figure 6.6: Results of the simultaneous fit to the number of pions, kaons and protons for the momentum slices a) $0.74 < p < 0.82$ b) $1.35 < p < 1.65$ c) $3.67 < p < 4.94$ and d) $4.94 < p < 6.69$.

Monte Carlo

Figure 6.7: Monte Carlo data for the distribution of R in the momentum slices a) $0.74 < p < 0.82$ b) $1.35 < p < 1.65$ c) $3.67 < p < 4.94$ and d) $4.94 < p < 6.69$. The curves show the result of a simultaneous fit to the number of pions, kaons and protons.

Standard Results:	F_{π} 0.64 $\pm$.007	F_K 0.266 $\pm$.009	F_p 0.094 $\pm$.006
Change in Fit	ΔF_{π}	ΔF_K	ΔF_p
Vary $\sigma_{dE/dx}$ (4.12% $\rightarrow$ 4.32%)	.001	.003	.004
Gaussian dE/dx distribution	.002	.002	.004
$\sigma_{dE/dx} \propto dE/dx$	.005	.011	.006
$\sigma_{dE/dx} = \text{constant}$	.006	.013	.007
Vary π -K separation by $\pm 0.3\%$	.003	.014	.017
Vary overall scale $\pm 0.3\%$	.008	.004	.012

Table 6.2: The stability of the fitted pion, kaon and proton fractions with respect to changes in the assumptions of the fit. These results are for the momentum range $3.7 < p < 6.7$ GeV/c.

		$P_{\text{vertex}} - P_{\text{TPC}}$ (GeV/c)		
p (GeV/c)	p_{min} (GeV/c)	$\pi^{\pm}$	$K^{\pm}$	$p, \bar{p}$
0.20 - 0.25	0.22	0.012	-	-
0.25 - 0.30	0.27	0.012	-	-
0.30 - 0.41	0.35	0.012	0.046	-
0.41 - 0.50	0.45	0.011	0.028	-
0.50 - 0.61	0.55	0.011	0.019	0.061
0.61 - 0.74	0.67	0.011	0.016	0.045
0.74 - 0.82	0.78	0.011	0.016	0.037
0.82 - 0.90	0.86	0.011	0.015	0.032
0.90 - 1.00	0.95	0.011	0.013	0.027
1.00 - 1.35	1.16	0.000	0.000	0.000
1.35 - 1.65	1.49	0.011	0.011	0.014
1.65 - 2.01	1.82	0.011	0.011	0.012
2.01 - 2.23	2.12	0.011	0.011	0.011
2.23 - 2.72	2.46	0.011	0.011	0.011
2.72 - 3.67	3.16	0.011	0.011	0.011
3.67 - 4.95	4.18	0.011	0.011	0.011
4.95 - 6.69	5.47	0.011	0.011	0.011

Table 6.3: The mean shift in momentum between the event vertex and the TPC due to energy loss in the beampipe and TPC pressure wall. Since energy loss is a function of velocity rather than momentum, different corrections are needed for pions, kaons and protons. We estimate a systematic error of 20% on each number in the above table.

		Momentum Shift (GeV/c)		
P (GeV/c)	P _{bin} (GeV/c)	$\pi^\pm$	K $^\pm$	γ
0.20 - 0.25	0.22	0.002	-	-
0.25 - 0.30	0.27	0.002	-	-
0.30 - 0.41	0.35	0.002	0.031	-
0.41 - 0.50	0.45	0.002	0.014	-
0.50 - 0.61	0.55	0.002	0.011	0.034
0.61 - 0.74	0.67	0.002	0.007	0.017
0.74 - 0.82	0.78	0.002	0.009	0.005
0.82 - 0.90	0.86	0.002	0.006	0.003
0.90 - 1.00	0.95	0.002	0.003	0.006
1.00 - 1.35	1.16	0.002	0.006	0.005
1.35 - 1.65	1.49	0.004	0.022	0.016
1.65 - 2.01	1.82	0.010	0.036	0.034
2.01 - 2.23	2.12	0.027	0.058	0.063
2.23 - 2.72	2.46	0.026	0.174	0.090
2.72 - 3.67	3.16	0.045	0.115	0.158
3.67 - 4.95	4.18	0.086	0.144	0.221
4.95 - 6.69	5.47	0.044	0.076	0.189

Table 6.4: The difference between the mean momentum for all tracks found in a bin and the momentum at which the cross sections are quoted. These values have been calculated via Monte Carlo and include the effects of momentum smearing between bins. We estimate a systematic error of 20% on each value quoted above.

of our results to changes in the assumed behavior of the dE/dx distributions by refitting the data for the momentum range $3.7 < p < 6.7$ GeV/c under a variety of assumptions. The results, quoted in terms of the fraction of charged hadrons measured to be pions, kaons or protons, are summarized in Table 6.2. The fit is not very sensitive to the shape of the distribution of truncated means. Replacing the form in Eq. 6.2 with a gaussian, or changing the value of α introduces only small variations in the fitted fractions. Changing the parameterization of the variation of resolution with average dE/dx produces a slightly larger difference, since it affects the particle species differently. The biggest sources of systematic error, however, are uncertainties in normalization and in the behavior of the average dE/dx as a function of $\beta\gamma$. Based on the studies of Section 5.4, we quote an uncertainty of $\pm 0.3\%$ in our estimate of the average dE/dx . Combining these sources of uncertainty, we estimate the total systematic error on the number of particles of each species as $\pm 2\%$ of the total number of hadrons in the momentum bin.

Because a track's dE/dx depends on its velocity in the medium being ionized, the data are fit as a function of the momentum measured in the TPC. This momentum differs from the true value at the beam interaction point for two reasons. First, energy loss in the material before the TPC systematically reduces the track's measured momentum. Second, the combination of a rapidly falling spectrum and resolution smearing between bins can bias the measured momentum distribution. The shift in momentum between the event vertex and the TPC has been evaluated using the detector simulation Monte Carlo. The momentum and species dependent corrections are given in Table 6.3. The effects of momentum smearing have been investigated using two techniques, an analytic approximation and a Monte Carlo simulation. The first method involves using the measured cross sections and the measured form of the TPC momentum resolution $((dp/p)^2 = \{0.06\}^2 + \{0.035p\}^2)$, for p in GeV/c) to estimate the amount of smearing between

p (GeV/c)	$\pi^\pm$	$K^\pm$	$p, \bar{p}$
0.20 - 0.25	4763 $\pm$ 70	-	-
0.25 - 0.30	6189 $\pm$ 79	-	-
0.30 - 0.41	12078 $\pm$ 110	171 $\pm$ 41	-
0.41 - 0.50	9635 $\pm$ 98	363 $\pm$ 20	-
0.50 - 0.61	10507 $\pm$ 102	577 $\pm$ 25	94 $\pm$ 21
0.61 - 0.74	10479 $\pm$ 103	526 $\pm$ 30	241 $\pm$ 23
0.74 - 0.82	5186 $\pm$ 72	440 $\pm$ 22	149 $\pm$ 17
0.82 - 0.90	4956 $\pm$ 72	484 $\pm$ 28	189 $\pm$ 20
0.90 - 1.00	4871 $\pm$ 77	515 $\pm$ 40	237 $\pm$ 22
1.00 - 1.35	-	-	746 $\pm$ 70
1.35 - 1.65	7355 $\pm$ 101	1408 $\pm$ 68	495 $\pm$ 31
1.65 - 2.01	-	1226 $\pm$ 52	-
2.01 - 2.23	2761 $\pm$ 70	672 $\pm$ 50	389 $\pm$ 81
2.23 - 2.72	4560 $\pm$ 75	-	-
2.72 - 3.67	5026 $\pm$ 81	-	-
3.67 - 4.95	3196 $\pm$ 75	1192 $\pm$ 53	519 $\pm$ 40
4.95 - 6.69	1579 $\pm$ 43	753 $\pm$ 37	329 $\pm$ 26

Table 6.5: The number of pions, kaons and protons per bin after corrections for momentum smearing and energy loss before the TPC have been applied.

bins. For the second method, the smearing is calculated using the detector simulation Monte Carlo. This simulation, in conjunction with the Lund event generator, does a reasonable job of reproducing the measured momentum spectra and therefore should provide a reasonable estimate of the smearing effects. We estimate a systematic uncertainty of 20% on the size of the bias due to smearing between bins. The two methods agree within this uncertainty for all momentum bins quoted. The studies of momentum smearing indicate:

- For momenta below 6 GeV/c, a given bin contains roughly equal numbers of tracks feeding down from higher momenta and feeding up from lower momenta. Thus, the overall momentum bias in these bins is small. The size of this bias for each momentum bin used in this analysis is summarized in Table 6.4.
- Above 6 GeV/c, the measured distribution consists largely of tracks that rightly belong in lower momentum bins. Because the contamination is so large, we cannot quote cross section measurements in this momentum region.
- Although the overall bias in the mean value of p is small for momenta below 6 GeV/c, resolution smearing introduces correlations between the measured cross sections in neighboring bins. In calculating the effect of these correlations on the systematic uncertainties in the cross sections, we have made the assumption that the inclusive spectra are smooth functions of momentum. Without such an assumption, the quoted errors would increase by roughly 50%.

Corrections for the momentum smearing bias and for energy loss in the material before the TPC are made by performing a linear interpolation between raw data bins. The number of pions, kaons and protons remaining in each bin after this interpolation is done are given in Table 6.5

p (GeV/c)	Acceptance (%)		
	$\pi^\pm$	$K^\pm$	p, $\bar{p}$
0.20 - 0.25	48.6 ± 3.04	—	—
0.25 - 0.30	43.1 ± 3.03	—	—
0.30 - 0.41	42.6 ± 3.02	17.8 ± 5.47	—
0.41 - 0.50	46.3 ± 2.03	29.5 ± 4.29	—
0.50 - 0.61	48.4 ± 2.03	37.4 ± 2.84	9.8 ± 6.02
0.61 - 0.74	48.5 ± 2.04	43.3 ± 2.96	25.5 ± 5.17
0.74 - 0.82	47.3 ± 2.08	43.8 ± 2.90	30.2 ± 3.34
0.82 - 0.90	46.8 ± 2.07	44.0 ± 3.04	33.3 ± 3.45
0.90 - 1.00	46.3 ± 2.08	41.3 ± 3.35	34.4 ± 3.45
1.00 - 1.35	—	—	41.1 ± 3.42
1.35 - 1.65	45.4 ± 2.08	41.4 ± 3.14	40.0 ± 4.07
1.65 - 2.01	—	41.9 ± 2.95	—
2.01 - 2.23	45.0 ± 2.18	42.8 ± 3.57	41.5 ± 11.1
2.23 - 2.72	43.3 ± 2.08	—	—
2.72 - 3.67	43.1 ± 2.06	—	—
3.67 - 4.95	43.5 ± 2.08	43.2 ± 3.02	41.9 ± 3.47
4.95 - 6.69	40.9 ± 2.13	42.5 ± 3.14	39.5 ± 3.30

Table 6.6: The pion, kaon and proton acceptances as a function of momentum. Below 2 GeV/c, the p+ $\bar{p}$ cross section is calculated using twice the measured value for p.

6.3 Acceptance Corrections

The charged particle acceptance is calculated using the TPC Fast Monte Carlo. Data summary files containing generated events are analyzed using an identical procedure as for the real data. Fig. 6.7 shows the R distributions and fits obtained from the Monte Carlo simulation for the same momentum bins as Fig. 6.6. The fitted numbers of pions, kaons and protons per accepted Monte Carlo event are compared to the number of particles generated per event when the physics Monte Carlo is run without radiative effects and without detector simulation. We include in the generated numbers the decay products of all particles with lifetimes less than 5×10^{-10} . The ratio of found tracks per event to generated tracks per event is defined as the acceptance:

$$A = \frac{\left(\frac{1}{N_{ev}} \frac{dN}{dp} \right)_{\text{Found}}}{\left(\frac{1}{N_{ev}} \frac{dN}{dp} \right)_{\text{Generated}}} \quad (6.4)$$

The results of this procedure are shown in Fig. 6.8. At low momenta, the acceptances are dominated by interactions in the material before the TPC. Electromagnetic energy loss produces sharp thresholds for seeing pions, kaons and protons at 120, 280 and 410 MeV/c. Above these thresholds, the acceptances vary due to the momentum dependence of the nuclear interaction cross sections. These efficiencies flatten out above 1 GeV/c; the final values are dominated by the geometric cuts and the requirement of 80 dE/dx samples.

To estimate the systematic uncertainties associated with the Monte Carlo modeling of energy loss and nuclear interactions, we have examined the size of these effects by generating sets of events with each of these processes turned off in the Monte Carlo. In addition, to study the systematic uncertainty in the probability for a decaying pion or kaon to be accepted as a good track by the pattern recognition software, we have varied the minimum angle at which the Monte

Figure 6.8: The total acceptance for a) pions b) kaons and c) protons with the event and track selection cuts outlined in the text. These numbers include the effects of decays in flight, energy loss, nuclear interactions and initial state radiation.

Carlo pattern recognition simulation recognizes kinks due to particle decays. Based on these studies we have assigned momentum and species dependent systematic errors to the acceptance. These errors range from 5% (for low momentum kaons and protons, where the acceptance is changing rapidly) to 2% (for high momentum pions).

The sensitivity of the results to our track and event selection criteria was determined by refitting the cross sections with a variety of different cuts. Changes that were made include lowering the minimum wire cut to 40 and raising it to 120, changing the cut on the angle λ to $0^\circ < \lambda < 60^\circ$, and changing the event selection criteria so that only events with more than 7 tracks were used. None of these variations caused significant changes in the acceptance corrected cross sections. Results for the momentum and species dependent acceptances are summarized in Table 6.6. Since the uncertainties in these numbers are dominated by our ability to estimate the pattern recognition efficiency, nuclear interaction and energy loss probabilities and distribution of number of dE/dx sample per track, the systematic errors quoted here are highly correlated between momentum bins.

6.4 Results

Fig. 6.9 shows the fraction of charged hadrons that are pions, kaons and protons as a function of momentum. At low momentum, over 90% of the charged particles are pions. This fraction decreases with momentum, while both the kaon and proton contributions rise. One might ask whether the momentum dependence of the K and p hadron fractions indicates that these particles are produced in a different manner than pions. Monte Carlo studies indicate that such an interpretation is not necessary. In both the extended Feynman-Field and the Lund models, the rate of baryon production is controlled by a momentum independent constant: the relative diquark to quark production rate. Yet, these models are able to reproduce the observed momentum dependence of the particle fractions. The mechanism responsible for this variation is the decay of resonances. Particles as the ρ and the ω produce large numbers of low momentum pions and therefore enhance the $\pi^\pm$ cross section at low p . This effect is illustrated in Fig. 6.10, where the observed proton fraction is compared to the predictions of the Lund Monte Carlo, both before and after resonance decays. A similar effect is observed in the $p_\perp$ dependence of the proton fraction⁷¹.

We have chosen to normalize the π , K and p differential cross sections to the total annihilation cross section. This quantity is equivalent to the fragmentation function. The results for $1/\sigma_{TOT} d\sigma/dx$, where $x = E/E_{beam}$, are shown in Fig. 6.11 and are summarized in Table 6.8. Our measurements of the proton cross section are in good agreement with results obtained by the TASSO detector. Below 1 GeV/c, however, the TASSO and TPC results for $K^\pm$ disagree significantly, with the TPC values falling below the TASSO measurement by as much as 30% (see insert in Fig. 6.11). The $\pi^\pm$ cross sections for the two experiments are consistent within the quoted errors, but the TPC data are systematically higher. A recent measurement of the $K^\pm$ cross section in the region $0.04 < x < 0.07$ by the MARK II collaboration is slightly higher, but consistent within systematic uncertainties with the TPC values⁷². It is not consistent with the TASSO results.

The charged hadron cross sections, along with the predictions of several Monte Carlo models, are shown in Fig. 6.12 through Fig. 6.14. The models all reproduce the measured $\pi^\pm$ cross section adequately. This agreement is not surprising. All of the Monte Carlos contain tunable parameters that have been determined using PEP and PETRA data. The pion inclusive distribution can be reproduced by any model that generates momenta via longitudinal phase space. In general, the kaon cross section is also well represented by the models, although the Webber Monte Carlo produces too soft a K spectrum. The present version of that model does not treat the fragmentation of heavy quarks properly. (The model decays charm and bottom clusters using a spectator model, where the c or b quark undergoes free-particle β decay.) It is quite likely that an improved treatment of these decays will change the predicted shape (but not the normalization) of the kaon

Figure 6.9: The fraction of charged hadrons that are pions, kaons and protons as a function of momentum. For the two lowest momentum points, the proton fraction was assumed to be 1%.

spectrum. The Lund and Webber models reproduce the baryon spectrum, while the Gottschalk (CIT) cluster model produces too few protons. For Lund, the agreement in rate is again a result of parameter tuning (the diquark to quark production rate was set using PETRA data). In the QCD cluster models the relative rate of baryon to meson production is determined solely by phase space and spin factors in the cluster decays. However, this means that the total baryon rate is a strong function of the maximum cluster mass where phase space decays are permitted, $W_{max}^{cluster}$. The fact that the Webber model reproduces the proton data well is a direct result of tuning $W_{max}^{cluster}$ to 3.5 GeV. Lowering the value by 20% substantially reduced the predicted rate of proton production. In the CIT model, clusters are split via a string mechanism until their masses are quite low. The version of the Monte Carlo used here does not provide any mechanism for creating baryons during this cluster splitting. Thus, baryon production is kinematically suppressed during the phase space decays. The difficulty this model has in reproducing the proton cross section can be overcome by introducing a diquark mechanism during string breaking⁷³.

The mean number of $\pi^\pm$, $K^\pm$ and $(p + \bar{p})$ produced per event is determined by fitting the differential cross sections to the sum of 2 exponentials to extrapolate to momenta where no measurements are possible. Since the systematic uncertainty on the acceptance is correlated between bins, we remove this contribution to the point-to-point errors before fitting. The final uncertainties on the π , K and p multiplicities are then obtained by combining the error obtained by the fit and the estimated uncertainty in acceptance in quadrature. On average, a multihadron event contains $10.7 \pm 0.6 \pi^\pm$, $1.35 \pm 0.13 K^\pm$ and $0.60 \pm 0.08 (p + \bar{p})$.

Figure 6.10: The proton fraction F_1 , as a function of momentum. The curves show the predictions of the Lund Monte Carlo for primary particles (dashed line) and after resonance decays (solid line).

Figure 6.11: The π , K and p differential cross sections $(1/\beta\sigma_{T,\pi}) d\sigma/dx$, where $x = E/E_{beam}$. The insert shows a comparison of our results for the kaon cross section with earlier measurements made by the TASSO group at $\sqrt{s} = 34$ GeV.

Figure 6.12: The pion differential cross section $(1/\beta\sigma_{Tot}) d\sigma_\pi/dx$, where $x = E_\pi/E_{beam}$. Included as well are the predictions of the Lund, Webber and Gottschalk Monte Carlos.

Figure 6.13: The kaon differential cross section $(1/\sigma_{\text{Tot}}) d\sigma_K/dx$, where $x = E_K/E_{\text{beam}}$. Included as well are the predictions of the Lund, Webber and Gottschalk Monte Carlos.

Figure 6.14: The proton differential cross section $(1/\beta\sigma_{T,t})d\sigma_t/dx$, where $x = E_p/E_{\text{beam}}$. Included as well are the predictions of the Lund, Webber and Gottschalk Monte Carlos.

P (GeV/c)	$\langle p \rangle$ (GeV/c)	F_π (%)	F_K (%)	F_p (%)
0.30 - 0.41	0.35	95.7 ± 1.2	3.3 ± 1.2	—
0.41 - 0.50	0.45	93.0 ± 0.8	5.5 ± 0.9	—
0.50 - 0.61	0.55	91.0 ± 0.6	6.5 ± 0.6	2.5 ± 0.5
0.61 - 0.74	0.67	88.3 ± 1.0	7.8 ± 0.6	3.9 ± 0.8
0.74 - 0.82	0.78	88.0 ± 1.3	8.1 ± 4.0	3.9 ± 0.5
0.82 - 0.90	0.86	86.4 ± 1.4	9.0 ± 1.3	4.6 ± 0.6
0.90 - 1.00	0.95	94.5 ± 1.6	10.0 ± 1.4	5.5 ± 0.7
1.00 - 1.35	1.18	—	—	5.3 ± 0.7
1.35 - 1.65	1.50	77.7 ± 2.5	16.3 ± 2.4	5.9 ± 2.1
1.65 - 2.01	1.83	—	16.5 ± 2.7	—
2.01 - 2.23	2.12	71.0 ± 3.4	18.2 ± 2.6	10.8 ± 3.8
2.23 - 2.72	2.46	73.0 ± 3.3	—	—
2.72 - 3.67	3.16	69.1 ± 2.7	—	—
3.67 - 4.95	4.18	64.8 ± 2.7	24.3 ± 2.7	10.9 ± 2.3
4.95 - 6.69	5.47	59.7 ± 2.9	27.4 ± 2.8	12.9 ± 2.4

Table 6.7: The fraction of charged particles that are pions, kaons and protons as a function of momentum.

P (GeV/c)	$e^+e^- \rightarrow \pi^\pm + X$		$e^+e^- \rightarrow K^\pm + X$		$e^+e^- \rightarrow p^\pm + X$	
	$\langle x \rangle$	$\left(\frac{1}{\beta\sigma_{T,\pi}} \right) \frac{d\sigma}{dx}$	$\langle x \rangle$	$\left(\frac{1}{\beta\sigma_{T,K}} \right) \frac{d\sigma}{dx}$	$\langle x \rangle$	$\left(\frac{1}{\beta\sigma_{T,p}} \right) \frac{d\sigma}{dx}$
0.20 - 0.25	0.018	$176. \pm 11.$	—	—	—	—
0.25 - 0.30	0.021	$192. \pm 14.$	—	—	—	—
0.30 - 0.41	0.026	$181. \pm 13.$	0.042	16.2 ± 3.0	—	—
0.41 - 0.50	0.032	$147. \pm 7.$	0.046	17.5 ± 2.6	—	—
0.50 - 0.61	0.039	$122. \pm 5.$	0.051	14.7 ± 1.2	0.075	12.4 ± 2.4
0.61 - 0.74	0.047	97.4 ± 4.2	0.057	12.7 ± 0.9	0.080	12.1 ± 2.6
0.74 - 0.82	0.055	84.1 ± 4.0	0.064	10.5 ± 1.5	0.084	9.0 ± 1.2
0.82 - 0.90	0.060	73.1 ± 3.5	0.068	9.8 ± 1.4	0.088	8.4 ± 1.1
0.90 - 1.00	0.066	65.6 ± 3.2	0.074	9.6 ± 1.4	0.092	8.3 ± 1.0
1.00 - 1.35	—	—	—	—	0.103	5.0 ± 0.6
1.35 - 1.65	0.103	31.7 ± 1.7	0.108	7.3 ± 1.1	0.122	3.4 ± 1.2
1.65 - 2.01	—	—	0.130	5.0 ± 0.7	—	—
2.01 - 2.23	0.146	19.8 ± 1.0	0.150	4.5 ± 0.7	0.160	3.1 ± 1.2
2.23 - 2.72	0.170	12.4 ± 0.7	—	—	—	—
2.72 - 3.67	0.218	$7.14 \pm .41$	—	—	—	—
3.67 - 4.95	0.288	$3.40 \pm .21$	0.290	$1.29 \pm .15$	0.295	$0.60 \pm .13$
4.95 - 6.69	0.378	$1.36 \pm .09$	0.379	$0.63 \pm .07$	0.383	$0.30 \pm .06$

Table 6.8: The differential cross sections for inclusive pion, kaon and proton production.

Chapter 7

Rapidity Distributions and Correlations

The QCD improved parton model successfully describes the process $e^+e^- \rightarrow \text{hadrons}$. The observation of jet structure provides evidence for the dominance of annihilation via two partons. The existence of planar events supports the QCD prediction of gluon bremsstrahlung. The mechanism by which the initial partons turn into hadrons is, however, not well understood. The majority of e^+e^- data can be reproduced by a variety of Monte Carlo models that have widely different approaches to hadronization. Even the most basic tenets of the standard models, such as local conservation of quantum numbers and of $p_\perp$, are difficult to verify experimentally. In this chapter we attempt to study the hadronization process by measuring two-particle correlations in rapidity space. The aims of the analysis are to demonstrate that the process $e^+e^- \rightarrow \text{hadrons}$ is mediated by an initial $q\bar{q}$ pair and to provide evidence for local conservation of quantum numbers during hadronization.

We choose to study correlations as a function of rapidity for several reasons. First, since longitudinal phase space models predict a uniform y distribution, the presence of any structure in rapidity indicates non-trivial physics. Second, rapidity distributions transform simply under Lorentz boosts:

$$y' = y + \xi \quad \text{where} \quad \xi = \text{arctanh}(\beta) \quad (7.1)$$

Rapidity differences are therefore independent of Lorentz frame. Finally, rapidity information is not destroyed by resonance decays. If a particle with rapidity y decays, its decay products are distributed about y .

The techniques used to measure rapidity correlations were first developed by groups studying hadronic interactions^{74,75}. More recently, the TASSO and PLUTO groups at DESY have applied these methods to e^+e^- annihilations^{52,63}. In these studies, a particle (called the "test" particle) is chosen in a specified rapidity range and the behavior of all other particles in the event is examined, either by calculating a correlation function or by finding an associated density.

7.1 Acceptance Corrections

Two algorithms, sphericity and thrust, are commonly used to define the jet axis. The sphericity algorithm has already been described in Section 2.3. This method finds the axis that minimizes the sum over all tracks of the $p_\perp^2$. The thrust axis is the direction that maximizes the sum of $|p_{\parallel}|$. We have chosen to estimate the initial parton direction using the sphericity axis calculated from charged tracks (including conversion electrons). For rapidities between .1 and 3.5, differences in $d\sigma/dy$ between the sphericity and thrust methods are less than 10%. The value of y used to generate Monte Carlo predictions is also calculated using the sphericity axis found from charged particles.

The rapidity distributions are measured using the event and track selection criteria described in Sections 4.5 and 6.1, with two exceptions. First, to insure adequate reconstruction of the jet axis, we require that the cosine of the angle between the sphericity axis and the beamline be less than 0.8. Second, the cut on the minimum number of dE/dx samples has been reduced to 40.

Figure 7.1: The effect of holes in the forward acceptance on the measured rapidity distributions is demonstrated a) for events where the jet axis is perpendicular to the beamline and b) for events where the jet axis is close to the beamline. In the first case tracks at low rapidity (perpendicular to the jet axis) are preferentially lost and in the second case high rapidity tracks (those close to the jet direction) are missed.

Acceptance corrections are made as a function both of y and of the angle of the jet axis with respect to the beamline, θ_{jet} . For a given jet direction, the forward holes due to the beampipe result in a localized dip in the acceptance (See Fig. 7.1 and Fig. 7.2). Such dips are especially dangerous in two particle studies, where they appear as spurious correlations. For distributions where particle identification is not used, the acceptances are calculated in the same manner as for the inclusive distributions. Lund Monte Carlo events, generated with radiative effects, are passed through a detector simulation and analyzed using the same cuts as for the real data. The number of tracks per bin per event is then compared to the number generated by the Lund Monte Carlo without radiative effects.

To find the rapidity distributions for pions, kaons and protons, we could in principle fit the particle content in each rapidity bin using the techniques describe in the previous chapter. This method, however, is not applicable to the correlation studies that follow. We therefore choose to use the identification algorithm described in Section 5.5 and to correct the measured distributions for particle misidentification. For studies involving identified pions, kaons and protons, we include in the distribution of measured particles of species a all tracks that meet the usual track selection

Figure 7.2: The pion acceptance as a function of the angle of the jet axis with the beamline for several values of y

cuts and in addition have

$$\begin{aligned} W_a(\chi^2) &\geq 0.7 & \text{and} \\ W_r(\chi^2) &< 0.15 \end{aligned} \quad (7.2)$$

where $W(\chi^2)$ is defined in Eq. 5.10. The effect of these cuts is demonstrated in Fig. 7.3. The plot shows the dE/dx distribution and the corresponding fit to the number of pions, kaons and protons for tracks with $3.7 \leq p < 6$ GeV/c. The cross hatched areas represent those particles that meet the above $W(\chi^2)$ cuts. The method produces a clean pion sample (purity greater than 90%, acceptance greater than 0.5) at all rapidities. However, due to the large number of pions relative to kaons and protons, and due to the small dE/dx separation between kaons and protons, highly pure K and p samples cannot be reached in the high momentum (and high rapidity) region. The kaon acceptance in the range $2 < y < 4$ is only about 25% and the sample purity is approximately 75%. The corresponding proton purity and acceptance are similar. The combination of the $W_r(\chi^2)$ cuts and the requirement of at least 40 usable dE/dx samples reduces the electron contamination to less than 2% at all rapidities.

The purity of the observed K sample depends not only on the number of kaons, but on the number of pions and protons as well. Purity corrections therefore cannot be handled as multiplicative factors. The true particle distributions instead must be unfolded from the simultaneously measured π , K and p distributions. We define the misidentification probability,

$$\epsilon_{ij}(y, \theta_{Sph}) = \begin{array}{l} \text{the probability that a track in the TPC of} \\ \text{species } j \text{ is identified as being species } i \end{array}$$

Figure 7.3: The distribution of dE/dx (normalized to the expected value for pions of the same momentum) for tracks with momenta between 3.7 and 6 GeV/c. The cross hatched regions represent those particles that would be called pions, kaons and protons with the criteria given in Eq. 7.2.

This ϵ matrix includes acceptance losses due to track and dE/dx selection criteria, as well as misidentification probabilities. Its coefficients depend on the dE/dx separation between species and on the dE/dx and momentum resolutions. The results for ϵ can be calculated via Monte Carlo and are not sensitive to the details of the event generator. The overall multiplicity, particle fractions and degree of jet collimation affect ϵ only to the extent that the number of usable dE/dx samples (which determines the dE/dx width) and the distribution of track measurement errors (which determines the momentum resolution) are effected by these quantities. The measured distribution of particles of type i can be expressed in terms of ϵ and the true particle distributions in the TPC:

$$\rho_i^{\text{Meas}} = \sum_j \epsilon_{ij}(y, \theta_{\text{TPC}}) \rho_j^{\text{True,TPC}} \quad (7.3)$$

This system of simultaneous equations can be solved by inverting the ϵ matrix:

$$\rho_i^{\text{True,TPC}} = \sum_j \epsilon_{ij}^{-1}(y, \theta_{\text{TPC}}) \rho_j^{\text{Meas}} \quad (7.4)$$

The misidentification probabilities relate the true and observed track distributions in the TPC. An overall acceptance correction is then applied to determine the number of particles of species i at the event vertex from the number of particles seen in the TPC:

$$\rho_i^{\text{True,Vertex}} = A_i \rho_i^{\text{True,TPC}} \quad (7.5)$$

This correction factor takes into account energy loss in the beampipe and pressure wall, decays in flight, detector geometry and radiative effects.

Since changing a particle's mass hypothesis systematically shifts the measured rapidity, separate ϵ values must be computed under the assumptions that all particles are pions, kaons or protons. For example, to determine the kaon rapidity distribution, we calculate the rapidity y_K for all tracks using the kaon mass to determine each particle's energy. The inversion is then performed on the distributions $d\sigma_\pi/dy_K$, $d\sigma_K/dy_K$ and $d\sigma_p/dy_K$, using the appropriate ϵ coefficients.

The extension of this technique to two particle correlations is straightforward. The probability that particle a with identity k and particle b with identity ℓ are measured as having identities m and n , respectively is

$$\bar{\epsilon}_{ij}(y_1, y_2, \theta_{Sph}) = \epsilon_{mk}(y_a, \theta_{Sph}) \epsilon_{n\ell}(y_b, \theta_{Sph}) \quad (7.6)$$

where y_a and y_b are the rapidities of particles a and b , j is the two particle state where tracks 1 and 2 have identities k and ℓ , respectively, and i is the state where the tracks 1 and 2 have identities m and n . The single particle misidentification probabilities ϵ_{mk} and $\epsilon_{n\ell}$ are 3×3 matrices (each index can be a π , K or p). The two particle probability, $\bar{\epsilon}_{ij}$, is a 9×9 matrix (i and j can be $\pi\pi$, πK , πp , $K\pi$, etc.). Thus, $\bar{\epsilon}_{ij}^{-1}$ relates the nine true two particle densities $\rho_{kr}^{True,TPC}$ to the nine measured two particle densities ρ_{mn}^{Meas} :

$$\rho_{kr}^{True,TPC} = \bar{\epsilon}_{ij}^{-1}(y, \theta_{Sph}) \rho_{mn}^{Meas} \quad (7.7)$$

The Monte Carlo determination of ϵ involves systematic uncertainties which must be included in the quoted errors on the rapidity distributions. As with the fitting of the inclusive distributions, a major source of uncertainty is the approximation of the shape of the dE/dx curve vs. $\beta\gamma$. Applying the same technique as was used to determine the systematic uncertainty in the dE/dx fits (see Table 6.2), we have determined the systematic errors on the elements of ϵ . In addition to these species-dependent errors, we estimate a systematic error on the overall normalization of the rapidity distributions of $\pm 5\%$.

Muons and pions cannot be separated using the dE/dx method because their masses are so similar. Our "pion" sample therefore includes a small ($\sim 1\%$) contamination of direct muons. We have chosen to include these muons rather than subtract out this contribution because a subtraction based on Monte Carlo calculations can not be done reliably for the two particle flavor correlation studies to follow.

7.2 Measurement of Inclusive Rapidity Distributions

To allow comparison with other experiments, we present the inclusive rapidity distribution calculated for all particles using the pion mass to determine y . This distribution is shown in Fig. 7.4, along with similar results from the TASSO⁷⁶ experiment and the predictions of several Monte Carlo models. The data exhibit the expected "rapidity plateau" for $-2 < y < 2$ and then falls off. A small decrease in $d\sigma/dy_\pi$ is seen at $y_\pi = 0$. All of the models, except for the independent fragmentation Monte Carlo, produce some $y = 0$ dip. None of the Monte Carlos fit the data well, however. Both Lund and Webber underestimate the size of the dip. CIT gives a dip of about the right magnitude, but predicts an overall spectrum that is too hard. We have not attempted to tune the models to improve the agreement.

Proponents of different hadronization schemes attribute the $y = 0$ dip to different causes. In the CIT model, the kinematics of heavy-cluster decays can cause an excess of high rapidity particles. Because perturbative radiation is cut off at large Q^2 in this model, such heavy clusters are common⁷³. In the Lund and Webber models, a dip can result from the treatment of gluons^{76,48}. In Lund, the depopulation results from string dynamics⁷⁷ and in Webber it is the result of angle ordering⁴⁸. There are, however, more mundane explanations of the rapidity dip. At least part of this structure is the result of assigning the pion mass to all particles. This procedure systematically shifts kaons and protons to higher values of rapidity. The size of that effect can be measured using the matrix inversion technique described above. The π , K and p

Figure 7.4: The normalized differential cross section $1/\sigma_{TOT} d\sigma/dy_\pi$, as measured by this experiment ($\sqrt{s} = 29$ GeV) and by the TASSO collaboration⁷⁶ ($\sqrt{s} = 34$ GeV). The TPC distributions have been calculated relative to the sphericity axis. The TASSO analysis uses the thrust axis. This difference accounts for the disagreement at small y . The predictions of several Monte Carlo models are included as well.

contributions to $d\sigma/dy_\pi$ are shown in Fig. 7.5. Approximately half of the “dip” is actually a bump due to kaons and protons. Heavy quark decays can also cause a bump at moderate y .

The rapidity distributions for pions, kaons and protons are presented in Table 7.1 and in Figures 7.6 through 7.8. In all three plots, we have folded the distributions about $y = 0$. The error bars include the systematic uncertainties in the matrix inversion method described above, but do not include the overall $\pm 5\%$ uncertainty in normalization. The proton distributions have been calculated using only negatively charged tracks to prevent contamination from nuclear interactions in the beampipe. Because the proton acceptance in the K-p dE/dx cross-over region is so low, the proton yield in the rapidity bin $1.33 < |y| < 1.67$ has an error of more than 100%. We therefore do not quote a result in this region.

In general, the agreement between data and Monte Carlo is acceptable. The two exceptions to this rule are the predictions of the Webber model for the kaon distribution, which are too soft, and the Goitschalk model predictions for proton cross section, which are too small and are peaked at high y . Both of these pathologies have been discussed in Section 6.4.

Integrating the rapidity distributions, we obtain $\pi^\pm$, $K^\pm$ and $(p+\bar{p})$ multiplicities of 10.7 ± 0.6 , 1.43 ± 0.13 and 0.64 ± 0.09 respectively. These numbers are in good agreement with the results obtained by extrapolating the inclusive cross sections (see Section 6.4).

Figure 7.5: The contribution of pions, kaons and protons to $d\sigma/dy_\pi$. About half of the observed dip is in fact a bump at $y = 1$ caused by kaons and protons.

Figure 7.6: The normalized differential cross section $1/\sigma_{\text{Tot}} d\sigma_\pi/d|y|$ for charged pions. Included are the predictions of several Monte Carlo models.

Figure 7.7: The normalized differential cross section $1/\sigma_{T,\text{tot}} d\sigma_K/d|y|$ for charged kaons. Included are the predictions of several Monte Carlo models.

Figure 7.8: The normalized differential cross section $1/\sigma_{T,\text{tot}} d\sigma_P/d|y|$ for protons plus antiprotons. Included are the predictions of several Monte Carlo models.

Rapidity Bin	$\pi^\pm$	$K^\pm$	$(p + \bar{p})$
0.00- 0.33	4.078 ± 0.035	0.456 ± 0.057	0.269 ± 0.050
0.33- 0.67	4.238 ± 0.034	0.429 ± 0.040	0.287 ± 0.025
0.67- 1.00	4.312 ± 0.033	0.512 ± 0.038	0.290 ± 0.023
1.00- 1.33	4.290 ± 0.033	0.506 ± 0.049	0.250 ± 0.057
1.33- 1.67	3.980 ± 0.033	0.610 ± 0.049	-
1.67- 2.00	3.531 ± 0.032	0.588 ± 0.056	0.263 ± 0.162
2.00- 2.33	2.808 ± 0.029	0.479 ± 0.116	0.165 ± 0.033
2.33- 2.67	2.033 ± 0.025	0.334 ± 0.029	0.102 ± 0.015
2.67- 3.00	1.284 ± 0.020	0.196 ± 0.015	0.045 ± 0.010
3.00- 3.33	0.770 ± 0.015	0.112 ± 0.009	-
3.33- 3.67	0.435 ± 0.011	0.056 ± 0.007	-
3.67- 4.00	0.265 ± 0.008	0.023 ± 0.004	-

Table 7.1: The normalized differential cross section $1/\sigma_{T,\text{tot}} d\sigma/d|y|$ for pions, kaons and protons. The correct particle masses are used to determine y and the jet axis is found by the sphericity method. No measurements are quoted for bins where the error is greater than 100%.

7.3 Charge Correlations

Correlation studies provide a means of examining the two time scales governing hadronization. The large momentum transfer involved in the creation of the initial parton system results in a large separation in phase space of the hadrons containing those partons. Since the primary quarks carry opposite quantum numbers, this process gives rise to long-range quantum number correlations (LRC). On the other hand, the final formation of hadrons is usually assumed to be governed by small momentum transfers. Such processes generate local, short-range correlations (SRC). Previous LRC studies by the TASSO and PLUTO groups have provided evidence for the existence of charged primary partons (quarks)^{52,53}. We present in this section similar results from the TPC. Since the measurements here can be made in the absence of particle identification, we use the pion mass to calculate all rapidities in this section. For the flavor correlation studies presented in Section 7.4, however, we use the correct mass. The event and track selection criteria for this analysis are identical to those used for the inclusive rapidity measurements. We have not introduced any cuts to remove three jet candidates from the sample, since such cuts would introduce model dependent systematic uncertainties in the results.

Let us begin by choosing a “test” particle at rapidity y_{test} . Charge conservation requires that the total charge of the remaining particles in the event exactly cancel that of the test particle. To examine the distribution of this compensating charge in rapidity, we introduce the associated charge density⁷⁴, i.e. the probability that the test particle’s charge is compensated by an oppositely charged particle at rapidity y :

$$q(y) \equiv \rho^+(y) - \rho^-(y) \quad (7.8)$$

where $\rho^+(y)$ and $\rho^-(y)$ are the densities ($1/\sigma_{T,\text{tot}} d\sigma/dy$) of particles at rapidity y with the opposite and same charge as the test particle, respectively. The charge conservation constraint is

$$\int q(y) dy = 1 \quad (7.9)$$

Fig. 7.9 shows the associated charge density $q(y)$ for four y_{test} bins: a) $0 \leq y_{\text{test}} < 0.75$, b) $0.75 \leq y_{\text{test}} < 1.5$, c) $1.5 \leq y_{\text{test}} < 2.5$ and d) $2.5 \leq y_{\text{test}} < 4.0$. The predictions of the

Figure 7.9: The associated charge density $q(y)$ as a function of rapidity for four ranges of test particle rapidity y_{test} : a) $0 \leq y_{test} < 0.75$, b) $0.75 \leq y_{test} < 1.5$, c) $1.5 \leq y_{test} < 2.5$ and d) $2.5 \leq y_{test} < 4.0$. The solid and dotted lines show the predictions of the Lund and Webber models, respectively.

Figure 7.10: The charge asymmetry A for the y_{test} range $1.5 \leq y_{test} < 4.0$. The solid curve shows the predictions of the Lund model. The dashed curve represents a model with neutral primary partons.

Lund and Webber Monte Carlos are also included. In the lowest test bin, $q(y)$ is dominated by a peak at $y = 0$. As the y_{test} range is moved, this signal moves as well. For large y_{test} , the peak is no longer centered about the test range. This asymmetry is a consequence of energy-momentum conservation; it is not possible to produce two particles at high rapidity in the same jet. Although the data exhibit a clear short-range correlation, the interpretation of this result is not straightforward. The models described in Chapter 2 all predict local charge conservation in the hadronization process. However, the presence of large numbers of resonance decays such as $\rho^0 \rightarrow \pi^+ \pi^-$ and $\omega \rightarrow \pi^+ \pi^- \pi^0$ provides an independent source of such correlations. Because we cannot accurately determine the total resonance contribution to the short-range correlation signal, we cannot use these data to test the principle of local charge compensation at the parton level.

For the two highest y_{test} intervals, $q(y)$ exhibits a small but statistically significant opposite-jet peak. Both Monte Carlo models reproduce this feature qualitatively, although they slightly underestimate its size. We interpret this long-range correlation as evidence that the primary partons are charged. To support this viewpoint we compare in Fig. 7.10 the charge asymmetry⁵²

$$A = \frac{\rho^O(y) - \rho^S(y)}{\rho^O(y) + \rho^S(y)} \quad (7.10)$$

for the y_{test} range $1.5 \leq y_{test} < 4.0$ to the predictions of the Lund Monte Carlo and of a model with neutral initial partons. The solid curve shows the asymmetry found with the Lund model. The dashed curve is the result of a Monte Carlo calculation where the initial quark-antiquark pair is replaced by a pair of gluons. The Lund model predicts a long-range asymmetry, similar to that seen in the data, while the neutral parton model exhibits a much smaller effect. (The remaining asymmetry in the neutral parton case is a result of energy-momentum conservation,

coupled to charge conservation. If a particle is produced at large rapidity, it uses a large fraction of the jet energy. Thus, there will be few other particles in that jet and it is more likely that its charge will be compensated in the other jet.) Because the data do not fit the neutral parton model predictions, we conclude that the observed long-range correlation provides evidence that the primary partons carry charge.

7.4 Flavor Correlations

Due to the small number of s quarks produced in a typical event, strangeness is a much cleaner label of a specific $q\bar{q}$ pair than charge. In this section we extend the above correlation studies to identified pions and kaons. The goals of the analysis are twofold. First, by studying short-range strangeness correlations, we hope to test the picture of color screening via vacuum polarization. Second, by examining long-range KK and πK correlations, we expect to demonstrate that the primary partons carry opposite flavors.

In analogy to Eq. 7.8, we define the flavor-tagged associated charge density $q_a^b(y)$ as the compensating charge density seen at rapidity y in particles of species b when a test particle of species a is placed at rapidity y_{test} . The charge conservation condition (the analog of Eq. 7.9) now becomes

$$\sum_{\substack{b=a,b,\bar{b} \\ s_{\text{test}}, y_{\text{test}}}} q_a^b(y) dy = 1 \quad (7.11)$$

In the case where a and b are kaons, $q_K^K(y)$ is equivalent to an associated strangeness density. This density does not integrate to 1, however, because it includes only that fraction of the strangeness carried by charged kaons.

The mechanisms responsible for various flavor correlations are shown in Fig. 7.11. Both short-range and long-range opposite-charge correlations exist between pions. The short-range correlations are due to a combination of resonance decays and primary hadronization. The long-range correlations reflect conservation of charge between the primary quarks. Short-range KK correlations result from the vacuum production of $s\bar{s}$ pairs. A small contribution (estimated below) from resonance decays is also present. Long-range KK correlations come from two sources. First, if an initial $s\bar{s}$ pair is created from the virtual photon, kaons will be created in both jets. Second, since charm and bottom quarks almost always decay into kaons, correlations between high rapidity kaons occur as a result of heavy quark decays. In addition to producing KK correlations, charm decays provide a source of long-range πK correlations. Because the charm and anticharm quarks in opposite jets undergo the decays $c \rightarrow \pi^+ K^- X$ and $\bar{c} \rightarrow \pi^- K^+ X$, this signal has the peculiar feature of being a same-charge rather than opposite-charge correlation (see Fig. 7.11c). Short-range πK correlations result primarily from the decay of such resonances as the K^{*0} and the D .

The flavor-tagged associated charge densities for the rapidity ranges $0 \leq y_{\text{test}} < 1.5$ and $1.5 \leq y_{\text{test}} < 4$ are shown in Fig. 7.12 through Fig. 7.14. The first region corresponds to the central rapidity plateau; the second is enriched in hadrons containing the primary quarks. These figures also include the predictions of the Lund and Webber models. These models represent philosophically opposite approaches to hadronization. The extent to which they agree is an indication of our lack of sensitivity to the details of the hadronization scheme. To demonstrate the influence of heavy quark decays on the shape of the long-range correlation, we present the results obtained with the Lund model when heavy quark production is suppressed.

The observed $\pi\pi$ correlations look quite similar to the charge correlations discussed in the previous section. They are adequately represented by both Lund and Webber. Integrating $q_\pi^\pi(y)$, we find that $(88 \pm 3)\%$ of the time the charge of a pion is compensated by another pion.

In both y_{test} bins, $q_K^K(y)$ exhibits a small, but statistically significant short-range correlation. (Note that the overall normalization of $q_K^K(y)$ is a factor of five below $q_\pi^\pi(y)$ and $q_K^\pi(y)$.) In addition, the high y_{test} region shows the same-charge long-range correlation expected from heavy quark decays. Both the Lund and Webber Monte Carlo reproduce this correlation, while the

Figure 7.11: The mechanisms responsible for long-range and short-range flavor correlations. a) Long-range pion correlations reflecting the initial $u\bar{u}$ quark flavors and short-range correlations due to resonance decays and to vacuum production of a $d\bar{d}$ pair. b) Long-range kaon correlations due to production of a primary $s\bar{s}$ pair and short-range correlations from the decay of ϕ mesons. c) Long-range KK and πK correlations due to initial charm production and KK short-range correlations due to strange production from the sea.

Figure 7.12: The associated charge density $q_{\pi}^{\pi}(y)$ for the rapidity ranges a) $0 \leq y_{\text{Test}} < 1.5$ and b) $1.5 \leq y_{\text{Test}} < 4.0$. The solid and dotted lines show the predictions of the Lund and Webber Monte Carlos. The dashed line gives the prediction of Lund in the absence of c and b quarks.

Figure 7.13: The associated charge density $q_{\pi}^K(y)$ for the rapidity ranges a) $0 \leq y_{\text{Test}} < 1.5$ and b) $1.5 \leq y_{\text{Test}} < 4.0$. The solid and dotted lines show the predictions of the Lund and Webber Monte Carlos. The dashed line gives the prediction of Lund in the absence of c and b quarks.

Figure 7.14: The associated charge density $q_K^K(y)$ for the rapidity ranges a) $0 \leq y_{\text{Test}} < 1.5$ and b) $1.5 \leq y_{\text{Test}} < 4.0$. The solid and dotted lines show the predictions of the Lund and Webber Monte Carlos. The dashed line gives the prediction of Lund in the absence of c and b quarks.

Figure 7.15: The normalized correlation function $R^{KK}(y, y_{test})$ for a) opposite charge and b) same charge kaons.

model without heavy quarks fails to do so. Although the high y_{test} distribution predicted by the Webber model is too soft, this problem is most likely a result of the inadequate treatment of heavy quark fragmentation and does not represent a fundamental difficulty with the model.

The KK densities exhibit short-range correlations for both y_{test} ranges. The high y_{test} bin shows a large long-range correlation as well. Integrating $g_K^a(y)$ gives a value of $(54 \pm 6)\%$ for Fig. 7.14a and $(57 \pm 7\%)$ for Fig. 7.14b. These numbers are consistent with the fact that approximately half of the strangeness is carried by neutral particles.

The short-range KK correlation observed in Fig. 7.14 could in principle result either from an increase in the number of oppositely charged kaons or a reduction in the number of same charge kaons for events where a test K has been chosen. To distinguish between these possibilities, we introduce the normalized correlation function, $R^{ab}(y, y_{test})$. This function relates the two particle density, $\rho^{ab}(y, y_{test})$, to the single particle uncorrelated densities $\rho^a(y)$ and $\rho^b(y_{test})$:

$$R^{ab}(y, y_{test}) = \frac{\rho^{ab}(y, y_{test})}{\rho^a(y)\rho^b(y_{test})} - 1 \quad (7.12)$$

Fig. 7.15 shows the normalized correlation function for same-charge and for opposite-charge KK combinations, averaged over the y_{test} region $0 \leq y_{test} < 1.5$. A positive correlation is clearly present in the opposite-charge data. No correlation is seen in the same-charge plot. (We do not present R for the large y_{test} region since the constraint of energy-momentum conservation produces a negative correlation that overpowers any dynamic correlations in the data.)

The data presented above demonstrates that strangeness is locally conserved in e^+e^- annihilation events. In order to interpret this result as evidence for local flavor compensation at the parton level, we must show that the observed correlation is not merely the byproduct of normal resonance decays. We present here an estimate of the contribution of resonance decays to our observed K^+K^- signal. If the primary mesons produced during hadronization are predominantly in the ground state scalar and vector nonets, the ϕ meson is the only resonance that produces a significant number of KK pairs at in the low y_{test} bin. At high y_{test} , the F meson contributes as well. Fig. 7.16 shows the invariant mass distribution for all K^+K^- pairs contributing to our cor-

Figure 7.16: The invariant mass distribution for all K^+K^- pairs that contribute to $q_K^K(y)$.

relation signal. A peak is seen in the ϕ mass region⁷⁸. We estimate the ϕ contribution to our KK signal to be approximately 10% of the observed SRC for the small y_{test} range and approximately 15% for the large y_{test} range. The cross section for F production cannot be measured using the presently available TPC data. We therefore have estimated the F contribution to our signal, using the Lund Monte Carlo. We find that F mesons (including the $F \rightarrow \phi\pi$ mode) contribute 5% of the SRC signal in Fig. 7.14b. It appears, therefore, that the observed KK short-range correlation is too large to be explained by resonance decays and is in fact evidence that color screening is a soft process and that flavor is locally conserved during hadronization.

Because the kaon rapidity spectrum has a significant dip at $y = 0$, it is important to demonstrate that the apparent division of $q_K^K(y)$ in Fig. 7.14b into long-range and short-range correlations is not merely a reflection of this dip. We therefore present in Fig. 7.17 for pions and for kaons the analog of the charge asymmetry defined in Eq. 7.10. The data indicate that when a test particle is selected at large y , the asymmetry goes through a minimum at $y = 0$ both for pions and for kaons. Thus, there is a separation between the long-range and short-range flavor compensation signals.

The relative strengths of the long-range and short-range correlations depends on the probability that a particle at rapidity y_{test} is the decay product of a meson that contained the primary quark or antiquark. To examine size of the long-range correlation for pions and for kaons, we have parameterized $q(y)$ as a sum of two gaussians, corresponding to the long-range and short-range components of the signal. The relative strength of the long-range correlation is defined to be the ratio of the area of the long-range correlation gaussian to the total area of the two gaussians together:

$$F_{LR} \equiv \frac{\int_{LR} q(y)}{\int_{SRC} q(y) + \int_{LR} q(y)} \quad (7.13)$$

For $q_\pi^K(y)$ and $q_K^K(y)$, the fitted positions (y_{SRC} and y_{LR}) and widths (σ_{SRC} and σ_{LR}) of these gaussians and the value of F_{LR} are given in Table 7.2. To interpret these results, let us examine the pattern of correlations we expect if test particle a is a meson that contains quark-antiquark

Figure 7.17: The charge asymmetry for a) pions and b) kaons with $1.5 \leq y_{\text{cent}} < 4$. This test region is enriched in hadrons that contain the primary quark or antiquark. The curve shows the predictions of the Lund Monte Carlo.

pair $q_1 \bar{q}_2$. If a is not a decay product of the meson containing the initial quark, its flavor must be compensated locally. Therefore, it produces only a short-range correlation signal. If, on the other hand, particle a does contain a decay product of the initial quark, both long-range and short-range correlations are found. The presence of the primary antiquark in the other jet produces a long-range flavor compensation signal. Because antiquark $\bar{q}_2$ was pulled from the vacuum along with a companion quark q_2 , a short-range correlation is present as well. In a model where pions are the only particles produced during hadronization and where u and d quarks are the only initial partons (and they are produced at the same rate), the relative strength of these short-range and long-range correlations must be equal. For a realistic model, however, the strength of the short-range correlation is influenced by decays, the fact that the number of s quarks pulled from the sea is smaller than the number of u and d quarks and the fact that the relative rates of primary u , d , s , c and b production are not equal. We therefore cannot interpret our measurements of F_{LRC} in model independent manner. If, however, we use the Lund Monte Carlo to guide our expectations, we find the following. Averaging over all initial quark species, Lund predicts that kaons resulting from the decay of hadrons containing a primary quark exhibit a long-range correlation strength of 53%. In the case of pions, the corresponding number is only

Table 7.2: The results of fitting $q_a^{\pi}(y)$ and $q_b^K(y)$ to the sum of two gaussians, representing the short-range and long-range components of the correlation signal. F_{LRC} is defined in Eq. 7.13.

	σ_{SRC}	y_{SRC}	σ_{LRC}	y_{LRC}	F_{LRC}
$\pi\pi$	$0.95 \pm .02$	$1.64 \pm .02$	1.1 ± 0.1	-2.1 ± 0.1	$6.16 \pm .02$
KK	$0.82 \pm .13$	$1.58 \pm .17$	0.8 ± 0.2	-1.7 ± 0.3	$0.37 \pm .06$

Figure 7.18: The measured value of F_{LRC} as a function of $x_F = p_{||}/E_{\text{beam}}$ for pions and for kaons.

26% because resonance decays enhance the size of the observed short-range correlation. Although these values are model dependent (they will change with the vector to pseudoscalar production ratio, for example), the implication is clear. The data are consistent with the primary mechanisms for π and K long-range correlations being the same, with resonance decays significantly distorting the observed $\pi\pi$ spectrum.

We have investigated this conclusion as a function of $x_F = p_{||}/E_{\text{beam}}$ as well. The technique is identical to that used above, except that the test particle is chosen to fall within a fixed x_F range rather than within a given y_{cent} bin. Because the effect of dE/dx cross-overs is more pronounced in this variable, the bins do not cover the complete x_F range. Due to the limited statistics in each x_F bin, the fitted results exhibited a large correlation between the width σ_{LRC} and the center y_{LRC} of the long-range correlation gaussian. During fitting we therefore have fixed one of these quantities. The quoted errors include the systematic difference between the two methods of fitting. The results are shown in Fig. 7.18. Although the error bars are quite large, the data suggest that the pion and kaon decay products of primary partons have similar x_F distributions.

In conclusion, we observe both short-range and long-range flavor correlations in e^+e^- an-

annihilation. Short-range correlations support the hypothesis that quantum numbers are locally conserved during hadronization, as expected in models where $q\bar{q}$ pairs are created in the color force field. Long-range correlations provide additional evidence that jets are produced by a pair of flavor-carrying partons (quarks).

Chapter 8

Summary and Conclusions

The PEP-4 Time Projection Chamber provides charged particle tracking over close to 4π steradian solid angle. In addition, the detector provides mass identification for stable charged hadrons over most of the available momentum range by simultaneously measuring each track's momentum and ionization energy loss. A dE/dx resolution of better than 4% has been achieved in multihadron final states. This resolution makes it possible to select samples of relativistic protons and charged kaons, with overall purities of greater than 70%. The corresponding sample of charged pions has a purity of more than 90%.

Using the superior particle identification capabilities of the TPC, we have studied the inclusive production of charged hadrons. Our measurements of the $\pi^\pm$, $K^\pm$ and $(p + \bar{p})$ differential cross sections indicate that similar mechanisms are responsible for producing all particle species. A typical multihadronic final state at a center-of-mass energy of 29 GeV contains 10.7 ± 0.6 charged pions, $1.3^\pm \pm 0.13$ charged kaons and 0.6 ± 0.08 protons plus antiprotons.

We have also studied flavor correlations in rapidity space. The data are consistent with the assumption that the reaction $e^+e^- \rightarrow \text{hadrons}$ is governed by two distinct processes. The large momentum transfer interaction can be described by the production of a quark-antiquark pair, plus some number of hard gluons. Such an interaction results in long-range flavor correlations. The transformation of the initial partons into observable hadrons occurs on a much longer time scale. It is characterized by limited transverse momentum relative to the initial parton direction and by local neutralization of quantum numbers.

Several phenomenological descriptions of hadronization have been developed in recent years. These models all begin with the principles discussed above, but differ widely in their approach to the hadronization process. Despite these differences, the models all do a reasonable job of describing the data. Although there are a few cases where agreement between a given model and the data is poor, these instances are most likely a result of inadequate tuning of the Monte Carlo models than evidence of fundamental difficulties.

One of the challenges of the next generation of hadronization studies is to find reliable techniques for differentiating between models (or alternately, showing that the seemingly different approaches are in fact identical). We had originally hoped that the correlation studies presented in the previous chapter could provide such a technique. A glance at the similarities of the predictions of the various Monte Carlos, however, indicates that this approach is not likely to be fruitful. However, the KK correlation studies have perhaps given us some hints of an approach for the future. In Section 7.4 we argued that the advantage of kaons over pions for correlation studies is that the kaon spectrum is not as sensitive to resonance decays. That argument is even stronger in the case of baryons. In addition, the mechanisms responsible for baryon production in the various models are quite different. Although the TPC group has already begun to examine baryon-antibaryon correlations⁷⁹, the data suffer from limited statistics. We hope in the future to have enough data to study the baryon-antibaryon correlation function as a function of rapidity and of $p_\perp$.

The study of the mechanisms by which quarks and gluons are transformed into hadrons is

still in its earliest stage. The results presented in this dissertation support the basic belief that hadronization occurs via local color neutralization and thus through the production of low mass color singlet states. A great deal of work is still necessary in order to understand the nature of these color singlets and how they are formed.

Appendix A

Energy Loss in Argon-Methane

A quantitative treatment of ionization energy loss requires an understanding of the atomic structure of the gaseous medium. We present here a calculation of the dE/dx spectrum for Argon-Methane gas. Our approach is based heavily on the work of Lynch⁸⁰, Hadley⁸¹, Lapique and Piuz⁶⁹ and Allison and Cobb⁵⁹.

The calculation of the energy loss distribution can be divided into four steps. First, the oscillator strength $f(E)$ is approximated by a sum of analytic terms corresponding to the atomic energy levels and the complex dielectric constant ϵ is calculated in terms of the approximate $f(E)$. Second, the mean number of resonance and Rutherford type collisions that occur per atomic level per unit distance is determined. Next, the distribution of observed dE/dx per collision is derived. Finally, the results of the previous steps are folded together either by Monte Carlo or by analytic approximation.

Following the method of Lapique and Piuz⁶⁴, we will approximate the oscillator strength by the sum

$$f(E) = \sum_i w_i f_i(E) \quad (A.1)$$

where w_i is the probability that an electron in the gas is found in atomic level i and $f_i(E)$ has the form:

$$f_i(E) = \begin{cases} \frac{s_i}{E_i} \left(\frac{E}{E_i} \right)^{s_i}, & \text{for } E > E_i; \\ 0, & \text{otherwise.} \end{cases} \quad (A.2)$$

Argon is represented by four atomic levels and methane by two. The values of s_i and E_i are given in Table A.1.

The complex index of refraction $\epsilon = \epsilon_1 + i\epsilon_2$ can be expressed in terms of the oscillator strength $f(E)$, the electron density ρ and ω_p , the plasma frequency⁶⁰:

$$\epsilon_2 = \frac{2\pi^2 \rho e^2}{m_r E} f(E) \quad (A.3)$$

$$\epsilon_1 = 1 + w_p^2 \int \frac{f(\omega')}{\omega'^2 - \omega^2} d\omega' \simeq 1 + \frac{\omega_p^2}{\omega^2} \int_0^\omega f(\omega') d\omega'$$

The equation for ϵ_2 is just the definition of the complex dielectric constant. The form of ϵ_1 is derived from the Kramers-Kronig dispersion relations. Substituting the above expressions for $f(E)$ and ϵ into Eq. 5.4 and normalizing to the density of electrons, we find:

$$\begin{aligned} \frac{d^2 N}{dx dE} = \frac{A}{\beta^2 E} \sum_i f_i(E) & \left[\ln \left(\frac{2m_r c^2 \beta^2 \gamma^2}{E |1 + \beta^2 \gamma^2 (1 - \epsilon(E))|} \right) + \beta^2 \epsilon_1(E) \right] \\ & + \frac{A}{\beta^2 E^2} \int_0^E \sum_i f_i(\mathcal{E}) d\mathcal{E} \end{aligned} \quad (A.4)$$

Atom	Level	w_i	E_i	s_i	E_{eff}
Ar	K	.111	3206.	2.75	5341.
Ar	L	.444	248.	2.29	490.
Ar	M	.133	52.	3.20	82.
Ar	M	.311	15.8	3.20	25.
CH ₄	-	.8	11.5	2.15	27.
CH ₄	K	.2	283.	2.52	506.

Table A.1: Oscillator Strengths for Argon and Methane. The values are from Ref. 80.

where $A = 2\pi e^4 \rho / m_e c^2$. We wish to integrate Eq. A.4 to find the mean number of collisions per unit sample thickness. To do so, we make the approximation that we can ignore the energy dependence within the logarithm and replace E with the average value of the ionization potential, defined by:

$$\ln E' = \int_0^\infty \ln(\mathcal{E}) f(\mathcal{E}) d\mathcal{E} \quad (A.5)$$

giving:

$$E'_i = E_i \epsilon_i^{-1/s_i}$$

Thus, integrating Eq. A.4 gives:

$$\begin{aligned} \frac{dN}{dx} = & \frac{A}{\beta^2} \sum_i \frac{w_i}{E_i} \frac{s_i - 1}{s_i} \left[\ln \left(\frac{2m_e c^2 \beta^2 \gamma^2}{E_i [1 + \beta^2 \gamma^2 (1 - \epsilon(E'_i))]} \right) + \beta^2 \epsilon_1(E'_i) \right] \\ & + \frac{A}{\beta^2} \sum_i \frac{w_i}{E_i} \frac{s_i - 1}{s_i} \end{aligned} \quad (A.6)$$

where the term on the first line gives the sum of the mean number of resonance type collisions per cm per atomic level, N_i^{res} and the term on the second line gives the sum of the mean number of Rutherford type scatters per cm per level, N_i^{Ruth} .

The next step in our calculation is to find the distribution of measured ionization per collision. Although we refer to the observed ionization in the TPC as dE/dx , this quantity is not strictly speaking an energy loss. What is truly measured is the number of electrons reaching the sense wire. To a moderately good approximation, this number is proportional to the energy deposited. We shall here derive a more accurate formula for the effective energy seen in an ionization detector.

When an atom is excited by collision with a fast moving charge, its response is a function of the energy transfer in the collision. In Table A.2 we summarize the most important decay modes for argon and methane. The electrons released by these processes then interact in the gas producing secondary ionization. The number of such secondaries can be approximated by the form⁶²:

$$N_s(k) = \frac{k - k_0}{W} \quad (A.7)$$

where k is the kinetic energy of the initial electron and the values of k_0 and W depend upon the composition of the gas. We have been unable to find an experimental measurement of k_0 for CH₄ and therefore have used the argon value of 16 eV⁸⁰. W is calculated from the fraction of methane in the gas using the formula:

$$\frac{1}{W} = \frac{1 - f_{CH_4}}{W_{Argon}} + \frac{f_{CH_4}}{W_{CH_4}} \quad (A.8)$$

The values of W for argon and methane are 26.4 eV. and 28 eV respectively⁸⁰.

We have defined the effective energy in atomic level i as the number of electrons released multiplied by the mean energy per electron:

$$(E_{eff})_i = \sum_{\substack{\text{primary electron} \\ \text{secondary electrons}}} \left(1 + \frac{k - k_0}{W}\right) W \quad (A.9)$$

The first term within the parenthesis corresponds to the primary electron produced and the next to the secondary electrons. The values of E_{eff} are summarized in Table A.1.

We now wish to combine Eq. A.6 and Eq. A.9 to determine the distribution of observed dE/dx per sample. The most straightforward approach is to use a simple Monte Carlo. The procedure is:

- Calculate separately for the resonance and the Rutherford regimes the mean number of collisions per atomic level per sample using Eq. A.6. For each sample, pick the number of collisions in each level using to a Poisson distribution.
- For collisions in the resonance region, make the approximation that the width of the distribution of energy loss per sample is narrow when compared to variations in deposited energy due to the width of the dN/dx distribution. In this case, the energy transfer in every collision can be set equal to is E_{eff} . (Note that one could instead pick the energy of each collision according to the distribution defined in Eq. A.2. We have chosen not to do so merely to decrease computer time.)
- For Rutherford type collisions, individually generate the energy deposited in each collision. (Since the number of such hard scatters is small, this is not very time consuming.) The collisions are generated according to a $1/E^2$ distribution that is cut off for energies below the ionization potential E_i for level i .
- The total value of E for the sample is then the sum of the energy deposited in Rutherford and resonance collisions for each atomic level.

If one is only interested in calculating the variation of measured dE/dx with $\beta\gamma$ rather than the distribution of dE/dx per sample, it is not necessary to use a Monte Carlo technique. We shall make the assumption that the 65% truncated mean scales with the most probable energy deposited. Since the most probable value for a Poisson distribution with mean N is $N - \frac{1}{2}$, we can express the most probable energy loss in resonance collisions as:

$$(MPE)_{\text{Resonance}} = \sum_{N_i^{R_{\text{res}}} > \frac{1}{2}} (E_{eff})_i (N_i^{R_{\text{res}}} - \frac{1}{2}) \quad (A.10)$$

The most probable value for the Rutherford scattering term was first derived by Landau and has been more accurately calculated by Maccabee and Papworth:⁸⁵

$$(MPE)_{\text{Rutherford}} = \frac{Ax}{\beta^2} \left[\ln \left(\frac{Ax}{\beta^2 \sum_i \frac{w_i}{E_i} \frac{x-1}{x_i}} \right) + 0.198 \right] \quad (A.11)$$

It is not in general possible to calculate the most probable value of a distribution from the most probable values of its constituents. If the width of this Landau term is large in comparison to the resonance term width, however, we can describe the most probable energy loss as the sum of Eq. A.10 and Eq. A.11. The Landau term has a fwhm:⁸⁵

$$(fwhm)_{\text{Rutherford}} = 4.02 \frac{Ax}{\beta^2} \quad (A.12)$$

Atom	Energy Transfer (eV.)	Probability	# e^-	Energy Released (eV.)
Ar	15.8-248	1	1	$k_1 = E - 15.8$
Ar	248-280	1	2	$k_1 = E - 248$ $k_2 = 200$
Ar	280-326	1.	2	$k_1 = E - 248$ $k_2 = 200$
		.3	2	$k_1 + k_2 = E - 265$ $k_3 = 200$
Ar	326-3206	.65	2	$k_1 = E - 248$ $k_2 = 200$
		.35	3	$k_1 = E - 326$ $k_2 = 200$ $k_3 = 50$
Ar	> 3206	.8	3	$k_1 = E - 3206$ $k_2 = 2670$ $k_3 = 200$
		.2	2	$k_1 = E - 2870$ $k_2 = 3206$ (photon)
CH ₄	11.5-283	1	1	$k_1 = E - 11.5$
CH ₄	> 283.	1	2	$k_1 = E - 283$ $k_2 = 243$

Table A.2: Ionization Decay Modes of Argon and Methane. The data for argon is from Ref. 64. The methane numbers are from Ref. 83 and Ref. 84.

Since the number of resonance collisions is large, we will approximate it as a gaussian distribution. The width therefore is:

$$(\text{fwhm})_{\text{resonance}} = \sqrt{\sum_i (E_{\text{eff}})_i^2 N_i^{R_{\text{res}}}} \quad (A.13)$$

For a path length of 5 mm and a $\beta\gamma$ of 4, the width of the Landau term is twice that of the resonance term. To the extent that this difference is large enough to make the assumption of adding most probable values reasonable, our full expression for the most probable energy loss as a function of $\beta\gamma$ becomes:

$$\begin{aligned} \text{MPE} = \frac{A}{\beta^2} \sum_i (E_{\text{eff}})_i, \max \left(\frac{w_i}{E_i} \frac{s_i - 1}{s_i} \left[\ln \left(\frac{2m_e c^2 \beta^2 \gamma^2}{E_i |1 + \beta^2 \gamma^2 (1 - \epsilon(E'_i))|} \right) + \beta^2 \epsilon_1(E'_i) - \frac{1}{2} \right], 0 \right) \\ + \frac{Ax}{\beta^2} \left[\ln \left(\frac{Ax}{\beta^2 \sum_i \frac{w_i}{E_i} \frac{s_i - 1}{s_i}} \right) + 0.198 \right] \end{aligned} \quad (A.14)$$

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