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CERENKOV RING IMAGING AND SPECTROSCOPY OF
CHARGED KSTAR CHARGED KSTAR
INTERACTIONS AT 11 GeV/c

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Abstract

A thesis in two independent halves.

Part I: A scientific proof-of-principle prototype was constructed and tested in order to evaluate the performance capabilities of a ring imaging Cerenkov detector. The detector measures the velocity of relativistic particles, and is the first Cerenkov counter to take advantage of the most fundamental properties inherent to Cerenkov radiation. Single photons from the Cerenkov cone of a radiating particle in a gas or liquid Cerenkov medium are focussed into a ring. Those Cerenkov photons with far ultraviolet wavelengths of the electromagnetic spectrum are converted to electrons, with high efficiency, in the uniform electrostatic field of an electron drift tube. The electrons are drift in the field to a novel proportional wire chamber which is sensitive to single electrons. Analysis of the reconstructed Cerenkov rings from an 11 GeV/c test beam show that an excellent velocity measurement can be obtained, allowing clean π/K particle identification at a level of ~ 15 standard deviations to be achieved.

The physics and technology of this new Cerenkov detector are discussed, including materials studies, construction techniques, and resolution measurements. Sources of resolution error are individually identified and measured where possible. The results of all studies indicate that the measurement resolution is understood. This work has led to the adoption of a large scale ring imaging detector as part of a new high energy physics spectrometer, the SLD, at the Stanford Linear Accelerator Center.

Part II: Results from an amplitude analysis of strange meson final states in $K^-\pi \rightarrow \bar{K}_S^0 \pi^-\pi^+$ interactions are presented. The data derive from a 4 event/nb exposure of the LASS (Large Aperture Superconducting Solenoid) spectrometer to an 11 GeV/c K^- beam. The data sample consists of $\sim 100,000$ events distributed over the Dalitz plot of the channel.

The process is observed to be dominated by the production and decay of natural spin-parity ($J^P = 1^-, 2^+, 3^-, \dots$) strange meson states. The data can be understood in terms of a simple model in which the resonant $\bar{K}^{*-}$ are produced predominantly via natural parity exchange in the t channel. The leading $K^*(890)$, $K_2^*(1430)$, and $K^*(1780)$ resonances are clearly observed and measured, and the underlying spec-

troscopy is also extracted. Indications of higher mass resonance production are also shown.

The observed properties of these states are used to confront current models of quark spectroscopy in strange meson systems.

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all along. My parents have been solidly supportive through these trying years, and that is in addition to sending the occasional check. David and Jennifer, my brother and sister, have been very close to me. I am fortunate to know two such fun, intelligent, and compassionate people.

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Dedicated to my Mother and Father

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I. Introduction to Cerenkov Ring Imaging

The exploration and discovery of natural phenomena has always motivated the work of scientists, authors, and philosophers. Such discoveries have been achieved through the efforts of the men and women of history to learn, invent, study, and create new tools and techniques with which to probe the unknown. In the specific case of particle physics, the unknown has required that new machines, the particle accelerators, be created which enable the study of ever higher energy regimes, and hence ever smaller distances. These machines accelerate charged particles of a variety of types, e.g., electrons and protons, to interact with other matter to produce new interesting physical systems. In parallel with the accelerator development, the detectors which enable the scientist to study these new systems have also evolved.

With the promise of the rich physics expected from the decay of the Z_0 produced in e^+e^- collisions from new accelerators, the Linear Collider at Stanford and the Large Electron Positron project at CERN, new detectors are being constructed to permit detailed studies in this frontier energy regime. The center-of-mass energy of ≈ 100 GeV, will allow scientists to make systematic studies at a mass scale which was hitherto unknown. This new energy scale will call for new detection techniques, one of which is a novel device for identifying the type, or flavor, of charged particles as they fly away from the interaction point in the e^+e^- collisions. This device, the Cerenkov Ring Imaging Detector, pushes to new levels the application of Cerenkov radiation as a tool for doing high energy physics.

The Cerenkov effect was first studied by P.A. Cerenkov in 1937, where he observed two peaks in the angular distribution of radiation produced by the beta decay of radium B in carbon disulphate. Analysis showed that these peaks corresponded

to two different velocities of electrons in the emission^[1]. It was at that time that he first discussed the possibility of using the effect to determine the velocities of particles. It was ten years later before Cerenkov counters were constructed as practical devices.

A Cerenkov detector is designed to measure velocities in a specified momentum range, using Cerenkov radiators with refractive indices chosen so that the Cerenkov angle varies with the velocity, from threshold to the highest anticipated momentum. To date, most Cerenkov detectors have not been used in this way. Traditionally, Cerenkov devices are used as threshold devices, where the simple detection or absence of light from a track in the radiator, constituted the signal. While this method has proven useful in a wide range of experimental applications, it makes poor use of the information intrinsically available from the Cerenkov radiation. The geometrical properties of the Cerenkov cone, could they be taken advantage of, would enormously increase the separation power of a device. The difficulty of making large scale devices which use the full Cerenkov information, has been the reason that threshold counters have been so successful for so long in high energy physics.

Cerenkov counters are inherently delicate instruments, in that they must be sensitive to single photon signals in the most extreme cases, often requiring expensive and difficult photomultiplier systems which have sometimes proven to be less robust than would be desired, in that signal-to-noise problems are often severe. One physicist was overheard to have quipped, "...particle identification is the largest waste of good manpower in high energy science today...", which reflects upon the frustration that Cerenkov counters have caused in the past. Also, costs have been too prohibitive to allow good Cerenkov coverage in an experiment, so that the segmentation is often coarse, which leads to problems in high multiplicity environments.

The ring imaging idea was originally conceived by Roberts^[2], though never adopted on a large scale in that form. The contemporary excitement for the technology was generated by Ypsilantis and Seguinot^[2-4]. They recognized that in this day of 4π spectrometers, any particle identification device must be uniformly sensitive over the full solid angle, with excellent segmentation, high quantum efficiency, and good time resolution. They first studied the potential photocathodes, the requisite chamber gases, and introduced the essential long drift technology. In

collaboration with others^[7], suitable electron counters were soon developed which enabled efficient readout of the Cerenkov devices. Soon, other teams worldwide became involved^[8,9], with the results leading to the construction of the SLD and LEP detectors.

This thesis will detail the scientific principles, construction details, and the results of tests made with a prototype Cerenkov ring imaging detector, promising to provide particle identification for experiments at new levels. In the ring imaging device, the Cerenkov photons are detected and the power of imaging optics is used to actually take 'photographs' of the Cerenkov light. The details of the reconstructed image have simple interpretations in terms of the particle velocities and give unprecedented particle identification performance.

Chapter II introduces the ring imaging concept, motivating the design essentials which the physics of radiating relativistic particles imposes upon such a detector. Some of the associated physical pitfalls are also introduced. Chapter III is essentially a recipe for constructing a ring imaging counter, taking advantage of the limited opportunities which Mother Nature has provided in this far ultraviolet corner of the electromagnetic spectrum. Chapter IV details some new and exciting work using carbon monofilament fibers as the sensitive elements for the very heart of the counter, while Chapter V reports upon the results obtained with the prototype ring imaging counter which was constructed as a result of the understanding of the science described in the previous chapters. Chapter VI will conclude this ring imaging portion of the complete thesis, summarizing the convincing evidence presented in favor of the conclusion that the ring imaging detector is clearly a new and important tool for the high energy physicist.

II. Principles of Cerenkov Ring Imaging Detectors

II.1 Properties of Cerenkov Radiation

Consider a relativistic particle of mass m moving with velocity β in a medium with index of refraction n . The fundamental relationship for the angle of emission of a Cerenkov photon from the particle is given by

$$\cos\theta_c = \frac{1}{n\beta} \quad (\text{II.1})$$

Since $|\cos\theta_c| \leq 1$, there is no Cerenkov emission until $\beta_1 = n^{-1}$, where β_1 is the threshold velocity in units of the speed of light. The complete nature of the spectrum of Cerenkov light was understood theoretically with the work of Frank and Tamm^[10]. Introducing the index of refraction into Maxwell's equations led them to conclude that Cerenkov radiation was essentially an electromagnetic analog of a sonic shock wave. In effect, the charged particle is moving with a velocity greater than the phase velocity of light in the medium. The intense electric field polarizes the medium, which relaxes back to equilibrium by emission of an electromagnetic wave. Their treatment showed that the effect leads to production of a coherent wavefront, in addition to many other details and interesting predictions which were amenable to direct tests. In particular, the energy radiated per unit frequency interval per unit path length is given by^[11]

$$\frac{d^2E}{dx d\omega} = \frac{e^2 \omega}{c^2} \left(1 - \frac{1}{\beta^2 n^2(\omega)} \right) \quad (\text{II.2})$$

where n is a function of the frequency of radiation, i.e., there can be chromatic dispersion. When integrated over the radiating path length L ,

$$\frac{dE}{d\omega} = \frac{L\epsilon^2\omega}{c^2} \left(1 - \frac{1}{\beta^2 n^2(\omega)}\right) \quad (11.3)$$

From this, the number of individual Cerenkov quanta emitted can be calculated. The total energy emitted for N radiated photons is

$$E = \sum_{i=1}^N \hbar\omega_i = N\hbar(\omega), \quad (11.4)$$

so that

$$N = \frac{L\alpha_e}{c} \int \left(1 - \frac{1}{\beta^2 n^2(\omega)}\right) d\omega \quad (11.5)$$

or, perhaps in a more illuminating fashion,

$$N = 2\pi L\alpha_e \int_{\beta n > 1} \left(1 - \left(\frac{\beta n(\lambda)}{\beta}\right)^2\right) \frac{d\lambda}{\lambda^2} \quad (11.6)$$

where α_e is the fine structure constant. The total number of photons emitted depends upon the wavelength integration, but in general, as (11.5) indicates, the number of photons emitted per unit length per unit energy is a constant, when there are no absorption bands close to the interesting frequency range. This constant is just the mean of a Poisson distribution in the number of photons, because it is an inherently statistical process.

When there are absorption bands nearby, the calculation becomes more tedious. There are several models, dating from the late 19th century, which approximate the problem of anomalous dispersion. If the first absorption band is the dominant one for a given radiator, the Sellmeyer equation

$$n^2(\omega) = 1 + \frac{A}{1 - (\omega/\omega_0)^2} \quad (11.7)$$

is a good approximation, where $2\pi c/\omega_0$ is the wavelength of the first absorption line. If the absorption is very strong, more complicated dispersion relations must

be used. These equations are important here, because they can be used to estimate the amount of chromatic dispersion expected from a Cerenkov radiator, the limiting source of error in practice.

11.2 Cerenkov Detectors

This calculation for the ideal case must be modified to account for the inefficiencies of any practical detector. Because the detection efficiency, $\epsilon(\omega)$, is a function of frequency,

$$N = \frac{L\alpha_e}{c} \int \epsilon(\omega) \left(1 - \frac{1}{\beta^2 n^2(\omega)}\right) d\omega \quad (11.8)$$

where N is now the number of photons detected, and $0 \leq \epsilon(\omega) \leq 1$ for all ω . If a device is designed to detect photons in a spectral region far from where the radiating medium has absorption bands, n can be taken to be independent of frequency and the traditional equation used to describe Cerenkov counters results,

$$N = N_0 L \sin^2 \theta_c \quad (11.9)$$

where

$$N_0 = \frac{\alpha_e}{c} \int \epsilon(\omega) d\omega. \quad (11.10)$$

The quality factor, N_0 , is the figure of merit for the performance of a given Cerenkov counter. The larger N_0 , the better the device. It is important to note here that for the Cerenkov ring imaging device, the Cerenkov light is in the far ultraviolet where the refractive index of any suitable UV transparent radiator is very frequency dependent, i.e., all radiators are substantially dispersive. The full integral relationship should be used to evaluate N_0 .

Cerenkov detectors are used in high energy physics to distinguish particle types. A particle type is distinguished by its mass, and from simple relativistic kinematics, that mass m is related to β through the momentum,

$$|\vec{p}| = m\gamma|\vec{\beta}| \quad (11.11)$$

in units where $c=1$. This means that β and p must be known before m can be determined. p is determined by measuring the trajectory of the particle as it moves in a magnetic field. This is typically determined to a higher precision than other measurements and will be assumed to be so here. The Cerenkov counter is used to find β through (11.1). At issue here is the accuracy to which m can be determined, and more importantly, the degree to which different masses can be distinguished. The accuracy is a function of the Cerenkov angle and the number of detected photons, so that to the extent that the momentum measurement is accurate, the errors in the Cerenkov angle measurement need to be understood and controlled to the highest degree in order to optimize the separation potential of a Cerenkov device.

Most of the Cerenkov light is emitted in the far ultraviolet region of the electromagnetic spectrum, as indicated by (11.6), where the λ^{-2} behavior is explicitly given. It is the nature of work in the ultraviolet that there are only a handful of substances sufficiently transparent to permit good Cerenkov measurements. Low levels of contamination in the form of more abundant materials like water and oxygen, which absorb strongly in the UV, can frustrate the collection of the meager amount of light radiated. This means that extraordinary care must be taken to insure that Cerenkov radiator volumes are pure at the parts per million level, which can become a formidable task over large volumes. Just the nature of Cerenkov light can make it hard to correlate a signal with a mass, in that the Poisson distribution of the number of photons for a particle just above threshold leads to ambiguities in the identification. An even more severe constraint is the lack of sensitive photocathodes in the far ultraviolet. Typically, to obtain quantum efficiencies of just 19% requires the use of difficult materials such as CsI and RbTe. Not only must extreme care be taken with the preparation, handling, and maintenance of such materials, the consequent expense makes them unsuitable for large surface area detectors. It is in fact the recent development of highly efficient ultraviolet photocathodes which makes Cerenkov ring imaging possible. Nonetheless, Cerenkov counters have taken many shapes and are now considered a bread-and-butter tool for the high energy particle experimentalist. The usefulness of the technique could be enormously increased if the shortcomings could be overcome; improvements that the Cerenkov ring imaging technology will make practical.

The Cerenkov ring imaging technique evolved from considerations of the simple

optical system shown in Figure 11.1. A relativistic particle is normally incident on a spherical mirror of radius R . The focal length f of the mirror is $f = \frac{1}{2}R$. The mirror is immersed in a volume of radiator, a gas or liquid, with index n , so that the emitted Cerenkov radiation is focussed to a ring of radius

$$r = f \tan \theta_c \quad (11.12)$$

at f . A photosensitive detector at f converts the Cerenkov light into signals which are electronically sensed and recorded, preserving the image information. By accurately determining r , θ_c and hence β can be calculated. This is the ring imaging technique in its simplest form. In principle, the mirror need not be spherical, and in general, the incident particle will not be normally incident at the center of the mirror. These conditions give rise to 'rings' which are not necessarily circles, but nonetheless form unique shapes in the focal plane. While these shapes may have complicated descriptions, they can be analyzed, and so do not affect the power of the method in principle.

Figure 11.1: A simple ring imaging system

In practice, the images do not form smooth curves in the focal plane. Because the Cerenkov radiators have finite lengths, the number of photons emitted

is limited, so that the Cerenkov angle measurement becomes an average of several independent angle measurements of the few points which comprise the image. Due to the Poisson nature of the photon emission, there are fluctuations in the number of points detected in the image. The higher the number of detected photons the better the measurement, so that the mean $\bar{n}$ of the Poisson distribution for $\beta=1$ particles should be as large as possible. It is important to understand and optimize all the contributions to N_0 .

The Cerenkov radiators used in this imaging system are in the gas phase. A typical gas radiator will have an index of refraction of the order of 1.0016, which implies $\beta_{th} \approx 0.9984$. These radiators work well for separating very relativistic particles, where $\beta \approx 1$. For particle separation at slower velocities, an alternative technique has been developed.

11.3 Proximity Focusing Detectors

The basic ring imaging system just discussed employs a gas radiator as the Cerenkov medium. Gases work well for relatively high momentum tracks, where, since $(n-1)$ is small, β is close to 1. For slower particles, the same ring imaging technology can in principle be applied, with the use of the proximity focussing technique. The basic detector is schematically shown in Figure 11.2. A relativistic particle is normally incident (for simplicity) upon a liquid Cerenkov radiator. Such liquids typically have indices on the order of $n \approx 1.3$, so that $\beta_c \approx 0.7$. The Cerenkov angle in such radiators is substantial, so that the photons are quickly well separated from the track. The liquid must transmit ultraviolet, which severely limits the choices of radiator, so that the photons reach the boundary window and escape into the surrounding environment. The environment is a low index transparent gas, so the photons are severely refracted as they exit the boundary window. As in the gaseous radiator case, the photons pass into a photon detector, where they are converted into signal electrons. Because of the large refraction, the circle of Cerenkov photons from the radiator make a large circle of signal electrons in the photodetector, if the spacing between the radiator and the detector is large. The distance is chosen to be long enough so that the error on the radius of the circle of signal electrons is small in comparison to the mean radius. An additional contribution to the error not

present in the gas case is due to the finite thickness of the radiator. It contributes a "flat-top" error, proportional to that thickness.

Figure 11.2: The proximity focussing technique. The vertical line segment at the bottom of the figure is the particle trajectory through the liquid radiator volume. The drift volume of the photodetector is at the top. The band represents the allowed photon trajectories, for a specific radiator.

The resolving power of such a device does not approach that of the gaseous devices, but it is very useful nonetheless for it enables good particle identification in a momentum range where gaseous radiators are insensitive. Also, it is easily installed into the same detector volume with the gaseous device, so that the two integrated, but complementary, instruments cover most the available phase space of interesting particle identification.

11.4 The Photon Detector

Once the photons have been generated, they must be detected and measured. A schematic drawing of the photodetector which will be discussed here is given in

Figure II.3. The detector volume proper in the sketch, with the boundary windows in the planes normal to the photons, is embedded in an electrostatic cage formed of thin wires, shown as the uniformly spaced lines in Figure II.3. The detector is sensitive to ultraviolet photons from in front of or behind it, as indicated by the rays in the figure. The photons are converted to single electrons, which must be transported in the gas-filled volume to an electron counter at one end. The electrons are constrained to move along the electric field lines which are defined by a plane of high voltage at the opposite end of the detector, and maintained by the field cage wires. The cage is on the order of one-meter in length. From the drift properties of the detector gas, as will be described, the electric field must be kept maintained at about 400 V/cm , implying a maximum potential at the high voltage cathode of about 40 kV .

Figure II.3: A simplified schematic of the photodetector. Incident Cerenkov photons must make it into the drift volume and to the electron counter, preserving all of the information about their precise point of photoconversion.

The field wires encircle the volume, on the inside and outside, and are constrained to maintain a uniform potential drop, i.e., electric field, along the volume length. The two sets of exterior and interior wires are essential, as they conspire to prevent sources of field lines exterior to the volume (or box) from distorting the highly uniform field inside. Also, and in the same sense, the wires are kept in intimate contact with the windows, to prevent any polarization or ion build-up from occurring.

Inside the drift box, a vapor of gaseous photocathode is carried in a UV transparent base gas. The photocathode is present at a relatively low level, consistent with its vapor pressure. Photons which enter the volume through the UV transparent face of the drift box are absorbed by the photocathode, which may respond by releasing a bound electron. The quantum efficiency of the photocathode used is high, at least $1/3$ of the incident photons are so converted. The photoelectrons drift to the electron detector under the action of the electrostatic field set up by the equipotential wires. At the photodetector, the electrons are electrostatically guided to a sensing plane, a plane of anode wires, the sense wires, which detect the arrival of single electrons. The signal is electronically amplified, and converted into the coordinates of the point of photoconversion, including the effects of any measurement errors. It is important that the acceptance for drifting electrons in any coordinate be uniform and give a unique measurement of the initial position of the incident photoelectron.

II.5 The Efficiency of Cerenkov Photon Detection

The efficiency of photon detection for the ring imaging counter is a the direct product of the independent component efficiencies. To understand the effects of the various efficiencies it is helpful to examine their relative contributions for the idealized example of Figure II.4. In the figure, there is the proximity focussing system consisting of the liquid radiator, a boundary window, a gas-filled gap, another boundary window, and the photodetector, all in front of the gas system and first encountered by a relativistic particle. It is discussed later. The gas system consists of the gas volume, the mirror, a boundary window, and the photodetector. The photodetector is divided into two regions, the drift or photoconversion volume, and the electron counter. The Cerenkov photons can be lost due to absorption in the

radiator gas, due to absorption or scattering at the reflecting surface of the mirror, due to absorption (without photoelectric conversion) in the photodetector, or due to scattering or absorption in the boundary windows. If the photons are photoconverted, then the signal electrons are subject to losses due to recombination in the photodetector, due to bad electrostatics in the photodetector, or due to inefficiencies in the electron counter. Briefly then, the efficiency, $\epsilon(\omega)$, can be succinctly written as

$$\epsilon(\omega) = \frac{\alpha_e}{c} \int T_{\text{rad}}(\omega) R_{\text{mir}}(\omega) T_{\text{win}}(\omega) Q E_{\text{pc}}(\omega) \lambda_d \epsilon_{\text{det}} d\omega \quad (11.13)$$

where

Figure 11.4: Schematic of a prototype ring imaging detector. The proximity focusing system is placed before the gaseous radiator system.

T_{rad} - is the probability of photon transmission from the radiator volume to the detector,

R_{mir} - is the reflectivity of the mirror,

T_{win} - is the transparency of the detector window,

$Q E_{\text{pc}}$ - is the quantum efficiency of the photocathode for conversion of photons to drifting electrons,

λ_d - is the probability for a drifting electron to reach the detector,

ϵ_{det} - is the probability for an electron to be counted.

Each of these probabilities can be determined by independent tests. The composite of the results gives N_0 , effectively the acceptance of the ring imaging counter for photons from the full Cerenkov spectrum. The results of such tests and the determination of N_0 for a particular ring imaging counter design are described in subsequent chapters.

11.6 The Resolution of a Ring Imaging Detector

It is not enough to ask that there be a large number of photons per measurement of the Cerenkov angle, it is important that each individual photon be well measured. There are several factors contributing to the ultimate resolution in a ring imaging counter, and just as for the N_0 estimate, they can be separately identified and studied. The expected sources of error are:

- Chromatic - the frequency dependence of the index of refraction,
- Measurement - the coordinate measurement errors in all dimensions, including those due to distortion and segmentation.
- Diffusion - the random thermal motion of the drifting electrons,
- Optical - mirror aberrations and focal astigmatism,
- Tracking - momentum measurement and multiple scattering

Some of these contributions can be directly measured and others are extrapolated by indirect measurements. The effects of the errors can be calculated, since the contributions to the total Cerenkov angle are related in known ways.

From (11.1), the error on the velocity is related to the index and measured Cerenkov angle by

$$\frac{\delta\beta}{\beta} = -\frac{\delta n}{n} + \tan\theta_c \delta\theta_c. \quad (11.14)$$

Use (II.12) to write it in terms of experimental parameters as

$$\frac{\delta\beta}{\beta} = -\frac{\delta n}{n} + \frac{(\frac{r}{f})^2}{1 + (\frac{r}{f})^2} \left(\frac{\delta r}{r} - \frac{\delta f}{f} \right). \quad (\text{II.15})$$

As discussed previously, the optical errors are small, for a very good mirror, so that for this example, $\delta f = 0$. Also, the Cerenkov angles are small, but the focal lengths of the mirrors are long, reflecting the lengthy gas radiator volumes, so that $(\frac{r}{f})^2 \ll 1$. Considering the most illustrative case of a highly relativistic particle, $\beta \approx 1$. Also, the index for the gaseous radiators is very close to one. Using these approximations in (II.15) gives

$$\sigma_\beta^2 = \sigma_n^2 + \left(\frac{r}{f}\right)^4 \left(\frac{\sigma_r}{r}\right)^2. \quad (\text{II.16})$$

The point is that the resolution on the gas rings is dominated by chromatic and measurement errors. The simplifying approximations used above are only approximately true in a practical counter, where the geometry of the counter and the momentum distribution of the tracks gives rise to a variety of circumstances where other errors cannot be ignored. For example, when the momentum of the radiating particle is low, (II.11) implies that the momentum error can become comparable to the other errors, due to multiple scattering. Also, for off-axis, oblique Cerenkov light rays, the optics of the spherical mirror are such that the focal length can vary appreciably over the ring image, i.e., astigmatic focussing errors can become significant.

There is another source of error which will degrade significantly the low momentum resolution properties of practical counters in the colliding beam spectrometers. The magnetic field required for tracking in the spectrometers, means that all tracks are bent as they traverse the Cerenkov radiators. Low energy electrons, in particular, will suffer severe bending, so that the Cerenkov light emitted from the track will be smeared across the face of the gas radiator mirror, smearing the ring image. This contribution to the measurement error goes like

$$\sigma_{mag} = \frac{30}{p} \text{ mrad} \quad (\text{II.17})$$

and in fact dominates the error at low momenta. Unfortunately, this parameter could not be studied here, and so will not be discussed further. However, as it is a well understood phenomena, independent of construction technique and materials issues that a prototype device is built to address, the loss to the prototype studies was felt to be small.

II.6.a Chromatic Errors

Given $n(\omega)$, the chromatic error can be estimated. Unfortunately, it is difficult, and expensive, to accurately determine the frequency dependence of n in the far ultraviolet and the instrument for making the exact measurement was not available for these studies. This measurement requires a UV refractometer, a device that measures the refraction of incident UV rays as a function of wavelength and angle in a given material. However, measurements can be found in the literature for any of the radiators that are interesting here. From a measurement of the Cerenkov angle from beam tests, the average index can be determined, which serves as a consistency check on the assumed values of n . The relevant error criterion is just the average of the chromatic dispersion over the Cerenkov photon acceptance of a specific counter design.

The measurements which do exist indicate a variation in $n-1$ of about 1 in 10 for the interesting gas radiators, and about 1 in 27 for the liquid radiators. In that case, the derivative of (II.1) gives

$$\sigma_c = \frac{1}{n\sqrt{n^2 - 1}} \sigma_n \quad (\text{II.18})$$

as an estimate of the chromatic error, when averaged over the photon acceptance of the counter. For the gas radiator used here, isobutane, $\sigma_c^{gas} \approx 1.2 \text{ mrad}$ per photon, and for the liquid, $\sigma_c^{liq} \approx 3.1 \text{ mrad}$ per photon, at a momentum of 5 GeV/c.

II.6.b Measurement Errors

The measurement errors are the errors incurred in the reconstruction of the Cerenkov angle due to the electrical, mechanical, and thermodynamical hardware of which the Cerenkov counter is made. These are the errors over which the experimenter has the most control. The largest range of potential error sources fall

under this heading. All contributions to the error due to relative misalignments in the system components, due to manufacturing imprecision, and due to electrical or electronic problems are included here. The greatest impact here is in the design and construction of the photodetector.

The measurement errors affect the three coordinate measurements in distinct ways. One coordinate is determined by the absolute position of a sense wire. An electron incident upon the electron counter will avalanche in the neighborhood of a single sense wire, i.e. ionize many gas atoms and liberate many additional signal electrons by the Townsend multiplication process in a locally intense electric field. This in contrast to conventional wire chambers where the signal is distributed over several wires. The electrical detection of the avalanche generates a coordinate given by the location of the sense wire. This coordinate is uniformly sampled over the breadth of the volume due to the constant spacing of the sense wires. As such, the error introduced by this spacing is just the "flat-top" error defined as $w/12$, where w is the wire spacing. Measurements indicate that the error on this coordinate is consistent with being due to the integral wire spacing. σ_x represents the error in this coordinate, where

$$\sigma_x^2 = w^2/12. \quad (11.19)$$

A second coordinate is determined as a function of the time between the passage of the charged particle and the arrival of a sensed photoelectron, the electron drift time. The photoelectrons will have drifted in the gas-filled box for distances of as much as one meter or more. In order to reconstruct where the photoelectron originated, the time-of-arrival relative to the particle passage must be known to good precision, at least as good as the wire spacing measurement if this error is not to dominate. The proper measurement of this coordinate requires the establishment a relationship between time-of-arrival and initial position, that is, requires that the drift velocity of the electrons in the drift volume be known to high accuracy. The drift velocity depends upon the composition of the gaseous environment, the local electric field, the pressure and temperature. The drift field must be uniform over the drift volume, and if not, a precise drift-map of the volume must be determined. Non-linearities in either the time or space measurements will cause distortions in the drift path, which will confuse the direct relationship between position and time-of-arrival at the detector.

Measurements of the detected electron position in this coordinate were made both as a function of initial position in this coordinate and in the transverse wire coordinate. Similar measurements taken by varying the electric field and gas composition were shown to be consistent with other measurements of the drift velocity in the same gases. Once understood, the drift velocity serves as a calibration for the system, to monitor the stability of the CRID gas environment. This allows any variations in the data to be corrected. The drift time is measured electronically by dividing the drift box in discrete time slices. The slices had a fixed width, on the order of tens of nanoseconds, so that a simple count of the number of slices between a detected avalanche on a sense wire and an event trigger gave the measurement of initial photoconversion in this coordinate, when the drift velocity is known. The error in this coordinate, σ_y , is just then the statistical error from the width of the time slice interval, and the systematic error from the drift velocity measurement, so that

$$\sigma_y^2 = \sigma_{v_d}^2 t^2 + v_d^2 \sigma_t^2. \quad (11.20)$$

In practice, this error is small enough to be an insignificant contributor to the overall error.

The third coordinate measurement needed is that along the length of the sense wire, that is, in the coordinate orthogonal to the wire position and the drift path. It also depends upon the electron avalanche to generate the signal. The technique developed here depends upon the charge division of the induced signal between the two ends of the sense wires. The requisite precision, compatible with the measurements on the other two coordinates, might be obtained by measuring the induced signal on the cathode of the sense chamber, and in fact this method has been adopted by others [12]. However, Monte Carlo studies indicated that a substantial signal loss due to dead time was inherent to that method, in addition to the fact that the already severe signal-to-noise requirements on the anode wires would be exacerbated for such a cathode chamber. Given that a suitable technology could be found, the charge division technique was the most appealing.

When an avalanche occurs along a resistive wire, the current in the wire is a function of the position of the avalanche along the wire. This current pulse may be integrated, and the resulting charge measurement can be simply related to the position of the avalanche. In the practical device, the resolution required is on the order of 1% of the wire length. A 1% percent charge division accuracy is common

in high energy physics experiments where the wire length is on the order of meters, not centimeters. This means that conventional sense wire materials will not offer the necessary performance to the CRID in a charge division mode. A thorough study of the problem resulted in the successful implementation of a new and precise charge division technique, the details of which are describe in a chapter III. The error in this coordinate, σ_x , then is

$$\sigma_x^2 = \sigma_{cd}^2. \quad (11.21)$$

11.6.e Diffusion Effects

The diffusion of the photoelectrons during the drift contributes an independent gaussian smearing to the error in each coordinate. While the amount of diffusion error depends upon the coordinate, whether along or transverse to the drift direction, the value in either case is amenable to direct measurement. Diffusion is a function of the time that an electron spends in the drift volume, and in principle, also a function of the local magnetic field strength. When the magnetic field is perpendicular to the drift, the Lorentz force damps the diffusion. The diffusion is given by

$$\sigma_d = \sigma_0 L, \quad (11.22)$$

where L is the drift length, $L = v_d t$, and σ_0 the diffusion coefficient in units of $\mu/\sqrt{\text{cm}}$. A careful study of the correlation between the known initial positions of some test electrons and the detected position of the same electrons provides information about the diffusion error. A method was developed which employed an ultraviolet laser and optics to produce the distributions from which the diffusion errors were determined, which were on the order of $\sigma_0 = 200 \mu/\sqrt{\text{cm}}$.

11.6.d Optical Errors

The optical error contribution to the resolution is also amenable to direct measurement. The focal properties and surface quality of mirrors which are suitable for ring imaging work are not very demanding in the technical sense. Any astronomical telescope mirror of the right focal length will be far more accurate than the Cerenkov detector requires. Nonetheless, the mirrors used must be understood so

that proper account of any focussing mistakes may be taken. Also important, the relative alignment of the mirrors and drift box. Errors from either source implies that photons will be displaced from their predicted locations in the focal plane of the mirror. Straightforward geometrical optics can be used to predict the resolution contribution from these problems.

11.6.e Momentum Effects

The tracking errors are due to multiple scattering of the incident particle and divergence in the tracks. The amount of multiple scattering is calculated from a knowledge of the amount and composition of the total material traversed by the radiating particle. The multiple scattering errors go like $\frac{1}{\sqrt{L}}$, and are small in the gas ring case, but dominate the resolution of the liquid rings. There can be calamitous results in the instance of severe kinking in the material between any tracking devices and the ring imaging counter. The beam divergence error, which is relevant for the test beam results presented here, is determined from the track angles as found by fitting to track data from a system of tracking chambers which in the beam line near the ring imaging detector.

11.6.f Total Resolution

The total resolution measurement can be written as the quadrature sum of all of the independent errors, neglecting the magnetic effects as they are irrelevant to these tests. Since there are N independent measured points on the ring, the Cerenkov angle error is improved by $N^{-1/2}$ over the individual point error. The actual error for the full Cerenkov angle measurement then is

$$\sigma_c = \sqrt{\frac{1}{N} \sum_1^N (\langle \sigma_{\text{chrom}} \rangle^2 + \langle \sigma_{\text{meas}} \rangle^2 + \langle \sigma_{\text{diff}} \rangle^2 + \langle \sigma_{\text{optic}} \rangle^2 + \langle \sigma_{\text{track}} \rangle^2)} \quad (11.23)$$

The resolving power of the ring imaging detector, in relation to other Cerenkov devices, is due to the fact that there are several measurements of the same Cerenkov angle per track. The best resolution will require that the errors from the individual contributors to σ_c be kept small.

11.7 Backgrounds

In a full scale high energy physics application, the Cerenkov ring imaging technique will become an exercise in hypothesis testing, where information from the tracking will be used to help distinguish between e, μ, π, K and p guesses. From a small distribution of roughly correlated hits in the photodetector, which appear as pixels to the analysis, a conclusion must be drawn about the velocity and hence the mass of the track which produced the hits. As the number of hits per track is on the order of 10 for a typical track, it is very important that there be few additional hits which would be presented to the analysis. A goal for a useful detector then would be the virtual elimination of sources of background hits. Three distinct sources of background hit generation were identified at an early stage in the evolution of ring imaging counters.

First, and the source of most concern, is the track itself. In this era of 4 π detectors, the radiating tracks necessarily pass through the photodetector. As they pass through, the ionization trail that they leave as a result of dE/dx loss, gives rise to backgrounds. The gases used in the photodetector are heavy, so that the amount of ionization per track can be as great as 1000 times as large as the single electrons from the Cerenkov photons.

Second, it is common detector lore that there is a simultaneous emission of UV radiation from an avalanche around a wire. The precise nature of the radiation depends upon the quantum mechanics of the gases involved, so that the spectral characteristics of the emission can be very complex. Because the photocathode used in the ring imaging detectors is very sensitive to UV radiation, the very beauty of it, any photoconversions of these regenerated UV photons, the "photon feedback", will make photoelectrons which then become signals in the electron detector.

Third, as might be garnered from the brief words about dE/dx , the dynamic range of signals that the counter must process is on the order of 1000. This huge number places severe constraints on the electronics which are used to sense the wire chamber signals. Large pulses can generate crosstalk over many wires unless special precautions are taken to prevent the problem. The crosstalk can give rise to false pulses which are subsequently recorded as hits. The investigation and treatment of these sources of background are essential, prior to the practical implementation of

a working ring imaging counter.

11.7.1 The dE/dx Problem

The problems with the ionization trails from tracks are three-fold. First, the majority of the time a track will not pass through a photodetector at right angles where the ionization left would be at a minimum. As a result, the track will spread itself out over a longer drift space and over several wires. In addition, diffusion of the charges will increase the width of the charge distribution which arrives at the electron detector. The large molecules of the detector gas then implies that a large amount of ionization will give rise to hits over several time slices, expanding over several wires. This cannot be completely avoided in the current schemes. Possible alternatives are to invent ways to decrease the intensity of the ionization, by decreasing the density of the gas or by choosing alternative lighter gases as the photodetector gas, so as to decrease the amount of generated dE/dx . However, the alternative gas must meet all of the other criteria of a good electron counting gas, and to date, no good candidate exists.

Also, when the dE/dx ionization arrives at the electron detector, it will undergo avalanche multiplication, and so generate feedback photons. As there is much more initial charge per avalanche than for the Cerenkov photoelectrons, the amount of UV emission will be proportionately larger. In that sense, the feedback problem from the track is the most severe source of extra "salt-and-pepper" background. Feedback is discussed in more detail in the next section. Also, the larger pulse height generated in the avalanche of the dE/dx implies that the crosstalk problem is most severe for the track.

11.7.2 The Photon Feedback Issue

Photon feedback is a particularly troublesome phenomena in the Cerenkov ring imaging environment. It places severe constraints on the design and implementation of a counter. The nature of the regeneration is clear, as discussed briefly above. The generation of UV light in an avalanche is an interesting physical phenomena in its own right. There have been proposals for detectors which use this light as a means to detect an avalanche, and in particular, as a possible means of gathering third coordinate information^[13]. The amount and spectrum of the light is very dependent

on the local environment of the avalanche. The very strong non-uniform electric fields which define the avalanche region, coupled with the dynamics of the gases, means that not only is UV emission a function of the atomic quantum mechanics of the constituents, but also a property of the molecular interactions which may take place during an avalanche. A complete description of the UV generation in an avalanche environment will not be offered here, because it is peripheral to the immediate problem. It is important to note that it exists and that in the context of a viable ring-imaging system, must be controlled.

The choice of a detector gas which is suitable for ring imaging work is limited to the saturated hydrocarbons. Unfortunately, these gases radiate strongly in the UV, and to date the only pragmatic to deal with this feedback problem is the straightforward mechanical one. A practical chamber design must optically isolate the avalanche from the photodetector proper. This is achieved by so-called "blinding" schemes. Figure 11.5 provides a simple illustration. The lines in the figure are legitimate electron drift paths as defined by the electrostatics of the photodetector. The lines end on sense wires, where charge multiplication occurs, and feedback UV generated. Photons which are emitted in the direction of the photodetector proper from the avalanche are optically shadowed by the blinding grids shown, the thick, round wires. The blinding grids cannot be randomly inserted into the chamber, however. The electrostatics of any blinding structure must insure that all of the true Cerenkov photoelectrons reach a sense wire. Otherwise, precious Cerenkov signal is lost. Thus, the electrostatics of the blinding structure becomes an integral part of the Cerenkov counter design. Any blinding scheme adopted must also isolate sense wires from neighboring sense wires. Because the feedback is in the UV, care must also be taken to insure that the metals in the blinds and chamber cathodes are chosen to minimize any photoelectric generation of charges which would drift to the sense wire and generate false hits.

The design used here was studied by simulating the electrostatics of the problem with the aid of computer programs which solve Maxwell's equations in two dimensions, which is valid at points well away from the end of the sense wires. The results and experience with the chosen design are presented here.

11.7.c Crosstalk

As indicated, the chamber signals from a ring imaging proportional wire cham-

Figure 11.5: Principle of "blinding" Photons emitted from the avalanche escape to the photodetector proper where they convert to create false Cerenkov signals

ber vary in size over as much as three orders of magnitude, but only the first decade is interesting. That is, the signal processing requirements slightly in that linearity is only important for the smallest, but most important signals. In fact, a non-linear pulse saturation would be desired, which limits the effects of largest signals. Nonetheless, the extraordinary dynamic range must be handled without adversely affecting the signal shape over as much of the range as possible. It turns out that the range of useful pulse heights from Cerenkov photoelectrons is close to 10, so that a very good signal must be measured for this range. If the crosstalk from neighboring simultaneous hits is too severe, the measurement is poor and the consequent resolution for the particular hit suffers. More serious, the large pulse heights from tracks must not induce signals on neighboring sense wires. For the work to be shown here, this particular difficulty was avoided by arranging the optics so that any Cerenkov photoelectrons would arrive at the electron counter well ahead of the primary ionization from the beam particle. The primary emphasis of

this work is to show proof that a good N_0 can be achieved for these chambers, so that the difficulties of the dE/dx from the track were not interesting. Later work, with improved custom electronics, has been done which indicates that in fact the difficulty can be minimized [14].

III. Prototype of a Cerenkov Ring Imaging Detector

III.1 The Cerenkov Detector for the SLD

Cerenkov detectors are designed for specific applications. In the case of the ring imaging counter discussed here, the device is being built to serve as the tool for identifying particles at the SLC accelerator^[15] facility at SLAC, as part of and in conjunction with the SLD^[16] detector. The SLC machine is being built for the express purpose of studying the physics of the Z_0 , the neutral intermediate vector boson of the Weinberg-Salam-Glashow theory. The Z_0 , with mass $m_Z \approx 92$ GeV, was discovered at CERN in 1983. It is responsible for such effects as the observation of a left-right asymmetry in the scattering of polarized electrons from a deuteron target, as observed at SLAC^[17]. The accelerator is designed to allow a complete energy scan across the Z_0 resonance, enabling high precision studies of the mass, width, and decay properties of the particle. It is thought to couple to all quarks and leptons with mass less than $m_Z/2$. The importance of studying the spectroscopy of all of the decays of the particle motivates the implementation a state-of-the-art particle detector like the SLD. The need to be able to identify the types of particles produced in the decay of the Z_0 's motivates the ring imaging counter design. From that point of view, the counter must be able to identify particles well for individual particle energies of $m_Z/2$, i.e., to approximately 50 GeV.

The ring-imaging counter must be built to cover as much of the complete solid angle around the interaction point of the colliding beams from the SLC as possible.

As such, barrel and endcap modules for the detector are being constructed. Figure III.1 gives an artists view of the counter as it will be installed in the SLD. The barrel and endcap sections are necessarily very different in design, because of the 0.6kG magnetic field produced by the SLD solenoid. In the barrel region, the magnetic field is parallel to the drift of photoelectrons in the drift volume, while in the endcap it is perpendicular. While the ring imaging principle is the same in both regions, the implementation of the counter is quite different.

Figure III.1: Cutaway view of one quarter of the SLD. The counter is at a 1m radius from the interaction point.

The geometry and resolution requirements of the SLD detector have driven the design choices for the ring imaging counter. It is necessary to have two distinct

radiators in order to identify particles over the full momentum range of outgoing particles, a liquid radiator to separate low momentum particles, and a gaseous radiator for the high momentum ones. The required resolution is on the order of 1mm in each dimension of the detector. This criterion was chosen to optimize performance at a reasonable cost. The most economical counter would be one where the number of electronics channels used in the chamber readout is smallest. Early on, it was recognized that this could be best accomplished if the photoelectrons were drifted along the barrel, parallel to the magnetic field. If the proper combination of gases and materials were found which would permit *SINGLE* electrons to drift unperturbed by any external forces over distances on the order of 1 meter, the detector could be realized. However, not only must the electrons drift unperturbed, they must also survive the drift, that is, there can be no electronegative components in the drift gas which would steal precious signal electrons from the chamber by capture. A successful implementation of this "long drift" detector was the challenge.

III.2 The Choice of a Photocathode

Certainly the most important element of the ring imaging detector is the photocathode. It is crucial that the quantum efficiency of the photocathode be very high, $\geq 30\%$, in order to get good particle identification. Because the Cerenkov light is emitted in the far ultraviolet, the list of candidate photocathodes was small. While solid photocathodes like CsI and NaI are sensitive in the far ultraviolet, they can at best do as well as 10% quantum efficiency, and that requires extraordinary care with the handling of the substances. There were liquid alternatives. Benzene, acetone, and triethylamine (TEA) have been used as photocathodes. Their quantum efficiency, while high, begins in the extreme UV, near 1300Å or less, where the choice of window materials is limited to the difficult and expensive alkali-halides. Also, early studies^[10] showed that the materials were too sensitive to O_2 and H_2O contamination. Figure III.2 illustrates the frequency of the quantum efficiency of several possible photocathodes, and the wavelength cutoff of some potential window materials.

The curve in Figure III.2 labeled TMAE reflected the quantum efficiency of the latest entry on the list of interesting photocathodes at the time. TMAE (tetrakis-dimethylamino-ethylene) as might be garnered from Figure III.2, has the

Figure III.2: The quantum efficiency of some potential photocathodes. The frequency cutoff of several window materials is also shown.

lowest photoionization threshold of any known liquid, at 5.36eV^[9]. The TMAE molecule, $C_2(N(CH_3)_2)_4$, illustrated schematically in Figure III.3 is one example of a tetraaminoethylene, which are fully substituted ethylene molecules, and its proclivity to give up electrons is reported to be due to the strong electron donor properties of the attached dialkylamino groups. By inspection of Figure III.2, the fortuitous property of TMAE is that it has good quantum efficiency in a spectral region where the window material could be fused silica (SiO_2). In fact, this overlap of material properties is what makes possible the cost effective implementation of a large-scale ring imaging. Fused silica is very robust and relatively inexpensive, and has an excellent spectral match with TMAE and the transmission region of potential ring imaging detector gases. While there are other TMAE-like molecules and hence potential alternatives, TMAE is the only one which is readily amenable to practical use.

The proper means of introduction of TMAE into the detector is achieved by bubbling the carrier gas through a container of TMAE. TMAE has the property that its vapor pressure at room temperature is very low, on the order of 1/2 Torr. The

TMAE Molecular Structure

Figure III.3: The chemical structure of TMAE. The exterior methyl groups are responsible for its remarkable properties.

vapor pressure determines the quantity of TMAE available to be picked up by the carrier gas stream, and so determines the absorption properties of the photodetector gas. The absorption length of ultraviolet photons in the TMAE environment of the photodetector must be short enough so that a large fraction of incident Cerenkov photons convert. Fortunately, as will be shown, the TMAE properties are such that it has an absorption length on the order of 2cm at 28°C, which is short enough to maximize the number of photoconversions in a shallow photodetector volume of a large scale counter, where it is important to maximize the Cerenkov radiator length.

The absorption length also depends upon the purity of the TMAE, which must be kept clean to maximize the counter performance. TMAE, as it is delivered contains several other chemicals dissolved in it, some of which are very UV absorptive, some of which are very electronegative. The gas delivery system must remove the significant fraction of these pollutants. It was important, then, to monitor the absorption length in the TMAE-laden photodetector gas during the course of this work. The measurement is principally dependent on the understanding of the characteristics of an ultraviolet monochromator, and since this instrument is at the heart of this and several other measurements which will be presented, it is necessary to digress from the principal topic at this point to discuss the UV monitoring system

III.2.a Precision Measurements with an ultraviolet monochromator.

The instrument used to make measurements of relevant UV quantities was a dual port monochromator, the Acton Model 502 vacuum monochromator, fitted with a Hamamatsu HTV-C704 power supply and a model L879-01 deuterium lamp for the UV light source. The source is bright from 1100Å to more than 4000Å, though the emission spectrum consists of many sharp resonance lines, which implies that a reference measurement must be made for every interesting sample measurement. The monochromator is designed to work in an evacuated environment but it was never used in that mode. Instead, it was fixtured to allow a continuous purge with a UV transparent gas, such as argon or the very pure gas from liquid nitrogen boil-off. Inside the vacuum chamber is a diffraction grating on a movable mount which allows it to be swiveled to reflect the source intensity to a chosen port. Each port was terminated by a phototube, the RCA 8850 quantacon, in front of which was a glass plate which had been coated with p-terphenyl wavelength shifter. Each port had independent UV intensity control, as did the lamp, by means of micrometer controlled slits. The pulses from the quantacon tube were amplified, discriminated, and scaled in a nearby NIM crate. A sample chamber, which varied with application, was machined to insert between a phototube and the monochromator and installed in such a way as to isolate the system from the exterior environment. The lamp and central monochromator were fitted with CaF_2 windows, so that UV of wavelength as short as $\approx 1450\text{\AA}$ could be transmitted to either port.

The connection to the monochromator of the sample stream was by Swagelok fittings to a series of stainless steel tubes. It was found that clean stainless steel tubing was necessary to maintain purity levels, and so was used throughout not only the monochromator station, but the entire gas system. The monochromator station was built so that a complete variety of input streams could be sampled. It was necessary during the course of the ring imaging tests to be able to sample gases as they flowed into/out of the detector, to bypass the detector altogether in order to monitor transparency and/or absorption of all kinds of combinations of gas streams, and to permit the calibration argon or nitrogen to be sampled. Also, as an integral part of the system, the water vapor and O_2 levels in the gas stream were monitored in the TMAE free parts of the gas system. Water absorbs UV strongly, and oxygen is very electronegative. The oxygen level was monitored by a Teledyne Model 310

Oxygen Trace Analyzer which is accurate to 1 ppm, and the water content was observed by a Panametrics System 5 Hygrometer, also good to 1 ppm.

In general, any measurement with this monochromator system consists of a few straightforward tasks. The system was used to make transmission and absorption measurements of window materials and various gas mixtures, mirror reflectances, and TMAE absorption. First, the sample volume was purged with the calibration gas, and then baseline transmission measurements were made over the wavelength range of interest. Typically, the interesting region was from 1600Å to 2500Å. When this was completed, the sample gas was purged through the volume, until the oxygen and water meters showed that all traces of those two contaminants were at tolerable levels. Measurements of the properties of the sample were taken, usually several times in order that the repeatability of the measurement established a stable result. The sample measurement was then corrected by the calibration. Experiments showed definite correlations of measurements with the amounts of O_2 and H_2O observed on the respective meters, so that the meters were useful and often indicative of leaks and other problems in the gas delivery system. The individual experimental techniques and results are described in detail where appropriate.

III.2.b Measurements of TMAE Properties

The quantum efficiency of TMAE has been measured with high statistics^[20]. This measurement is difficult and cannot be made in situ, but the UV absorption measurement of TMAE-laden detector gases could be monitored locally. The absorption measurement became a sensitive indicator of the cleanliness of the TMAE used in the experiment. Problems would frequently arise with the TMAE samples used. Were the TMAE to be used without treatment as delivered, these pollutants would enter the gas stream and the photodetector volume, robbing either Cerenkov photons or the generated photoelectrons, thereby reducing N_0 . Fortunately, standard chemical processing techniques were developed which were both simple and effective, so that the delivered TMAE samples could be purified^[21]. This purification had to be done in an anaerobic environment, as TMAE reacts strongly with O_2 in the environment. If any substantial amount of O_2 is present, the reaction takes place spontaneously, easily recognizable by a beautiful green fluorescence from the reactants. So, a dry box was used for all transfers, where the box could be purged

with a transparent inert gas and the internal O_2 content reduced to less than 100 ppm, typically.

As mentioned briefly previously, the TMAE was introduced into the gas stream by a bubbling technique. The purified liquid TMAE is placed in a clean, Pyrex receptacle which was fitted with input and output tubes. The jar is submerged in a temperature controlled bath of ethylene glycol. The bath, the Alfa RTE-4 Recirculating Bath, can maintain a temperature tolerance of 0.1 degrees. The temperature is also monitored with a thermometer. As the gas bubbles up through the liquid, the bubbles become saturated with TMAE vapor. Then, as long as the temperature away from the bath is maintained a few degrees above the temperature of the bath, the TMAE will remain in the gaseous phase as it is transported through the ring imaging detector. The typical operating temperatures are above room temperature, which meant that the gas plumbing had to be heated, which was done by wrapping the pipes with heating tapes. The temperatures were monitored at various places along the piping, so that TMAE condensation was avoided. This must be done carefully, as even a small, localized cold spot can cause condensation, which leaves TMAE liquid in the plumbing, a condition to be avoided. There was also some concern about hot spots which could possibly breakup the TMAE molecule, and perhaps introduce pollutants into the gas stream.

The absorption length of the TMAE saturated detector gas has an important temperature dependence which must be understood. The temperature characteristics will determine such parameters as the depth of the long drift box. The box should be as thin as possible, so that the largest proportion of the available detector volume can be used for radiator, yet not so thin that an appreciable number of Cerenkov photons pass through the photodetector volume without being converted to electrons. The measurement of the absorption length as a function of temperature is made with the monochromator apparatus, with some knowledge of a few simple facts of liquid vapor pressure behavior. Given a beam of photons of intensity I_0 , the intensity remaining after traversing a path of length ℓ is given by

$$I(\lambda) = I_0(\lambda)e^{-x/\ell} = I_0(\lambda)e^{-\rho\sigma(\lambda)x}, \quad (III.1)$$

where ρ is the density of the absorber and σ is the wavelength dependent absorption cross section. The concentration, and hence the absorption length, of the TMAE

vapor depends upon the partial pressure, which in turn depends upon the bubbler temperature as [22]

$$P \propto T^n e^{-L_0/RT} \quad (III.2)$$

where L_0 is the latent heat of vaporization, R is the gas constant, and n is near one. In fact, in the region of experimental interest to the ring imaging detector, the power law temperature dependence is negligible and the exponential term dominates. Then, if the absorption length at 300K is known, it follows that

$$\mu = \mu(300K)e^{-\frac{L_0}{R}(\frac{1}{T} - \frac{1}{300})} \equiv \mu_0 e^{-\frac{L_0}{R}(\frac{1}{T} - \frac{1}{300})} \quad (III.3)$$

The test of this is a measurement of the absorption dependence over the relevant wavelengths, as described previously, over the modest temperature range of potential detector operation. The interesting quantity is

$$T(\lambda) = I_{\text{sample}}(\lambda)/I_{\text{ref}}(\lambda) = e^{-\mu(\lambda)\ell} \quad (III.4)$$

so that

$$\mu(\lambda) = -\ln T(\lambda)/\ell. \quad (III.5)$$

Included in the numerator and denominator of (III.4) is a background intensity subtraction, to allow for noise in the photomultiplier tube, a problem of very low level which stayed low throughout the experiment. The results of the tests are shown in Figure III.4. The measurement of the TMAE absorption was often repeated, nominally at the 28° temperature. The wavelength dependence of the temperature measurement was also reproducible. If the curves were not reproduced, it was an indication that there were leaks in the plumbing, or contaminants in the TMAE. Often, it was this measurement which warned of trouble.

III.3 Choice of Cerenkov Radiators

There are three physical considerations for a radiator of Cerenkov light. First and foremost, the radiator must be transparent to the photons which are emitted. Second, the index of refraction should be as high as possible in order to get the highest possible emitted photon yield in the spectral window of the device. Third,

Figure III.4: TMAE absorption coefficient as a function of temperature. The equation is a fit to the data with the Clausius-Clapeyron equation.

the dispersion of the emitted photons should be small enough that the spread of the Cerenkov cone does not severely compromise the separation performance. Note that this third criterion is an essential ingredient for a ring imaging counter, and in fact will turn out to limit the resolution of the liquid radiator. In a threshold counter, the dispersion of the radiation is not such a crucial parameter.

The majority of Cerenkov radiation is emitted in the ultraviolet region of the electromagnetic spectrum where the choice of radiator materials is limited, particularly for the liquid radiator. The choices for both types of radiators are listed in Table III.1 and Table III.2.

Ideally, complete particle separation for momenta as large as 50 GeV would be achieved in the perfect Cerenkov counter. However, at the SLD, Monte Carlo studies of heavy quark decays indicate that exceptional coverage of the 1 to 10 GeV momentum range is desirable. Appropriately, the liquid Cerenkov should provide particle separation at least the 3σ level out to a few GeV, where the gas coverage

Gaseous Radiators			
Gas	Index $(n-1) \times 10^{-4}$	E_0^π (GeV/c ²)	E_0^K (GeV/c ²)
Helium (He)	35	17.0	59.0
Argon (Ar)	283	5.67	20.9
Nitrogen (N ₂)	297	5.5	20.3
Methane (CH ₄)	441	4.5	16.7
Carbon Dioxide (CO ₂)	450	4.7	17.4
Ethane (C ₂ H ₆)	706	4.5	16.7
Isobutane (C ₄ H ₁₀)	1550	2.50	8.9
Neopentane (C ₅ H ₁₂)	1710	2.38	8.5
Perfluoropentane (C ₅ F ₁₂)	1725	2.37	8.4

Table III.1: Potential gaseous radiators for the ring imaging counter. The long wavelength index and the threshold energy for π 's and K's are also shown.

would then become useful at about the same level. Potential combinations of the radiators shown in the above tables indicated that a choice of perfluorohexane as the liquid and either isobutane or perfluoropentane as the gas radiator would give good performance. They have ultraviolet transparencies which overlap the region of TMAE quantum efficiency well, so that the optimum performance of the counter can be achieved. Figure III.5 is a sketch of the predicted spectral acceptance for the prototype Cerenkov ring imaging counter, where isobutane was chosen to be the gas radiator^{*}.

III.4 Choice of Window Material

The radiator volumes must be isolated from the photodetector, in the focal

^{*} Similar figures exist in the case where perfluoropentane is the radiator, as that gas will be used in the actual SLD and LEP counters. See, for example, the SLD Design Report

Liquid Radiators				
Liquid	Index	E_0^{π} (GeV/c ²)	E_0^K (GeV/c ²)	B.P. (°C)
Liquid Helium (He)	1.03	0.56	2.07	-269
Liquid Neon (Ne)	1.10	0.32	1.19	-246
Liquid Nitrogen (N ₂)	1.205	0.24	0.89	-190
Liquid Argon (Ar)	1.23	0.23	0.85	-186
Perfluorohexane (C ₆ F ₁₄)	1.277	0.22	0.8	56
Water (H ₂ O)	1.33	0.20	0.75	100

Table III.2: Potential liquid radiators for the ring imaging counter. The boiling points for the liquids are also shown.

plane of the spherical mirror, to confine all the photoconversions there. At the same time, the photons must be able to travel across the boundary with little loss. These two conditions mean that the chosen material must be transparent in the ultraviolet, while at the same time being non-porous and strong. The material must be inert, so that no undesirable chemical reactions occur at the interface. The necessity for ultraviolet transparency constrains the choice of window materials to a choice of lithium fluoride, magnesium fluoride, calcium fluoride, barium fluoride, natural quartz, or fused silica. The fused silica option is essentially the only practical one for the ring imaging systems, because it is not significantly hygroscopic, has a high tolerance to radiation, and primarily, can be fashioned into large flat sheets. There was some concern that it may be porous to water vapor and oxygen at very small levels, and that stresses introduced during construction may lead to development of microcracks. These questions remain unresolved as of this writing.

III.5 Gas Radiator Mirror

As explained in chapter I, the gas radiator volume is terminated by a UV reflective mirror. Such a mirror built of ground and polished, which is coated with

Figure III.5: Spectral acceptance of the gas and liquid systems. The TMAE quantum efficiency and the transparency for a quartz window selection are also shown.

a sputtered aluminum surface, and coated with a transparent UV layer, typically

MgF₂. The coating is necessary to prevent oxidation of the aluminum and hence maintain the high level of reflectivity. The magnesium fluoride is insensitive to water vapor and doesn't react to degrade its transmission when exposed to the Cerenkov radiator environment, whether it be isobutane or the perfluoropentane of the SLD design.

The requirements on the surface precision and focal properties of a mirror suitable for a ring imaging Cerenkov counter are very relaxed in comparison with even a modest astronomical mirror. The error introduced by any mirror imperfections must just effect the resulting ring image at levels smaller than other errors, typically no more than 0.5 mm. The mirror chosen for the beam tests with the prototype detector was an $f = 50\text{cm}$ Acton telescope mirror. Its properties surpassed what was required, with focal plane errors $\leq 0.5\text{ mm}$. The mirror was delivered with a coating of MgF₂. The monochromator was used to measure the UV reflectivity which was found to be $\geq 95\%$ over the interesting region.

III.6 Choice of Photodetector Drift Gas

Here, again, the list of potential candidates for the gas which fills the interior of the quartz drift box is a short one⁽²³⁾. First, any candidate drift gas must have the following properties:

1. UV transparent. Clearly, the radiator gas must not absorb Cerenkov photons. Any such effect is a direct loss to the counter.
2. Inert in the presence of the photocathode. There cannot be any poisoning effects when the TMAE and the detector gas occupy the same volume. Because the TMAE is mixed with the gas by a bubbling technique, there must not be any compounds produced which are either ultraviolet opaque or are electronegative. Nature has provided several candidate detector gases which are TMAE-friendly.
3. Non-electronegative. The Cerenkov photoelectrons must be able to migrate from the point at which they are produced to the electron detector without loss. This property of the gas is reflected in the measurement of electron lifetime. If a detector gas is to be a candidate for the counter, a single electron should be able to travel the full length of the box at least 80% of the time.

This corresponds to an electron lifetime time in the gas of approximately 5 times the box length. This is a very stringent constraint on the purity of the gas, and while there are examples, e.g., the noble gases and several hydrocarbons which in their pure form meet this requirement, there must be no electronegative contaminants in the gas stream. This figure of merit, the electron lifetime will, in the end, be the measure of performance for the entire counter.

4. Dielectric strength. The drift gases are also the avalanche multiplication gases for the electron detector. When single electron detection is necessary, it turns out that few gases can maintain a Townsend process prior to breakdown before a single electron counting plateau is reached. Since breakdown is thought to be associated with the onset of strong UV emission, the drift gas must also be able to absorb the UV which is generated in the avalanche, i.e., it must be self-quenching. This ultraviolet absorption must not be strong at the Cerenkov wavelengths though, or else signal photons are lost. Also, a high dielectric strength will aid in the suppression of spurious photons from points of UV corona emission inside the drift box. While special care is taken to protect against corona emission points, additional protection is always welcome.
5. Photon Feedback. It has been shown that several potentially interesting combinations of noble gases and hydrocarbons don't work, because they produce too much avalanche induced UV in the Cerenkov photon spectrum⁽²⁴⁾. There are no obvious or perfect choices, but careful chamber designs using just hydrocarbon mixtures have brought this problem under control.
6. Drift velocity. The drift velocity of electrons in the drift gas is an important resolution parameter. For example, to meet the 1mm resolution criterion, the drift velocity must not vary by more than 0.1% over the full drift volume of about one meter. This stringent requirement would ideally be met by a drift gas whose properties were such that the drift velocity was independent of electric field, a saturated drift gas. While candidates exist which behave like this at the few hundred volts per centimeter of electric field which is appropriate to the Cerenkov ring imaging detector, they do not meet the other criteria mentioned above. It is unfortunate that the hydrocarbon gases

have a nearly linear behavior of the drift velocity with electric field, because a careful calibration of the drift velocity is then essential for the performance of the counter. Such calibrations are routinely made however with the use of a N_2 laser. This laser has become an important tool for research in the ring imaging field, and more will follow about it.

7. Diffusion. Drifting signal electrons will suffer from diffusion effects due to the massive molecules of the drift gas. It is desirable that the effects upon the resolution due to diffusion be on the same order of magnitude, or less, as the other individual resolution contributions, i.e., order of 1mm or less. The hydrocarbons meet this requirement, especially the heavier ones, while the light (inexpensive) noble gases are too diffusive. For the large counters, the magnetic field effects on the diffusion should be considered.

Clearly, a hydrocarbon drift gas was favored. Pure or mixed methane, ethane, and isobutane environments were considered. Pure methane was quickly ruled out. Preliminary tests had shown that methane is very unstable by itself, i.e., it is subject to premature voltage breakdown. Also, the pulse height behavior was unsatisfactory. Ethane, at that time, was expensive and without a local supplier in large, relatively pure ($\geq 99\%$) quantities*. Pure isobutane required relatively higher chamber voltages in order to work well as a chamber gas. Also, indications were that it did not yield good electron lifetimes¹²⁹. In addition, the UV cutoff of isobutane was just at the edge of acceptability, and required special care in order to keep the transmission high.

The combinations of 70/30 methane/isobutane or 80/20 methane/ethane were used. Methane was readily available in very pure form in quantity, a small quantity of research grade ethane was procured, and less pure isobutane, as it came from the manufacturer was used. The isobutane could be purified, based on monochromator tests of the UV transmission, by venting off 10% of the filled vessel weight of the gas upon delivery. All of the results shown here use one of these mixtures.

III.7 The Gas Delivery System

A major focal point of effort involved the detailed construction of the gas

* The problems with ethane no longer exist as of this writing. In fact, it is now the preferred choice for the SLD ring imaging detector

delivery system. The system adopted must serve several functions.

1. The system must be flexible enough to work with different gas mixtures. Separate flowmeters were installed in the input gas lines to permit a tuning of the flow rates. The constituent gases then entered a small mixing volume and allowed to come to equilibrium before coming in contact with TMAE.
2. The severe cleanliness requirements required the use of electro-polished stainless steel tubing, all of which were internally cleaned with acetone and alcohol, prior to assembly. The pipes were connected with Swagelok fittings, except in the most critical assemblies, where Kajan VCR fittings were used. All valves used were high quality stainless steel valves with stainless steel bellows. These valves are expensive, but tests showed that the level of contamination introduced into the drift box would be minimized by the adoption of a fully stainless steel delivery system.
3. The input gases were to be of the highest quality available. The gas bottles came from the manufacturer with less than 10ppm of contamination of impurities. Nonetheless, additional safeguards were introduced to prevent any particulates from within the bottles from being allowed to enter the drift box stream. 13X filters were installed at the output of the low-pressure regulators for this purpose. Also, Oxisorb O_2 and H_2O filters were installed in the gas stream, to reduce these contaminants to the smallest level.
4. The TMAE-laden photodetector gas must be delivered to the drift box at a specified temperature. As the temperature is to be a variable for diagnostic tests, the control of the temperature at all positions along the gas system was essential. This was accomplished by the use of temperature-controlled baths at the TMAE bubbler, and with the aid of Variac-controlled heating tapes which were wrapped around the stainless steel pipes for the full delivery length.
5. The pressure control of the apparatus was another critical issue. A moderate overpressure between the intake and exhaust manifolds of the drift box would mean that the quartz windows would be the sole source of restraint against instant pressure relief in the system, potentially a disastrous situation in the making. Not only the problem of large cracks, or worse, developing, but the pressure overhead must not be allowed to open even the smallest

of cracks, which would then serve as doors for contaminants to enter the box by diffusion. Equally important, the escape of the hydrocarbon/TMAE mixture to the external environment must be prohibited. To this end, a sensitive system of pressure regulation by remotely monitored capacitance manometers were installed in parallel with the drift box exhausts.

6. The safety problems with TMAE are not well understood, but they are potentially severe. With that in mind, a refrigeration system was built at the exhaust of the complete gas line which served as a condenser of the entrained vapor. Kept slightly above the freezing point of water, to prevent water vapor condensation and potential ice blockages in the piping (a learned thing), the TMAE was extracted from the gas stream, potentially to be reclaimed and used for subsequent experiments.
7. As discussed briefly in the section on monitoring, the gas system was also designed to deliver a multitude of alternate gas streams to the monitoring equipment. Enough valves and pipe were installed to allow virtually any combination of input/output detector gas with/without TMAE to be delivered to any combination of the drift box/monochromator/oxygen meters/water vapor measurement apparatus. Also, a source of liquid nitrogen was plumbed in as the monochromator calibration.

The gas system is shown schematically in Figure III.6.

III.8 Photodetector Volume, or Drift Box

The photodetector volume which supports the TMAE-laden detector gas mixture was constructed entirely of fiberglass G-10 composite, fused silica sheets, and Shell Epon 826 and Versamid 140 epoxy adhesives. Preliminary tests had confirmed that the three structural members were not subject to corrosion in a TMAE atmosphere. Also, the construction must be of these good insulators, because the box must withstand a ≈ 60 kV voltage drop along its 80cm length. The box was 20cm wide, and tapered from 4cm in depth at the high voltage cathode end to 4.8 cm at the electron detector. This was to protect against diffusion losses near the window surfaces.

Figure III.6: The gas system of the prototype tests. The pressure control devices and monitoring equipment are also shown.

The sidewall G-10 pieces were fashioned with printed circuit tinned copper traces at 2.5mm inch intervals. The traces were connected to one another by 1% 1M Ω resistors. These resistors defined the electric potential along the length of the box. At the high voltage end, the resistor chain was soldered to a high voltage cathode, a flat, rectangular copper plate to define the initial voltage plane. At the detector end, the drift box was fitted with a G-10 throat, with the field gradient properly continued by traces on it, before fitting flush on the electron detector mounting. The termination of the resistor chain and the field are discussed in chapter III.

III.8.a Electrostatics of the Long Drift Prototype

As has been mentioned, the electric field in which the electrons drift must be very uniform in order that events be properly reconstructed. An electrostatic field cage was built into the long drift box of the prototype to accomplish. The cage consists of CuBe wires, strung at a tension of 100g, on the inner and outer surfaces of the front and back quartz windows. There are wires inside and out in order to prevent polarization of the fused silica window material, which might not only cause local field distortions outright, but also serve as a getter and store of ions in the drift

box, another potential source of undesirable electric fields. The wires, 75 μ copper-beryllium stretched to 100 grams, were soldered to traces on the inner and outer surfaces of the G10 side-pieces. The work was done in a very modest clean room, under a microscope to help insure that the inner and outer equipotential wires were well aligned. In addition, another external field cage of much looser wire pitch, one inch, was constructed on a G10 frame. The same potential gradient existed on this outer field cage, which served as more robust isolation of the long drift box from external ground sources. The high voltage connection for the field cage/drift box structure was made to a large flat copper plate at the bottom of the exterior field cage. This was positioned 3cm from the end of the drift box, and the field properly graded in between. The copper plate was fit into a three-quarter inch thick G-10 slab, which was the insulation from nearby grounds. The slab had G-10 posts glued to it, set in a pattern to define five concentric sets of smooth cornered rectangles of slowly increasing cross-section, on which thick copper wires were connected. Each wire was connected to a potential on the exterior field cage to properly degrade the high voltage to grounds at the edge of the slab. Electrostatic simulations of this geometry, using a program which solves Maxwell's equations in 2 dimensions for an arbitrary set of conducting surfaces [26], showed that the resulting field shape inside the drift box was extremely robust against penetration of any field lines from sources exterior to the it.

The most direct way of measuring any field distortions would be to map the distributions of wire hits at the electron detector due to charges which were created from sources at known positions along the length of the drift box. Such distributions should show deviations from linearity if there are field effects. This is the most basic test for any drift box design, and the measurement involved the use of the ultraviolet laser. Since this laser has become an essential tool for the development and testing of the detector, another aside to describe the laser apparatus and operation is warranted.

III.8.b Diagnostic Testing with an UV Laser.

The laser is an ideal tool for diagnostic studies of gas proportional chambers. The ability to direct the laser beam, and control the light intensity with optics and filters, allows the experimenter the means and flexibility to study many aspects of

drift chambers. In particular, for the ring imaging detector, potential uses include the study of electrostatic distortions, the parameterization of the drift properties of electrons in interesting gas environments, the study of the properties of potential electron detectors, and the ability to simulate the effects due to different initial charge depositions.

The laser used for this work was the Photochemical Research Associates Incorporated Model LN100 Nitrogen laser. The laser operates at a wavelength of 337nm, at a nominal output energy of 70 μ J per 300 psec pulse. The pulse is initiated by a spark gap, whose gap width and voltage were externally adjustable. The option is available to externally trigger the laser, although the internal trigger was most often used for this work. The nitrogen supply was delivered through a low pressure regulator from the same boiloff supply used for monitoring purposes. These few parameters were sufficient to provide a stable, reliable light source, provided the entire system was cleaned at regular intervals.

It remains a controversial subject, the precise mechanism of charge deposition in the gas of a chamber by laser light. The current lore is that a multiple photon absorption occurs in the gas when exposed to the high intensity light at the waist of the focussed beam. If this is the case, it should follow that the amount of charge deposited in a chamber by the laser vary as some power law of the laser intensity. The charge variation with intensity might be measured by integrating the output signal from the electron detector. Indications from simple tests here indicated that the charge varied like the square of the intensity.

The primary uses of the laser are for measuring lifetimes, studying electrostatics, and measuring drift velocities. Because the electron counter is sensitive to single photoelectrons, the raw laser intensity must be heavily attenuated. This was done with an assortment of Melles-Griot UV interference filters of differing optical powers. In order to determine when the laser light was most likely generating single electrons, a counting scheme was used. In principle, if the laser is tuned to provide a single electron in the drift chamber per laser pulse on average, Poisson statistics dictates that only 1/e of the time will a single electron actually be produced, for a Poisson process of mean 1. By the counting of the number of null hits in a large number of triggers, an estimate of the Poisson parameter can be made. By observing how this parameter changes with distance, voltage, or location, conclusions can

be drawn as to the sources of losses or distortions.

This requires that a stable, reproducible environment be maintained, both inside the photodetector gas volume and for the laser itself. The laser output was monitored by a photodiode which sampled a fraction of the UV from a beam splitter at the output coupler of the laser. The photodiode signal, which also served as the trigger for the laser tests, was sampled by an ADC which could readout on-line by the control system. Checks were made often that the pulse height distribution was sharply peaked and approximately Gaussian. The stability of the gaseous environment was not easy to monitor. In fact, when the laser ADC was stable, any difficulties were attributed to the ring imaging photodetector. However, many hours of working with gas system and readout systems finally showed that a consistent behavior could be obtained.

III.9 The Electron Detector

In order to use the laser to make measurements, or more importantly, to measure Cerenkov rings, an electron detector (or counter) sensitive to single photoelectrons had to be designed and constructed. The electron detector is the business end of the drift box. Here, electrons from the photodetector proper are electrostatically focussed onto small radius wires. Due to the potential structure in the neighborhood of the wires, an electron avalanche is precipitated there. It is the job of the electron detector to sense and amplify the signal locally, for subsequent electronic processing. The electron detector must be able to measure two of the three coordinates with good resolution, it must couple with good efficiency to the lower electric field of the photodetector proper so that all of the photoelectrons make it in to the electron detector, it must be designed with proper electrostatics to guarantee that the electrons end up on sense wires, and in these chambers, must also blind the photodetector proper from the UV photons from the avalanche. Thus a careful design of this apparatus is a must. As such, the whole of chapter III is devoted to the electron detector, and a further description will be reserved until then. A full description includes the electronics which is responsible for amplifying and distributing the signals to the external world.

III.10 The Womb

As has been stressed, the environment around the TMAE filled detector box must be temperature controlled. A full range of control from room temperature to greater than 50°C was important for the experiment. An aluminum enclosure was built to house the prototype and the radiator volumes. The temperature inside was controlled with the aid of platinum thermistors which allowed feedback-controlled heating tapes inside the volume to be turned on and off, to a precision of one degree. Also, the gas pressure inside the womb was held constant with respect to atmosphere, and more importantly, with respect to the pressure inside the enclosed drift box, by set of sensitive capacitance manometers. The end pieces from the womb box could be removed so that the drift box, gas radiator mirror, liquid radiator vessel, and external electric field cage could be serviced. The exterior of the womb was insulated with 2in thick styrofoam.

The drift box was hung vertically so that the high voltage cathode was nearest the womb floor. Gas was delivered to the top of the drift box and exhausted at the bottom. The outlet gas from the drift box was vented to the outside world through a glass rod which was fit into the G-10 block supporting the high voltage cathode of the drift box. The use of the glass tube was mandatory, because a non-conducting pipe had to be made to get the gas from the high voltage cathode to the external world. The high voltage was brought in near the mirror by high voltage feedthrough, which was easily disconnected from the outside. The high voltage supply was a Spellman Model RHR-100-N-100, which could deliver up to 100kV, although it was only taken to 90kV once during these experiments. Inside the womb, a 100kV insulating cable carried the potential to the cathode where it was carefully soldered.

The external field cage was mounted around the drift box on the G10 slab which simply rested on the womb floor. The work required to accommodate the high voltage proved tedious, but when sufficient care was taken to insure that all solder connections were rounded and that all wires were smooth and dust free, an electrically quiet, corona-free high voltage environment was created.

A mirror was mounted at a precise position on the downstream door, and was mounted so that the tilt of the mirror relative to an incoming particle beam could be

adjusted and known. A liquid radiator cell, which consisted of 2 transparent quartz windows mounted in a flange, was mounted on a non-conducting G10 frame which could be externally rotated into/out of the particle beam. The liquid radiator was connected to a recirculating pump, so that the perfluorohexane could be recycled, cleaned, and monitored by the UV monochromator. A schematic view of the womb interior is shown in Figure III.7. A view into the womb interior from the upstream end is shown in Figure III.8.

Figure III.7: The interior of the womb. The relative locations of the drift box, radiators, mirror, electron detector, and external field cage is shown.

Figure III.8: A view of the ring imaging prototype. The liquid radiator cell (the small circular glass vial on the G-10 arm), the spherical mirror (mounted on the downstream door), the long drift box (the two sets of narrowly spaced printed circuit traces extending from top to bottom on left and right define the sides of the quartz windows), the external field cage, and the womb vessel are all shown.

IV. The Carbon Fiber Electron Detector

The resolution requirements of the ring imaging proportional counter are very demanding, if the highest levels of performance are to be attained. The particle identification capability of a given design is in direct proportion to the resolving power of the instrument, which in this case amounts to the degree to which single photoelectrons can be detected, analyzed, and their point of origin in the photodetector determined. If all three spatial coordinates of an electron's origin are found with good accuracy, the resolution will be optimized. Also, if no coordinate is to dominate the resolution, all three coordinate measurements should be of like accuracy. In the ring imaging chamber, two coordinates are given in a relatively straightforward manner, as described in chapter II. One coordinate is given by the position of the sensing anode wire, the other by the drift time to that wire. The position of the electron along the wire must be also determined. As the wires are short, on the order of 10 cm, it will be shown that conventional wire chamber technology cannot provide the requisite resolution in this coordinate, so that a new technique had to be found and developed into a viable, practical, and most importantly, functional part of the ring imaging apparatus.

IV.1 Properties of Proportional Wire Chambers

The science of the proportional wire chamber has been studied for almost 25 years, since Charpak first built one in 1967^[27]. The study of wire chambers and the behavior of electrons in gases under the influence of large electric fields of course precedes Charpak, back to the time of Townsend in the early part of the twentieth

century. An excellent compendium of the of contemporary knowledge about wire chambers is found in Sauli^[24]. The details of the process of electron multiplication around a sense by the avalanche mechanism can be found there. Only the salient points of the process will be described in any detail here. In particular, most of Sauli's discussion is focussed on the problem of charge multiplication for large initial electron concentrations, which is very different from the interesting case of single electrons in the ring imaging detector.

The freshman physics problem of a charged wire along the axis of a concentric cylinder defines the electrostatics of a simple proportional wire chamber. Between the wire, the anode, and the cylinder, the cathode, there is a uniform density of gas. If a potential difference is applied between the anode and cathode, simple electrostatics means that an electric field is set up between the conductors, a field which varies as the inverse of the distance from the anode to the cathode. At this point, if an electron is introduced into the gas by some means, it will move in the electric field to the most positive potential, defined to be the wire here. If the potential is increased, eventually the electric field imparts sufficient energy to the electron so that in a collision with a gas molecule, it can give up some of its kinetic energy to the molecule, liberating bound electrons, which then propagate under the field, gain energy, and liberate even more bound charges. This is the Townsend avalanche. In most gases, as the potential difference is further increased, the amount of charge liberated is proportional to the voltage, hence the name proportional chamber. At still higher voltages, so much charge can be liberated, that a continuous discharge is developed in the gas, leading to a spark. The region of operation of interest to the ring imaging device is the proportional region, for two reasons. First, as the name implies, a spark emits copious amounts of radiation, and again, generation of ultraviolet light other than that from Cerenkov radiation is to be avoided in the ring imaging counter. Second, a spark also emits radiation in the radio frequencies, which leads to a noisy electronic environment which must also be avoided. Other modes of operation of a wire chamber could in principle be utilized, for example, the limited streamer mode, but the ultraviolet background problem would still be immense. In fact, even when operated in the proportional mode, the UV emission is considerable. Efforts must be made to keep this extraneous source of ultraviolet to a minimum.

Before any electron counter is deemed acceptable for large scale application, it

must be shown to be efficient for electron counting. Full efficiency implies that every charge incident upon the counter generates a detectable signal. At that point, the chamber is said to be plateaued. The beginning of the plateau region, the knee, is strictly a function of the electronic sensitivity of the proportional chamber readout. The length of the plateau is a property of the detector gas. For single electron sensitive counters, however, the plateau is not so clearly defined. It has been known for some time that the pulse height spectrum from single electron chambers are very broad, with output pulses varying from zero to 4 or 5 times the mean. For a given electronic threshold then, some fraction of pulse heights is always below the threshold, and full counting efficiency is unattainable. For a purely Townsend process, one of an exponential pulse height behavior, a significant number of signals can be lost. Fortunately, nature has provided some gases which appear to avoid this problem in that under the extreme high field conditions in the single electron wire chambers, the statistical process changes. For reasons that are not completely understood, the heavier hydrocarbon gases exhibit peaked pulse height distributions prior to breakdown. One explanation for this behavior is that one tenet of the Townsend model, that the probability of electron ionization is independent of the distance from the last ionization, is no longer valid⁽²⁰⁾. In any case, the observation of a peaked pulse height distribution from single electrons can be regarded as the knee of the plateau for a single electron counter.

IV.2 Electron Detector Options

The goals for the electron detector have been previously defined. An efficient photoelectron detector must be developed which will provide the means for satisfying all of the resolution goals of the ring imaging detector, i.e., space point resolution on the order of 1mm in all coordinates, while at the same time minimizing the amount of spurious photoelectron generation due to the conversion of ultraviolet photons which are emitted in the routine process of avalanche multiplication. The electron detector must be robust against failure from breakdown, yet be capable of operation over a huge dynamic range of the initial charge which will be incident upon it. It must be amenable to mass production techniques as would be suitable for a practical particle detector.

Consider then the proportional chamber shown schematically in Figure IV.1. In the figure, the anode wires are represented by the points in the middle of the open cells, between the fingers of the scalloped cathode which extend out beyond the plane of the anode wires, so as to serve as optical isolation between adjacent wires. The cathode is a single piece of aluminum which was machined to specification on a digitally controlled mill. As aluminum has a relatively low work function, on the order of 4ev, it was necessary to coat it with nickel, work function about 5.5ev, in order to minimize oxidation of the aluminum, and to reduce the secondary emission of electrons from the cathode by the photoelectric effect. The cathode-to-anode potential defines the avalanche properties of the chamber.

In between the cathode and the photodetector proper is the blinding grid, or affectionately, the "cowfence". The cowfence is made from 1mm diameter tinned copper welding rods, which were work-hardened by pulling the rods to their elastic limit. They were cut to length, and soldered to traces on G10 support pieces at the ends of the rods, so that a high voltage could be connected to each of the three planes of wires. The purpose of the cowfence is to prevent the ultraviolet photons emitted from the avalanche from propagating back into the photodetector. As such, the dimensions of the grid were chosen so that only 6% of 2π did not have an opaque barrier in the photon line-of-flight. Also, if some of the ultraviolet photons did convert inside the grid structure, the electrostatics were designed so that those photoelectrons drift to either the cathode or to another part of the grid, so that no false signals are made. As well, the potentials on the cowfence are defined so that all charges which drift into the structure from the photodetector proper are electrostatically focussed to the anode wires. That is, all the field lines from the exterior drift box end on a sense wire. This construction was first simulated with the aid of the electrostatics program. An important series of tests were performed which proved that the design was correct for the actual prototype detector. These tests are discussed in ch.IV. There is a finite probability of regenerated photons being converted and hence detected as real signals; it is simply unavoidable at some level. The geometry of the detector is such that the complete assembly is less than one TMAE absorption length in dimension from the anode wire to the outermost cowfence blind, so that the majority of the dangerous feedback photons are stopped. However, the chance of emission and conversion along the wire is still a potential problem, but the solid angle for that process is a small fraction of the total. An

Figure IV.1: The prototype proportional wire chamber with "cowfence". The true distances are indicated by the dimensions on the right, the high voltage connections are indicated at left.

empirical upper limit of the effectiveness of the blinding scheme is derived from the tests to be described later.

The sense wires were made of carbon monofilaments, a new and important innovation which has shown to very effective and very robust in these chambers. The choice of carbon monofilaments as the sense wires was entirely motivated by the need to have 1mm millimeter resolution in the coordinate along the anode wires. There were two candidate means for obtaining the required measurement precision in this coordinate: the induced cathode pick-up and the charge division. The cathode

pick-up method is, at first glance, the preferred method, since it does not require the development of any new technology. As detailed in Sauli, the potentials inside of the wire chamber are maintained so that the negative charges avalanche onto the anode wires. The large, slow positive ions drift slowly away from the anode wire, in the process inducing an opposite charge on the anode wire, and at the same time, a like charge on the cathode surfaces. If the cathode surfaces were segmented in such a way that the induced charge could be detected electronically, a determination of the position of the avalanche could be made that way. However, in contrast to the anode wire which feels the full effect of all of the field lines from the drifting ion cloud, the entire cathode surface shares the field lines from the same. Hence, any cathode segmentation implies that only a small portion of the available signal can be detected on any one cathode part. In conventional chambers, where signal levels are typically much larger than noise, this is not a problem, because the induced signals are still relatively easy to detect. At the single electron level, though, where the anode wire signal is at the very limit of detection, the cathode signal is that much more difficult. Such a design has been demonstrated for a ring imaging detector, where great pains have been taken to optimize the coupling to individual cathode segments by placing the anode wires very close to the cathode at very close tolerances^[30]. Small tolerance errors mean large gain variations, the gain is exponential in the separation, so that an extraordinary amount of care must be exercised in the construction of such a chamber. The alternative method of charge division then becomes much more attractive, provided a means of doing it can be found.

IV.3 Elementary Charge Division

When this work began, carbon monofilaments had not been used in proportional chambers. It turns out that independently and concurrently, a group of Japanese nuclear physicists had also recognized the imaging potential of carbon monofilaments^[31]. Their measurements and experiences with the fibers were very similar to ours.

Now, the TMAE absorption length, λ , was found to be $\lambda \approx 2\text{cm}$ at the temperatures of interest. This implies that for a photon normally incident upon the photodetector, the box must be at least 6 cm deep if the majority of photons are

to be converted inside the box. If a precision of 1mm is desired, then a position measurement must be accurate to $\approx 1\%$ of the wire length. The accuracy attainable is related to the signal to noise performance of the electronics. Preliminary work had indicated that the mean charge which was to be detected by the electronics for single electrons was $\approx 2 \times 10^5$ or less. That is, the onset of the peaking behavior of the pulse height spectrum from photoelectrons corresponded to the that charge gain. It was also known that electronic pre-amplifiers were available which had noise levels on the order of 1000e's. When connected for a charge division measurement, however, that noise level increased. The situation is illustrated schematically in Figure IV.2.

Figure IV.2: The charge division technique. The ends of the fiber are connected to the inputs of the operational amplifiers as shown. The input impedance is represented by the resistor at the inputs.

Operational amplifiers have the property that the inputs are almost at ground potential, so called virtual ground. Therefore, the situation is that of a resistance connected to the same potential at both ends. The thermodynamic phenomena of Johnson noise then serves as a current source for the detectors at each end. Because there is no external source, the total current in the wire at any time is zero, but depending upon the point of generation, each detector will see different amounts of charge. If these detectors are charge sensitive, the currents can be integrated onto capacitors, shown as C_f in the figure. From simple noise theory, the r.m.s. noise

charge, called the ENC (Equivalent Noise Charge)^[22] is

$$ENC = 2.7 \times \sqrt{\frac{kT\tau}{2R_w} + \frac{kTR_sC_f^2}{2\tau}} \quad (IV.1)$$

for the circuit in Figure IV.2. Here, R_s is the noise resistance in the first transistor in the amplifier, that is, the first amplifying element, τ is the shaping time of the circuit, which depends upon the values of the individual elements, and kT is the thermal energy. This noise must be added in quadrature with the intrinsic, i.e., unloaded electronic noise. This last contribution is a function of the avalanche location. To evaluate it, define

$$A_Q = \frac{Q_R - Q_L}{Q} = \frac{Q_R - Q_L}{Q_R + Q_L} \quad (IV.2)$$

as the charge asymmetry for any charge division. For a charge incident at z as in the figure, the desired measurement would relate this position to the asymmetry by

$$z = \frac{L}{2}(1 + A_Q). \quad (IV.3)$$

If the input impedance of the amplifiers can be neglected, as is the case here because the wire resistance is so high, A_Q varies between -1 and 1. Using all of the relations gives that the total charge error is

$$\sigma_Q = \sqrt{(ENC)^2 + \left(\frac{Q_L}{Q}\right)^2 \sigma_R^2 + \left(\frac{Q_R}{Q}\right)^2 \sigma_L^2} \quad (IV.4)$$

and the resolution is

$$\frac{\delta_z}{L} = \frac{\sigma_A}{2} = \frac{\sigma_Q}{Q}. \quad (IV.5)$$

From (IV.5), the noise must not exceed 2000 electrons for the required 1% to be obtained. Typical values for the value of the R_s term in (IV.1) is $\approx 1000e$'s which leads to a lower limit on R_w of 15k Ω s.

Recall that the sense wires are to be only of the order of 10 μ m or less. A 5cm long 20 μ m wire would have to be made of material with a resistivity of 10k $\mu\Omega$ -cm.

There is no such wire. The same wire made from stainless steel would have only 785 Ω resistance, so that it is ruled out. Carbon, as graphite, has a resistivity of 1600 $\mu\Omega$ -cm. Simple arithmetic has that a 7 μ m filament of carbon would have $R_w=20k\Omega$, just what is required. Fortunately, just that size of carbon filament was being manufactured by a few companies in America and Japan. The filaments are manufactured by a process of reducing long hydrocarbon polymer molecules in a high temperature furnace. Long strands of filaments can be routinely produced, which when tied together, make an extremely strong, yet lightweight conductor. These strands are imbedded in epoxies and used as very lightweight Faraday cages. The 7 μ m size was just what was being used in industry for this purpose. However, the thought of using single strands of the stuff was uncertain, for the process control required to maintain a uniform diameter was of concern.

A sample of filaments, more than enough to string all of the chambers for the SLD detector, was obtained from American Cyanamid corporation. These fibers had a nickel coating with the properties listed in Table IV.1 The nickel coating made working with the fibers very simple, since they could be individually soldered to tinned traces on printed circuit boards. The nickel could be removed from the rest of the wire in nitric acid, without attacking the underlying carbon.

Carbon Filament Properties	
Diameter	7 μ m
Breaking Strength	450,000 psi
Elongation	1.5 %
Uncoated Resistivity	1600 $\mu\Omega$ -cm
Coated Resistivity	26 $\mu\Omega$ -cm

Table IV.1: Properties of the American Cyanamid Carbon Monofilaments

IV.4 Carbon Fiber Chamber Tests

A single wire chamber was built to approximate the one cell of the final CRID

electron detector, shown in Figure IV.3. An aluminum cathode, 5 cm in length with a 2 mm wire-to-cathode spacing, was fit into a G-10 carrier equipped with Lemo connectors, and enclosed in a Faraday cage. Cables routed the signal to the amplifiers mounted in a nearby NIM crate. The amplifiers were not mounted right at the sense wire for mechanical reasons, not an optimal design, because the capacity seen at the preamplifier input in this scheme is higher than would be desired, although in practice it closely resembled the configuration of the final prototype electronics. The chamber gas was a 70/30 mixture of methane/isobutane, based on a preliminary sample of pulse heights as observed directly on an oscilloscope. The 4.5 cm long carbon fiber used in the test cell was well within the 90 cm critical length limit^[23] calculated at the breaking tension of 13 gm (based upon the fiber geometry at the breaking point in Table IV.1), so that the fiber would not be subject to electrostatic instability. The carbon fiber was strung to 11 grams of tension and its resistance was measured to be $R_w=19k\Omega$.

The single photoelectrons were produced by illumination of the cathode through an 0.7mm slit, which was placed on top of a UV transparent Aclar cell cover. The light source was a commercial DC EPROM erasing lamp. A difficulty with this method was that the aluminum was reflective and produced some background to the signal; to be discussed later.

The low noise amplifier used was the HQV810 from LeCroy Research^[24]. The schematic of the test electronics is shown in Figure IV.4. The signal was shaped to have a 26 nsec rise time and a 60 nsec decay time. From the LeCroy specification sheet, and using $\tau=45$ nsec., the mean time of the shaping filters, the channel resistance of the input FET is calculated to be $R_i \approx 40k\Omega$. The capacitance of the input channels, including effects from the chamber and cabling, were measured to be 23 pF. This quantity was determined by isolating the HQV electronics from the chamber and measuring the noise voltage on an oscilloscope as a function the size of a capacitor which was directly connected between ground and the signal inputs. The noise voltage measured with the chamber attached implied the 23pf. Assuming 20 pF of on chip capacitance, (IV.1) gives ENC = 1588 electrons per channel.

The output of each channel was input to an ADC and to a low gain amplifier stage (a 733 stage with a gain of 7). The outputs of the two low gain stages were summed and became a total pulse height trigger for the data acquisition system. The

Figure IV.3: The single wire bench test cell. The ultraviolet source lamp is indicated by the bulb.

bench data were taken with a LeCroy 2249A ADC, and sent to an Apple II/Plus computer through a CAMAPPLE interface [13]. The data analysis and graphics were done with the Apple.

The gain of the electronics was measured by attenuating NIM level pulses and presenting that voltage to a 5 pF capacitor, connected to the amplifier input. As the amount of attenuation was changed, the input charge varied and calibration curves were generated for both signal channels independently. The calibration points were fit with a quadratic for low ADC counts and a line for larger pulse heights, (> 85

Figure IV.4: Carbon Fiber Bench Electronics. The fiber is connected directly to the input of the first amplifier

counts). The charge gain of the circuit was found to be 2500 fC/fC. Then the two channels were given known charges simultaneously, so that charge division could be simulated. From the width of the measured A_Q distribution, which should be ≈ 0 , the noise per channel was found to be about 1600 electrons per channel, consistent with the expectations based on the LeCroy specifications.

IV.5 Charge Division Performance

With the light source positioned above the slit and over the wire, single photoelectron pulse height spectra were measured at several chamber voltages, an example of which is shown in Figure IV.5. The spectra obtained were fit by a Polya distribution,

$$P(n) = \frac{b^{\bar{n}}}{\bar{n} \Gamma(b)} \left(\frac{n}{\bar{n}} \right)^{\bar{n}-1} e^{-\frac{bn}{\bar{n}}} \quad (\text{IV.6})$$

where n is the pulse height, $\bar{n}$ is the mean of the pulse height, and b is a parameter which was typically about 2.5, first measured here and subsequently confirmed by others. It must be emphasized that the choice of fitting function was motivated by the theoretical work of Alkhazov and others. Polya fits to the data here gave $\chi^2/\text{dof} \geq 1.5$, which implies that the data are not completely consistent with the Polya hypothesis. However, the discrepancy is not inconsistent with Alkhazov, in that the geometry of the proportional chamber used here is certainly an intermediate

case in the sense of his model, where a Polya distribution is expected for a cylindrical geometry. The mean values of these distributions were defined to be the charge gains of the chamber. For low voltages, a simple exponential pulse height distribution was observed. As the voltage was increased, the distribution of pulse heights below the mean fell near the electronic threshold until a slightly peaked distribution was observed. There were always a number of pulse heights which were too small to measure because of electronic thresholds. A peaked distribution indicated that a majority of the pulse heights were being measured, i.e., a peaked distribution implied a "plateaued" chamber. The higher the peak, the more pulses measured, but once peaked, the number of additional electrons detected as the peak of the spectrum increased was at the few-percent level, until breakdown conditions were reached. A plot of mean charge* versus chamber voltage is shown in Figure IV.6. Since the Johnson noise is a constant, independent of pulse height, the S/N ratio is expected to get worse as the voltage on the chamber is lowered. Furthermore, at a fixed voltage the charge division precision should vary. Figure IV.7 shows a scatterplot of pulse height versus AQ . The distribution is broad for small pulse heights and very narrow for large pulses. For the largest pulses, the signals are saturated and the no measurement is possible. The points in the scatterplot away from the central distribution were attributed to photons which were reflected from the cathode and converted at some other point in the cell, away from the collimation slit. There may also have been some multiple electron avalanches at different locations on the wire. However, the ratio of the number of noise hits to signal hits did not vary as a function of pulse height, implying that they were just an artifact of the technique.

For each measurement then, a mean standard deviation, (σ_A) is computed. This number is just the weighted quadrature sum of the widths of five pulse height regions of the scatterplot, from threshold to just before saturation. At a mean gain of 2×10^5 the mean error was calculated to be 1.7% in the wire length, the contribution to the error growing quickly as the pulse height decreased, as expected. Figure IV.8 shows the distribution of (σ_A) as a function of chamber gain. The behavior is as expected, i.e., $\sigma_A \propto \frac{1}{Q}$. A complete understanding of the plot requires a knowledge

* The definition of charge gain used here reflects the measured mean of the observed distributions. Other authors correct the signal to reflect the total chamber gain, which includes the additional signal from the long tail due to the slowly drifting positive ions. However, in this predominantly methane mixture, the correction is only on the order of 1.5. Thanks to J. Vavra for pointing this out.

Figure IV.5: An example of a single electron pulse height spectrum. The mean charge is about 3×10^5 electrons at a chamber voltage of 1700 volts.

of the contributing error sources. There are three sources:

1. The Johnson noise of the resistive carbon filament.
2. The electronic noise from the amplifiers.
3. The distribution of avalanches due to imperfect source collimation. This spread arises from the fact that a 1 mm slit introduces a constant σ_x and hence a constant σ_A . This is indicated by σ_A^0 in the figure.

The contributions of the independent error sources, their quadrature sum, and the data are shown as a function of pulse height in Figure IV.9. The errors have been converted into equivalent position errors along the wire. The plot is for the case of illumination of the central section of the wire, and the total expected noise. The measurement agrees well with the expectation. The errors in this plot are primarily due to uncertainties in the value of the total charge (especially at low total charges), and to uncertainties in the extrapolated value of R_s and in the "averaged" τ value

Figure IV.6: Chamber gain as a function of voltage. The gain is calculated relative to the mean value of the Polya fit.

of the shaping used.

Figure IV.10 exhibits the excellent performance graphically. The upper plot displays the peaks in the A_Q distributions as the source is scanned along the wire, the lower is the linearity measurement for a more complete set of source positions.

Clearly, the carbon monofilaments can be utilized as sense wires in the Cerenkov ring imaging application. The excellent linearity of measurements prove that when the fibers are used in conjunction with low noise electronics, the potential to achieve good resolution accuracy in the coordinate along the sense wire in the ring imaging chamber electron detectors can be realized. Based upon these results, a full-size prototype electron detector for the SLD detector was constructed with the carbon monofilaments.

Figure IV.7: Measured asymmetry as a function of collected charge. The trombone-shaped distribution reflects the variation in resolution with signal.

IV.6 The Prototype Electron Detector

A plane of sixty-four wires was strung with carbon monofilaments. The wire plane formed the sense plane of the "cowfence" described in section III.2. This chamber was used to make the beam test measurements for the final ring imaging tests.

The fabrication procedure was arrived at after several weeks of experimentation. A G-10 frame with tinned copper traces etched at the requisite cathode pitch served

Figure IV.8: Error on the asymmetry as a function of charge. The charge is determined from the mean of the Poisson fit.

as the mounting platform. A fiber was selected and stretched across the frame gap at a tension of 6.5 grams. A conventional lead-tin solder was used to make the electrical contact between the fiber and the traces. Once attached, a hypodermic needle was used to form a bubble of nitric acid at the tip, which was moved by hand along the length of the exposed fiber, removing the nickel coating. The entire procedure was completed under a microscope. The frame was washed with methanol to clean it.

The sense plane was mounted to the cathode piece, the blinding grid above it, and finally, the guide wire plane. The spacing of the various layers determined the necessary voltage drop which would generate a constant electric field inside the "confence" structure. Note that the field inside the electron detector was kept at least $4\times$ that of the photodetector proper. This is essential in order to guarantee that all electrons drifting into the electron counter are collected. The assembly was then mounted on the womb volume above the throat of the drift box.

There were three high voltage (perhaps intermediate voltage is more appropriate in this application) supplies required to define the potentials properly, as shown in Figure IV.1. The guide wires, the last "fence rung", and the cathode all had

Figure IV.9: Measurement error as a function of charge. The contributions are described in the text.

independently controlled voltages. The anode wires were held at ground potential, in order to prevent the necessity of capacitive decoupling at the front end of the charge sensitive electronics. A kapton multi-conductor strip removed the sense wire signals to the electronics which sat at the top of the womb. There was about $2''$ of the conductor prior to the first stage of electronic amplification, corresponding to about 20pf of input capacitance, the same as measured on the bench tests.

The readout electronics for the chamber used the HQV810 chips, one at each end of the wire. Only the first stage of amplification was at the womb. The signal was removed to some electronic racks where the second stage amplification, also the HQV810, was mounted. This shaper stage output was the input signal for the waveform digitizer, the LeCroy model 2241 Image Chamber Analyzer, or ICA. The schematic is drawn in Figure IV.11. There was only enough electronics to instrument twenty wires, which provided 6.35cm in chamber coverage.

Figure IV.10: Linearity as a function of source position. The scale on the upper plot is arbitrary.

With this system then, systematic studies of the prototype ring imaging counter could proceed. Chapter IV details the tests and results of those studies, from simple gas property studies, to the eventual ring imaging beam work.

Figure IV.11: The long drift prototype electronics. The detector sense wire is the input signal for the first stage at left. There are two such circuits per fiber, one at either end

V. Tests with the Prototype CRID

V.1 Introduction

With the advent of commercial ultraviolet lasers, it became clear that a majority of the interesting properties of the long drift scheme were amenable to test in a controlled and systematic way, barring tests of the Cerenkov performance per se. The laser measurements would permit a more complete quantitative understanding of the Cerenkov results, since the interesting properties of the detector which contribute to the Cerenkov resolution, could be thoroughly studied in the absence of Cerenkov light, i.e., in the absence of particle beams. The set of laser tests which established the understanding of the detector are detailed here, culminating with the beam test results.

V.2 Data Acquisition

The battery of tests made with the prototype CRID were completed while tests of other prototype particle detectors for SLD were also in progress. The ensemble of experiments used the common trigger and data processing electronics of the SLAC Test Beam Facility (TBF). Each individual group defined their own trigger locally. The signal was sent to the TBF master trigger, which handled the trigger arbitration for all the groups. When a trigger decision was made, the set of CAMAC crates appropriate to a trigger were readout into a VAX 750. From there, the data was

logged to a tape drive unit or sent to another VAX 780 which supported the on-line event analysis*.

V.2.a ICA Signal Analysis

The LeCroy ICA module was relied upon for the bulk of the work done in the test beam^[25]. The module is built around a high speed CCD chip, which permits real-time voltage sampling of an input signal. For purposes here, the module was used in the uniphase mode, with separate ends of a given fiber serving as the inputs to the two common phase ports per module. In this way, there were two fibers instrumented per ICA. The signals from the chamber electronics vary from 0-2 volts, so that ICA counts were in the 2000-4000 region. The ICA pedestal was set to be commensurate with this. A clock signal was generated externally, using a NIM pulse generator. The frequency was a variable which was set by a dial, and monitored by a frequency counter. The modules require a common STOP signal, which was generated relative to the trigger from the CRID trigger electronics. Because the ICAs provide just 256 time buckets, it was often the case that the ICA time slice was moved by a global offset to a time when the signals were expected. That is, in very low fields, or for very long drift distances, the STOP was moved to provide greater resolution of the signal within the ICA slice. Upon receipt of an event trigger, the generation of further STOP signals was inhibited, and the ICAs read out by the data acquisition VAX over CAMAC for analysis and/or logging.

V.2.b On-line Event Analysis

The on-line event analysis was particularly user-friendly. Histograms of any interesting quantity could be made and reviewed easily, on the fly. The "touch-panel" technique proved most natural, and very flexible. In addition, local hardcopy was available at a "touch".

The algorithm for analyzing the information took two distinct forms. In either case, the signal from the electronics had to occupy at least 4 ICA time buckets. First, for the majority of the laser tests and the Cerenkov work, a simple threshold

* The results presented here represent the successful collaboration of the groups in the TBF, whose cooperation is very much appreciated. In particular, thanks to David Sherden of SLAC, for his unselfish help, at any hour of the day or night.

peak finding technique was used. By defining an ICA count level as a threshold, any voltage signal which surpassed the level was counted as a pulse. The threshold could be set from the "touch" panel. The threshold is defined relative to the pedestal level. The pedestal level had to be defined first. By sampling a few buckets at the beginning and end of the 256 time slices per trigger, or at "quiet" intervals in the 256, a least-squares line fit to the bin values established a pedestal level. The same pedestal technique was used in conjunction with the alternate peak finding algorithm, the derivative method. This method proved slightly more efficient (98% versus 95%) at finding peaks, based on studies with calibration data. The method places constraints on the pulse shape, by taking the local derivative at every time bucket, and from studying the resulting characteristic pulse shapes from calibration data. Multiple peak finding was enhanced with this method, however, and in fact, peaks only 1 time bin apart could be resolved. Single and double hit ICA outputs from the ring imaging data are shown in Figure V.1. The ICA signals from isolated pulses, as in (a), can be used to resolve the close pair of pulses in (c) for pulse height information. Note how the derivative technique finds the pulse pair accurately. The smooth pedestal behavior permits accurate pulse height reconstruction.

The gain calibration for each channel was crucial to the proper event reconstruction. Because of electronic variations from wire-to-wire, and from one wire end to the other, a careful gain calibration was made at repeated intervals. Every electronic channel was instrumented with a precision 5pf capacitor at the input of the front-end HQV810, in parallel to the fiber connection. The single calibration signal, which was a NIM transition that was attenuated with in-line coaxial 50 Ω attenuators, was distributed between four calibration lines on the womb electronics board. The topology was such that even-odd channels on each side of the board could be calibrated independently. The calibration consisted of taking about 100 attenuated NIM pulses per attenuation level. The input voltage per calibration could be easily calculated, but was always measured by oscilloscope at the calibration input as a precaution against non-linearities in the attenuators. Using the threshold technique, the mean and error for the 100 pulses was calculated on-line, and recorded. The resulting data points were fit with lines, with the slope and intercept then kept as a file on-line, so that subsequent data sets were gain corrected. The typical channel gain was 4 V/pC, measured from the peak of the pulse. The effect of the calibration is shown in Figure V.2. The distribution of the asymmetry parameter as a function

Figure V.1: ICA displays for Cerenkov data. The plots (a),(c) are actual event snapshots, the plots (b),(d) the are derivatives. In (a), a single peak is found, indicated by the small arrow. In (c), two peaks are resolved. The large negative signal is indicative of the beam crosstalk.

of pulse height should be reflect the noise in the system only, if the calibration is done properly, i.e., there should be no systematic non-zero asymmetry values on a wire.

Figure V.2: Asymmetry before/after calibration. There should be no variation of the asymmetry with pulse height if the calibration is correct.

V.3 Chamber Efficiency

A measurement of the chamber efficiency for single electrons is given by determining the plateau behavior of it. This is the simplest test for the prototype. The laser spot is focussed to a point in the center of the instrumented wires. The focus is adjusted so that the hits for a number of triggers all occur on a single wire. The focus could be checked by a knife edge test, where a razor blade, mounted on a micrometer, at a fixed position relative to the last focussing element is slowly moved across the laser beam. The number of observed hits is tallied as a function of blade distance. As the blade passes through the beam, the number of observed hits drops. The distance through which the blade is required to move to completely block the beam is extrapolated to the focus, based upon the known focal properties of the optics. Spots as small as $\approx 250\mu$ were obtained, and spot sizes on the order of 0.5mm were the minimum requirement.

The efficiency is calculated based upon the assumption that the laser is depositing charge in the chamber according to a Poisson distribution. The light is attenuated until, on average, there are no observed hits. Then, the intensity is increased until some number for $\bar{n}$ less than one is obtained, where now $\bar{n}$ represents the mean of the Poisson distribution in the limit of large statistics. For these purposes, $\bar{n}$ is calculated by counting the number of null events in N triggers. From

simple Poisson statistics,

$$\bar{n} = -\ln(1 - \frac{n}{N}) \quad (\text{V.1})$$

with an error given approximately by

$$\sigma_{\bar{n}} = \frac{\sqrt{\bar{n}}}{N - n}, \quad (\text{V.2})$$

and n is just the number of triggers in N where any number of hits is recorded. Then, for a plateau study, $\bar{n}$ is plotted as a function of cathode-to-anode voltage, keeping the absolute value of the blind-to-cathode voltages fixed. Such a plateau curve is plotted in Figure V.3. From the curve, which was made without TMAE in the chamber, the knee of the plateau is at $\approx 1750\text{V}$. With the addition of TMAE, the knee is observed to begin earlier by 50V , so that 1700V was used as the cathode-anode potential difference for the majority of the experiments.

Figure V.3: A typical plateau curve for the carbon fiber chamber. The efficiency, $\bar{n}$, is the mean of the Poisson distribution as determined by counting the number of null events in a large number of triggers.

V.4 Drift Velocity Determination

The drift velocity measurement is the first, and simplest, measurement. It requires relating the drift time for produced electrons to the electric field in the

drift box. The electric field is given simply by the voltage difference between the Spellman supply and the potential at the guide wires, divided by the 330M Ω resistor chain value. The drift time was determined by simply counting the number of ICA buckets between the laser trigger and the peak of the signal distribution. Slightly better resolution can be obtained if the prior knowledge of the shape of the pulses is invoked. It is independent of the efficiency. A plot of the drift velocity for the gases used in the prototype is shown in Figure V.4. Using the same technique, the drift velocity in pure CH_4 was measured and found to be consistent with other similar measurements.

Figure V.4. The drift velocity in 70/30 CH_4/C_4H_{10} . The drift velocity was also measured in pure CH_4 and found to be consistent with other measurements.

V.5 Electron Lifetime

Certainly the most important measurement of the drift box environment at any time is the measurement of the electron lifetime. This measurement indicates

the relative purity of the drift gas, and is the staple diet of the ring imaging experimentalist. The measurement is sensitive to the presence of any electronegative pollutants in the gas stream. The source of any such difficulties is not given by the lifetime determination, but the measurement can be made expediently enough to allow systematic tests of the gas system hardware to be performed in an effort to locate the problems. The lifetime was always the first parameter monitored when changes to the plumbing were made, or when new gas mixtures or TMAE supplies were introduced.

The measurement only makes sense when the chamber is stable. The technique is straightforward, simply determine the efficiency either at a given laser location as a function of electric field or as a function of laser position at a fixed electric field. The observed $\bar{n}$ should lie along an exponential curve whose parameter is the lifetime. Then, the parameter can be used to extrapolate to zero time, to extract the actual initial $\bar{n}$, $\bar{n}(0)$. This corrected efficiency is important to the calculation of actual ring imaging detector performance.

A lifetime measurement by electric field measurement is shown in Figure V.5, while the same measurement with laser position is shown in Figure V.6. The extrapolated lifetimes from the two figures is on the order of 50 μ sec at a drift field of 400 V/cm.

The history of the electron lifetime measurements reflects the progress in the understanding of the detector gas environment. In the early days of this work, lifetimes were on the order of microseconds. As time progressed, various calamities occurred which shed light on polluting sources in the gas stream, in addition to direct TMAE immersion studies which helped point to contributing evils. As new things were learned, the lifetime reflected the improvement. As an aside, then, Figure V.7 illustrates the improvement in understanding with time.

V.6 Longitudinal Diffusion

The determination of the contribution of diffusion to the uncertainty in the arrival time of a photoelectron is made by fine-tuning the technique used for measuring the drift velocity. The measurement requires an unfolding of the laser spot size from the distribution of arrival times of photoelectrons generated by the laser

Figure V.5: An electron lifetime by field variation. The electric field values are appropriate to the SLD CRID.

Figure V.6: An electron lifetime by source variation. The distances are relative to the electron detector sense plane.

as a function of drift distance or electric field. In effect, a very high statistics drift velocity measurement is made. With the laser positioned very close to the electron detector (about 5cm was the closest distance of approach), the leading edge of the

Figure V.7: The history of CRID electron measurements. The continued improvement of the lifetime over the years reflects the knowledge gained along the way.

drift velocity distribution is measured. Attributing the shape of the leading edge to spot size and diffusion implies that the intercept should reflect the spot size and the slope is attributed to diffusion. In principle, this measurement is also independent of the efficiency. The results from this measurement are shown in Figure V.8. The best fit line to the points gives a longitudinal diffusion coefficient of $\sigma_l = 180 \mu / \sqrt{\text{cm}}$.

Figure V.8: The longitudinal diffusion measurement. The result is consistent with similar measurements.

V.7 Transverse Diffusion

The measurement of the transverse diffusion coefficient for the operating parameters of the ring imaging counter is a much more tedious one. It requires a good understanding of the electrostatics of the chamber. The measurement of this quantity was made with a simpler proportional chamber, where the efficiency is more easily understood, and more with primitive electronics. But the gas mixture was identical, and the properties of the standard gold-tungsten wire chamber used were understood.

The box geometry chamber, blindless except for the conducting copper walls between anode wires, had a particularly easy electrostatic geometry to simulate. As shown in Figure V.9, the electrostatics of the chamber guarantees that some fraction of the incident photoelectrons are focussed onto the blnds and never get recorded as signals. Such an effect must be folded in with the transverse diffusion effects in order to account for the observed hit distributions. The laser must be scanned in small, $\approx 250\mu$, steps across several detector cells, and the efficiency of each cell measured as a function of laser position. An $x-y$ scanning table mount

was built for the laser which gave the proper precision.

Figure V.9: The contributions to the transverse diffusion measurement. The electrostatic loss is easily modeled.

The scan requires high statistics at each point so that the system was required to be stable over periods on the scale of hours. When this stability was established, the results of scans made across the same cells for the 70/30 mixture at two different drift distances are shown in Figure V.10. The distances correspond to drift times of

4.7 μ s and 17.4 μ s, and these imply that diffusion in the transverse dimensions, σ_t , is $\sigma_t = 211\mu\text{cm}$, for the simulated 80% flat top cell efficiency of Figure V.9. The shift in the mean values of the signal distributions from these tests indicated that the long distance electrostatic distortions amounted to less than a wire spacing.

Figure V.10: Transverse diffusion in the 70/30 mixture. The two curves are averages over several cells at each distance.

V.8 Third Coordinate Tests

With the other contributing factors measured, i.e., the drift velocity, the trans-

verse diffusion, the laser spot size, and the channel gains, the performance of the third coordinate can be studied. For this measurement, a small mirror was mounted on an insulating G-10 arm which was secured to the exterior field cage. The mirror was adjusted to reflect the focussed laser light across the wide dimension of the photodetector volume, parallel to the plane of the electron detector sense fibers. The plane of the beam was 56cm from the fibers. This implied that there would be non-negligible effects due to diffusion which would necessarily smear any z-distributions. However, the diffusion contribution was known, and can be accounted for in the resulting measurements.

The calibration of the fibers was performed as discussed in chapter IV.2.b. When all channels were calibrated carefully, the laser beam was again used to generate ionization in the photodetector volume. Because the spot was extended over several centimeters inside the box, the average spot size was $\leq 1.5\text{mm}$. This could be estimated from the known focal properties of the optics and the observed laser spot on the far womb wall. Ten wires were instrumented for this test, whose results are shown in Figure V.11. The peaks of the hit distributions lay along a line. A typical wire asymmetry is shown in Figure V.11b. When the diffusion and laser spot size are accounted for, it is found that $\sigma_{\text{eq}} \leq 1.7\text{mm}$. This represents an upper limit, and is about a factor of 1.5 worse than bench tests would indicate. It reflects the more noisy environment of the prototype environment. Nonetheless, the knowledge of the conversion point in the photodetector volume is clearly improved, indicating that parallax errors in the ring imaging problem would be greatly reduced.

V.9 Beam Tests

The acid test of a Cerenkov counter is to make and measure Cerenkov light. This requires a source of relativistic particles. The momentum and direction of the particles must be well known, in order that any contribution to the Cerenkov measurement be minimized and understood. Fortunately, at SLAC, the Test Beam Facility was built for just this purpose.

V.9.a Beamline Modifications

For these purposes, SLAC beamline 20-21 was used as the source of ν s and K s, as described in Chapter VIII. There were some notable modifications to the

Figure V.11: The observed hit distribution for 12 wires. The test is of the third coordinate performance for a complete prototype carbon monofilament chamber.

beamline however. The RF separators were not used. As a result, the polarity of the beam chosen to be positive, because the percentage of K^+ s was known to be higher in the + beam than in the - beam. The LASS magnets and chambers were not used in the tests to be described here, but instead, the beam chamber package (as described in Chapter VIII) was used for additional tracking by resurveying them into positions on the beam axis, forward and aft of the womb.

The C_W and C_K Cerenkov counters and the SE scintillator were also used for triggering and tagging the beam particles. The beam flux was controlled by adjusting the lead collimators at the entrance to the experimental hall. The flux was adjusted to count about 0.3 particle/pulse from the 10Hz spill. This helped prevent multiple beam track events from confusing the issue. There were additional scintillators installed forward and aft of the womb, 1 of which served as trigger scintillators, 1 of which was a halo veto. Of the trigger scintillators, one, the A counter, was mounted on the front of the womb just before the thin aluminum beam window on the front door of the womb box. The B.C. and D counters were set up just downstream of the box. The D counter, a 1×1 in² paddle, was used in coincidence with the A counter, a 1×1 in² device, as the main trigger. The B and

C counters, both like A, could be included in the trigger if a tighter collimation was required. They were seldom used.

With this configuration, the best performance was obtained when only the forward set of beam chambers was used to define the candidate track. There was sufficient material inside the womb box to cause perceptible multiple scattering effects which led to relatively poor track parameters if the hits from the aft set of beam chambers were included in the beam fitting. With this procedure, the beam divergence was known to 1mrad.

It was important to insure that only single beam tracks generated triggers. A study of the efficiency of the beam chambers in coincidence with the SE counter produced good single track events when the beam chambers indicated that there was only a single track, and the pulse height from $SE \leq 180$. If the SE ADC cut was not incorporated, there were still several multiple track events.

V.9.b Gas Ring Images

With an understanding of the beam, the essential ingredients of the ring imaging problem are known. The error sources have been determined. What remains is a demonstration of the resolution and separation properties of the complete prototype detector. The beam is predominantly pions, which are used to study the systematics of the gas ring reconstruction. The few kaons in the beam are studied to determine if the expected separation properties are attained. There were 650 good single track events of which 11 were identified kaons, as determined from the beam Cerenkov truth table (table VIII.1).

For these tests, the expected ring from the gas radiator was positioned at 25cm from the electron detector, which is 50cm from the beam. This was done by tilting the spherical $f = 50$ cm focal length Acton mirror, which is mounted on a support fixture 44.7cm downstream from the center of the driftbox, toward the electron detector by 12.5°. This rotation introduces an astigmatism into the focal properties of the mirror, which introduces an error on the order of 0.5mm into the ring reconstruction. Also, the tilt focusses the image in the photodetector where the depth of the drift box is 4.6cm.

The gas mixture consisted of 80/20 CH_4/C_2H_6 for the events shown here. The chamber properties of this mixture are very similar to the 70/30 CH_4/C_4H_{10}

mixture discussed previously. The cathode voltage was -1.7 kV with a drift field of 400 V/cm . The drift velocity was found to be $5.4\text{ cm}/\mu\text{sec}$. The TMAE temperature for these runs was 28° , which from figure III.4, implies an UV absorption length λ of $\lambda \approx 1.9\text{ mm}$. The corresponding electron lifetime was measured to be $\tau_e \approx 50\mu\text{sec}$. There were 20 wires instrumented with electronics, which meant that 5% of the ring was not observed, on average. The ICAs were externally clocked at 25 MHz .

The raw event data are corrected for calibration, using the procedure discussed earlier. The local pedestal was determined from the ICA bin values in the intermediate region between the beam and the gas ring. Using the algorithms discussed, the photoelectron signals for events consistent with pions are found and analyzed for wire number, timing, and third coordinate. The scatterplots shown in Figure V.12 indicate that the third coordinate information has already improved the resolution.

Figure V.12: Third coordinate improvement. In (a), the raw x and raw y values are shown. In (b), the raw x against the third coordinate corrected y values are plotted.

The data were then corrected for the tilt of the mirror, which simply rotates the image into the focal plane of the mirror. The center of the image must be defined

before any resolution parameters can be extracted. The center is determined by translating to a local coordinate system near the center of the Cerenkov ring. The corrected scatterplot is then unfolded in azimuth. The distribution of the local ring radius is histogrammed, and the process is iterated until the narrowest ring distribution is obtained*. The ring radius distribution is plotted in Figure V.13.

Figure V.13: The radius of the Cerenkov ring. The best fit gives $\sigma_r = 2.09 \pm 0.22\text{ cm}$.

The radius distribution shown in Figure V.13 is for the ensemble of gas points. Of interest is the event averaged ring radius, as it will constitute the Cerenkov ring measurement in a working detector. The event averaged radius is obtained by

* In the SLD device, the mirror positions will be aligned to high accuracy beforehand, by laser sighting techniques, in an effort to reduce the amount of time spent with this procedure.

making a $\pm 3\sigma$ fiducial cut about the radius distribution. Then, all photons which fall inside the cut on an event basis are fit with the best circle, taking into account the errors from the ring radius distribution and an error for the center of the ring based upon the ensemble average. The results are shown in figure V.15. The event radius is

$$r = 27.3 \pm 0.71 \text{ cm.} \quad (\text{V.3})$$

which corresponds to an index of refraction in isobutane of

$$n_{C_4H_{10}} = 1.00157 \pm 0.00008, \quad (\text{V.4})$$

and a Cerenkov angle of

$$\theta_c = 54.5 \pm 1.5 \text{ mrad.} \quad (\text{V.5})$$

consistent with the measurements made by a similar group at CERN^[27].

The relative contributions of known error sources can be compared against the experimental result. The expected values are given in Table V.1. The diffusion contributions have been discussed, but in fact, the CH_4/C_2H_6 mixture may be cooler than the CH_4/C_4H_{10} mixture actually measured previously. The wire spacing and bucket errors are just simple "top-hat" errors. The charge division is estimated from the measured characteristics. The multiple scattering and beam divergence errors are found from the tracking. The focal errors are based upon the known mirror parameters. One point should be clarified. The contribution to σ_r from the measured σ_z errors is relaxed because the tilt weights the z -contribution by the tangent of twice the mirror tilt angle.

The total expected error estimated is worse than the observed number. It may indicate that the diffusion is overestimated, or that several of the smaller errors, which are calculated, are pessimistic. An error of no more than 15% is assigned to the calculated value, which is clearly not inconsistent with the measurements.

These π rings should manifest other Cerenkov features. The distribution of the number of hits on a ring should be Poisson distributed about some mean value. The distribution is shown in Figure V.14. Since the mean is high, a Gaussian approximates the distribution well. The fit finds an average of 8.6 photons per ring. The number of observed hits can be corrected for known losses in order to calculate the observed N_0 , as is done in Table V.2. The lifetime correction is for the

Error source	Parameter	Error contribution
X Diffusion	$\sigma_x = 200 \frac{\mu}{\text{cm}}$	1 mm
Wire Spacing	3.175 mm cells	0.92 mm
Bucket Spacing	40 nsec bucket @ 5.4 cm/μsec	0.62 mm
Y Diffusion	$\sigma_y = 200 \frac{\mu}{\text{cm}}$	1 mm
Charge Division	2%	0.6 mm
Z Diffusion	$\sigma_z = 200 \frac{\mu}{\text{cm}}$	0.5 mm
Multiple Scattering	0.2 rad length	0.5 mm
Beam Divergence	1 mrad	0.5 mm
Mirror Aberrations	50 cm focus	0.5 mm
Astigmatism	$\lambda = 2.2 \text{ cm}, 4.6 \text{ cm image depth}$	0.5 mm
Total Point Error	Quadrature Sum	2.2 mm

Table V.1: Contributions to the single point errors from known sources. The expected is larger than the observed.

25cm drift, the wire loss accounts for the incomplete instrumentation of the ring image area, and the photoconversion loss is for the photon loss through the drift box at this TMAE temperature. There is probably an additional correction for the plateau efficiency, because the mixture used is not the best choice, but was the best at the time. The corrected N_0 then is $\approx 85 \text{ cm}^{-1}$ for the 44.5cm radiation length in isobutane.

The kaon events are treated identically. Kaons are selected for single beam tracks according to the truth table of Chapter VIII. There are only a few positively identified K^+ events. The event averaged radius for these events is plotted alongside the same quantity for the π events in Figure V.15. The average radius for K^+ 's was 16.6mm, while 16.4mm is expected for the index of refraction determined from the π measurements. The average number of kaon hits on a ring was measured to be

Figure V.14: The number of hits on the gas rings. The distribution is fit with a Gaussian, a good approximation for a mean of 8.6.

Observed photon yield	-	8.6 ± 0.1
Electron Lifetime	10%	9.5 ± 0.1
Area Loss (no wires)	5%	10 ± 0.1
Photoconversion Leakage	10%	11.1 ± 0.2

Table V.2: Corrections to the observed N_p .

3.3 ± 0.3 , where the expected number would be 3.1. A pair of events, one pion and one kaon, are shown in Figure V.16.

Figure V.16: Event averaged radii for K 's and π 's. Using the measured π resolution implies a 15σ π/K separation power at 11 GeV/c.

As a check on the power of the third coordinate measurement, the z -distribution of hits is shown in Figure V.17. The asymmetry is defined so that values near one indicate photon hits nearest the downstream window of the photodetector volume, the one nearest the radiator mirror. The observed asymmetry does not extend to ± 1 because of the small opening cone of the drift box. The aperture is 4.8 cm at the electron detector, only 4.6 cm at the gas ring location. If the electrostatics is good, the cutoff should reflect this difference, as it does. The fast falloff of the distribution towards $A=-1$ is due to the 10% photon loss through the upstream quartz face of the drift box.

A simple Monte Carlo program was written which generated Cerenkov rings according to the geometry appropriate here. The TMAE absorption length and the smearing in this coordinate were excluded. When compared with the measurements, the best fit value of the TMAE absorption length was $\lambda = 20.5$ mm, and the Gaussian smearing of $\sigma_z = 2.4$ mm. The absorption length compares very well with the 1.9 cm expected. The smearing includes diffusion, charge division, and any

Figure V.16: Two Cerenkov ring images. The upper plot is for a pion, the lower for a kaon. The fiducial curves and the fitted circle are shown.

electrostatic effects. However, the electrostatic effects are unknown. Ignoring it implies a charge division error of 2mm. This is slightly larger than expected. Some residual calibration error probably remains. Nonetheless, the performance is very satisfactory.

Figure V.17: The distribution of $L-R$ measurements for the gas ring data. The Monte Carlo fits a distribution of hits for the present geometry, extracting the parameters shown as the best fit quantities.

V.9.c Liquid Ring Images

The liquid radiator was mounted upstream of the drift box, some 13cm from the quartz face. At this distance, and assuming an index of refraction of ≈ 1.277 , the proximity focussed ring radius would be 17cm. For the instrumented area then, $< 5\%$ of the liquid arc is observed in the region before the beam. Liquid ring data were taken under these rather extreme circumstances, subject to exactly the same detector environment as in gas ring case. The distributions of the reconstructed angles made by reconstructed Cerenkov photons, relative to the center of the liquid radiator cell are shown in Figure V.18. The difference in the raw and third coordinate data show the resolution improvement gained when the third coordinate

information is used. The error on the raw liquid angle (b) of 2.3° is reduced to 0.91° (a), with a mean value of $\approx 52^\circ$. This angle is related to the Cerenkov angle by Snell's law and some simple trigonometry, and gives $\theta_c^{\text{liquid}} = 38.4^\circ$, or an index of $n_{C_6H_{14}} = 1.276$, in remarkably good agreement with the advertised value. A similar calculation finds an expected error of 0.86° on the raw angle, where the largest contribution is due to the multiplicative effect of refraction on the already large chromatic dispersion, which goes like the square of the index. The extrapolation to a full ring from the number of Cerenkov hits observed is 34, which corresponds to an N_0 of 86 cm^{-1} . These liquid ring data give comfortable answers, but are considered to suffer so from the acceptance that firm conclusions will require further tests.

Figure V.18: Angle distributions for liquid rings. In (a) the z corrected data is shown. In (b), the raw data.

Nonetheless, with the accumulated results, the expected separation power of the counter built and tested here is shown in Figure V.19. This resolution curve is, on the scale of particle identification devices, spectacular.

Figure V.19: The resolving power of the prototype. The top plot reflects the high momentum gas performance, the lower the liquid.

V.10 Backgrounds

The question of background phenomena was first raised in chapter III. The problems of dE/dx , photon feedback, and crosstalk were identified as potentially dangerous difficulties worthy of special attention.

The dE/dx problem is primarily responsible for the other two. As mentioned, the difficulties were essentially avoided in these tests, for good reasons of the more pressing interest in N_0 . Some comments can be made here, however, based on prior work from earlier generations of this same prototype. The blinding scheme, as it relates to the beam, appears to have had a dramatic effect on the suppression of photon feedback^[28]. A reduction of 4-5 in the population of feedback hits from the beam was attained with the blinding in place. The improvement was noted by a simple counting of the observed chamber signals with/without the blinds. As mentioned, the issue of beam crosstalk has been addressed, in the form of electronic improvements to the front-end amplifier circuitry. Also, additional capacitive decoupling of the chamber cathodes to ground is being implemented in the SLD detectors. Great improvements are expected.

Using a similar counting scheme, the contribution of feedback photons generated by Cerenkov photoelectrons can be estimated. By examining the scatterplot of the gas ring radius as a function of azimuth, as was done in order to find the image center, a comparison of the number of hits observed in the counter in equal area wire-time slices in regions just prior to and just after the gas ring region can be made. Such a calculation gives an upper limit of 0.07 photons/photoelectron in the present blinding scheme. This is given as an upper limit, because the precise nature of the observed hits is not understood.

The crosstalk was estimated to be $\leq 2\%$. The estimate was obtained by studying ICA plots, looking for correlated signals on neighboring wires from ring signals. The correlations were always negative, with peak values on the order of the number given.

VI. Conclusions and Prospects

The design, construction, and successful operation of a Cerenkov ring imaging counter has been demonstrated. The physical principles of Cerenkov light in all of its aspects, have been used and observed, and shown to be amenable to large scale construction. The intrinsic information in Cerenkov radiation can, for the first time, be utilized to realize a very high performance particle identification device.

The properties of natural materials appropriate for such a device have been studied in order to best take advantage of the full capability. The transparencies, quantum efficiencies, electronegativities, and other properties of a host of materials have been quantified and assimilated. A new proportional chamber strung with carbon monofilaments was successfully tested, yielding enhanced resolution properties for the detector. Finally, test results confirm the expected superior properties of this particle identification technique.

In fact, the results shown here constitute the end of the proof-of-principle stage, the proof that the physical potential can be realized, of the ring imaging detector development work. Since this work was completed, in the spring of 1986, much more has been done. There exists now a complete catalog of the contaminants known to be the primary causes of electron and ultraviolet loss in the device. New technologies for constructing the electrostatic field cages, the blinding grids, and the carbon monofilament chambers are being used for the final SLD construction. Better understanding of quartz, radiator, and chamber gas properties have been gained so that the final product will be even more robust in the difficult, high radiation environment of a colliding beam accelerator experiment. All of these new developments are important and essential extensions of the groundbreaking scientific

work described here.

With the CRID at the SLD, the RICH (Ring Imaging Detector) for the DELPHI LEP experiment at CERN, the E-605 experiment's ring imaging counter at Fermilab, and the Cerenkov detector at the Omega spectrometer also at CERN, the ring imaging technology is mature. There have been proposals made to use such detectors to study atmospheric cosmic ray showers, for CP violation experiments, and in other 4π experiments around the world. Clearly, the high energy physics community recognizes the promising potential of this technique, and all indications are that Cerenkov ring imaging counters have now become an important tool of the experimental physicist.

VII. Introduction to Strange Meson Physics

Studies in light quark spectroscopy had revealed a qualitatively satisfactory picture of the mass spectrum of the light mesons and baryons by the early 1970s. The known states were understood in terms of the allowed quantum numbers predicted by the naive quark model involving up, down, and strange quarks. The states could be arranged in SU(3) and SU(6) multiplets, some of which were complete, and the observed decay branching ratios could be described in terms of a small number of coupling constants and mixing angles. In some instances, simple phenomenology was able to predict the masses and quantum numbers of undiscovered states^[39-43], but no comprehensive model of quark-quark interactions existed that quantitatively reproduced the complete level structure observed in the hadrons.

The "November Revolution" of 1974, i.e., the discovery of the J/ψ particle, and the later discovery of the heavier Υ resonance, opened new windows onto meson spectroscopy^[14,45]. These narrow particles were identified as the ground states of $c\bar{c}$ and $b\bar{b}$ quarks^[16,47], and subsequent investigation revealed entire families of such states, each with the quantum numbers of the photon, $J^{PC} = 1^{--}$. The spectroscopy of the lowest-lying vector states was soon determined and found to be in good agreement with quark model predictions that employed a naive non-relativistic potential to model the forces between the quark and anti-quark^[48]. The potential model is only an approximation, yet it is quite successful in describing the phenomenology of the heavy $q\bar{q}$ systems. Such potential descriptions are very

useful in helping to fold the phenomenological descriptions into a more fundamental description of the strong interaction (QCD).

The systematic study of the heavy-quark mesons in the J/ψ and Υ families at e^+e^- storage rings has heavily impacted the understanding of meson physics^[49]. Pure vector states are directly formed in e^+e^- collisions, and many important new discoveries have been made in this way, e.g., the χ , ψ , D^* , D_S mesons. Studies of the decays of these vector mesons, particularly the radiative decays, have also proved interesting, and have resulted in the observation of the $\eta(1440)$, $f_2(1720)$, and $\chi(2220)$ mesons. Other states are not readily studied in this context. For example, the expected meson states corresponding to large $q\bar{q}$ orbital angular momentum excitations are rarely produced and none of the light-quark radial recurrences has been observed. Such systems are more readily produced in fixed target experiments, which therefore play a complementary role in the study of quark spectroscopy.

Knowledge of the masses and widths of the excited mesons provides several important clues to the nature of the forces binding the constituent quarks into hadrons. The hyperfine structure and radial excitation splittings reflect the relative strengths of the terms in the $q\bar{q}$ Lagrangian, so it is important to understand these splittings as a function of both the radial and orbital quantum numbers. Studies of the decays of these excited states are also valuable. These decays are believed to be strongly affected by the $q\bar{q}$ wave function, so that their variation with the $q\bar{q}$ excited state should provide insight into the dynamics of these systems.

There are other reasons to study light-quark meson systems aside from their intrinsic interest. The speculation that exotic states such as "glueballs" may exist^[50] with masses below 2 GeV/ c^2 , has increased the need to better understand the spectroscopy of the conventional states. One such glueball candidate, the $\eta(1440)$, may have been mistakenly identified as such^[51], but without a comprehensive understanding of the conventional spectroscopy, such ambiguities are difficult to resolve.

The spectroscopy of the light-quark mesons has been experiencing a Renaissance of sorts in the last few years. Several light-quark nonets can now be considered complete with well established states: the pseudoscalar nonet comprising the π , K , η , and η' ; the vector nonet consisting of the $\rho(770)$, $K^*(892)$, $\phi(1020)$, and $\omega(783)$; the tensor states, the $a_2(1320)$, $K_2^*(1430)$, $f_2(1270)$, and $f_2'(1525)$; and more re-

cently the 3D_3 nonet, with the $\rho_3(1690)$, $\omega_3(1670)$, $K_3^*(1780)$, and $\phi_3(1850)$. It is important to continue the pursuit of an experimental understanding of mesons so that the systematics can be better understood and new directions for theoretical work established.

The dense backgrounds in the busy environment of hadronically-produced systems compounds the problem of extracting resonance parameters from scattering data. These backgrounds must be kept small, while the event samples should be as large as possible. The resonance structure can only be extracted by partial wave analysis, which requires that the event sample be unbiased by acceptance effects. These observational difficulties are compounded by the lack of theoretical knowledge about of the systematics of the excited meson states. The non-relativistic potential models are very successful in heavy-quark systems (for example, in charmonium where $(v/c)^2 \approx 0.3$), and yet predicted levels are in good agreement with observation. Naively, these models might not be expected to work as well in light-quark systems, although in fact, they do surprisingly well.

Apace with the latest experimental results, impressive theoretical progress has recently been made in the light-quark meson sector. Several authors have formulated models that insert the dynamical features expected from QCD into a potential, which they then modify to take into account the relativistic effects in light-quark systems [52]. In particular, Godfrey and Isgur (GI) [53], have shown that it is possible to account for the established light and heavy-quark mesons in a theoretical framework that is based on such an approach. Besides reproducing the spectra of the well-established states, their model makes predictions for the masses and decay rates of the underlying spectroscopy.

The strange mesons are an ideal place to study this spectroscopy. The states are observed to be narrower than their non-strange partners, yielding better signal-to-noise ratios. The spectroscopy is expected to be simpler for several reasons: (a) the strange mesons have fewer open decay channels compared with the isovector and isoscalar states, which makes the phenomenological interpretation of the data simpler; (b) pure glueball states are not possible; and (c) the neutral states are free from the ambiguities inherent in the isoscalar-isovector sectors. Most importantly, the excited strange meson states are readily produced in collisions between high-energy kaon beams and stationary proton targets. Attempts to unravel the strange

meson spectrum have had limited success, however. The spectrum is sufficiently complex that it requires sophisticated analyses of large data samples to probe beyond the lowest-lying, leading excited K^* states, so that only a handful of decay channels have been studied.

The primary sources of data on strange mesons have been the two-body $K\pi$ and three-body $K\pi\pi$ final states. Each of these channels has been observed in either the non-charge-exchange reaction $K^-p \rightarrow K^{*-}p$ or in the charge-exchange reaction $K^-p \rightarrow \bar{K}^{*0}n$. Recently, data from the two-body $K\eta$ final state have been analyzed as well.

The non-charge-exchange reactions that have been studied with the highest statistics are $K^-p \rightarrow \bar{K}_0^0\pi^-p$ [54-51], $K^\pm p \rightarrow K^\pm\pi^+\pi^-p$ [52-57], and $K^-p \rightarrow K^-\eta p$ [58]. The two-body $K\pi$ and $K\eta$ channels are allowed decay modes for the natural spin-parity mesons ($J^P = 0^+, 1^-, \dots$), the spectroscopy of which is illustrated in Figure VII.1. To date, the non-charge exchange reactions have provided information on the leading K^* resonances, the $1^- K^*(892)$, the $2^+ K^*(1430)$, the $3^- K_3^*(1780)$, and the $4^+ K_4^*(2060)$, with hints of underlying S -wave. The $K^\pm\pi^+\pi^-p$ final states have provided detailed information on the properties of the unnatural parity strang mesons.

The charge-exchange reactions are dominated by leading K^* production in the final state. Studies of the two body reaction $K^-p \rightarrow K^-\pi^+n$ [59-71] have yielded extensive information on the leading natural spin-parity states up to $J^P = 5^-$, have established the existence of the underlying $J^P = 0^+ K_0^*(1350)$, the $J^P = 1^- K_1^*(1700)$, have resulted in the observation of a $J^P = 1^- K_1^*(1410)$ signal. There have been several studies of the three-body charge-exchange reaction $K^-p \rightarrow \bar{K}_0^0\pi^-\pi^+n$ in which the natural spin-parity $2^+ K_2^*(1430)$, the $3^- K_3^*(1780)$, and the $4^+ K_4^*(2040)$ have been observed. The three-body analyses have also clearly observed strong resonance production in the underlying 1^- amplitude, and have also yielded evidence for a higher lying 2^+ state in the 1.8-2.0 GeV/c² region.

This thesis will report the results from a new study of the $K^-p \rightarrow \bar{K}_0^0\pi^-p$ reaction using data from experiment E-135 obtained with the Large Aperture Superconducting Solenoid (LASS) spectrometer at the Stanford Linear Accelerator Center. The experiment was a high-statistics study of inelastic $K^\pm$ interactions at 11 GeV/c. The entire K^- data set consists of $\sim 113,000,000$ events corresponding

Figure VII.1: The natural-parity spectroscopy. All of the states shown have an allowed two-body decay mode, with the exception of the 1^+ at 1400 MeV/c².

to a sensitivity of ~ 4 events/nb. The experiment yields $\sim 100,000$ $\bar{K}_0\pi^-p$ events in which the $\bar{K}_0$ is observed to decay weakly into two charged π 's.

The analysis of this reaction is performed by selecting the sample of candidate events from the complete data set. This event sample is analyzed in the context of a one-particle exchange model, where the results of this analysis give a description of the $\bar{K}_0\pi^-$ system in terms of the spin-parity amplitudes for the intermediate resonant states.

This half of the full thesis opens with an abridged description of the LASS facility in Chapter XIII, where the hardware and and software of the experiment are

discussed. The selection process that isolates the $\bar{K}_0\pi^-p$ events from the total event sample is described in Chapter IX. The features of the event sample are shown, and the quality of the data discussed. The efficiency for the detection of the $\bar{K}_0\pi^-p$ final state and the subsequent acceptance correction are discussed in Chapter X. Chapter XI presents the acceptance corrected moments of the $K\pi$ angular distribution and interprets their significant features. Chapter XII describes the amplitude analysis, and interprets the results in terms of parameters characteristic of the observed resonant states. The conclusions of this study are summarized in Chapter XIII.

VIII. The LASS Facility

VIII.1 LASS

The Large Aperture Superconducting Solenoid spectrometer was built and operated at the site of the two-mile long Stanford Linear Accelerator in the late 1970s/early 1980s, primarily for the purpose of studying the strong interactions. The apparatus consists of two distinctly different magnetic spectrometers, and is the first of this kind in high energy physics (it provides complete coverage of the full 4π solid angle in the laboratory frame for particle interactions occurring within it). Although it is a forward spectrometer, (i.e., the nuclear target is at rest in the laboratory frame) and in this sense it may be thought of as an electronic bubble chamber. But, where bubble chambers are relatively slow and cumbersome, LASS has all of the advantages of the flexibility and speed of contemporary electronics, and in addition has rather extensive particle identification capability.

The instrument records to a high precision the trajectories of charged particles emitted from interactions in the liquid hydrogen target located near the upstream end of the LASS spectrometer, where a secondary beam of highly relativistic kaons from the Stanford Linear Accelerator is focussed. (Figure VIII.1). The emerging charged particles follow helical paths in the 22.4kG solenoidal field, depositing minute but detectable amounts of energy as they traverse the cylindrical and/or planar proportional chambers which instrument the solenoid volume. Further downstream, additional wire chambers record the particle tracks as they move into the

region of the dipole spectrometer. The dipole serves to momentum analyze low-angle high momentum tracks which are not well-measured in the solenoid. The two spectrometers provide excellent momentum resolution and very high angular acceptance over the nearly all of the available phase space. In addition, two Cerenkov counters, a time-of-flight system, and dE/dx information from the cylindrical chambers, permit a high degree of particle identification in the experiment.

Figure VIII.1: The LASS Spectrometer. The incident beam direction, the solenoid/target, the Cerenkov counters, and the dipole chambers are shown.

There is extensive computing power and event analysis software which supports the spectrometer. A high speed data acquisition network allowed an accumulation of nearly 140 million events to be kept for analysis. A sophisticated tracking program, a thorough Monte Carlo simulation of the detector, and a multiple vertex fitting program were developed which permitted the successful, highly efficient reconstruction of the events from LASS.

A comprehensive review of the LASS facility can be found elsewhere^[2a]. For continuity and convenient reference, a few of the details of the spectrometer which

are particularly relevant to the present physics will be reviewed here. It should be noted here that no prior experiment attempt in a to study the spectroscopy of particle states in hadronic interactions can approach the LASS in experiment acceptance and statistical power.

VIII.2 The LASS Beamline

The 11-GeV/c kaons which comprise the incident beam for this experiment, experiment E-135 at LASS, originate from the secondary debris generated when 21 GeV electrons from the SLAC linac are focussed onto a production target. The spills are 1.6 μ sec long, as frequently as 130 pulses/second. The water-cooled target is composed of 0.25 radiation length of copper backed by 1 radiation length of beryllium. Electromagnetic showers in the target produce a host of secondaries, of which those within a 17.45 ± 3.85 mrad production angle are collected to provide SLAC beamline 20-21[79]. The beam line is shown schematically in Figure VIII.2

Figure VIII.2 SLAC Beam Line 20-21-22 The relative positions of the components are greatly exaggerated.

The optics is such that the beam is focussed four times along the line, at points labeled as F1-F4 in the figure. F1 limits the momentum of the secondaries by collimation to 2.5% in $\delta p/p$. Between F1 and F2, there is a short 0.56 radiation lengths of Pb to help remove any electron contamination, and then two radio frequency (RF) separators. The separators generate electric fields, the strength of which is dependent upon the RF phase at the time of arrival of a charged particle. This momentum dependent dispersion is organized such that by the time the particles arrive at the achromatic focus F2, a clear particle separation has formed. A collimator selects a single species of beam particle. Only one separator was necessary to achieve the beam purity levels essential to the experiment.

A hodoscope of scintillation counters was installed at F3. The instrument consisted of six lengths of scintillating plastic which measured the momentum of individual particles. The counters were overlapped to prevent any dead regions. With this construction, the incident beam momentum per particle was known to 0.5%. At the nominal 11 GeV/c momentum setting of the beam, the individual particle momentum is

$$p_0 = 10.600 + 0.0495n,$$

where n is one of the 11 counter regions. Any uncertainty in the beam tuning also contributed to the beam momentum error. This effect was found to be detectable, and will be detailed later. The error on the beam momentum is estimated to be ≈ 20 MeV/c.

At focus F4, another downstream hodoscope, the $\Theta - \Phi$ hodoscope, measures the spatial position of each beam particle. It consists of the horizontal (Θ) and vertical (Φ) counters, each an array of 12 $1/2''$ foil-wrapped scintillators. This hodoscope also generates a signal for the experiment's trigger logic, in an effort to isolate single beam particles.

Between the $\Theta - \Phi$ hodoscope and the entrance to the LASS experimental hall proper, are 2 threshold Cerenkov counters. Each counter is gas-filled and fitted with a parabolic focussing mirror at the downstream end of the device. The mirror focuses the ultraviolet light from the Cerenkov emission in the gas onto a photomultiplier tube, where a photocathode generates a signal. The C_K device is filled with 40 psia of hydrogen gas, in which only 11 GeV/c pions can radiate. The C_π counter is filled with carbon dioxide, and both pions and kaons produce light. The beam particle is

identified according to Table VIII.1 it was found the RF separation achieved a K/π ratio of 70-80 for the 11 GeV/c K^- beam. Other contaminations were negligible.

Cerenkov Identification			
Counter	π, e	K	p
C_{π}	✓	✓	
C_K	✓		

Table VIII.1: The beam Cerenkov truth table. Possible track hypotheses are indicated

Just prior to the LASS solenoid a set of proportional counters were installed. This beam-chamber package consisted of ten planes of proportional wire chambers (PWC's), each built with sixty-four 20μ diameter gold-tungsten sense wires spaced at 1.016mm. The set was divided into two distinct groups 1 meter apart. The upstream four planes measured x, y, p and e coordinates, where $e \equiv (x - y)/\sqrt{2}$ and $p \equiv (x + y)/\sqrt{2}$. The six downstream chambers measured x, y, x', y', e and p , where the primed coordinates are offset by 1/2 of a wire-to-wire separation relative to the unprimed ones. The beam package provided a beam vector at the $300\mu\text{m}$ level in x, y and 0.3mrad in direction.

A set of 3 scintillators form the last pre-target device. The first, SE, located 5 meters from the solenoid, serves to count both beam particles and halo particles, where halo is any in-time off-axis particle. It generates the start pulse for the time of flight system, and is made especially thick (0.952 cm) to reduce time jitter in the output pulse. The SY and the RING counter comprise the final beam tag. Both counters are aligned on the beam axis, but RING has a 2.54cm hole at its center. RING serves as a beam halo veto. to help reduce stray triggers.

The final beam focus, F4, is an achromatic focus at the end of the liquid hydrogen target which provides for a beam in the target of 0.5cm r.m.s. in each dimension.

The 20-21-22 beamline typically ran at a 90 Hz repetition rate during data

taking, yielding a trigger rate for the experiment of 3-5 beams/pulse randomly distributed in time over the 1.6 μs spill.

VIII.3 The Solenoid Region

The solenoid magnet supports the liquid hydrogen target cell, the inner cylindrical wire chambers, the gap and bore chambers, and is capped on the downstream end by the C_1 Cerenkov counter and the time-of-flight (TOF) hodoscope, as shown in Figure VIII.3 The magnet itself is constructed of four separate superconducting coils, where the first segment holds the cylinder package, and the downstream three pieces support tracking chambers. The operating current of 1600 Amps generates an axial field of 22.4kG, which requires 5000 liters of liquid He from two local refrigeration units at 60 l/hr in normal operation for cooling. The axial magnetic component was measured to be uniform to 1% over the full inner diameter of 185cm for the entire 465cm length. The radial component is negligible on axis, becoming significant near 300cm downstream, because of the downstream opening into C_1 . A fiducial region was defined where the average field variation was less than 3.1%, with a radius of 75cm and length of 320cm.

VIII.3.a The Liquid Hydrogen Target

The liquid hydrogen target cell provides the initial state protons for the experiment. At room temperature, the cell is 85.04cm long and 2.54cm in radius, and shrinks to 84.06cm when full of liquid H_2 . The hydrogen is captured in a mylar bag, which is surrounded by mylar super-insulation, which is inside of an 0.71mm thick aluminum vacuum jacket. It is supported by steel bracing to the upstream mirror plate of the solenoid. The liquid H_2 condition is monitored for temperature and pressure, as the target density is important for determination of the experimental interaction cross section. The 2.54cm of liquid H_2 implies that protons must have $p \geq 225 \text{ MeV}/c^2$ to escape to the vacuum jacket, if they start on axis and propagate radially; this is a consequence of the $0.0716 \pm 0.0004 \text{ g}/\text{cm}^3$ mean liquid H_2 density measured over the course of the experiment.

Figure VIII.3: The solenoid magnet and interior components. The cross-hatched areas represent the solenoid itself, with the target and tracking chambers shown in appropriate relation. The K beam is incident from the left (upstream) end.

VIII.3.b Proportional Wire Chambers

The detectors inside the solenoid are all proportional wire chambers. That is, a particle traversing the device will deposit some of its energy in the form of ionization in the gas; this will be electrostatically focussed to an anode wire, and, in the presence of a high electric field, will generate an electrical signal proportional to the initial charge. These wire chambers are essential tools of the high energy experimentalist. In brief, the electrostatics are defined by the cathodes which form the walls of the chambers. An important consequence of the pulse formation near the anode wire, is that there is also an induced charge on the cathodes. This occurs when the fast, mobile electrons left in the wake of the track arrive near the anode wire and multiply by Townsend multiplication. The remaining collection of positive ions drift away from the anode, giving rise to the cathode signal. By segmenting the cathode, an additional measurement of the amount of charge deposited can be

obtained. Using the induced cathode signals, which typically cover 2-5 strips in the LASS chambers, and the anode signal as well, a three-dimensional reconstruction of the charge deposition locale is generated, and this constitutes a track measurement.

VIII.3.c The Cylindrical Vertex Detector

The heart of the spectrometer is in the vertex package which surrounds the target. Because of the close coupling there, the interaction target is essentially instrumented over the complete solid angle. Albeit that the fastest tracks go forward and are only approximately momentum-measured, if at all, in these chambers, the central and backward regions in the lab are well analyzed with this arrangement. This makes LASS unique in forward electronic spectrometers, providing bubble chamber like reconstruction in most cases.

There are six cylindrical proportional wire chambers, arranged to be coaxial with the target, with central radii measuring from 6.05cm to 49.02cm. The chambers consist of a central anode wire plane of 20μ diameter gold-plated tungsten wires, sandwiched between inner and outer layers of cathode strips. The strips make $\pm 10^\circ$ angles to the anode wires over 100cm of active length in the first four PWC's and $\pm 15^\circ$ angles for 87cm in the outer pair of chambers. The outer pair of chambers are recessed 7cm from the upstream solenoid mirror.

The cathode strips are made of metal foil upon an interstitial support of a double cylinder structure of aluminum-mylar laminate and paper honeycomb sheet. The ends are covered with fiberglass printed circuit boards, where the anode wires are anchored and readout. There is an additional lucite support ring halfway along the cylinder length.⁽⁶⁰⁾

The result of this delicate construction is to reduce the material density in the chamber region to ≈ 0.11 gm/cm³. This is excellent for many of the interesting physics channels, as it reduces the multiple scattering and energy loss contributions to resolution, and hence improves the tracking resolution. However, the physics to be reported here requires good resolution even for the slowest particles which emerge from the target, in particular the slow recoil protons. The density of material in the chamber region will play a significant role in some kinematic corrections which will be detailed in chapter VIII. Suffice it to say here though, that it is impossible to understand the exact consequences of the complicated chamber construction on

a per track basis. That is, a perfect simulation is not realizable. The effects of this will become clear, but as will be demonstrated, the problem can be circumvented in a most satisfying way.

The cylindrical chamber package was not intended to be used to measure energy loss of slow moving particles within them. It is a tribute to the care of the scientists who built and studied the system, that in the end, a very beautiful dE/dx measurement capability was wrested from the cylinder signals, as reviewed in chapter VIII. The signal strength from the cathode foils was used to measure energy loss for slow tracks of 600 MeV/c or less. There were noticeable effects from pulse saturation in the gas-filled chambers, but none the less, correlating the measured ionization from several cylinder measurements per track provided useful particle identification information.

VIII.3.d The Gap, Bore, Plug, and Trigger Chambers

Each of the three downstream segments of the solenoid is instrumented with a gap, bore, plug, and trigger chamber ensemble. These planar tracking chambers generate space points from ionization deposited by particles moving in the helical orbits which they follow in the solenoidal field.

The gap chambers are single planes which cover the breadth of the solenoid inner diameter. There are 758 anode wires (the wires are of the same manufacture as for the cylinder, and all other, chambers) strung 2.032 mm apart. The cathodes are of aluminum-mylar laminate, with 6.86 mm etched strips 1.27 mm apart defining the readout structure. The cathode-anode spacing is 0.508 cm, and each cathode is oriented at 45° to the anode wires, clockwise on the upstream plane, counterclockwise downstream. A central 16.51 cm diameter is covered with a styrofoam backed mylar barrier, which prevents any signal generation in the region. This is essential in order to prevent the busy axial region from confusing the off-axis tracking in these chambers. The plug chambers cover this region of solid angle.

There are actually 5 plug chambers in the spectrometer. There are the three in the solenoid and two additional ones downstream. The plugs are small, only 26 cm by 28 cm, but the 256 anodes at a 1.016 mm spacing, provide good resolution in the beam region. Each chamber has three such anode planes, oriented to provide x , y , and z coordinates, where z is at 35° to the x wires. These chambers worked well

in the intense 5-10 MHz charged particle instantaneous rate environment near the spectrometer axis.

The bore chambers instrument the inner radius of the solenoid. Again, there are x , y , and z planes, with the z wires at 45° to the other two. There are 640 sense wires at a 2.032 mm pitch, with a 0.406 cm cathode-anode spacing. These chambers have octagonal frames, to maximize the tracking coverage inside the magnet sections.

The trigger chambers were originally intended to measure the track multiplicities for the event trigger, in the end, they primarily provided more redundancy in the tracking. Each chamber consists of an anode and segmented cathode planes, where the anodes just provide the electric field here, i.e., they do not provide signals, and as such are widely spaced at 4.064 mm. The upstream, or tracking, cathode, has two concentric rings of pads defined on it, 3 pads in the interior and 128 in the outer ring. The downstream, or trigger, cathode is more coarsely divided into 16 pads in the outer ring, with the like 8 inside. The trigger foil of the most downstream chamber was not read out.

VIII.3.e The C_1 Cerenkov Counter and the TOF hodoscope

All of the physics amenable to study with the spectrometer depends upon the knowledge of the mass of the particles which leave tracks in the spectrometer. In some cases, the kinematics and topology of an event are enough to allow an unambiguous determination of the masses involved. However, there are also important classes of events where additional particle identification is required. For example, very high momentum tracks are difficult to identify without additional information. To this end, a great deal of particle identification hardware was built into the spectrometer. The C_1 Cerenkov counter and the time-of-flight (TOF) hodoscope are two examples. It is important to note that the final states studied here did not require any particle identification information from the dedicated devices. However, for completeness, a brief description of the hardware is included here.

The C_1 Cerenkov counter is inserted into the last of the four solenoid coils, filling the coil aperture, and extending down the beamline. The counter is a conventional gas threshold device, where Freon 114 is the fill gas. In this gas, high momentum pions begin to radiate at ≈ 2.6 GeV/c, and kaon threshold is ≈ 9.2 GeV/c. The total path length in the counter can be 180 cm, which implies that a fast pion

could generate a cone of Cerenkov light of nearly 2cm in radius at the end of the radiating volume.

The volume itself is segmented into 38 cells, arranged into 2 small semi-cylindrical cells on axis, and three concentric layers of 12 azimuthally symmetric cells. The cells do not extend over the full length of C_1 , beginning some 40cm from the upstream end. This design means that a radiating particle may emit into more than one cell, even though the extrapolated track occupies only a single cell. Each cell roughly focuses the incident light onto a mirror which reflects the light at right angles through a Fresnel lens, a collection horn, and onto a photomultiplier cathode. The phototubes are RCA 8850 Quantacons, magnetically shielded with μ -metal and iron jackets, and a set of bucking coils, so as to correct for the strong fringing magnetic fields, on the order of 2-4 kG, of the solenoid.

The performance of the counter was studied using pions from K_0^0 decays. A particle trajectory is extrapolated through the Cerenkov counter in the presence of the solenoidal field. (More on the details of the tracking in sec.VII.8) The expected pattern of Cerenkov emission for the track is calculated and then compared with what is actually observed. A careful study has shown that the average efficiency of the device is 0.975 for particles above threshold, which improves to 0.995 for tracks which radiate exclusively into a single cell.

The TOF hodoscope is part of the time-of-flight system. It is used in conjunction with the SE upstream scintillator to determine the length of time a particle actually took to traverse the solenoid. The hodoscope consists of 24 15° wedge paddles surrounding a set of four on-axis 90° scintillator segments. Each scintillator is 1cm thick, and is optically connected to a lucite light guide consisting of a light pipe, an elbow, and a Winston cone. The guide terminates at the face of a photomultiplier tube, which is, again, magnetically shielded in a fashion very similar to the C_1 tubes. The tubes are readout by both an ADC and a time-to-digital converter (TDC).

In practice, a signal from the SE counter starts the TDCs. When the scintillation light from a passing particle is sensed in a phototube, the attached TDC is stopped and the pulse height recorded by the ADC. The pulse height information is used to correct the observed time (T_{SE-TDC}) for "time walk" in the readout electronics. The resulting resolution on the TOF measurement was ≈ 0.5 ns. The

expected time for a mass hypothesis of π, K, p is calculated based upon the track momentum and point-of-origin, is compared to the measured time, both of which have been corrected for transit time of the light in the scintillator and dispersion. A probability reflecting a given hypothesis is determined. In this fashion, $\approx 3\sigma$ separation is found for π/K up to momenta of ≈ 1.1 GeV/c, and π/p to ≈ 2.5 GeV/c.

VIII.4 The Dipole Region

The downstream half of LASS consists of the 18kG dipole magnet and sets of magnetostrictive and proportional chambers. This spectrometer measures the momentum of particles within 50-100 mrad of the beam axis which have $p > 1.5$ GeV/c. The dipole region is shown in Figure VIII.4.

VIII.4.a The Magnet

The magnet is a water-cooled aluminum pancake, built of four stacked coils, which draws up to 7050A at full field. At this excitation, an 11 GeV/c particle will develop about 900 MeV/c of transverse momentum, due to the 30.1kG-m field integral along the beam-axis. The magnet was operated at full field for better than half of the experiment, and about 2/3 strength for the remainder.

VIII.4.b The Dipole Chambers

The spectrometer consists of are a combination of magnetostrictive chambers and proportional counters located as shown in Figure VIII.4. The magnetostrictive chambers generate the majority of the coordinates for the track finding code in LASS, and the technology provided a cost-effective means for obtaining good accuracy and high efficiency over a large area.

A magnetostrictive chamber is a spark chamber readout with special magnetostrictive wire which efficiently converts the electrical energy of a spark into a travelling acoustic pulse. The spark chamber consists of two planes of woven wire cloth separated by 1cm. One plane is kept at high voltage, the other at ground potential. When an event trigger is received, a high-voltage pulse is generated by a

Figure VIII.4: The Dipole Region. The magnet, chambers, and scintillation hodoscopes are shown

thyratron amplifier, so that a spark is created by any ionizing radiation in the gap. The induced current pulse on the wire cloth transferred as an acoustical signal on a magnetostriuctive wire wand. The wand is connected electrically to a TDC, which records the arrival time of the full set of sparks from the trigger and stores it for subsequent readout.

In detail, the first five chambers are $150 \times 300 \text{ cm}^2$, and the most downstream pair of "super-chambers" are $200 \times 400 \text{ cm}^2$. The upstream wire plane in each chamber determines the x and y coordinates of each spark, while the downstream plane is readout at 30° to the normal, generating $e \equiv x \cos \theta - y \sin \theta$ and $p \equiv x \sin \theta + y \cos \theta$ points. There are polyurethane plugs installed on axis in order to deaden the spark

gaps, otherwise the high beam flux would seriously degrade chamber performance. The downstream chambers are displaced so that the dipole-bent beam will pass through the plugs.

There are two readout wands per chamber, reading out opposite sides of the chamber, except for MSD1 and MSD2 which have only one wand each. Also, each wire plane has two "fiducial" wires placed along the side of the plane. A voltage induced on the wires is detected by the wands. A voltage pulse is generated whenever a spark occurs across the chamber gap. The fiducial wires are also pulsed on the same trigger and are used to calibrate the time scale for each event. The full pulse train for the event is digitized and stored in ANNA modules,^[a] a set of sample-and-hold electronics which is subsequently readout by CAMAC. The arrival time of each pulse is determined relative to a 20 MHz clock pulse train, which translates into roughly 0.27mm spatial sensitivity, for the x and y chambers.

The magnetostriuctive readout scheme must be compensated for stray magnetic fields, which can be large. It was necessary to magnetically bias the wands of the upstream magnetostriuctive chambers. A separate coil of wire is wrapped around each wand to help alleviate the stray field problem. However, it was still observed that the upstream chambers are less efficient than the downstream ones, as might be expected.

VIII.4.c The Johns Hopkins Wire Chambers

The proportional chambers shown in Figure VIII.4 were built at Johns Hopkins University, and consist of a single plane of 512 vertical wires with a 4.23mm wire spacing. The aluminum and polyester mesh cathodes of JHUP and JHDN are built with a 1.27mm gap width, while the JHXY segmented aluminum-mylar foil cathode is of 128 2.62cm wide horizontal strips at a 0.16cm gap width. These counters provided in-time corroboration and position information for dipole tracks, though the JHXY cathode efficiency is only about 50% and of limited use.

VIII.4.d The Scintillator Hodoscopes

The HA and HB scintillator arrays are immediately downstream of the last magnetostriuctive chamber. Each counter consists of two rows of scintillator paddles, with 21 paddles per row in HA and 38 per row in HB. The center paddles

in each array are deadened in the middle so that non-interacting beam particles do not generate signals. Each paddle is 83.82cm long. These counters trigger the spectrometer electronics on events with dipole particles, and provide in-time corroboration of tracks for track reconstruction. A small circular counter of 9.84cm radius, the LP3 counter, over the central hole in the hodoscope array functions as an anti-coincidence counter, helping to ensure that there was no downstream beam tracks in a real event.

VIII.4.e The C_2 counter

The final apparatus in the 20-21 beamline is the C_2 Cerenkov counter. It is also a threshold Cerenkov device used to aid in π/K separation for fast dipole tracks. The counter is divided into 8 segments, and filled with 1 atmosphere of Freon 12 (CCl_2F_2) which serves as the Cerenkov radiator. Particles traverse a 2mm thick aluminum sheet and at least 175cm of gas. Any radiated photons are reflected from mirrors at the rear of the counter, two rows of four mirrors tilted at 10° from the vertical, through light pipes onto photomultipliers, located four on top of and four below the counter.

The pulse height from the tubes is proportional to the number of incident photons from the given mirror, which is of course dependent upon the momentum of the radiating particle. For Freon 112, π threshold is at 2.9 GeV/c and K 's begin to radiate at 10.3 GeV/c. However, the geometry and magnetic environment of the unprotected photomultipliers move the thresholds up, on the order of 10% for the π 's.

VIII.5 The E-135 Trigger

The event trigger for the experiment is minimal, so that the widest variety of interesting physics from target interactions is obtained. All that is required is that a beam particle interact in the target and produce at least two charged secondaries, thus effectively triggering on the full K^-p cross section into charged particles. This physics trigger is supplemented by a few special purpose triggers. The full trigger consists of two pieces, the beam trigger and the cluster logic, which respectively determine the trigger criteria just described.

VIII.5.a The Beam Trigger Logic

An acceptable beam particle is identified whenever the beam trigger logic, shown in Figure VIII.5, determines that an incident beam particle is a kaon, that the particle enters the target, and that there is only one such particle which enters the target within a 20 nsec window. A beam kaon meets the Cerenkov criteria as given in Table VIII.1. A target particle is defined when signals from SE and SY are coincident and they are anti-coincident with RING. The RING veto prevents off-axis "halo" particles from generating false triggers. The $\Theta - \Phi$ hodoscope signals are used as fast timing signals to maintain the 20 nsec trigger window.

In Boolean shorthand, the trigger logic definition is:

$$BT = \Theta\Phi \geq 1 \wedge \Theta\Phi \geq 2 \wedge SE \wedge SY \wedge \bar{R} \wedge C_K \wedge \bar{C}_\pi$$

VIII.5.b The Cluster Logic

The cluster logic determines when a viable target interaction is detected. By counting the number of hits on the anode wires in the innermost two cylindrical PWC's and the first two plug chambers, an event is recorded or discarded. The cluster count from each anode plane forms the logic signals of Table VIII.2. All but the 32 central wires of plug 1 are OR'ed to generate the cluster multiplicities of Table VIII.2, where the central wires are ignored as an additional safety against multiple beam triggers. Conversely, only the central wires of plug 2 are included. This allows for events with a single very forward track to be counted. Every new signal as given in the table includes the requirement that cylinder 2 have at least 1 hit, because cylinder 1 has a large accidental rate due to its trigger proximity which must be kept in check.

VIII.5.c A complete set of event triggers

The full set of physics triggers, $T0_i$, $i=1,4$, and four additional scaled down minority triggers, $T1, T2, T3$, and $T4$ form the complete set of possible event triggers. A minority trigger is one which causes an event to be recorded at only a fraction of the time it actually occurs. These triggers are important for calibration and diagnostic purposes. The trigger definition logic is shown in Figure VIII.6.

Figure VIII.5: The Beam Trigger Logic. This simplified schematic indicates the signals used when defining a valid beam trigger.

Briefly, the separate trigger topologies are:

1. $T0_1$ - ≥ 3 secondaries and 2 cylinder tracks
2. $T0_2$ - 2 cylinder and 0 dipole tracks
3. $T0_3$ - ≥ 1 forward and 1 or 2 transverse tracks
4. $T0_4$ - ≥ 2 forward tracks

Signal	Definition
$PLG_1 \geq 1$	at least 2 plug 1 planes have ≥ 1 hit
$PLG_1 \geq 2$	at least 2 plug 1 planes have ≥ 2 hits
$PLG_2 \geq 1$	at least 2 plug 2 planes have ≥ 1 hit
$PLGS \geq 1$	$PLG_1 \geq 1 \wedge PLG_2 \geq 1$
$CYL \geq 2$	$CYL_2 \geq 1 \wedge (CYL_1 \geq 2 \vee CYL_2 \geq 2)$
$CYL \geq 3$	$CYL_2 \geq 1 \wedge (CYL_1 \geq 3 \vee CYL_2 \geq 3)$
$CYL \geq 3$	$CYL_1 \geq 3 \wedge CYL_2 \geq 3$
$CYL = 2$	$CYL_1 \geq 2 \wedge CYL \geq 3$
$CYL = 1 \text{ or } 2$	$CYL_1 \geq 1 \wedge CYL \geq 3$

Table VIII.2: The cluster logic signal definition. PLG_i is the i th plug chamber, etc...

5. $T1$ - beam decay trigger. The 3-prong decays of the beam kaons, τ decays, serve as calibration monitors for the dipole magnet and the P-hodoscope. Only 1 of 10 recorded.
6. $T2$ - elastic trigger. K^-p elastic scattering events for momentum calibration. Only 1 of 100 recorded.
7. $T3$ - beam phase space trigger. Only 1 of 1000 recorded.
8. $T4$ - random interaction trigger. Only 1 of 50 recorded.

The trigger deadtime is driven by the time needed to clear the track ionization from the spark chamber gaps, about 15msec, but can be larger for very high multiplicity events. On average, the detector was taking data 60% of the time that the beam was being delivered.

Figure VIII.6: The trigger logic. The Boolean definition for each trigger is also indicated

VIII.6 Data Acquisition

The LASS data acquisition system can be resolved into five distinct parts. First and on the front-end, are five device controllers which monitor the five data branches from the spectrometer electronics. There is a DEC PDP-11/04 minicomputer that receives the controller output, an IBM System 7 computer which acts as the communications port between the PDP-11 and the SLAC central computer center, and

finally, the SLAC TRIPLEX computer system*, which logs the data to magnetic tape. The control program for the entire system is the SLAC REALTIME program running on the TRIPLEX.

The device controllers temporarily store the electronic data from the CAMAC data branches of the spectrometer, before being transferred to the PDP. Each of the five CAMAC branches (only four of which are used at any one time) consists of raw event processing electronics, a controller, and a direct memory access (DMA) unit to quickly transfer data to the PDP.

One CAMAC branch is dedicated to the hodoscope, TOF, Cerenkov counter and beamline electronics. Another handles the signals from the PWC anodes. A third controls the cathode foil strip signals, and the fourth services the magnetostrictive chambers through the ANNA system. The fifth branch serves to monitor the voltage settings of some counters and to record the contents of some beam counting scalars.

The PDP-11 reads out the device controllers through the DMA branch driver. On every event, the PDP memory is partitioned to accept the contents of four controllers, which write data to memory in a parallel fashion. When all the data have been transferred, the PDP formats the new data for transfer to the System 7 computer.

The System 7 computer shuttles the PDP data to the TRIPLEX system, over a 0.5km data link. It reads the PDP event buffer, sets up the correct communication protocols with the TRIPLEX, and transfers the information over a connecting coaxial cable.

The actual tape logging and the online monitoring of the spectrometer are at the control of the REALTIME operating system. The tape logging procedure has the highest priority of any task, and must be complete for a previous event before another event can be received from the spectrometer. A background analysis program samples the incoming event stream, unpacks the data, and a crude reconstruction program can generate event information as it arrives from the spectrometer. A host of kinematic quantities, e.g., missing mass, effective mass, etc., could be computed and observed on-line, with the aid of an IBM 2250 graphics display terminal. As

* The TRIPLEX consisted of a 2Mb IBM 360/91, 2 3Mb IBM 370/168s, 10 6250bpi tape drives, 4 IBM 5098 control units, and numerous disk drives and I/O devices.

well, the spectrometer vital signs could be monitored on-line. Various readout efficiencies, chamber efficiencies and multiplicities, etc., could be examined, so that the performance of any part of the spectrometer could be checked at any time.

VIII.7 Spectrometer Alignment and Calibration

The alignment and calibration of the spectrometer is one of the most important jobs of the experimental physicist. Consequently, it is also the most tedious, time-consuming, and difficult of work. Due to the variety of distinctly different devices which must be understood, great care must be taken to insure that the separate positions of each device, relative to one another and to the beamline, are known.

VIII.7.a Spectrometer Alignment

The measurement frame chosen for LASS is defined by the gap chambers, since they could be surveyed *in situ* most easily. The x -positions were recorded, and the anode wire locations were already known from precise bench measurements. The gap cathodes were aligned in two steps. During stretching of the cathode foils onto the support frames, small but measurable distortions of the foils get built into the final device. The distortion is mapped and recorded over the foil surface for every foil, fit to a polynomial, and used in conjunction with the pulse height data from real events to correct for the real $x-y$ position of a hit. Then, single anode signals are matched to the found x value. A chi-squared minimization is made for a large number of hit matchpoints, which allows for small anode displacements and global chamber rotations. The best-fit results from gap chambers 1 and 3 then define the measurement frame.

The other solenoid and dipole chambers are "found" relative to the measurement frame using straight line fits to tracks with the magnets turned off. The cylindrical chambers are located next. The cylinders have unique problems, because they are not exactly round. The small distortions can be understood, by studying the helical tracks found in the spectrometer with the magnets on. Tracks found in the planar chambers are projected back into the cylinders, residuals determined, and the entire process iterated upon until the alignment is satisfactory. The beam chambers are aligned using T3 trigger events. The tracks in the plug and

full-bore chambers are extrapolated into the beam chambers, and the beam chamber positions are adjusted until the residuals are acceptable. These alignments were stable over the course of the experiment, unless repairs to chamber assemblies were required, or the beam polarity was reversed requiring the moving of the downstream chambers to center the beam flux in the chamber dead regions. A set of calibration data was always taken in such cases.

The magnetostrictive chambers were more difficult, and it was found that their alignments could vary on the scale of hours. The affect was traced to instabilities in the wands, and so the wand variations were carefully watched with a re-calibration made whenever the degradation was too severe (≈ 50 times during the experiment).

VIII.7.b Spectrometer Calibrations

Proper event reconstruction requires that the calibrations of all of the experimental apparatus be understood. The readout electronics must accurately obtain true pulse information, the TOF system must give accurate timing measurements, the Cerenkov system must reflect the true momentum dependence of radiating particles, and the dE/dx signals should exhibit pulse height behavior consistent with known energy loss behavior from moving particles.

The ADC electronics used throughout the experiment was subject to run-to-run variations over the length of data taking. These problems were due to pedestal fluctuations in the cases of the Cerenkov and TOF systems, and to both pedestal and gain changes in the SHAM-BADC electronics of the cathode readouts. In the first case, 100 events at the beginning of each run were used to define the pedestals for the run. In the latter case, special calibration electronics was built into the channel amplifiers which permitted a set of calibration constants to be calculated at a greater frequency during runs.

The TOF calibration, while important to the whole of E-135, has no bearing on the present physics and other sources should be consulted for an informed look at the procedure^[78].

The momentum calibration of the spectrometer was also monitored on a run-by-run basis. By finding the absolute calibration in the dipole region, the solenoid spectrometer and P-hodoscope could be studied for consistency with the dipole measurement.

The dipole calibration is obtained by reconstructing K_s decays where at least one of the daughter π 's goes through the dipole magnet. The reconstructed $\pi^+\pi^-$ mass is compared to the real K_s mass, and the difference is converted to a dipole field scale factor, which is 1.007 for full field and 1.004 for half field. Also, elastic and τ triggers are sources independent event sets with good dipole tracks. The elastic events where the scattered K^- is dipole measured, provided a relative calibration between the P-hodoscope and the dipole at 11 GeV/c. The momentum dependence of the dipole calibration was checked using the slower π 's from the τ decays.

A good measure of the run-to-run variations is found by plotting $\delta p_z \approx p_{\text{final}} - p_{\text{beam}}$ versus run number, as seen in Figure VIII.7. Clearly, there are small but distinct variations, which coincide with the occasion of change of polarity of the incident kaon beam, so that the downstream dipole chambers were moved. The details of this plot in between the jumps indicate that there are some small drifts in the system, which were attributed to the P-hodoscope calibration because the $\bar{K}_0$ mass in the dipole reconstructions shows no variation.

An alternative approach to the fine scale δp_z offsets uses the 4-C (momentum and energy constrained) kinematic fit events from the present analysis. For the case where all of the final state tracks are seen, those consistent with the $\bar{K}_0\pi^-\pi$ channel, as discussed in chapter VIII, are fit. One fit is a simple vertex pointing fit (GEOFIT). Events which have acceptable GEOFITs are then kinematically fit. For events with good kinematic fits, the GEOFIT missing mass squared against the $\bar{K}_0\pi^-$ system is histogrammed, as in Figure VIII.8. The abscissa is the calculated missing mass squared minus the proton mass squared. Were the beam calibration exact, there should be no momentum offset. The observed offset corresponds to a 20 MeV/c beam momentum deficit for the plotted 1981 data (the 1982 data had only a 5 MeV effect). These shifts are just at or below the resolution limit on this parameter, but nonetheless these momentum shifts were used in a subsequent refitting of the data sample. It is worth noting here that a similar effect in the $\bar{K}_0\pi^+\pi^-\pi$ analysis was seen in recoil mass from the meson system. The neutron missing mass squared residual was not quite 0. Because the neutron is not seen, any error at the 20 MeV/c level was not detectable, i.e., the resolution on the missing mass for a 1C fit is not nearly so good as for the 4C case.

Figure 3

Figure VIII.7: The longitudinal beam calibration. δp_z against run number is plotted for τ decays

VIII.7.c The dE/dx Calibration

To extract dE/dx information from the cylinder chambers, the pulse height information from the cylinder cathodes must be understood. The seen pulse height, PH , is related to the ionization left by a track by

$$I = f_\theta PH,$$

where θ is the angle defined by the track's momentum vector and the normal to the cylinder surface, and f_θ is the path length correction. If the cylinder package were homogeneous in composition, f_θ should vary as $\cos\theta$. But as seen in Figure VIII.9a, there is an excess of signal as the tracks are more and more forward. There are saturation effects in the chamber which probably account for some of the increase, but the exact source of difficulty is not known. A polynomial quadratic in $\cos\theta$ is used to further correct the pulse height.

Figure VIII.8: Missing mass squared against the K_0K . The proton mass squared has been subtracted. The residual is consistent with an additional 20 MeV beam momentum uncertainty

As the gas gain in the chamber exhibits run-to-run variations, an offline calibration for each cylinder on a per run basis using fast negative tracks is made. Then the β dependence is measured using the protons from elastic events. Proton dE/dx loss in liquid hydrogen is known to exhibit a Landau distribution. However, it was observed here that the variable $(f_0 PH)^{0.3}$ was approximately Gaussian in nature. The β dependence of the average value of the quantity is shown in Figure VIII.9b, and was parametrized with a quadratic.

With these simple parametrizations, a likelihood function for the probability that a track is a π or p is constructed. In practice, the likelihood ratio for the respective probabilities is formed from the product of likelihood ratios of each cylinder hit. The result can be used to differentiate between the two track hypotheses. This procedure was very useful in the $K_0^+ \pi^+ \pi^- n$ analysis^[77]. Here, it is used as a consistency check on the cylinder protons, as will be discussed in chapter IX.

Figure VIII.9: The characteristics of the cylindrical PWC pulse heights. $f_0/\cos\theta$ is plotted in (a) for a typical run. The β dependence of $(PH f_0)^{0.3}$ is plotted in (b).

VIII.8 Event Reconstruction

The large number of signals from the spectrometer must be decoded and analyzed before physics studies may commence. Real charged particle trajectories in the complex magnetic field must be reconstructed, interaction and decay vertices localized, and correct corroboration with the particle identification devices must be determined. An event reconstruction program was written to perform these tasks, which generates a list of associated track parameters, track coordinates, and event topologies. The program as written has four clear procedures. It transforms the raw signals into coordinates, finds particle trajectories, determines the event topology (i.e., identifies vertex candidates), and stores the results on tape in concise event records.

The raw data record from the data tapes is organized into a fixed header and

up to four data blocks. The first block holds the particle identification signals, some strobe signals, and some online scaler counts, such as the elapsed time since the last beam track. The second block contains the PWC anode information, consisting of a list of anode hits and some relative wire timing signals which locate a particular hit with respect to the event trigger. As there may be a cluster of anodes, all "ON" for a given hit, the data were recoded as a single wire hit. The signals are translated into the lab coordinate system and assigned an error equal to $w/\sqrt{12}$, where w is the wire spacing in the particular chamber. The third block contains the cathode signals from the gap and cylinder chambers. The list consists of the strip number and the associated ADC count, which is then adjusted with channel pedestal, gain, and cross-talk corrections from prior calibrations. The strips are grouped into signal clusters, and the pulse height data used to find the charge centroid. Typically, resolutions are on the order of 200μ in the gap chamber y coordinate and 600μ in the cylinder z coordinate. The last block packs the magnetostrictive chamber data consisting of the pulse arrival time at the wands. The fiducial counts are also recorded here and are used to convert the pulse timing signals into spatial coordinates with errors which depend upon the given wand, but average about 600μ .

VIII.8.a Tracking

The beam particles are measured by the beam chamber package. Often, there are several candidate beam tracks per trigger, and care must be taken to insure that only the correct track is found. First, anode clusters are formed in the chambers and converted to spatial coordinates. Then, lists of possible matchpoints are formed for the upstream and downstream beam chambers, requiring that the time slot information for associated clusters be consistent. A beam track candidate is defined when a pair of consistent matchpoints from the upstream and downstream chambers is found. Candidates with the best timing consistency, a satisfactory number of coordinates along the track, and good agreement with the allowed beam phase space are kept. Finally, in-time corroboration with the $\Theta\Phi$ hodoscope and SY are required. Other found beam tracks in an event are usually due to non-interacting beam kaons which can also be found in the downstream dipole chambers. In that case, the relevant matchpoints are not used for any subsequent track finding. Typically, there are 1.5 beam tracks per event in the data.

Track finding in the solenoid region begins by defining precise matchpoints in the PWC chambers. Acceptable track candidates should trace out well-defined helices in the uniform magnetic field. In that vein, three matchpoints, each from a separate chamber, are chosen such that they are consistent with a helical trajectory. This track is used to look for corroborating matchpoints in the other solenoid chambers, where the search is done over areas defined by the estimated errors on the original matchpoints. Based on the total number of corroborating points found, a χ^2 fit to a helix is performed on the coordinates. The track is included on the list of solenoid tracks if the fit is acceptable. The fit allows for multiple scattering, assuming that the track is a π and that the scattering errors are uncorrelated. In the case that the fit is poor, coordinates which contribute more than 12 to the χ^2 can be dropped and the track re-fit. This process is iterated until either too many track points are dropped or the fit succeeds, where a successful fit has a confidence level $> 10^{-4}$. The track finding efficiency is better than 99% in the solenoid region, as determined from Monte Carlo studies.

Tracks in the dipole region are found by first finding line segments in the region downstream of the magnet. A pair of x -coordinates from two similar coordinate planes is found. If the associated spark gaps from each chamber corroborate the x -coordinates, a candidate line segment is defined. The line segment is extrapolated into the other chambers, and if the number of found hits is acceptable, a least-squares fit to a common line is calculated. The fit is kept if the confidence level is $> 10^{-4}$. Line segments in the chambers upstream of the dipole are found by the same process. The full set of line segments from each side of the dipole is then crossed through the center of the dipole magnet. The upstream-downstream pair which is the best match for any track is selected, a first order estimate for the track momentum based upon the angle between the segments is made, and the estimate used to swim the downstream segment back through the dipole to the upstream piece. The process is iterated until the best angle through the field, i.e., the most likely track momentum, is obtained. The momentum errors are based upon the chamber errors and the multiple scattering contribution.

Finally, the dipole and solenoid tracks are joined by finding a solenoid track which matches to a dipole track and refitting the solenoid helix so that its momentum agrees with the dipole measurement. If the fit is successful, the original helix fit results are replaced with the new one. About 95% of the refits are successful, the

failures arising primarily from "kinks" due to particle decay or large angle scatters.

VIII.8.b Topology Recognition

The topology of an event refers to the configuration of vertices reconstructed by the tracking. Usually, a single event can be described by two or more topological patterns, unless specific assumptions about the event are made. As such, care is taken to insure that events can overlap several topological types prior to the more detailed analysis procedures.

There are eight event topologies recognized by the E-135 data processing software. They are:

1. All tracks emanate from a single primary vertex.
2. A number of tracks and a V0 where the V0 is consistent with a K_s .
3. A number of tracks and a V0 where the V0 is consistent with a Λ or $\bar{\Lambda}$.
4. A number of tracks and a V0 where the V0 is consistent with a γ or e^+e^- pair.
5. A number of tracks from a primary vertex and a V^- , where the V^- decays into a negative and a neutral V0. The V^- must be consistent with a Ξ^- or Ω^- , and the V0 with a Λ .
6. A primary vertex and two secondary V0's.
7. A primary vertex with a V^- and a V0.
8. A primary vertex with 3 V0's.

The program also finds scattered tracks, and slow tracks which may spiral several times within the solenoid.

For the $\bar{K}_0\pi^-p$ final state, events are primarily topology 2. However, care is taken to save events which decay close to the primary vertex and are given the topology 1 hypothesis. The $\pi^+\pi^-$ pair which constitute the $\bar{K}_0$ must have $0.436 < m_{\pi\pi} < 0.557$ GeV/c². The primary vertex is found by fitting the extrapolated intersection point of the V0 and the other tracks in the event to a common point based upon the distance of closest approach. Tracks which cause the χ^2 to increase dramatically can be ignored and the event refit. The V0 must also meet a soft decay length criterion. However, any events which are cut are recovered as different

topological candidates, but most are eventually lost anyway with the imposition of tighter final decay length cuts.

VIII.8.c DST Production

Finally, candidate events are written onto magnetic tape as a data summary tape (DST). The compressed event information is stored in a data bank record format, which can be easily added to or subtracted from by subsequent processing steps. This structure allows great flexibility for subsequent physics analysis programs.

VIII.9 Monte Carlo Simulation

Contemporary spectrometers, and LASS is no exception, are very complicated devices. A relativistic particle moving through such a thick ensemble of magnets, liquid hydrogen, and wire chambers of varying composition can suffer a variety of ends so that real tracks, and hence real events, can be confused or lost to the event recognition programs. There are weak decay losses, nuclear absorption losses, losses due to particles being scattered abruptly from their helical orbits, and trigger losses. In order to estimate the magnitude of these problems in a spectrometer, and the resulting effects on the physics, computer programs which model the device in software and simulate real tracks interacting with it, are written. The tracks are generated by the Monte Carlo technique, i.e., large samples of random tracks which populate the software spectrometer in selected regions of phase space. Proper calculations of physics quantities will suffer unless a detailed, well-written Monte Carlo program exists.

More specifically, the $\bar{K}_0\pi^-p$ events relevant to the present analysis can be discussed in terms of four quantities which describe the physics. The mass of the meson final state, $m_{\bar{K}_0\pi^-}$, the four-momentum transfer to the final state proton, t' , and two angles.* These quantities describe the phase-space of the channel. Then, by generating a large number of simulated events which reproduce the behavior (roughly) of these quantities and essentially counting the number of events which

* See the discussion of the Gottfried-Jackson frame in Chapter IX, section 3.

are kept versus those which are input, a detailed schematic of the acceptance of the spectrometer as a function of the given parameters for the $K_0\pi^-p$ channel can be mapped out. This map, also called the efficiency, ϵ , is then used to predict the true number of real physics events which were produced, given those that were actually found. The efficiency is a very complicated function with no a priori analytic representation, and is mapped by using the Monte Carlo technique to generate coordinates corresponding to random events which occupy all of phase-space. These events are treated just as real beam events in every way, so that the same analysis techniques can be used to study them.

VIII.9.a Generating Monte Carlo events

The generated event information is channel specific, so that here the generated event consists of the primary vertex coordinates, the beam 4-vector, the two meson 4-vectors, and the proton recoil. The thrown phase space of the beam kaon follows from a study of T3 triggered events. The beam track is randomly interacted in the target, accounting for nuclear absorption effects in the liquid hydrogen. This first measured vector is then retracked back out of the target, corrected for energy loss and multiple scattering, and the found momentum smeared with the beam chamber errors. This method gives a good reproduction of the beam particles. The other tracks from the T3 events are propagated through the spectrometer to create hits in the downstream chambers. Smeared by the respective chamber errors, these tracks then are used to simulate the effects of secondary beam particles, so that effects in the track finding due to these signals are properly accounted for.

As the interaction products move away from the target, the Monte Carlo tracking simulates the effects of multiple scattering, energy loss, absorption, weak decays, and non-uniformities in the magnetic fields. The details of the spectrometer are precisely reproduced to insure that all simulated tracks closely approximate real ones. Every generated track is swum through the magnetic fields using a Runge-Kutta stepping algorithm to generate track coordinates in chambers until the track has interacted, decayed, ranged out, or exited the spectrometer.

Multiple scattering errors are added whenever material would be encountered in the real spectrometer. Energy loss is simulated, including an approximation to

the "effective" magnetic field felt by a particle as it loses energy over a step size.* Nuclear absorption and weak decays are modeled simply as exponential effects with known characteristic absorption lengths, decay lifetimes, and decay branching ratios. Finally, a field map of the magnetic measurements in the spectrometer was used to generate polynomials to correct for local field distortions.

The responses of the chambers, scintillators, and Cerenkov counters are also modeled. Using the generated track parameters, the resulting coordinates depend upon the angle of incidence of the track and the measured device efficiencies. Cathode simulation is more difficult than anode simulation, because the signal is more spread out and so the chance for overlap of two adjacent signals is greater. The cathode peak reconstruction accounts for this effect. The rest of the chambers and the scintillators are modeled to reflect the effects seen in the real data. The quality and performance of the simulation were checked against real data from n -prong, τ decay, and elastic K^-p events. The elastic events serve as a check on the normalization of the experiment, because the cross section for elastic scattering is well known. The three-pronged τ events from K^- decays were also studied and found to reproduce the kaon mass very well.

In summary, the Monte Carlo correctly simulates the kinematic resolution and reconstruction efficiency of the spectrometer.

* This approximation, while fine for most particles, turns out to be a poor one in the very difficult case of slow recoil protons appropriate for the present physics channel. The effects of this are covered in detail in chapter IX.

IX. The Data Sample

The final sample of candidate $\bar{K}_0^0 \pi^- p$ events was obtained from the much more extensive set of E135 K^- DST's. By a series of event filtering steps, where the topology appropriate to the present channel was the primary selection criterion, a group of $\approx 100,000$ events were eventually chosen from the initial $\approx 113,000,000$ E135 events as true $\bar{K}_0^0 \pi^- p$ events. There were 5 principal tape strips in the series:

1. The Meson Strip (MES). A coarse strip of the DSTs which isolates the low-multiplicity events where the missing momentum is small.
2. The "V0+2 prong" topology strip. Events consistent with one V0 and at most 2 additional prongs are isolated.
3. The FPB strip. Events are selected from the "V0+2" strip which pass a more restrictive set of kinematic and geometrical detector cuts.
4. The MVFIT sample. Constrained kinematic fits are made which select only those events conserving 4-momentum.
5. The Final sample. The final set of cuts which selects the final event sample from the MVFIT results.

These strips are detailed here, including discussions of the cuts used, and examples of the precision, sensitivity, and statistical power of the data are given.

IX.1 The Coarse Data Strips

IX.1.a The Meson Strip

The meson strip is the first reduction step from the full DST tapeset. This selects events which are characterized by a forward-produced multi-particle system accompanied by a seen or unseen recoil state. Events with $2 \leq n \leq 9$ charged tracks are selected, subject to the condition

$$(p_{\text{beam}} - \sum p_i) < 2.0 \text{ GeV}/c,$$

where p_{beam} is the beam momentum and the sum is over all of the found tracks. About 1/3 of the events from the original DSTs survived this strip. This set of strip meson strip tapes contains events with topologies appropriate to, e.g., the $K^- \pi^+ n$, $K^- \pi^+ \pi^- p$, $\bar{K}_0^0 \pi^+ \pi^- n$, and $\bar{K}_0^0 \pi^- p$ channels.

IX.1.b The "V0 + 2 prong" topology strip

This strip takes the meson strip events and selects events with at least 4 charged tracks that can be topologically interpreted as two decay prongs from a secondary V0 vertex downstream of the primary vertex where at least 2 additional charged prongs are found. In addition, the V0 mass is required to be consistent with the mass of the $\bar{K}_0^0$. The "V0 + 2 prong" strip is actually a misnomer, because events with only 1 prong at the primary vertex were also selected. This is to ensure that events in the present channel with a slow recoil proton which cannot escape the target volume due to energy loss are also saved.

This strip achieves a 6-to-1 reduction from the meson tapeset, yielding a set of 58 magnetic tapes of 4,486,000 events*.

* Thanks to Pekka Sinervo, who had the foresight to include the $\bar{K}_0^0 \pi^- p$ events for this analysis in his $\bar{K}_0^0 \pi^+ \pi^- n$ event stripping. At this point, however, the event analyses for the two channels diverge.

IX.1.c The FPB Strip

This is the first strip which defines the sample based upon spectrometer characteristics. Cuts are made to insure that events which interact in the target and which have V0's that are very likely to be $\bar{K}_0$ s are selected. All of the cuts are generous, so as to allow for as large a fraction of true $\bar{K}_0$ -p events as possible to survive. All of the cuts could be studied with a Monte Carlo event sample so that their effects were understood. The cuts are tabulated in Table IX.1. The final two rows indicate the percentage of events with the given topology that survive from the "V0 + 2 prong" strip. The track cut did not change from the input strip, so it had no effect. The primary vertex cut just insures that the interaction occurred in the target. The missing momentum cut is very loose, in order to keep the events with a slow recoil. Note that this cut applies strictly to the V0+1,2 prong topology, so that events with extra tracks which create a larger momentum imbalance are rejected. The decay length cut is to eliminate backgrounds due track pairs from the primary vertex which might fake a $\bar{K}_0$. The d^2 cut serves to remove very poorly measured $\bar{K}_0$ s. The cut is made upon the average of the distance of closest approach of each daughter π to a common $\bar{K}_0$ decay vertex. The missing mass squared cuts are to help ensure that the missing track is compatible with a proton in the V0+1 prong case, and to check that there is nothing missing in the V0+2 prongs. There were $\approx 1,110,000$ events remaining after this set of cuts.

IX.2 The MVFIT results

The fourth stage in the strip process is kinematic fitting. This event fitting is carried out using a multiple vertex least-squares fitting program, MVFIT, which minimizes the chi-squared determined from simultaneously fitting helix parameters to all of the track coordinates. The helices are constrained to pass through a common space-point with other helices in the same candidate vertex. Any secondary vertex is constrained such that the connecting track associates correctly with the primary vertex. This geometrical fit, or GEOFIT, is made for every event.

There are two more fits to every candidate event. The " $\bar{K}_0$ " constrained fit begins with GEOFIT results, but constrains the V0 mass to that of a $\bar{K}_0$ in addition to the pointing constraint. Finally, the overall kinematic fit imposes 4-momentum

Cut	Min	Max	% remaining
Number of tracks	3	9	100.
Primary z vertex (cm)	10.	105.	98.
Missing p_z (GeV/c)	-2.	2.	70.7
$\bar{K}_0$ decay length (cm)	0.5	-	42.
d^2 of $\bar{K}_0$ (cm^2)	0.	0.4	40.5
Missing mass V0+1 prong (GeV^2/c^4)	0.	2.	6.7
Missing mass V0+2 prong (GeV^2/c^4)	-0.5	0.5	21.0

Table IX.1: The cuts which define the FPB strip. The final column indicates the percentage of events which have survived from the initial V0 strip.

conservation on each event, starting with the " $\bar{K}_0$ " fit results. A mass hypothesis is assigned to each track on every fit, a π for the V0 daughters and the negative from the primary vertex, a p for the positive. The kinematic fit occurs in two forms, a 4-constraint (4C) fit which conserves energy and momentum of the V0+2 prong events, where all of the 3-momentum is seen, and a 1-constraint (1C) fit to the V0+1 prong events, where the missing particle is assigned the proton mass. Note that all 4C events are also fit as 1Cs, where the positive track from the primary vertex is ignored. This is useful for making background checks on the 1C sample. All of the fit results are appended to the initial event records.

The set of fits above requires a large amount of time on a mainframe computer system. In order to more cost-effectively meet the processing demands, a pair of 3081/E emulators were built.^[12] A subset of a more extensive emulator "farm", these computers processed all of the fits, using the SLAC 3081 mainframes only minimally as I/O units. This permitted exclusive, around-the-clock processing of the 1 million events, and in the end was a great time savings.

The MVFIT results reduced the sample to 5 tapes of $\bar{K}_0$ -p data. At this point in the analysis, there remained 155,347 events, of which 129,252 were converged 4-C event kinematic fits, and the rest were 1-C fits. The quality of a converged fit

is measured by the confidence level of the result, and the final analysis sample is selected from these events with the imposition of final cuts, including a cut on confidence level.

IX.3 A Frame of Reference

Prior to defining the final analysis sample, some difficulties with the data had to be resolved. As will become clear, there were problems with the MVFIT measurements of recoil momentum which had to be understood in order that the physics make sense. To formally address the problem, it is important to digress here to define the physical system pertinent to an understanding of K^* production in the $K^-p \rightarrow \bar{K}_0^*\pi^-p$ reaction.

The most appropriate reference frame for understanding the scattering of spinless projectiles from an unpolarized target was first discussed by Gottfried and Jackson in 1964^[93]. The $\bar{K}_0^*$ and π^- are considered to be the decay products of a resonant state, produced against a recoil proton in the present case. The incident beam K^- defines the z -direction in the $\bar{K}_0^*\pi^-$ rest frame. Four variables are required to completely describe the final state. The invariant mass of the $\bar{K}_0^*\pi^-$ meson system, $m_{\bar{K}_0^*\pi^-}$, is given by

$$m_{\bar{K}_0^*\pi^-}^2 = (p_K + p_\pi)^2 = (E_K + E_\pi)^2 - (\mathbf{p}_K + \mathbf{p}_\pi)^2, \quad (\text{IX.1})$$

where p is the 4-momentum, E is the energy in the frame, and $\mathbf{p}$ is the 3-momentum in the same frame. The 4-momentum transfer from the beam kaon to the $\bar{K}_0^*\pi^-$ system is given by

$$t_{BK\pi} = q^2 = (p_K - p_{\text{beam}})^2, \quad (\text{IX.2})$$

and $t_{\min}$ is the minimum value of this quantity for a given $m_{\bar{K}_0^*\pi^-}$. Two angles completely describe the spatial distribution of the decays of the $\bar{K}_0^*\pi^-$ system. The Jackson cosine, $\cos\theta_J$, is the polar angle of the $\bar{K}_0^*$ w.r.t. the z -axis. The azimuthal angle, the Treiman-Yang angle φ_{TY} , after S. Treiman and C.N. Yang^[94] who first discussed its relevance to similar problems in 1962, is the angle between the plane of the $\bar{K}_0^*$ and π^- and the plane of the beam and the recoil, also called the production plane. Their use of the angle differs slightly because their choice of frame was different, but in fact it can be shown to be the same azimuthal angle used here^[93]. Figure IX.1 is a graphical depiction of the GJ frame.

Figure IX.1: The Gottfried-Jackson frame. The scattering angles are shown and some relevant kinematic quantities defined.

IX.4 Slow Recoil Protons

For the events which pass the MVFIT selection, Figure IX.2a histograms the recoil proton transverse momentum in the lab frame. Here, the 4-momentum transfers are small, so that $t \approx p_t^2$. A smooth distribution is expected, as opposed to one with a sharp peak seen at ≈ 260 MeV/c. In order to understand the effect, a number of Monte Carlo events were generated, with $m_{\bar{K}_0^*\pi^-} \approx 0.9$ GeV/c², where the same effect was observed.

Protons with the lowest of these momenta typically do not escape the cylinder chamber region. In fact, they have already lost most of their energy to dE/dx loss in the liquid H₂ of the target, for the p_t plotted in the figure is actually the estimated value of the original p_t at the primary vertex. They are heavily ionizing at this momentum, where $\beta = p/E \leq 0.25/M_p \approx 1/4$. MVFIT attempts to account for the effect, assuming some average density for the cylinder chamber environment. But for a highly non-linear process like energy loss, small errors in the accounting of the

Figure IX.2: The recoil proton p_t for 4-constraint events. In (a), the shoulder at low p_t is an excess which indicates that too many protons are found with their true momentum underestimated. In (b), the pseudo-1C fit produces a more natural distribution.

material can lead to non-trivial errors in the reconstructed momentum. Again, the cylinder package is a complicated structure, constructed of aluminum and mylar layers which sandwich a cellular honeycomb of hex-cell plastic. As the protons exit the target, which is more easily modeled because it is of nearly homogenous construction, the protons have lost most of their original momentum. These protons spiral along a quickly decaying helical trajectory, depositing a greater fraction of their energy as they move.

Since the Monte Carlo events exhibited the same behavior, some ad hoc corrections to the cylinder density were made in software, which cured the problem. However, the Monte Carlo simulation incompletely describes the cylinder chamber region, so that a correction which alleviates the problem in Monte Carlo, only modestly improves the situation for real events. It became clear that an adequate correction for additional proton dE/dx loss outside of the target, which worked for

both Monte Carlo and experimental events, was intractable.

Instead, an approach based upon the assumption of a proton recoil was attempted. The transverse momentum of the beam kaon, the $\bar{K}_0$, and the π^- are well measured in the lab frame. Using these quantities to calculate $m_{\bar{K}_0\pi^-}$, a Lorentz boost is made to the center-of-mass frame along the beam direction. This boost preserves the mass and p_t values which were calculated in the lab frame. The p_t of the $\bar{K}_0\pi^-$ system is also that of the proton. The recoil energy is then known, given the proton recoil hypothesis. The energies of the mesons are then recalculated to conserve overall energy and momentum, keeping the angles between each meson and the beam axis the same. These modified momenta are boosted back into the laboratory frame, and the 4-momentum transfer t is recalculated. The process is essentially a 1C fit to the 4C events, or pseudo-1C fit. The scatterplots in Figure IX.3 show the improvement when using this procedure on Monte Carlo events. The changes in the internal angular distributions of the $\bar{K}_0\pi^-$ system due to this technique were below the level of the resolution in these quantities. The p_t distribution for experimental events corrected in this way is shown in Figure IX.2b.

With this improvement, it was decided that the kinematic fit quantities MVFIT would be used to arrive at a final sample of $\bar{K}_0\pi^-p$ data, except that the 4-momentum transfer to the final state would be determined by this pseudo-1C fit.

IX.5 The Monte Carlo Event Sample

Every experimenter would like to build the ideal spectrometer, one with perfect efficiency. Then, whatever secrets that nature would allow one to see could be studied to arbitrary accuracy and all would be clear. Such is not the case, and perhaps it would not be as interesting an undertaking if it were. Yet, it is still essential that any results from a given experiment not be biased by instrumental effects. So the efficiency of the spectrometer, i.e., the acceptance, must be understood.

The Monte Carlo technique is used throughout high energy physics to study the efficiency of any detector. By simulating the response of the device to a model of expected physics, the experimenter is able to learn about, and perhaps correct for, any biases in the instrument which may affect results. Prior to detailing the final event selection, the MC sample is described here. The results from the analysis of

Figure IX.3: The calculated 4-momentum transfer to the recoil proton. In (a), the MVFIT $\delta p_t = p_t^{\text{thrown}} - p_t^{\text{found}}$ is plotted versus p_t . The failure at low p_t 's is evident. In (b), the same quantities are plotted for the assumption of a recoil proton in the pseudo-1C fit described in the text.

these MC data will subsequently be compared to those from the final data sample, as evidence of the high quality of the physics events.

A cursory look at the GEOFIT results from the final data sample provided a basis for defining the distribution of the kinematic variables which are generated for the Monte Carlo events. By examining the $t' = |t - t_{\min}|$ dependence of the data in different, broad $m_{K^0\pi}$ slices, rough estimates for the t' dependence of the Monte Carlo events were obtained. The MC events were generated with exponential distributions in t' , according to

$$\frac{dN_{MC}}{d(t')} \propto e^{-\alpha(t')}, \quad (\text{IX.3})$$

where α , the t' slope, is mass dependent, as tabulated in Table IX.2. The mass dependence of the thrown events was picked to roughly exhibit the distribution in the data, so that the ratio of thrown to real events per mass slice was approximately constant over the full $m_{K^0\pi}$ mass range. Note, however, that this does not imply good sampling over the entire D^{*-} plot. In the first place, because the emphasis here is on the $K^*\pi$ resonance production, there was no MC generation at all for $m_{K^0\pi} > 2.5 \text{ GeV}/c^2$. Above this mass, the cross section for any interesting $K^*\pi$ physics is too small. More importantly, as is well known, the production of high spin K^* 's gives rise to sharply forward/backward peaked angular distributions in the GJ frame, which implies event concentrations near Dalitz plot boundaries. This means that a relatively larger number of MC events must be thrown at more forward and backward angles (in the GJ frame) so that adequate sampling of the acceptance phase space is maintained. Consequently, there were 5 times as many events thrown in the angular region where $\cos\theta_j > 0.7$ relative to the number at other values of $\cos\theta_j$, for events with thrown mass larger than $1.6 \text{ GeV}/c^2$. This increased the statistical precision of the sampling in that interval, but the events were deweighted by 5 to preserve the flat acceptance correction.

Also, slightly different numbers of MC events were generated for 1981 and 1982 data samples, corresponding to full and half field dipole configurations, because the 1982 data set is about 10% smaller than the 1981 set. The differences in the configuration of the spectrometer for the two samples are accounted for in the MC generation. The distribution of thrown MC events with mass, generated in $20 \text{ MeV}/c^2$ bins over the full range, is shown in Figure IX.4. A comparison with figure IX.14 shows the mass dependence to which the MC was tailored.

Mass Region (GeV/c ²)	t' slope	t' range
$0.64 < m_{\bar{K}_0^0 \pi^-} < 1.6$	4.5	$t' > 0.$
$1.6 < m_{\bar{K}_0^0 \pi^-} < 1.9$	3.8	$0. < t' < 0.7$
$1.9 < m_{\bar{K}_0^0 \pi^-} < 2.7$	1.5	$t' > 0.$

Table IX.2: The mass dependence of the t' slope. Figure IX.4: The number of generated MC events vs. $\bar{K}_0^0 \pi^-$ mass. The total sample, 1981 and 1982 simulation, is the unshaded histogram, with the 1982 contribution in the shaded.

IX.6 Trigger Acceptance

The high trigger acceptance of the spectrometer is due to the nearly 4π coverage of the device. The physics trigger, $T0$, acceptance is illustrated in Figure IX.5. The curves are shown for 3 different values of $m_{\bar{K}_0^0 \pi^-}$. The trigger acceptance is large,

$\approx 85\%$, and approximately uniform at all mass values.

Figure IX.5: The trigger acceptance for Monte Carlo $\bar{K}_0^0 \pi^-$ data. The acceptance is plotted as a function of $\cos \theta_j$.

IX.7 The Final Event Sample.

The final event sample consists of the two distinct event topologies, with/without (4C/1C) the seen proton recoil. The two samples are treated differently for the remainder of the event selection process as they are subject to differing backgrounds and resolutions.

IX.7.a The $V0$ s in the Sample

The $V0$'s in the data should be consistent with a K_s . The properties of the K_s are well known. Because it is a spin 0 object, the distribution of the decay angles

of the π 's in the rest frame of the K_s should be flat. More precisely, the angle (θ_h) between the π^+ and the axis (defined by the direction of motion of the K_s in the lab) should be uniformly distributed in $\cos\theta_h$, the helicity cosine. Deviations from the expected mass and helicity cosine distributions indicate very poorly measured events, or backgrounds in the V_0 sample. Also, the proper lifetime of the V_0 's should be consistent with the known K_s proper lifetime ($c\tau = 2.675\text{cm}$). Again, deviations from this value indicate poorly measured or non- K_s events. The 4C events are expected to be relatively background-free, because the 4C constraint is a very powerful one. The 1C events are more subject to backgrounds, as will be shown.

The GEOFIT V_0 mass distribution for all 4C events, and for those where the kinematic fit had a confidence level (CL) better than 10^{-6} , are shown in Figure IX.6a. The choice of $\text{CL} = 10^{-6}$ was based on Monte Carlo event studies where it was observed that 3% or less of the 4C Monte Carlo events were this poorly fit.

The resolution on the V_0 mass is shown in Figure IX.7. The points are the events from the final sample, binned in $1 \text{ MeV}/c^2$ intervals, with $0.84 \leq m_{K_S\pi} \leq 0.94$. The curve results from a fit to the points using two gaussians, whose widths of $4.75 \text{ MeV}/c^2$ and $11.14 \text{ MeV}/c^2$ were determined by identical fits to Monte Carlo events in the same mass region. There is a slight systematic shift of $-0.4 \text{ MeV}/c^2$ from the nominal K_s mass of $0.49768 \text{ GeV}/c^2$, but this $\sim 0.1\%$ effect is compatible with the uncertainty in the magnetic field measurements. A similar systematic effect was also observed in the $K_0^*\pi^+\pi^-$ analysis from this experiment^[77].

The helicity angle cosine of the daughter π 's is a very sensitive indicator of background problems. There is the possibility that a p from a Λ decay could be mistakenly identified as a π^+ , and there is phase space for the resulting invariant mass to reflect into the K_0^0 region. Such events reflect as a band across the $\pi\pi$ invariant mass plot, as seen in Figure IX.8. Also, if the π^- is actually a K^- from the decay of a $K^*(892)$ resonant pair at the vertex, then the kinematics are such that a broad band of events should be apparent in the figure which cross the K_0^0 mass at $\approx \cos\theta_h = 0.3$. There is also phase space for events from higher mass K^* events to reflect into the plot but no effect is seen, and likewise there is no evidence for $\bar{\Lambda}$ background from $\bar{\Lambda}pp$ events.

Figure IX.9 shows the distribution of the helicity cosine distributions prior to

Figure IX.6: The GEOFIT V_0 mass distributions. In (a), the unshaded histogram is for all converged 4C fits, the shaded for those events which pass the final cuts. In (b), the unshaded plot is for all 1C fits, the shaded for those events which pass the final cuts.

acceptance correction. The 4C events are flat after final cuts, with no indication of background contamination. The 1C events have been severely cut prior to acceptance correction. These events are most subject to background, as the GEOFIT distribution in the plot indicates. The 1C events with $\cos\theta_h > 0.8$ are discarded, as well as those 1C events with $t' < 0.02$ and $0.2 < \cos\theta_h < 0.5$. These are both flat cuts and easily simulated in Monte Carlo, so rather than use more sensitive, detailed cuts, these were adopted. Because the most interesting physics in this channel is due to natural parity exchange where most of the cross-section is at high t' , the 1C events are not as important as in the $K^-\pi^+\pi^-$ channel. In this case, statistics were sacrificed for cleanliness, a reasonable tradeoff.

There are cuts on the decay length and proper lifetime of the V_0 's. These cuts help to remove events where the V_0 is not a true K_s , but is something else which reconstructs with an invariant mass within the K_s mass window. Examples are

Figure IX.7: The V_0 mass resolution. The fit is to a the sum of two gaussians, whose widths were determined from identical fits to the Monte Carlo data in the same mass region.

events where a K_s was made from 2 π s from the primary vertex that the reconstruction software made into a downstream vertex, other $K\pi$ overlaps, and V_0 s from false tracks. The lifetime cut was used to help remove backgrounds in the 1C sample and has a very minor effect upon the 4C events. A comparison of the lifetime distribution from the final sample and the MC events is shown in Figure IX.10

IX.7.b Other Cuts

There is a substantial cut on 1C events with no seen dipole tracks. Figure IX.11 shows the missing mass squared distribution for events without (a) or with (b) at least 1 dipole measured π , either from the primary vertex or from the K_s decay; the proton mass squared has been subtracted. It is clear that events without a dipole track are so poorly measured that any Δ signal cannot be discerned, or estimated. The π 's from events of this type are either absorbed in the iron of the solenoid, or

Figure IX.8: $\cos\theta_h$ vs. $\pi^+\pi^-$ invariant mass. These are from a sample of uncut GEOfit tracks. The Λ events are the dark band at the top. There is no evidence for either misidentified $K^-\pi\gamma$ events from the primary vertex or $\bar{\Lambda}$ contamination.

have such a steep angle that they never make it through the dipole aperture. These are just geometrical effects in the spectrometer and are reproducible in the Monte Carlo. Since there is no information in the events, they are discarded.

There is a cut on the missing mass squared, mm^2 , on both 4C and 1C samples. This cut is also sensitive to both background and resolution problems in the channel. In the 4C case, the cut simply removes the most grossly mismeasured events. In the 1C case, there is a real background from $K^-p \rightarrow \bar{K}_0^0\pi^-\Delta^+$ events, where the Δ is not seen. As shown in Figure IX.12, the missing Δ events manifest themselves as a shoulder on the high missing mass side of the proton peak. A cut is made to keep only events with $-0.6 < mm^2 - M_p^2 < 0.2$, except when $t' < 0.02$, when the upper cut is made at 0.1. A fit to the MC 1C missing mass resolution is used

Figure IX.9: The helicity cosine, $\cos \delta_h$, distributions. The unshaded histograms show the uncut GEOTIT distribution, and the shaded are the final MVFIT distributions prior to acceptance correction.

to estimate the expected missing proton distribution, and that fit imposed on the figure as shown. The agreement between data and Monte Carlo is excellent. The Δ is fit as a p -wave Breit-Wigner, and that curve extrapolated under the missing proton peak to estimate the Δ background in the events*. The fit finds that a maximum of 5% Δ contamination is allowed for events with $t' < 0.02$. There is $\leq 3\%$ contamination for $0.02 \leq t' \leq 0.06$.

The set of final cuts for the 4C events is listed in Table IX.3. Similarly, the final 1C cuts are shown in Table IX.4

* The Δ fit is determined by fitting 1C GEOTIT events with $t' > 0.06$, where Δ production is dominant.

Figure IX.10: The lifetime of the V_0 . The abscissa is normalized to the known proper lifetime of the K_s , $c\tau=2.675$ cm.

The t' cut in the tables demarks the separation between the 4C and 1C samples. Any event with $m_{K_0\pi} < 1.6$ GeV/ c^2 and $t' < 0.06$ will be treated in the subsequent analysis as a 1C event. For $m_{K_0\pi} > 1.6$ GeV/ c^2 and $t' < 0.04$, the event is also defined to be a 1C event. All other events are 4Cs, and the 4C fit for the event is used.

One indication of the 4C sample purity can be seen in Figure IX.13. In the figure, the logarithm of the ratio of proton likelihood to pion likelihood is plotted as a function of recoil momentum. The likelihood is calculated from the dE/dx information from the cylinder package. The negligible amount of signal in the left half of the plot indicates that there is no appreciable contamination in the sample

Figure IX.11: The missing mass squared distribution for 1C events. In (a), there are no d -pole tracks in the event, whereas in (b), there is at least one dipole measurement. The proton mass squared has been subtracted.

due to misidentified slow recoil π s. This should be contrasted with the $\bar{K}_0^*\pi^+\pi^-\pi$ sample from this same experiment, when the proton recoil is actually a source of background and a cut had to be made on the likelihood parameter.

The $m_{\bar{K}_0^*}$ distribution after final cuts but prior to acceptance correction is shown in Figure IX.14. The mass distribution is shown here, because the acceptance correction developed for this analysis has a coarse dependence upon the shape of the observed data mass distribution*. The contribution to the total of the 4C and 1C samples is also shown. This first look at the mass distribution provokes two immediate observations. First, the 4C contribution clearly dominates the sample, and that the 1C sample more severely cut. Second, there are obvious concentrations of events in the region of the $K_1^*(892)$, another smaller "bump" near the

* This was done in order to correct for acceptance effects in the mass reconstruction of the reconstructed events. A further discussion of this is presented in chapter X.

Figure IX.12: Δ background in the 1C sample. The Δ is fit as a p-wave Breit-Wigner resonance, and extrapolated under the missing proton peak to estimate the remaining contamination after cuts.

$K_2^*(1430)$, and a hint of some additional structure in the 1800 MeV/c² region. The threshold, intermediate, and high mass ranges appear as featureless, slowly varying, in fact almost flat, distributions. The implication is that the cross section is dominated by resonance production at larger ($t' > 0.06$) 4-momentum transfers, particularly at low mass. As the mass increases, resonance production must decrease very rapidly, as there are few events at high mass relative to the 1 GeV/c² range. This is consistent with previous results in this channel, but is in contrast to the $K^-\pi^+\pi$ and $\bar{K}_0^*\pi^+\pi^-\pi$ data from this and other experiments. More definitive statements as to the nature of the observed mass structure must await the moments and amplitude analysis of chapters XI and XII.

IX.8 Normalization

For a fixed number of observed events, N_{obs} , the cross section is inversely pro-

Cut Description	Min	Max	# cut	% left
Input 4C Data Events	-	129295	-	100.0
4C Conf level	10^{-6}	-	21082	83.7
$\bar{K}_0$ mass (GeV/c ²)	0.478	0.518	7963	77.5
$\bar{K}_0$ lifetime	0	0.2	131	77.4
$\bar{K}_0$ decay length (cm)	0.5	-	2204	75.7
$\bar{K}_0$ d ² (cm ²)	0.	0.1	1185	74.8
Missing mass mm ² (GeV ² /c ⁴)	-0.15	0.15	2641	72.8
t'	0.06	30	1617	71.5
Output 4C Events	-	92472	-	71.5

Table IX.3: The 4C final cut list.

portional to the incident flux. The sensitivity, S , in units of nb⁻¹, of an experiment is an oft-quoted number which is related to the flux, and to N_{obs} by

$$N_{obs} = \epsilon \sigma S, \quad (IX.4)$$

where $S \equiv N_{beam} \rho_p L$, and N_{beam} is the total number of incident probe particles, ρ_p is the target density, and L the target length. The efficiency (acceptance), ϵ , is determined from the Monte Carlo. N_{beam} is the scaled number of beam particles, corrected for absorption in the target, beam decay between the upstream beam trigger package and the target, and various hardware and software inefficiencies. The greatest correction is due to event processing losses, estimated to be on the order of 18%. The net correction factor to N_{beam} is 0.768 ± 0.026 . The target density is just that of liquid H₂, averaged over the E-135 running, and gives $\rho_p = 4.728 \times 10^{22} \pm 0.5\% \text{ cm}^{-3}$. The target length is $81.6 \pm 0.2 \text{ cm}$, which is the cooled, filled length. The product gives

$$S = 4.082 \pm 2.7\% \text{ events/nb}, \quad (IX.5)$$

Cut Description	Min	Max	# cut	% left
Input 1C Data Events	-	36052	-	100.0
$\bar{K}_0$ mass (GeV/c ²)	0.478	0.518	5747	84.1
$\bar{K}_0 \cos \theta_h$	-1.	0.8	825	81.7
$\bar{K}_0$ lifetime	0	0.2	2878	73.8
$\bar{K}_0$ decay length (cm)	0.5	-	184	73.3
$\bar{K}_0$ d ² (cm ²)	0	0.1	473	72.0
No dipole tracks	-	-	6314	54.4
Missing mass mm ² (GeV ² /c ⁴)	-0.4	0.2	11224	23.3
t'	0	0.06	2387	16.7
$\cos \theta_h$ ($t' \leq 0.02$)	0.2	0.5	331	15.8
Output 1C Events	-	5884	-	15.8

Table IX.4: The 1C final cut list.

which is about a factor of 4 larger than any prior K^-p experiment with $\approx 4\pi$ acceptance.

Figure IX.13: The dE/dx signal for the recoil particle. There is a negligible signal on the negative half of the plot, where events consistent with a recoil π would be plotted.

Figure IX.14: The MVfit $\bar{K}_0\pi^-$ mass distribution. The events are not corrected for spectrometer acceptance.

X. The Acceptance Correction

The calculation of the acceptance of the spectrometer in this physics channel is eased due to the simple, but clean, topology of the events, at least in the 4C case. There are always at least three solenoid found tracks, two of which must make a good V0. Because the spectrometer has almost 4π angular coverage, a large fraction of these events are found. Also, three found tracks guarantees that there are almost no triggering losses, except for events with the primary vertex far downstream in the target, where the V0 tracks do not separate in space enough so that both pass through the deadend plug region. Fortunately, the MC mimics the effect so that the consequences are understood. An early estimate of the average efficiency, i.e., just the number of events found by the number thrown, indicated that the efficiency was nearly 50%. With such high average acceptance with all of phase space occupied, it was decided that a weighted events technique could be used to calculate the detailed acceptance.

X.1 Weighted Events

The weighted events method has several advantages over other options, e.g., the maximum likelihood method or the method of moments technique. It is more flexible when examining any interesting feature of the data, because the acceptance does not have to be recalculated for every parameter change. It is fast, because each observed data event has an associated weight, representing the inverse of the efficiency in the same corner of phase space. The weight enters as a simple multiplicative factor when used in any calculation.

While it can be shown that the errors determined with this method for any acceptance corrected quantity are not optimal as they would be in the maximum likelihood case, the fact that the efficiency is so high in the present channel was felt to be good reason to opt for the relative convenience of the weighting technique. *A posteriori*, the results of the full analysis indicate that the physics did not suffer.

The calculation of the weight is straightforward, relying upon known properties of the spherical harmonic functions, $Y_l^m(\theta_j, \varphi_{TY})$. These functions are given in terms of the Legendre polynomials, $P_l^m(\cos \theta_j)$, and exponentials as

$$Y_l^m(\theta_j, \varphi_{TY}) = \sqrt{\frac{2l+1}{4\pi} \frac{(l-m)!}{(l+m)!}} P_l^m(\cos \theta_j) e^{im\varphi_{TY}}. \quad (X.1)$$

Consider the relationship between N_i produced events distributed as $A_i(\theta_j, \varphi_{TY})$, and the distribution of N_f observed events $A_f(\theta_j, \varphi_{TY})$ which have survived the pass through the spectrometer model and data cuts. The produced events could be the generated MC events, or as well, the real physics events. With $d\Omega = \sin \theta_j d\theta_j d\varphi_{TY}$, expand A_f in spherical harmonics.

$$\frac{dA_f(\Omega)}{d\Omega} = N_f \sum_{l=0}^{l_{max}} \sum_{m=-l}^l \alpha_{lm} Y_l^m(\theta_j, \varphi_{TY}), \quad (X.2)$$

where the α_{lm} 's are constants, and l_{max} represents the largest value of the orbital quantum number consistent with the structure in the moments. In strong interaction physics, parity is a good quantum number, so that using the property of the spherical harmonic functions under reflection

$$Y_l^m(\theta_j, \varphi_{TY}) = Y_l^m(\pi - \theta_j, \pi + \varphi_{TY})^* = (-1)^m Y_l^m(\theta_j, \varphi_{TY}), \quad (X.3)$$

(X.2) can be rewritten as

$$\begin{aligned} \frac{dA_f(\Omega)}{d\Omega} = N_f \left\{ \sum_{l=0}^{l_{max}} \{ Y_l^0(\theta_j, \varphi_{TY}) \} Y_l^0(\theta_j, \varphi_{TY}) \right. \\ \left. + \sum_{l=1}^{l_{max}} \sum_{m=1}^l \{ \text{Re} Y_l^m(\theta_j, \varphi_{TY}) \} 2 \text{Re} Y_l^m(\theta_j, \varphi_{TY}) \right\}, \end{aligned} \quad (X.4)$$

with $\langle \text{Re}Y_l^m(\theta_j, \varphi_{TY}) \rangle = \alpha_{lm}$. For the purposes of acceptance calculation, the MC events serve as the produced distribution. The Monte Carlo events were generated flat in the GJ angles so that

$$\frac{dN_l}{d\Omega} = \frac{N_l}{4\pi} \quad (\text{X.5})$$

Then, in any solid angle element $d\Omega$, the acceptance $a(\theta_j, \varphi_{TY})$ is given by

$$a(\theta_j, \varphi_{TY}) = \frac{dA_f/d\Omega}{dA_f} = 4\pi \left(\frac{N_f}{N_l} \right) \left\{ \sum_{l=0}^{l_{\max}} \langle Y_l^0(\theta_j, \varphi_{TY}) \rangle Y_l^0(\theta_j, \varphi_{TY}) + \sum_{l=1}^{l_{\max}} \sum_{m=1}^l \langle \text{Re}Y_l^m(\theta_j, \varphi_{TY}) \rangle 2\text{Re}Y_l^m(\theta_j, \varphi_{TY}) \right\} \quad (\text{X.6})$$

Here, the Y_l^m 's are normalized so that

$$\int_{4\pi} Y_l^m(\theta_j, \varphi_{TY}) Y_l^m(\theta_j, \varphi_{TY})^* d\Omega = 1. \quad (\text{X.7})$$

Then, multiplying (X.4) by $Y_l^m(\theta_j, \varphi_{TY})^*$ and integrating over 4π in solid angle gives

$$\langle Y_l^0(\theta_j, \varphi_{TY}) \rangle = \frac{1}{N_f} \int_{4\pi} Y_l^0(\theta_j, \varphi_{TY})^* \left(\frac{dA_f}{d\Omega} \right) d\Omega = \frac{1}{N_f} \sum_{\text{events}} Y_l^0(\Omega_i), \quad (\text{X.8})$$

and

$$\langle \text{Re}Y_l^m(\theta_j, \varphi_{TY}) \rangle = \frac{1}{N_f} \int_{4\pi} Y_l^m(\theta_j, \varphi_{TY})^* \left(\frac{dA_f}{d\Omega} \right) d\Omega = \frac{1}{N_f} \sum_{\text{events}} \text{Re}Y_l^m(\Omega_i), \quad (\text{X.9})$$

where the sums are over all events in a large sample. These normalized moments, Y_l^m for brevity, characterize any distribution, and will reflect any biases in the spectrometer with the MC events, and physics in the data events, when the acceptance correction is done properly. A useful quantity is the average acceptance, α_{00} , given by

$$\alpha_{00} = \frac{N_f}{N_l}. \quad (\text{X.10})$$

and with this definition, (X.6) becomes

$$a(\theta_j, \varphi_{TY}) = \frac{dA_f/d\Omega}{dA_f} = 4\pi \alpha_{00} \left\{ \sum_{l=0}^{l_{\max}} \langle Y_l^0(\theta_j, \varphi_{TY}) \rangle Y_l^0(\theta_j, \varphi_{TY}) + \sum_{l=1}^{l_{\max}} \sum_{m=1}^l \langle \text{Re}Y_l^m(\theta_j, \varphi_{TY}) \rangle 2\text{Re}Y_l^m(\theta_j, \varphi_{TY}) \right\}. \quad (\text{X.11})$$

The weight assigned to any found event at some Ω_i then is

$$w_i = \frac{1}{a(\Omega_i)}. \quad (\text{X.12})$$

This weight is just a multiplicative factor attached to an event in that particular Ω_i of phase space. It is a measure of the ratio of the number of produced events in Ω_i to the number of found events. When calculating any acceptance corrected quantity, e.g., the number of events in a given mass or t' slice, the weight is an additional factor included in the sum. For example, for a quantity, z , the unnormalized moment $\langle z \rangle$ corresponding to a sample of events with individual values z_i and w_i is given by

$$\langle z \rangle = \sum_i w_i z_i, \quad (\text{X.13})$$

with error

$$\delta_{\langle z \rangle} = \sqrt{\sum_i w_i^2 z_i^2}. \quad (\text{X.14})$$

In the particular case that $\langle z \rangle$ is a spherical harmonic moment, an error matrix E is defined so that

$$E_{lm'l'm'} = \sum_i w_i^2 Y_l^m(\Omega_i) Y_{l'm'}^*(\Omega_i). \quad (\text{X.15})$$

X.2 Mass Weighting

Every found event has an attached acceptance weight, whether it be real physics data or Monte Carlo. In the case of the Monte Carlo data, there are two more multiplicative weights. There is the $\cos\theta_j$ weight which is either 1 or 1/5 as discussed

in the last chapter, and there is a mass weight. The mass weight is used to include the mass resolution of the spectrometer into the acceptance corrected distributions. To this point, only the angular degrees of freedom have been considered in this acceptance discussion. There are also the mass, $m_{K_0^*}$, and t' variables to consider in the complete problem. The complete acceptance of the spectrometer for the mass range of interest requires that the acceptance weights be determined in fine mass and t' bins over the available phase space. Again, the mass sampling is in 20 MeV/c² intervals. The t' distribution is divided in many different bin sizes depending upon the thrown statistics. It varies from bins of 0.02 GeV²/c⁴ in the region of the $K^*(892)$ to bins as large as 0.1 GeV²/c⁴ at high $m_{K_0^*}$ and large t' .

The $m_{K_0^*}$ spectrum of figure IX.14 has large, narrow bumps at low mass, becoming essentially flat at high mass. This implies that in the $K^*(892)$ mass region, the angular distribution of events in the wings of the peak can be noticeably affected by a contribution due to a subset of true peak region events which are relatively poorly measured and reconstruct with an error on the mass value. Because this is a spectrometer effect, it can be simulated and studied with the Monte Carlo sample. It is true that the real data are not flat and preferentially occupy certain regions of angular phase space, but with sufficient MC statistics, the consequences can be understood.

The procedure is straightforward. A mass distribution like that of figure IX.14 was generated in 4 MeV/c² bins. That histogram was fit with 3 Breit-Wigner resonances and a fourth-order polynomial, normalizing the function, $f(m_{K_0^*})$, to unity where the data peaks. The mass weights were calculated independently for the 1981 and 1982 data samples. The distribution of mass weights is shown in Figure X.1. While the thrown $m_{K_0^*}$ lies within ± 10 MeV/c² of the central value of the mass bin, the reconstructed mass will not necessarily do so. Every reconstructed mass value was given the value of the $f(m_{K_0^*})$ appropriate to the center of the 4 MeV/c² bin. Then, the complete set of unnormalized moments were calculated in the 20 MeV/c² $m_{K_0^*}$ and t' bins for all $Y_l^m(\theta_j, \varphi_{TY})$ with $l \leq 10$, $m \leq 4$, with the weight factor included. That is, for all reconstructed masses within 100 MeV/c² of the central throw, $m_{K_0^*}$, the quantities

$$\langle ReY_l^m(\theta_j, \varphi_{TY}) \rangle = \sum_i \frac{N_i}{N} w_i^m w_i^{\cos \theta_j} ReY_l^m(\theta_j, \varphi_{TY})_i \quad (X.16)$$

were accumulated in 20 MeV/c² bins. This was repeated for the entire thrown sample. So, for a given thrown mass bin, the contributions to the unnormalized moments from the events in the 10 nearest thrown bins which reconstruct with masses in the current bin are summed to give the new acceptance moments.

Mass MC Weight

Figure X.1: The mass weight distributions. The points are the weights in the 4 MeV/c² bins, the histogram the average value in the 20 MeV/c² thrown bins.

The mass resolution determined from the MC sample is shown in Figure X.2 as a function of $m_{K_0^*}$ mass. Clearly, the resolution is much finer than the sizes of any relevant physical structures in the respective mass domains, based on Particle Data Book values. The MC events in the figure are for the 1981 data only, but an

essentially identical plot can be made for the 1982 sample. As it happens, the new resolution is only slightly better than for the MVfit value alone. A great improvement is observed when the same technique is applied to the GEOFIT sample. This is taken to be evidence to support the claim that the fitting procedure is understood. Any advantage from the weighting becomes less important as the mass increases, because the strength of the production of resonances decreases there.

Figure X.2: The mass resolution. The points are for the combined 1C and 4C events.

A slight modification to the weighting must be made in the $m_{K_0\pi^-} \approx 1.6$ GeV/c² and $m_{K_0\pi^-} = 1.9$ GeV/c² regions where the t' dependence of the thrown distributions changes. For 100 MeV/c² on either side of the two breakpoints, a renormalization weight is included for events which flow across the boundary in one direction or another. For the same number of total events thrown, the occupancy of a given t' bin must be corrected for the t' slope difference across the respective boundary. That is, every event which flows across the boundary must be corrected to represent the number which would flow if the t' slopes were the same. The correction removes any discontinuities in the weighting spectrum.

X.3 The Acceptance Moments

The normalized acceptance moments from the generated MC sample, calculated according to (X.8) and (X.9) are the important quantities. It was observed that the t' dependence of the normalized moments was flat in the separate 1C and 4C intervals in each mass bin. Consequently, linear least-squared fits to the distinct 1C and 4C t' regions were made, and the results used to calculate the value of any normalized moment.

Differences between the calculated values for the 1C and 4C samples are expected, because of the disparate cut tables for each set. The differences are manifest in Figure X.3 and Figure X.4. These two figures are representative. The a_{00} plots for the 4C samples indicate the almost flat, $\approx 50\%$ acceptance of the spectrometer in this $K_0\pi^-p$ channel. The acceptance value varies slightly where the resonance structure in the data is prominent, dipping slightly at the peaks and rising in the wings of the peaks. This is consistent with the expectation that more events will flow out of the peak region than will flow in, due to the nature of the $K_0\pi^-p$ production. The 4C acceptance is nearly constant in mass. The 1C acceptance, on the other hand, decreases monotonically above ≈ 1.4 GeV/c². Also, the plot of $Y_{K_0}^{20}$, which attains the largest non-zero values of any acceptance moment, has a very clear mass dependence in the 1C case, whereas the 4C moment is essentially 0. until ≈ 1.6 GeV/c², when it develops a small negative value. The reasons for the variations are understood. In the 1C case, there is a depletion of events in the central $\cos\theta_s$ region as $m_{K_0\pi^-}$ increases, hence the positive $Y_{K_0}^{20}$. This is due to a loss of events in the transverse plane, i.e., events which have no found dipole tracks. The effect of this essential cut grows with mass, which is to be expected, because of the resulting larger q -values of the decay pseudoscalars. In fact, a hole in the 1C acceptance is created for $m_{K_0\pi^-} \geq 1.6$ GeV/c². Hence, the 1Cs are not used in the amplitude analysis above that region. This is considered to be unimportant for this analysis, because the interesting physics in this channel populates the phase space with $t' \geq 0.08$, whereas the 1C region is $t' \leq 0.04$ at these masses. In the 4C case, the decay angular distribution in the GJ frame is more strongly peaked in the forward/backward directions, as the mass increases. The cut on the decay length generates a slight asymmetry in the moments, because the backward K_0 s are more affected than the forward ones.

Figure X.3: The dependence of a_{00} on $m_{K^0\pi}$. The 4C calculation, at $t' = 0.15$, is shown in (a), with the 1C calculation, at $t' = 0.03$, in (b).

This technique can be checked for self-consistency using the Monte Carlo data again. If the set of normalized moments which have been calculated are sensible, then using them to correct the round Monte Carlo events should reproduce all of the original characteristics of the thrown sample. In other words, consider the found Monte Carlo events as data to be acceptance corrected. Insert the normalized moments in the acceptance calculation, using (X.8), and (X.9) in (X.11) to get

$$a(\theta_j, \varphi_{TY}) = 4\pi a_{00} \left\{ \sum_{l=0}^{l_{max}} Y_N^{l0} Y_l^0(\theta_j, \varphi_{TY}) + \sum_{l=1}^{l_{max}} \sum_{m=1}^l Y_N^{lm} 2Re Y_l^m(\theta_j, \varphi_{TY}) \right\}. \quad (X.17)$$

Then, sum over the same set of Monte Carlo events which were used to generate the

Figure X.4: The dependence of Y_N^{20} on $m_{K^0\pi}$. The 4C calculation, at $t' = 0.15$, is shown in (a), with the 1C calculation, at $t' = 0.03$, in (b).

normalized moments, evaluate the spherical harmonics, calculate $a(\Omega)$ and so the acceptance weight w . Again, at a given t' , the value of the normalized moments is based on the linear fit in the mass bin. The mass dependence of the acceptance is just an interpolation between central values of the normalized moments in contiguous 20 MeV/c² bins, until $m_{K^0\pi} = 1.6$ GeV/c², when the relative lack of sampling in the small binning requires that the averaging take more neighboring bins into account. Using the weights, the acceptance corrected kinematic distributions were compared with the thrown distributions, using the procedure described, until the thrown and acceptance corrected distributions agreed.

X.4 The Acceptance Corrected Data

The acceptance corrected data are generated from the MVfit sample using the acceptance moments. For every data event, the acceptance weight is calculated based upon the $m_{\bar{K}_0\pi^-}$, t' , $\cos\theta_{\pi^-}$, and φ_{π^-} exactly as for the Monte Carlo, i.e., according to (X.17). Any events with $t' \geq 0.7$ are attached the weight given as determined from extrapolating the linear fits to the higher t' values.

There is an additional correction to be made which applies for the data events only. Because of the inadequate representation of the cylindrical chambers in the spectrometer which gave rise to the t' problem in the recoil protons, the Monte Carlo does not reproduce the observed t' dependence of the event sample in the 1C/4C overlap region. The effect is clear in Figure X.5. The problem is the same, the Monte Carlo underestimates the amount of material in the cylinders, so that the 4C events will become important at smaller values of t' . The problem is very localized to the 1C/4C t' boundary at low mass, with a gentle diffusion in t' observed as the mass increases. There are no unusual acceptance effects in either 1C or 4C samples as the boundary is crossed. Then, since the reasons for the effect are clear, a simple correction is made to correct the weighted cross section with a polynomial in the overlap region, so that there is no discontinuity across the boundary in the cross-section. This was checked against the number of $\bar{K}_0\pi^-p$ events found in the 1C events with $t' \geq 0.06$ (in the $K_1(892)$ region) and there was no disagreement. These additional 1C events were kept in the sample for this purpose, but were excluded from the physics analysis because the background Δ production clearly dominates the signal in t' low events.

The acceptance corrected mass distributions are shown in Figure X.6 and Figure X.7. The $\bar{K}_0\pi^-$ histograms show prominent peaks at the 2 lowest mass K^* 's, and also indicate just how dominant the high t' production is. The baryon-meson mass plots indicate the prominent sources of background to the present physics analysis, the namely the three $p\pi$ structures and a low mass bump in the $p\bar{K}_0$ mass plot.

The acceptance corrected data are shown in the Dalitz plot of Figure X.8, the Chew-Low plot of Figure X.9, the $\cos\theta_{\pi^-}$ vs. $m_{\bar{K}_0\pi^-}$ plot of Figure X.10, and the φ_{π^-} vs. $m_{\bar{K}_0\pi^-}$ plot of Figure X.11.

1C/4C Overlap at $K_1(890)$

Figure X.5: The overlap region at $t' \approx 0.06$ in the $K_1(892)$ region. There are 4Cs in the MC sample produced at a lower t' value than in the data.

The Dalitz plot clearly shows the production of resonances in the $\bar{K}_0\pi^-$ system. Seen are the $K^*(892)$ and the $K_2^*(1430)$. There are hints of structure at the $K_2^*(1780)$, and perhaps in the $K_2^*(2050)$ region. Also, the density enhancements between the resonant bands at high $p\pi$ mass is indicative of interference behavior between amplitudes. There is also evidence of strong baryonic resonance production. The left-most vertical band is immediately identified with the $\Delta(1232)$. There are two clear higher mass N^* 's as well. Along the upper right boundary, there is evidence for production of a Y^* at a mass of $\approx 1.8 \text{ GeV}/c^2$. The obvious overlap of the baryonic bands with the mesonic bands has non-trivial implications for the present analysis. Prior analyses from the present experiment and others on related

Figure X.6: The acceptance corrected $\bar{K}_0\pi^-$ mass. The 4C and 1C samples are shown separately.

$K\pi$ production, have been forced to make hard cuts on the baryonic mass in order to extract the interesting meson spectroscopy. Unfortunately, such simple cuts in the baryon systems reflect non-linearly into the mass, ℓ' , and angular distributions of the meson system. In the present analysis, however, the good efficiency suggests alternative means of coping with this difficulty, as will be shown. On the other hand, the overlap with the baryon systems has been used as of late to extract physics

Figure X.7: The mass distributions of the proton-meson systems. The $p\pi^-$ mass distribution is shown in (a), the $p\bar{K}_0$ distribution in (b).

from the meson reactions. If the amplitudes for baryon and meson production are coherent, the total amplitude for reactions involves an interference term which has consequences that can be understood in terms of resonance production in the meson system^[66,67].

The Chew-Low plot has several interesting features. Again, the resonance bands are clear, with stronger hints of the $K_4(2050)$ than the Dalitz plot would

Figure X.8: The Dalitz plot. Strong bands indicative of resonant production are observed in the $\bar{K}_0\pi$ system (the horizontal bands), the $p\pi$ system (the vertical bands), and the $p\bar{K}_0$ system (the diagonal band at the upper right of the plot).

imply. Note the depletion of events along the boundary for $m_{\bar{K}_0\pi} \geq 2.5 \text{ GeV}/c^2$, which it actually sets in at smaller mass values. This loss of 1C events is indicative of the acceptance as previously discussed. There are real 1C events at these large momentum transfers, where the recoil is forward and consequently must travel longer paths in the target before it escapes. The resonance bands extend to large momentum transfers. These events possess the same structure in the angular distributions as the low t' events, however their flatter t' dependence indicates that a different production mechanism may be in effect.

The plot of $\cos \theta_j$ against $m_{\bar{K}_0\pi}$ shows the event enhancements at forward and backward $\cos \theta_j$ at high $m_{\bar{K}_0\pi}$, characteristic of baryonic overlap. Actually, first

Figure X.9: The Chew-Low plot. The $\bar{K}_0\pi$ resonance bands are seen. Note the long tails extending to large momentum transfers in the resonant bands.

effects of the baryons can be seen at $\approx 1.3 \text{ GeV}/c^2$.

The plot of φ_{TY} against $m_{\bar{K}_0\pi}$ also illustrates the classic effect of baryons. The higher event concentration near $\varphi_{TY} = 0$ becomes broader and less restricted as the $K\pi$ mass rises. More interesting, is the obvious appearance of the double peaking of the φ_{TY} dependence, in sharp contrast to the $K^-\pi^+n$ case. This $\cos 2\varphi_{TY}$ behavior is characteristic of natural parity exchange, and is the first clear indication that the physics in this channel is different from that in $K^-\pi^+n$.

Figure X.10: A plot of $\cos \theta_J$ vs. $m_{K_0^* \pi}$. The effects of the production of baryonic resonances are seen as the dense concentration of events in the forward and backward regions of the plot as $m_{K_0^* \pi}$ increases.

Figure X.11: A plot of ϕ_{TY} vs. $m_{K_0^* \pi}$. The effects of baryonic production are seen as the relatively dense band of events moving away from $\phi_{TY} \approx 0$ as $m_{K_0^* \pi}$ increases.

XI. Moments

The nature of the $\bar{K}_0^0 \pi^- p$ data sample is exhibited in the distributions of unnormalized moments of figures XI.1-XI.5. The moments and errors are calculated using equations Equation X.13 and Equation X.14 where the measured quantity is just the weighted value of the moment. The quantity t_{lm} is related to the spherical harmonic moment by

$$t_{lm} = \sqrt{4\pi} \sum_i w_i Y_{lm}(\Omega_i), \quad (\text{XI.1})$$

with w_i the acceptance weight. The data are presented in 20 MeV/c² bins for $0.5 < m_{\bar{K}_0^0 \pi^-} \leq 1.18$, 40 MeV/c² bins for $1.18 < m_{\bar{K}_0^0 \pi^-} \leq 1.86$, 60 MeV/c² bins for $1.86 < m_{\bar{K}_0^0 \pi^-} \leq 2.52$, and mass is given in GeV/c². These are acceptance corrected moments and, in particular, t_{00} is the corrected number of events. The raw moments, i.e. the moments prior to acceptance correction, are not significantly different in shape. The data shown are summed over $t \leq 0.8$.

By inspection, the $m = 3$ and $m = 4$ moments are consistent with 0. Consequently, they are not used in the analysis. Also, the $l \geq 9$ moments are observed to be structureless for all mass and t . There is clear mass dependence in the $m = 0, 1, 2$ moments. In particular, the very large structures in the $m = 2$ moments are in sharp contrast to the $K^- \pi^+ n$ channel. The physics in the two complementary channels is clearly different, as expected.

As discussed previously, there are effects in the $\bar{K}_0^0 \pi^-$ distributions due to the production of baryon resonances. Such effects exhibit themselves as monotonically increasing (decreasing) functions of $m_{\bar{K}_0^0 \pi^-}$ in the $m = 0$ ($m = 1$) moments. In order to extract the relevant K^* measurements for the purposes of the present analysis,

Figure XI.1 The unnormalized $m=0$ moments. Moments are shown for $l \leq 8$. Errors are statistical only.

this background must be removed. The effects of the baryons are in fact much more extensive than a simple moment plot would indicate. They reflect into the meson system as complicated functions of $m_{\bar{K}_0^* \pi}$, t , and the GJ angles, and there is no way to analytically distinguish these backgrounds from the $\bar{K}_0^* \pi^-$ data.

XI.1 Comparison with Other Experiments

As a check upon the normalization and mass dependence of the observed moments, it is important to compare these measurements with others. The highest statistics measurement of the present reaction in this beam momentum range is that from the 10 GeV/c experiment at the CERN/Geneva non-magnetic spectrometer^{[35][34][57]}. There are other less significant measurements^{[40][61][68]}, and there is also an experiment at higher beam momentum using the non-magnetic spectrometer^[59]. It follows that the only meaningful comparison is with the CERN 10 GeV/c data and this is shown in Figure XI.6. The comparison requires that the data be plotted in terms of cross section, using the sensitivity of the experiment as discussed in section IX.8. Corrections have been made for the branching ratio $\text{BR}(\bar{K}_0^* \rightarrow K_s \rightarrow \pi^+ \pi^-) = 0.344$. To be compatible with the CERN data, only events with $0.1 \leq t' \leq 0.4$ are used, and the E-135 data have been scaled by a factor of $(\frac{11}{10})^2$ account for the beam momentum difference. The agreement is excellent, except perhaps for the t_{42} moment in the $m_{\bar{K}_0^* \pi} \approx 1.4$ GeV/c² region. The difference could be due to the fact that the Dalitz plot shows that in the $m_{\bar{K}_0^* \pi} \approx 1.4$ GeV/c² region, there is significant N^* production in the E135 data. The CERN data are not as sensitive to N^* final states, due to acceptance.

XI.2 Scattering Formalism

The moment plots plainly show the production of the well known strange meson resonances $K^*(892)$, $K_2^*(1430)$, and perhaps even $K_3^*(1780)$, based on the t_{00} and t_{62} plots. In order to extract precise resonance parameters, the exact form of the relativistic Breit-Wigner fits must be defined.

The S -matrix element describing $1 + 2 \rightarrow (31 + 32) + 4$ scattering is given by

$$S_{fi} = b_{fi} - i(2\pi)^4 \delta^4(p_f - p_i) (2^5 E_1 E_2 E_{31} E_{32} E_4)^{-1/2} T_{fi}, \quad (\text{XI.2})$$

Figure XI.2: The unnormalized $m=1$ moments. Moments are shown for $l \leq 8$. Errors are statistical only.

where p are four-momenta as before, and T_{fi} is the Lorentz-invariant transition amplitude. The differential cross section is given by the transition rate multiplied by the Lorentz-invariant phase space, and divided by the probe flux, as follows:

$$d\sigma = \frac{(2\pi)^4 M_2 M_4}{2p^0 E^*} |T_{fi}|^2 \delta^4(p_f - p_i) \frac{d^3 p_{31}}{(2\pi)^{3/2} E_{31}} \frac{d^3 p_{32}}{(2\pi)^{3/2} E_{32}} \frac{d^3 p_4}{(2\pi)^3 E_4}, \quad (\text{X1.3})$$

where the p^* , E^* are the momentum and energy in the CMS (center-of-mass) frame. The appropriate "Feynman" diagram is shown in Figure X1.7, and ignoring final state interaction effects, the invariant amplitude can be written as

$$T_{fi} = V_i \Delta_x V_f, \quad (\text{X1.4})$$

where the V_i, V_f are invariant vertex amplitudes, and $\Delta_x \equiv \Delta_x(t_x, m_x^2)$ represents the propagator which connects the top and bottom vertices. Then, (X1.3) can be written as

$$d\sigma = \frac{(2\pi)^4}{2p^0 E^*} \left\{ |V_i|^2 \delta^4(p_{31} + p_{32} - p_1 - p_2) \frac{d^3 p_{31}}{(2\pi)^{3/2} E_{31}} \frac{d^3 p_{32}}{(2\pi)^{3/2} E_{32}} \right. \\ \left. \cdot \Delta_x^2 \left\{ \frac{M_2 M_4}{2} \sum_{\text{spins}} |V_f|^2 \frac{d^3 p_4}{(2\pi)^3 E_4} \right\} \right\}. \quad (\text{X1.5})$$

For the specific case of pion exchange, V_b is taken to be a point interaction of the form

$$V_b = G \bar{u}_4 \gamma_5 u_2. \quad (\text{X1.6})$$

Summing over spin states, the second set of curly brackets in X1.5 becomes

$$\frac{G^2}{(2\pi)^3} \left\{ t_x^2 + (M_1 - M_2)^2 \right\} \frac{d^3 p_1}{1E_1}, \quad (\text{X1.7})$$

with $\frac{G^2}{4\pi}$ the appropriate pion-nucleon coupling constant. A change of variable to $M^2 \equiv (p_{31} + p_{32})^2$ and $t \equiv t_x$, recognizing that $M_1 = M_2 = M_{\text{proton}} \equiv M_p$ and that $p^0 E^* \approx p_{\text{lab}} M_p$, makes (X1.7) become, after integrating out a trivial azimuthal dependence,

$$\frac{(G^2/4\pi)}{16\pi p_{\text{lab}} M_p} t dM^2 dt. \quad (\text{X1.8})$$

Figure X1.3: The unnormalized $m=2$ moments. Moments are shown for $l \leq 8$. Errors are statistical only.

The top vertex is related to the virtual 2-body scattering process $1+x \rightarrow 31+32$, where all particles are bosons. With the magnitude of 3-momentum of the constituents at V_t in the rest frame of M given by

$$q_i = \frac{1}{2M} \sqrt{(M^2 - M_i^2 + t)^2 + 4M_i^2 t}, \quad (\text{XI.9})$$

the differential cross section for that process is,

$$d\sigma_i = \frac{(2\pi)^4}{4q_i M} \left\{ |V_i|^2 \delta^4(p_{31} + p_{32} - p_1 - p_2) \frac{d^3 p_{31}}{(2\pi)^3 2E_{31}} \frac{d^3 p_{32}}{(2\pi)^3 2E_{32}} \right\}. \quad (\text{XI.10})$$

It follows that for one pion exchange, (XI.5) becomes

$$d\sigma = \frac{(G^2/4\pi)}{16\pi(p_{lab} M_p)^2 (t + m_\pi^2)^2} (2q_i M \frac{d\sigma_i}{d\Omega}) dM^2 dt d\Omega, \quad (\text{XI.11})$$

where $d\Omega = d(\cos \theta_f) d\phi_{TV}$ in the t -channel helicity frame, and Δ_s has been expressed as the pion propagator. In the language of Chew and Low^[10], the differential cross section, $\frac{d\sigma_i}{d\Omega}$, may then be considered to describe off-mass-shell $K\pi$ scattering. This may be expressed in terms of the corresponding on-mass-shell cross section by writing

$$V_i(t) = \frac{V_i(t)}{V_t^{on}} V_t^{on} \quad (\text{XI.12})$$

where V_t^{on} is the invariant amplitude describing the real $K\pi$ process at the same value of M . Equation XI.11 then becomes

$$d\sigma = \frac{(G^2/4\pi)}{16\pi(p_{lab} M_p)^2 (t + m_\pi^2)^2} F(t) \left(2q_i M \frac{d\sigma_i}{d\Omega} \right)^{on} dM^2 dt d\Omega, \quad (\text{XI.13})$$

where q_i is obtained from (XI.9) with $t = -m_\pi^2$, and

$$F(t) = \left| \frac{V_i(t)}{V_t^{on}} \right|^2 \rightarrow 1 \text{ as } t \rightarrow -m_\pi^2. \quad (\text{XI.14})$$

The form factor, $F(t)$, may be well-described by an exponential of the form

$$F(t) = e^{-2b(t+m_\pi^2)}. \quad (\text{XI.15})$$

Figure XI.14: The unnormalized $m=3$ moments. Moments are shown for $l \leq 8$. Errors are statistical only.

The on-shell $K\pi$ differential cross section may be written as

$$\frac{d\sigma_l^{on}}{d\Omega} = \left| c_{\frac{1}{2}} f^{\frac{1}{2}}(\Omega) + c_{\frac{3}{2}} f^{\frac{3}{2}}(\Omega) \right|^2 \quad (X1.16)$$

where the scattering amplitude for isospin l is given by

$$f^l(\cos\theta_j, \varphi_{TV}) = \frac{\sqrt{4\pi}}{q_i} \sum_{l,m} a_l Y_l^0(\theta_j, \varphi_{TV}) \quad (X1.17)$$

with

$$a_l = \sqrt{2l+1} \varepsilon_l \sin \delta_l^l e^{i\delta_l^l}, \quad (X1.18)$$

and ε_l is the elasticity of the l^{th} partial wave of phase shift δ_l^l ; $c_{\frac{1}{2}}$ and $c_{\frac{3}{2}}$ are isospin CG coefficients.

It has been shown that this pion-exchange process is the dominant contributor to the production of the $K^-\pi^+$ system in the final state $K^-\pi^+n$; other contributions due to absorptive cuts and $\rho-\omega$ exchange are much less significant. For $K\pi$ production in the $\bar{K}_0\pi^+p$ final state, all contributions corresponding to isospin 1 exchange in the t -channel are reduced by a factor 4 with respect to those in the $K^-\pi^+n$ state, due to isospin CG coefficients. Explicitly, for pion exchange the renormalized coupling constant values at the baryon vertex are related by

$$(G^2/4\pi)_{pp\pi^0} = \frac{1}{2} (G^2/4\pi)_{p\pi^+\pi^+}. \quad (X1.19)$$

The $K\pi$ scattering amplitude relevant to the $\bar{K}_0\pi^+p$ state is given by

$$f = \frac{\sqrt{2}}{3} \left(f^{\frac{1}{2}} - f^{\frac{3}{2}} \right) \quad (X1.20)$$

whereas that corresponding to the $K^-\pi^+n$ final state is

$$f' = \frac{2}{3} \left(f^{\frac{1}{2}} + \frac{1}{2} f^{\frac{3}{2}} \right). \quad (X1.21)$$

The CG coefficients outside the brackets result in a further factor of two reduction in intensity. However, for the $K\pi$ S wave, which has a sizeable isospin $3/2$ contribution,

Figure X1.5 The unnormalized $m=4$ moments. Moments are shown for $l \leq 8$. Errors are statistical only.

there are other significant effects due to the fact that the two isospin amplitudes combine differently for the $\bar{K}_0\pi^-p$ and $K^-\pi^+n$ final states; as will be discussed.

The Chew-Low formalism has proven very successful, due in large part to the close proximity of the pion pole to the physical region. In the present channel, the dominant exchanges are natural parity isoscalars. The lowest lying isoscalar mesons are so massive as to be far away from the physical region, and it is not obvious that the simple propagator can be used here. Instead, the t -dependence of the natural parity exchange contribution is simply parametrized by the expression

$$(t')^{\frac{n_f}{2}} e^{-\alpha t'} \quad (\text{XI.22})$$

where n_f is the net helicity flip in the production process. As will be shown, the different representations of the pion and natural parity exchange contributions describe the data well.

For a resonant $K\pi$ amplitude produced by pion-exchange, the corresponding mass spectrum, obtained by integrating (XI.11) over angles and t takes the form

$$\frac{d\sigma}{d\Omega} \propto \frac{(M^2/q)\Gamma^2(M)}{(M_R^2 - M^2)^2 + M_R^2\Gamma_T^2(M)}. \quad (\text{XI.23})$$

Here, $\Gamma_T(M)$ is the total width of the resonance of mass M_R , where

$$\Gamma_T(M) = \sum_i \Gamma_i(M), \quad (\text{XI.24})$$

and $\Gamma_i(M)$ are the partial widths of the separate open decay channels. The $K\pi$ partial width is parametrized by

$$\Gamma(M) = \frac{q}{M} BF(M), \quad (\text{XI.25})$$

where $BF(M)$ is the Blatt-Weiskopf barrier factor given by

$$BF(M) = \frac{(qR)^{2l}}{D_l(qR)}, \quad (\text{XI.26})$$

Figure XI.6. A comparison of experiments. The open circles are data from the CERN 10 GeV/c experiment, the solid circles are for the data of this analysis.

Figure XI.7: The scattering picture. This is for an arbitrary $1+2 \rightarrow (31+32)+4$ scattering, where the shaded areas represent vertex couplings.

with l the angular momentum of the resonance. The $D_l(x)$ are the polynomials defined in Table XI.1. Then, with $\Gamma(M_R)$ as the width at the resonance peak,

$$\Gamma(M) = \left(\frac{q}{q_R}\right)^{2l+1} \left(\frac{M_R}{M}\right) \frac{D_l(q_R R)}{D_l(q R)} \Gamma(M_R); \quad (\text{XI.27})$$

if it is assumed that Γ and Γ_T have the same mass dependence, the mass distribution given by (XI.23) is completely specified in terms of M_R , Γ_R , and R . In the case of natural parity exchange, no assumption is made about the M dependence at the resonance production vertex (cf. fig XI.7). The resulting mass spectrum is then

$$\frac{d\sigma}{dM} \propto \frac{M\Gamma(M)}{(M_R^2 - M^2)^2 + M_R^2\Gamma_T^2(M)}. \quad (\text{XI.28})$$

Because the cross section goes as the square of the amplitude, the largest l amplitude can produce a $2l$ moment, from the properties of the spherical harmonic functions. Conversely, in a particular mass region, observation of structure in an even L moment, with no structure in higher moments, implies a maximum l amplitude given by $l = L/2$. This fact is used to define the set of moments, and hence

l	$D_l(x)$
0	1
1	$1+x^2$
2	$9+3x^2+x^4$
3	$225+45x^2+6x^4+x^6$
4	$11025+1575x^2+135x^4+10x^6+x^8$

Table XI.1: Blatt-Weiskopf barrier factors. $x \equiv qR$, and the radius parameter R is usually given in GeV^{-1} .

maximum partial wave, used in the fits to extract amplitudes. The relationship between the moments and amplitudes is discussed in Appendix A.

XI.3 Baryon Resonance Correction

Prior to fitting the moments, an important digression must be taken. As has been mentioned, there is clear evidence for the production of baryon resonances in the data sample. This is a direct consequence of the large, uniform acceptance of the experiment. The Dalitz plot indicates that the N^* production* reflects into the meson system for $m_{K^0\pi^+} \geq 1.2 \text{ GeV}/c^2$. This reflection is observed in the moments as a monotonic increase in the absolute value of the $m = 0, 1$ moments as the $K\pi$ mass increases. It follows that the corresponding $K\pi$ amplitudes and phases will be affected by this overlap.

Studies have shown that the l' dependence of the N^* production is not localized, so that l' cuts do not eliminate the effect. However, as discussed previously, the N^* s do reflect into specific regions of the $\cos\theta_j$ and $\varphi_{K\pi}$ angles, as seen in the scatterplots. With the excellent acceptance here, a technique which takes advantage of the localized GJ angle properties of the N^* reflections was used to cut these

* In what follows, N^* will be used to refer to both N^* and Δ production.

background regions away. The technique is just the method of moments, applied to the weighted events.

Consider an arbitrary distribution of events, $I(\Omega)$, and let Ω_c be the solid angle remaining from the full solid angle after an N^* cut is applied. Outside the region of N^* overlap, $I(\Omega)$ is the true K_0^* angular distribution, which may be written in terms of spherical harmonic moments as

$$I(\Omega) = \sum_{l=0}^{l_{\max}} c_{l0} Y_l^0(\Omega) + \sum_{l=1}^{l_{\max}} \sum_{m=1}^{m=2} c_{lm} 2ReY_l^m(\Omega), \quad (XI.29)$$

where $l_{\max}$ is some cutoff value of the orbital quantum number, the c_{lm} s are the true K_0^* moments, and only the moments with $m \leq 2$ are considered, as higher m values are not required by the data. Then, in turn, multiply (XI.29) by each of the spherical harmonics which make up the expansion on the right hand side, and integrate over the range of solid angle, Ω_c , which is left after the N^* cut. Typically, the region will include all of the space, except the region near forward $\cos\theta$, with $\varphi_{TY} \approx 0$. If there were no N^* cut, the original moments would be recovered, because in fact, that is how they were determined. The set of $3l_{\max} \times 3l_{\max}$ equations which results yields the following matrix (with $L \equiv l_{\max}$):

$$\begin{aligned} \int_{\Omega} Y_{00} d\Omega &= c_{00} \int_{\Omega} Y_{00} Y_{00} d\Omega + \dots + c_{L1} \int_{\Omega} Y_{00} 2ReY_{L1} d\Omega + c_{L2} \int_{\Omega} Y_{00} 2ReY_{L2} d\Omega \\ \int_{\Omega} Y_{10} d\Omega &= c_{00} \int_{\Omega} Y_{10} Y_{00} d\Omega + \dots + c_{L1} \int_{\Omega} Y_{10} 2ReY_{L1} d\Omega + c_{L2} \int_{\Omega} Y_{10} 2ReY_{L2} d\Omega \\ &\vdots \\ \int_{\Omega} ReY_{L1} d\Omega &= c_{00} \int_{\Omega} ReY_{L1} Y_{00} d\Omega + \dots + c_{L1} \int_{\Omega} ReY_{L1} 2ReY_{L1} d\Omega + c_{L2} \int_{\Omega} ReY_{L1} 2ReY_{L2} d\Omega \\ \int_{\Omega} ReY_{L2} d\Omega &= c_{00} \int_{\Omega} ReY_{L2} Y_{00} d\Omega + \dots + c_{L1} \int_{\Omega} ReY_{L2} 2ReY_{L1} d\Omega + c_{L2} \int_{\Omega} ReY_{L2} 2ReY_{L2} d\Omega \end{aligned}$$

The left hand side is just the usual weighted event sum, except that only events within Ω_c are added. The integrals on the right hand side are constants which may be evaluated by Gaussian quadrature over Ω_c . The resulting square matrix is inverted to find the c_{lm} coefficients. That is,

$$C = Y^{-1} I, \quad (XI.30)$$

where C is the column vector of the c coefficients, Y is the matrix of integrations, and I is the event sum. Also, the error matrix appropriate to the coefficients must

be calculated. It is given by

$$(\delta C)^2 = Y^{-1} E (\widetilde{Y^{-1}}) \quad (XI.31)$$

where E is the error matrix which accompanies I in the usual way. The quoted errors on the coefficients are just the square roots of the diagonal terms.

The $K\pi$ moments obtained in this way naturally have larger errors than the direct moment sums, because the statistics in a given sum are reduced, and because of the correlation terms introduced by the N^* cut. It should be observed that the important moments for the present analysis are the $m = 2$ moments, which reflect the predominantly natural parity exchange. Inspection of the unnormalized $m = 2$ moments of Figure XI.3 shows that these moments are not as severely affected by N^* s as the $m = 0, 1$ moments, due of course to the fact the contamination is largest where the value of the $m = 2$ moments is small. As such, the actual effect of the cut is reduced for these moments.

There were three N^* cuts used in the present work, the choice depending upon the mass region of interest. The cuts were chosen based upon the amount of N^* removed, measured by simply examining the histograms of $m_{\pi\pi}$ and/or $m_{\pi K_0}$ for the cut. The actual cuts were made in θ , φ_{TY} and are tabulated in Table XI.2. The N^* cut used in the presentation of any fit will be identified.

XI.4 Breit-Wigner Fits

The leading moments in the data sample are used here to extract the Breit-Wigner parameters for the resonances observed. Also, the observed lineshapes can be compared with similar measurements from the $K^-\pi^+\pi$ results of the present experiment^[71], so as to understand more clearly the difference in the physical processes.

XI.4.a The $K^*(892)$ region

The $K^-\pi^+K^*(892)$ is certainly the most well understood object in $K\pi$ scattering. The mass and width are well known and the resonance mass difference between the charged and neutral modes is relatively well understood^[90]. However, the higher statistics of E-135 will permit some additional conclusions to be drawn.

N^π cut	$m_{K^0\pi}$ (GeV/c ²)		t' cut (GeV ²)		θ_f cut ($\times\pi$)		φ_{TV} cut ($\times\pi$)	
	Min	Max	Min	Max	Min	Max	Min	Max
K_2^0 cut	1.2	1.6	0.0	0.2	0.0	0.2	-0.5	0.5
	1.2	1.6	0.2	0.8	0.0	0.4	-0.3	0.3
K_3^0 cut	1.52	1.96	0.08	0.8	0.0	0.15	-0.6	0.6
	1.52	1.96	0.08	0.8	0.15	0.25	-0.5	0.5
K_4^0 cut	1.52	1.96	0.08	0.8	0.25	0.35	-0.3	0.3
	1.7	2.1	0.08	0.8	0.0	0.15	-0.7	0.7
	1.7	2.1	0.08	0.8	0.15	0.25	-0.5	0.5
	1.7	2.1	0.08	0.8	0.25	0.35	-0.35	0.35
	1.7	2.1	0.08	0.8	0.80	1.00	$0.7 \leq \varphi_{TV} $	

Table X1.2: The N^π cuts. There is overlap between the cuts to permit continuity checks.

For the mass region in the vicinity of the peak, the t' dependence of the moment combinations is plotted in Figure X1.8. The fitted forms are from the $F(t)$ parameterizations of (X1.15), and the helicity projections σ_0 , σ_+ , and σ_- are defined in appendix A. The dominance of the helicity 1, natural parity exchange is apparent. The slope parameters from the fits are in good agreement with other similar measurements^[58].

The mass dependence of the moments and projections are shown in Figure X1.9, for $0 \leq t' \leq 0.8$. The fits used were described in section X1.2. The S wave contribution has been removed, motivated by the results of the $K^-\pi^+\pi^0$ analysis, where it was found that the S wave, the $K_0^*(1350)$, could be parametrized by an effective range term and a Breit-Wigner. Note that the S wave background is expected to be small in the present channel. The factor of $1/3$ in the projection definitions, another factor of $1/3$ from the $\sqrt{2t'+1}$ normalization, and another factor of $\approx 1/2$ from the magnitude of the S amplitude implies that the S^2 contribution to

Figure X1.8 The t' dependence at the $K^*(892)$. The helicity projections are shown, overlaid with fits based upon the $F(t)$ parameterization.

$\sigma_0^{t'}$ is on the order of $1/18$. From the t' plot, the contribution to $\sigma_+^{t'}$ from S is then negligible. The extracted Breit-Wigner (BW) parameters are listed in Table X1.3. The t_{22} measurement is the most reliable, as it reflects the dominant natural parity exchange only. As supporting evidence for resonance behavior in the P waves, the t_{10} , t_{11} , and t_{21} moments exhibit the oscillatory behavior expected for the P wave beating against a slowly varying S wave.

Figure X1.9 The mass dependence of the helicity projections. Each curve is fit with a small, increasing background to account for the N wave background

Fit	Mass (MeV/c^2)	Width (MeV/c^2)	Radius (GeV^{-1})
t_{00}	$890.9 \pm 0.2 \pm 0.5$	$46.2 \pm 0.5 \pm 0.5$	$5.6 \pm 0.5 \pm 1.$
t_{20}	$889.6 \pm 0.5 \pm 0.5$	$50.7 \pm 1.4 \pm 0.5$	$100. \pm 50. \pm 5.$
t_{22}	$890.5 \pm 0.3 \pm 0.5$	$46.0 \pm 0.7 \pm 0.5$	$11.4 \pm 2.5 \pm 1.$
σ_0^P	$890.3 \pm 0.2 \pm 0.5$	$50.4 \pm 1.7 \pm 0.5$	$0.00 \pm 10. \pm 1.$
σ_+^P	$890.7 \pm 0.2 \pm 0.5$	$46.4 \pm 0.4 \pm 0.5$	$15.4 \pm 5. \pm 1.$
σ_-^P	$898.8 \pm 7.7 \pm 0.5$	$87.9 \pm 30. \pm 3.$	$100. \pm 50. \pm 5.$

Table X1.3: BW parameters for the large moments in the $K^*(892)$ region. The first error is statistical, the second systematic

The cross section for the production of the $K^*(892)$ is compared against the world data in Figure X1.10. The measurement is corrected for the tails of the $\bar{K}_0^0$, the unseen $K_s \rightarrow \pi^0 \pi^0$ decay, the unseen $K^*(892) \rightarrow K^- \pi^0$ mode, and the tails of the $K^*(892)$ BW outside of the $0.84 \leq m_{\bar{K}_0^0} \leq 0.94$ mass region selected. The measurement gives

$$\sigma(K^- p \rightarrow \bar{K}^{*0}(892)p) = 116 \pm 5 \pm 3 \text{ } \mu\text{b.} \quad (\text{X1.32})$$

in good agreement with the expected value as extrapolated from the other measurements.

As Figure X1.8 indicates, there is an apparent slope change for $t' \geq 0.8 \text{ GeV}^2$, indicating that perhaps there is a change in the production of the K^* at these large momentum transfers. The effect is also evident in the Chew-Low plot of Chapter X, where the resonance regions have clear bands extending to large t' . In fact, the moments have the same behavior, i.e., the production physics is very similar. The low statistics in these tails prohibit a full amplitude analysis, but the production is due to natural parity exchange, albeit with a flatter t' dependence. In fact the behavior observed is quite consistent with the $\rho - a_2$ contribution to $\bar{K}^{*0}(892)$ production observed for the $K^- \pi^+ n$ final state^[10].

Figure X1.10: The cross section for the reaction $K^- p \rightarrow \bar{K}^{*0} p$ (892).

X1.4.b The $K_2^*(1430)$ region

The moments which will be used for analysis purposes are obtained from the unnormalized moments, after removing the N^* 's according to the K_2^* cut of Table X1.2. The effect of the K_2^* cut on the px mass distribution is shown in Figure X1.11. The majority of N^* events in the sample have been removed. The leading moments, and BW fits to them, are shown in Figure X1.12, and tabulated in Table X1.4. Note that there is no structure in the $l = 6$ unnormalized moments in this mass region, indicating no significant leading F wave production. Again, the raw $t_{10}, t_{21}, t_{30}, t_{32}, t_{50}, t_{52}$ distributions exhibit interference behavior. Such behavior in this range of moments implies both PD and DF interference. Of special interest is the t_{20} moment, as the fact that it is positive with a slow rise to a peak near $m_{\bar{K}_0^*} \approx 1.4$ GeV/c² and then a rapid falloff, is also evidence for an S wave background which resonates in the region. The helicity projections shown in Figure X1.12 are derived in appendix A. The σ_4^0 plot reflects predominance of the natural parity exchange, but the σ_0^0 plot is strongly influenced by the underlying S wave, hence the fit is not expected

to reproduce the data quite as well. Note the similarity in shape to the t_{20} moment. A more complete understanding of the S wave mass dependence will require the full amplitude analysis.

Figure X1.11: The px mass distribution for $1.2 \leq m_{K\pi} \leq 1.6$. The shaded events are removed by the N^* cut.

The t' dependence of the cross section in this mass region is more difficult to know with great precision. However, with the N^* cut, a fit to the t' dependence of the t_{42} moment gives a slope parameter for the cross section of $\alpha_{42} = 8.92 \pm 0.27$, in agreement with the same parameter from another measurement^[36].

X1.4.c The $K_3^*(1780)$ region

The moments which will be used for analysis purposes are obtained from the unnormalized moments, after removing the N^* 's according to the K_2^* cut of Table X1.2. The effect of the K_2^* cut on the px mass distribution is shown in Figure X1.13a, and that on the $p\bar{K}_0^*$ mass distribution is shown in Figure X1.13b. As previously

Figure XI.12: The leading moments in the $K_2^*(1430)$ region. The uncut moments are shown as the unfilled dots.

discussed, the 1C events were not used in the analysis for $m_{K^0\pi^+} \geq 1.5 \text{ GeV}/c^2$; only the 4C events within the P' cut of Table XI.2 are included in the analysis sample.

Fit	Mass (MeV/c^2)	Width (MeV/c^2)	Radius (GeV^{-1})
t_{00}	$1419 \pm 0.8 \pm 1.$	$99.3 \pm 3. \pm 2.$	$100. \pm 78. \pm 10.$
t_{22}	$1424 \pm 2.6 \pm 1.$	$100.7 \pm 10.1 \pm 3.$	$0. \pm 53. \pm 5.$
t_{40}	$1424 \pm 3.5 \pm 1.$	$191 \pm 11.3 \pm 3.$	$5. \pm 5. \pm 3.$
t_{42}	$1420 \pm 2. \pm 1.$	$89.8 \pm 5. \pm 1.$	$3.1 \pm 1.2 \pm 1.$
σ_0^D	$1412 \pm 4.8 \pm 1.$	$100.7 \pm 14.7 \pm 3.$	$100. \pm 59. \pm 15.$
σ_+^D	$1420.5 \pm 1.1 \pm 1.$	$98.8 \pm 4.4 \pm 3.$	$12.5 \pm 54. \pm 10.$

Table XI.4: BW parameters for the large moments in the $K_2^*(1430)$ region. The first error is statistical, the second systematic.

The fits to the leading moments are shown in Figure XI.14, and tabulated in Table XI.5.

Figure XI.13: The effect of the K_S^0 cut, where $1.6 \leq m_{K_S^0} \leq 2.0$. The shaded events are removed by the N^* cut. The $p\pi$ mass distribution is shown in (a), with the $p\bar{K}_0^0$ mass in (b).

Figure XI.14: The leading moments in the $K_S^0(1780)$ region.

Fit	Mass (MeV/c ²)	Width (MeV/c ²)	Radius (GeV ⁻¹)
t_{00}	$1747 \pm 12 \pm 4$	$115 \pm 59 \pm 10$	100.
t_{22}	$1738 \pm 21 \pm 5$	$195 \pm 36 \pm 15$	100.
σ_0^F	$1784 \pm 11 \pm 10$	$233 \pm 36 \pm 50$	100.
σ_2^F	$1711 \pm 9 \pm 5$	$213 \pm 60 \pm 10$	100.

Table XI.5: BW parameters for the large moments in the $K_S^0(1780)$ region. The first error is statistical, the second systematic.

XII. Amplitudes

An understanding of the physics underlying the set of $\overline{K}_0\pi^-$ unnormalized moments for the reaction $K^-p \rightarrow \overline{K}_0\pi^-p$ requires a partial wave amplitude analysis. As discussed in chapter XI, and given explicitly in appendix A, the moments reflect the production amplitudes. In particular, the smaller, or underlying amplitudes which are dominated by the leading waves in each mass region and can only be extracted by means of a partial wave analysis.

XII.1 The Scattering Model

To extract the $\overline{K}_0\pi^-$ scattering amplitudes from the observed moments in the Gottfried-Jackson frame, it is convenient to define helicity amplitudes with definite exchange naturality. In the one particle exchange model, the naturality is given by

$$\eta = (-1)^J P, \quad (\text{XII.1})$$

where J, P are the spin and parity of the exchanged particle. For a helicity amplitude H_{λ}^L , the production amplitudes $L_{\lambda\eta}$ in a $\overline{K}_0\pi^-$ system of spin L and t -channel helicity λ by natural ($\eta = +1$) or unnatural ($\eta = -1$) parity exchange is written as

$$L_{\lambda\eta} = c_{\lambda} \left(H_{\lambda}^L - \eta (-1)^L H_{-\lambda}^L \right); \quad c_{\lambda} = \begin{cases} 1/2, & \lambda = 0 \\ 1/\sqrt{2}, & \lambda \neq 0 \end{cases} \quad (\text{XII.2})$$

where $\lambda = 0, 1$ here. The observed moments are expressed as sums of bilinear products of the production amplitudes, the relationship given by^[9]

$$t_{lm} = \sum_{LL'\lambda\lambda'} \left(\frac{(2L+1)(2L'+1)}{2l+1} \right)^{\frac{1}{2}} \langle LL'00 | l0 \rangle (-1)^{\lambda'} \langle LL'\lambda\lambda' | lm \rangle \cdot \frac{1}{4c_{\lambda}c_{\lambda'}} \text{Re}(L_{\lambda\eta} L_{\lambda'\eta}^*) \quad (\text{XII.3})$$

where $\langle \quad \rangle$ are Clebsch-Gordon coefficients. Here, a sum over nucleon helicity, and the property of spin coherence at the nucleon vertex are assumed.

The set of moment-amplitude relations used here is listed in appendix A. The mass and t' dependence of the individual amplitudes must both be determined, because the interesting physics has consequences in both parameters. Solutions in a given bin, in mass or t' , or both, are obtained by minimizing a χ^2 function given by an effective χ^2 , where

$$\chi^2 = \sum_{\mu\nu} (t_{\mu} - \bar{t}_{\mu}) E_{\mu\nu}^{-1} (t_{\nu} - \bar{t}_{\nu}), \quad (\text{XII.4})$$

where t_{μ} are the measured moments, $\bar{t}_{\mu}$ are the expected moment values as given in appendix A, and E is the error matrix for the t_{μ} s. A fit was deemed acceptable if the χ^2 was less than 3x the number of parameters in the fit. All 3 helicity states for a given L are included in the fits, in general, but the analysis program was designed with enough flexibility to allow individual moments or amplitudes to be fixed or ignored.

The complete mass range for the amplitude analysis was from threshold to 2 GeV/c². The upper limit of the mass range reflects the lack of statistics coupled with the increased complexity of the analysis at high mass. The goal was to arrive at the best understanding of the amplitude properties over the full mass range. The largest angular momentum amplitude required in the fits was the F wave, because there was no significant structure in the moments with $L \geq 7$. The mass intervals of the partial wave fits are listed in Table XII.1. The t' ranges of the fits were defined to be

$$\begin{aligned} 0.0 \leq t' \leq 0.8 \quad (\text{GeV}^2); \quad m_{\overline{K}_0\pi^-} \leq 1.6 \text{ GeV}/c^2, \\ 0.08 \leq t' \leq 0.8 \quad (\text{GeV}^2); \quad 1.6 \text{ GeV}/c^2 \leq m_{\overline{K}_0\pi^-}. \end{aligned} \quad (\text{XII.5})$$

Mass Interval	Wave Set
$0.64 \leq m_{K^0\pi} \leq 0.96$	S, P_0, P_-, P_+ (elastic S, P_0)
$0.96 \leq m_{K^0\pi} \leq 1.52$	$S, P_0, P_-, P_+, D_0, D_-, D_+$ (elastic S)
$1.52 \leq m_{K^0\pi} \leq 1.96$	$S, P_0, P_-, P_+, D_0, D_-, D_+, F_0, F_-, F_+$

Table XII.1. The mass dependence of the partial wave set

The analysis procedure was an iterative one. The division of the full mass range into the three regions reflects the dominance of a single natural parity amplitude, the leading amplitude, in each region. The characteristics of the leading amplitudes were established first. The amplitudes with the same L produced by unnatural parity exchange were also measured in each region. These amplitudes are due predominantly to π exchange, and have been measured with much higher statistical accuracy in the $K^-\pi^+\pi$ analysis of this same experiment. To take best advantage of this, these amplitudes were parametrized identically in both analyses; in particular the helicity 0 amplitudes in the physical region are expressed in terms of the corresponding on-mass-shell $K\pi$ scattering amplitudes. In all cases, the observed leading helicity 0 amplitudes were compatible with those obtained in the neutral K^* analysis, and hence could be fixed to those values to take advantage of that more precise measurement. This choice of constraint was checked against the observed moments of the present channel, to ensure that the resulting χ^2 remained reasonable. The observed moments were always well reproduced when the amplitudes were treated in this way.

The amplitude fixing permitted compatibility checks between the different leading mass regions. Tails of the leading amplitude fits from one mass region could be fixed in the adjacent lower mass region. The π tails were allowed into the moment fits, and the results checked for consistency with the results in the free case. The new answers improved the resolution on the underlying spectroscopy.

XII.2 Amplitudes and Phases

In general, all of the partial waves have $I = 1/2$ and $I = 3/2$ components

However, as discussed previously, only the S wave has been found to require any significant $I = 3/2$ contribution²⁴. The on shell S wave is known to be elastic for $m_{K^0\pi} \leq 1.5$ GeV/c², so that in principle, the elastic amplitudes can be written

$$a_L = \sqrt{2L+1} \sum_{k=1/2}^{3/2} C(G_k \epsilon_L^k \sin \delta_L^k e^{i\delta_L^k}) \quad (\text{XII.6})$$

with ϵ_L^k the elasticity, δ_L^k the phase shift for the L th wave, and CG_k are isospin Clebsch-Gordon coefficients. Elastic unitarity requires $\epsilon_L^k = 1$, so that the amplitude is given by δ_L^k . This should be contrasted with the inelastic case, where

$$a_L \equiv |a_L| e^{i\phi_L} \quad (\text{XII.7})$$

This inelastic parametrization is used for all amplitudes here except for the S wave and P_0 wave for $m_{K^0\pi} \leq 1.2$ GeV/c².

In the inelastic region, there is an overall phase in the amplitude solutions which is not determined and must be chosen. The choice is arbitrary, picked in order to make the interpretation of the results simple. The other phases in the amplitude fits are relative phases. That is, the absolute phase of any amplitude is not determined, and so some method of phase fixing must be adopted.

XII.3 Normalization

The absolute normalization of the present experiment can be calculated in the context of the Chew-Low model for $K-\pi$ scattering at the pion pole. The helicity 0 piece of the unnatural parity exchange cross section at the $K^*(892)$ is used to test the normalization, because the resonance is known to be elastic, i.e., there are no other decay channels open.

The Chew-Low differential cross section is given by the (XII.11). In general, the scattering amplitude must include the piece due to production of isospin $3/2$, although for the P_0 wave, no significant contribution has been measured. This implies, e.g. for $K^-\pi^+$ scattering, that

$$f(\Omega) = \frac{2}{3} f^{1/2}(\Omega) + \frac{1}{3} f^{3/2}(\Omega) = \frac{2}{3} f^{1/2}(\Omega) \quad (\text{XII.8})$$

Using an exponential to describe the vertex form factor dependence, the Chew-Low equation may be written as

$$\frac{d^4\sigma}{dM dt d\Omega} = \frac{(G^2/4\pi)}{p_{lab}^2 M_p^2} \cdot \frac{te^{-\alpha(1+m_\pi^2)}}{(1+m_\pi^2)^2} \cdot \frac{2M^2}{9q} |a_S Y_0^0 + a_P Y_1^0|^2, \quad (XII.9)$$

where another factor of 2 has been divided out to account for the CG of the $\bar{K}\pi$ final state. This can be related to the measured four-fold differential cross section, given by

$$\frac{d^4\sigma_m}{dM dt d\Omega} = \frac{S}{\Delta t \Delta M} \left(\frac{2.56814 \text{ GeV}^{-2}}{1000 \mu\text{b}} \right) \sum_{l0} Y_N^{l0} Y_{l0}(\Omega), \quad (XII.10)$$

where S is the experimental sensitivity of Chapter IX, and $\Delta t, \Delta M$ are the 4-momentum transfer and mass intervals in the region where the comparison is to be made. The amplitude program fits a form

$$|FN|^2 \frac{M^2}{q} |a_S Y_0^0 + a_P Y_1^0|^2 \quad (XII.11)$$

to the measured quantities, with the same t dependence, so that if unitarity at the pion pole is to hold, FN should be given by

$$FN = \frac{1000 \Delta t \Delta M}{2.56814 p_{lab}^2 M_p^2} \frac{2 G^2}{9 4\pi}. \quad (XII.12)$$

The fits are done in 20 MeV/c² mass intervals at the $K^*(892)$, and the t' interval is chosen to be 0-0.2 GeV², as for the σ_0 plot of figure XI.8. The arithmetic gives $FN=18.725$, using the latest published measurement for the value of the $pp\pi^0$ coupling constant^[93], $\frac{G^2}{4\pi} = 13.1 \pm 0.1$. With this value for the FN parameter, the observed peak value of the P_0 amplitude should be consistent with one. The helicity 0 P wave amplitude is shown in Figure XII.1. Clearly, the unitarity bound is satisfied. If the old value for the coupling constant is used, the extrapolated pole amplitude in the figure drops to ≈ 0.92 , consistent with the change in the coupling constant. The curve in the figure is a BW fit to the data where a radius parameter of 3. GeV⁻¹ was used, as in the $K^-\pi^+\pi^-$ analysis. The fit found $m_{K\pi}(P_0) = 0.893 \pm 0.002$ and $F_{P_0} = 52.7 \pm 6 \text{ MeV}/c^2$.

Figure XII.1: The P_0 amplitude in the $K^*(892)$ region. Unitarity would imply that this amplitude approach 1 at the peak of the resonance.

XII.4 The Natural Parity Exchange Amplitudes

The leading natural parity amplitudes are shown in Figure XII.2. These waves are responsible for the structure in the moments with the largest l value in the respective mass slices. The curves are fits to the data points using the parametrization of (XI.28). The fit results are shown in Table XII.2. As discussed earlier, there are mathematical ambiguities, and this implies that in each bin there can be several different acceptable fits, the number of which increases with the number of partial waves allowed into the fit. It was observed that in the present channel, the leading natural parity amplitude solutions were approximately unique, i.e., while there might be several distinct solutions for the unnatural parity and non-leading natural parity amplitudes, these leading waves varied only slightly for the different solutions.

XII.4.a The $1^- K^*(892)$

The Particle Data Group lists $0.89211 \pm 0.00032 \text{ GeV}/c^2$ as the mass of the charged $K^*(892)$, somewhat above the value reported here. The amplitude however is quite consistent with the mass as determined from the moments (see Chapter

Figure XII.2: The leading natural amplitudes. The scale is arbitrary, as discussed in the text, but the errors indicate the relative statistical precision of each measurement.

XI). The discrepancy is not considered significant. The width is also smaller than the world average of 51.13 ± 9.34 GeV/c². Recall that this analysis has removed the resolution smearing due to the spectrometer, and in that sense, this narrower width may be more representative of the actual value. In any case, the resonance is clearly the $K^*(892)$.

The value of the neutral $K^*(892)$ mass measured in this experiment was $0.8948 \pm 0.6 \pm 0.2$ with a width of $0.0494 \pm 0.9 \pm 0.2^{[1]}$. This gives a neutral/charged mass difference of 5.4 ± 0.6 MeV/c², consistent with the PDG's 6.7 ± 1.2 MeV/c².

Resonance	Mass (GeV/c ²)	Width (GeV/c ²)	Radius (GeV ⁻¹)
$K^*(892)$	$0.8904 \pm 0.0002 \pm 0.0005$	$0.0452 \pm 0.001 \pm 0.002$	$12.1 \pm 3.2 \pm 3.0$
$K_2^*(1430)$	$1.4234 \pm 0.002 \pm 0.003$	$0.098 \pm 0.004 \pm 0.005$	$4.80 \pm 2.3 \pm 3.$
$K_3^*(1780)$	$1.720 \pm 10. \pm 15.$	$0.187 \pm 0.031 \pm 0.020$	$8.5 \pm 3. \pm 10.$

Table XII.2: The leading natural parity amplitudes. A single Breit-Wigner resonance is assumed for each wave.

XII.4.b The $2^+ K_2^*(1430)$

The value for the $K_2^*(1430)$ mass is consistent with the PDG's value of 1.4256 ± 1.5 GeV/c² and width of 0.99 ± 0.0024 GeV/c², and is consistent with the measurements of Chapter XI. This peak is clearly identified. Using the value of the neutral $K_2^*(1430)$ mass from this experiment, $1.4316 \pm 0.0018 \pm 0.0007$ and width $0.1165 \pm 0.006 \pm 0.002$, gives a value for the neutral/charged mass difference of 8.4 ± 3.0 MeV/c², consistent with the $K^*(892)$ mass difference.

XII.4.c The $3^- K_3^*(1780)$

The mass value derived from the amplitude analysis is below the PDG value for the $K_3^*(1780)$ of 1.780 ± 1.6 GeV/c². It is also lower than any of the moment measurements of Chapter XI. This discrepancy is due to the fact that this mass region, as will be shown, is not completely dominated by production of the leading resonance. There is significant 1^- resonance production in this region, as well as 2^+ background. There is also the set of unnatural parity amplitudes contributing to the cross section in this region. For these reasons, the values of the 3^- parameters are much more difficult to determine. The N^* cut necessary to extract the signal, also reduces the a priori statistical power of the data. Most importantly, with respect to the neutral K^* system, the cross section due to leading amplitude production is greatly decreased in this mass region. The measurement of the t_{52} moment is not dependent upon the complicated correlations which are introduced by the amplitude analysis, so that the corresponding resonance parameters are inherently better defined. The mass value of this resonance will subsequently be defined accordingly.

The curve plotted in Figure XII.2 is for the $1.74 \text{ GeV}/c^2$ mass value. The χ^2 change is $6.5/8$ instead of $3.4/7$. The mass value is still lower than that of the FDG, but is consistent with the recent measurements from the present experiment [68][71]. A systematic error of $15 \text{ MeV}/c^2$ is estimated on the basis of the sensitivity of the fitted mass value to variations on the N^* cut, binning changes, and choice of amplitude solutions. This wave was well defined in all cases, and is clearly due to production of $K^*(1780)$. The width was also subject to the same problems, but it too was also well defined and consistent with the value determined by fitting the t_{62} moment.

XII.5 The Fit to the Natural Amplitudes

With the parameters of the leading amplitudes established, the underlying spectroscopy could be examined. The mass dependence of the natural parity amplitudes from threshold to $2 \text{ GeV}/c^2$ is shown in Figure XII.3. The plotted points are the results from the amplitude analysis when as many as 30 random starts are made in a mass bin, and the amplitudes which are allowed to contribute to a fit in a mass bin are as given in Table XII.1. There were no ambiguities found in the natural parity waves for mass values below $1.6 \text{ GeV}/c^2$. Above this mass, there were two distinct choices per bin in the 1^- and 2^+ waves. The alternate solutions are shown as the open dots at high mass in the figures. These points, while good solutions to the moment equations, were not used in the mass dependent fits responsible for the overlaid curves in the figure. Such fits were attempted but never gave reasonable behavior in terms of the simple model of resonance production used here. While it is certainly mathematically inappropriate to ignore these solutions, the fact that the first choice of points showed smooth, reasonable behavior which provided results consistent with other measurements of resonance production, was thought to justify this choice as the more probable physical solution.

The lack of obvious structure in other than the leading amplitudes of Figure XII.3 is offset by the elegant structure observed in the natural parity relative phase behavior. The measured phases, $\varphi_P - \varphi_D$ and $\varphi_D - \varphi_F$, as required by the moment-amplitude relations of Appendix A, are plotted in Figure XII.4.

The model used to describe the data of figures XII.3 and XII.4 involves the superposition of interfering BW amplitudes with arbitrary relative production phases and no background. The fit requires 6 Breit-Wigner resonances, each with a nor-

Figure XII.3: The natural parity amplitudes. In (a), the 1^- amplitude, here the P_+ wave is shown, in (b), the 2^+ amplitude, or the D_+ wave, and (c) the 3^- amplitude, or F_+ wave. The fitted curves are the results of a simultaneous BW fit to the data.

malization, a mass, and a width parameter, 5 additional production phases, and a common radius parameter in the relevant Blatt-Weisskopf barrier factor. The results are sensitive to this radius, as it has a significant effect on the relative phase behavior of the interfering amplitudes, due to its effect on resonance tails. The fit results prefer a value of $R \approx 10 \text{ GeV}^{-1}$, approximately 2 fermi. The model incor-

Figure XII.4: The two natural parity relative phases. In (a) the $\phi_P - \phi_D$ phase, and in (b) the $\phi_D - \phi_P$ phase. The dashed curve indicates the expected phase behavior if the second P wave BW is excluded. The dotted curve indicates the expected phase behavior if the second D wave BW is excluded.

porates three 1^- BWs, two 2^+ BWs, and the single 3^- BW. A given amplitude is parametrized as

$$a_L \propto \sum_i BW_L^i e^{i\phi_L^i}, \quad (\text{XII.13})$$

where ϕ_L^i is the production phase of the k^{th} BW. This phase was set to 0 for the $K^*(892)$, consistent with the overall phase freedom described earlier. The fit shown gave a $\chi^2=148/86$. The mass and width of the leading amplitudes were fixed at the values tabulated in Table XII.2. The production phases of the $K^*(890)$ and the $K_2^*(1780)$ were set to 0, as suggested by $SU(3)$ coupling arguments, and by results from preliminary fits.

XII.5.a The P_+ Amplitude

The P_+ wave is clearly dominated by the production of the $K^*(892)$. At higher mass, the distribution of the amplitude is smooth and featureless, but clearly non-zero. The phase behavior is clearly indicative of additional resonance behavior, particularly near $1700 \text{ MeV}/c^2$. The need for a third resonance near $1400 \text{ MeV}/c^2$ is indicated by the phase behavior as well. The dashed curve represents the expected phase behavior for the coherent addition of just the $K^*(892)$ and 1700 resonances. Clearly, the phase of the P_+ wave in the 1400 region is systematically lower than would be expected. There is a small "bump" in the distribution of the magnitude in the same region. The addition of a third resonance describes the data well. The parameters of the resonances as determined from the fit are shown in Table XII.3.

Fit to the $1^- P_+$ Amplitude				
Wave	Normalization	Mass (GeV/c^2)	Width (GeV/c^2)	Phase (deg)
P_+^1	1.677 ± 0.010	0.8905 (fixed)	0.045 (fixed)	0 (fixed)
P_+^2	0.122 ± 0.019	1.367 ± 0.054	0.114 ± 0.101	69 ± 7
P_+^3	0.190 ± 0.008	1.678 ± 0.061	0.454 ± 0.270	-92 ± 17

Table XII.3: The P_+ resonance parameters. The subscript refers to the resonances discussed in the text in order of ascending mass

XII.5.b The D_+ Amplitude

The D_+ wave is clearly dominated by the production of the $K_2^*(1430)$. At higher mass, the distribution of the amplitude drops relative to the behavior expected for the high mass tail of the $K_2^*(1430)$, shown as the dotted line in the D_+ wave of Figure XII.3. In addition, the $D_+ - P_+$ relative phase (Fig. XII.4(b)) is quite different from that expected for the $K_2^*(1430)$ and $K_2^*(1780)$ resonances alone (dotted curve in Fig. XII.4(b)). These effects require the inclusion of an additional resonance D wave resonance near $1950 \text{ MeV}/c^2$. The parameters found for the pair of resonances required are shown in Table XII.4.

Fit to the $2^+ D_+$ Amplitude				
Wave	Normalization	Mass (GeV/c ²)	Width (GeV/c ²)	Phase (deg)
D_+^1	0.780 ± 0.012	1.425 (fixed)	0.1 (fixed)	$40. \pm 4.$
D_+^2	0.124 ± 0.009	1.978 ± 0.040	0.398 ± 0.047	$2. \pm 5.$

Table XII.4: The D_+ resonance parameters. The subscript refers to the resonances discussed in the text in order of ascending mass.

XII.5.c The F_+ Amplitude

The F_+ wave is specified using the parameter values obtained from the t_{62} moment, i.e., $m_F = 1.74$ GeV/c² and $\Gamma_F = 0.195$ GeV/c². The fit found that a normalization value of 0.434 ± 0.017 was required.

XII.5.d Discussion

The parameters of the 2 additional P wave resonances are consistent with those seen in the $\bar{K}_0^+ \pi^+ \pi^- n$ and $K^- \pi^+ n$ analyses from the present experiment. That the widths of the pair are not well determined is a reflection of the lack of an obvious "bump" structure in the amplitude plot of Figure XII.3. The higher lying recurrence is most naturally identified with the $1^3 D_1$ member of the $L = 2$ triplet to which the $K_2^*(1780)$ belongs. The interpretation of the second resonance as the $2^3 S_1$ state, the radial excitation of the $K^*(890)$ is attractive.

The parameters of the second D -wave resonance are in excellent agreement with those obtained in the partial wave analysis of the $\bar{K}_0^+ \pi^+ \pi^- n$ system produced in the reaction $K^- p \rightarrow \bar{K}_0^+ \pi^+ \pi^- n$ in the present experiment^[17]. The interpretation of the resonance in the context of the quark model is somewhat uncertain. It may be the radial excitation of the $K_2^*(1430)$ or it may be the 2^+ member of the $1^3 F$ ground state triplet in the quark model, spin 4 family. The mass splitting would make it more compatible with the first interpretation, as does the mass value, which is predicted to be in this range^[53].

The mass dependence of the individual amplitudes and phases are shown in the Argand plots of Figure XII.5.

Figure XII.5: The Argand plots. The observed mass dependence of the natural parity amplitudes and their phases are plotted. The production phase has been suppressed.

XII.6 The Unnatural Parity Exchange Amplitudes

Contributions due to the unnatural parity exchange amplitudes, i.e., the L_0, L_- waves, are also present in the data sample. One example has already been shown, namely the P_0 amplitude in the $K^*(890)$ region, which establishes the normalization of the experiment. There are 3 other helicity 0 amplitudes, the S, D_0 , and F_0 waves. The latter two were used to help constrain the possible solutions in the D wave and F wave regions, so they are not independently measured. However, the S wave is determined. It is interesting because of the $I = 3/2$ component which leads to differences in the total S wave shape between the $K^- \pi^+ n$ and $\bar{K}_0^+ \pi^- p$ channels.

Also, both the S and P_0 waves have been observed to resonate again at high mass^[71], and it is of interest to know if such effects can be detected in the present channel.

The helicity 1 unnatural parity amplitudes are not shown. They are very small here, reaching values on the order of $1/2$ of the helicity 0 amplitudes at the resonance positions.

XII.6.a The Low Mass S Wave

It is important to realize that the measured S wave in either the $K^-\pi^+n$ or the $\bar{K}_0^-\pi^+p$ channel is the total S wave amplitude. That is, as was observed in $K^+p \rightarrow K^+\pi^+n$ and $K^-p \rightarrow K^-\pi^-\Delta^{++}$, there is a significant $I = 3/2$ component contributing to the S wave. The $I = 3/2$ wave enters differently in the $\bar{K}_0^-\pi^+p$ and $K^-\pi^+n$ channels and this results in dramatic effects in the interference moments, which contrast sharply between the two final states. The $I = 3/2$ amplitude is well described by an effective range parameterization^[52],

$$\cot \delta = \frac{1}{(a_S)q} + \frac{1}{2}(r_S)q, \quad (\text{XII.14})$$

where the best fits gave $a_S = -1.03 \text{ (GeV/c)}^{-1}$ and $r_S = -0.94 \text{ (GeV/c)}^{-1}$. The contribution to the total S wave of the $I = 3/2$ part is different between $K^-\pi^+n$ and $\bar{K}_0^-\pi^+p$ because the isospin Clebsch-Gordon coefficients are different. Explicitly,

$$a_S^T(K^-\pi^+) = \frac{2}{3} \left(a_S^{I=1/2}(K^-\pi^+) + \frac{1}{2}a_S^{I=3/2}(K^-\pi^+) \right), \quad (\text{XII.15})$$

and

$$a_S^T(\bar{K}_0^-\pi^+) = \frac{\sqrt{2}}{3} \left(a_S^{I=3/2}(\bar{K}_0^-\pi^+) - a_S^{I=1/2}(\bar{K}_0^-\pi^+) \right). \quad (\text{XII.16})$$

From the $K^-\pi^+n$ analysis where the S wave is well measured, and using the effective range parametrization for the $I = 3/2$, the behavior of the S wave in the present channel can be calculated. A comparison between this calculation and the S wave for $m_{\bar{K}_0^-\pi^+} \leq 1.6 \text{ GeV}/c^2$ found in the present analysis is shown in Figure XII.6, where elastic unitarity has been assumed for $m_{\bar{K}_0^-\pi^+} \leq 1.5 \text{ GeV}/c^2$. The agreement is good, and again demonstrates the consistency between the helicity 0 amplitudes in the $K^-\pi^+n$ and the $\bar{K}_0^-\pi^+p$ final states.

* In principle, a new, more precise measurement of these parameters could be extracted from the E-135 data sample, but this is not attempted here.

Unnatural Parity Amplitudes

Figure XII.6: The total S wave amplitude and phase.

XII.6.b High Mass S, P_0

In the high mass region it is even more difficult to extract the unnatural parity amplitudes, due to the t' cut at 0.08 GeV^2 . The S, P_0 amplitudes are nonetheless shown in Figure XII.7. The $K^-\pi^+n$ analysis found evidence for resonance production in P_0 at $\approx 1700 \text{ MeV}/c^2$ with an elasticity of ≈ 0.4 , and in the S wave at $\approx 1900 \text{ MeV}/c^2$ with an elasticity on the order of 0.5. These amplitudes, although poorly measured here, are seen to be consistent with the values expected on the basis of the $K^-\pi^+n$ analysis^[71].

Figure XII.7: S, P_0 at high mass. In (a) the S wave, and (b) the P_0 wave.

XII.7 t' Dependence

In the present analysis, the only t' information which can be reliably extracted is for the leading helicity 1, natural parity amplitudes. Also, only the behavior of the leading amplitudes in a given mass slice is well measured. When the underlying natural parity amplitudes are significant, the t' dependence of the leading amplitude is assumed.

The t' behavior of the leading amplitudes is shown in Figure XII.8, and the mass windows and fit parameters are given in Table XII.5. The fit to the data is of the form

$$a_L(t') \propto (t')^{\frac{n}{2}} e^{-\alpha t'}. \quad (\text{XII.17})$$

This form is the consequence of angular momentum conservation in the forward direction, (i.e., $t' \rightarrow 0$), and n is the number of units of net helicity flip in the reaction. Since the predominant production mechanism involves isoscalar t -channel exchange (i.e., ω, f exchange), n is expected to take the value 1, i.e., the natural

parity exchange amplitudes would be expected to be $\sim \sqrt{t'}$ at small t' . The results for the D_+ and F_+ amplitudes are consistent with this. However, the value for the P_+ amplitude, although close to 0.5, is substantially different from it. This may be due to the contributions from $\rho - a_2$ exchange as observed for the $K^- \pi^+ n$ final state^[71].

Figure XII.8: The t' behavior of the leading amplitudes. The fitting function is described in the text.

Mass	Amplitude	Power	Slope	χ^2/dof
$0.84 \leq m_{\bar{K}_0^*} \leq 0.94$	P_+	0.465 ± 0.003	1.678 ± 0.02	45/7
		0.5	1.531 ± 0.01	99/8
$1.3 \leq m_{\bar{K}_0^*} \leq 1.54$	D_+	0.535 ± 0.05	1.481 ± 0.21	6.42/7
		0.5	1.34 ± 0.08	7/8
$1.6 \leq m_{\bar{K}_0^*} \leq 1.95$	F_+	0.854 ± 0.29	1.66 ± 0.92	2.3/5
		0.5	3.58 ± 0.26	3.7/6

Table XII.5: Results from the t' fits. The parameters when the power is free/fixed are shown.

XII.8 Cross Sections

From the masses, widths, and t' dependence of the amplitudes, the individual contributions to the total observed intensity can be calculated. Here, the mass dependence specified by the Breit-Wigner is integrated from threshold to $3 \text{ GeV}/c^2$, and the t' dependence is integrated from 0 to 10 GeV^2 . These account for the unseen tails of the cross sections. The cross sections are tabulated in Table XII.6.

XII.8.a The $K\eta/K\pi$ Branching Ratio

A result of some interest involves the determination of the $K^* \rightarrow K\eta$ to $K^* \rightarrow K\pi$ branching ratio at the $K_2^*(1430)$ and $K_3^*(1780)$ resonances. The mass spectrum in $\bar{K}^{*0} \rightarrow K^-\eta$ has been seen to be dominated by the production of $K_3^*(1780)$, with no indication of $K_2^*(1430)$ production. This result is expected in flavor SU(3) symmetry, if the η/η' mixing angle is on the order of -20° [94].

The branching ratio has already been reported [95], but it was based on a much earlier measurement of $K^* \rightarrow K\pi$ from a different experiment. With the present analysis, a new branching ratio value can be determined. The branching ratio calculation relates the ratio of produced events between the two channels of interest. In order to treat the data as democratically as possible, the $K\eta$ data

Resonance	Produced Events	Cross Section (μb)
$K^*(890) \eta_+$	71714 ± 718	105 ± 1
$K^*(890) \eta_-$	7995 ± 297	12 ± 0.5
$K^*(890)_{\text{tot}}$	79709 ± 801	117.2 ± 1.2
$K_1^*(1400) \eta_+$	1100 ± 367	1.6 ± 0.4
$K_2^*(1430) \eta_+$	18739 ± 606	27.5 ± 1
$K_2^*(1430) \eta_-$	6070 ± 432	8.9 ± 0.6
$K_2^*(1430)_{\text{tot}}$	24809 ± 820	36.4 ± 1.5
$K_1^*(1700) \eta_+$	1450 ± 167	2.1 ± 0.2
$K_3^*(1780) \eta_+$	3467 ± 423	5.1 ± 0.6
$K_3^*(1780) \eta_-$	2644 ± 727	3.9 ± 1
$K_3^*(1780)_{\text{tot}}$	6111 ± 750	9.0 ± 1.1
$K_2^*(1950) \eta_+$	241 ± 47	0.4 ± 0.1

Table XII.6: Resonance production cross sections. The unnatural parity cross sections were determined using the elasticity values from the $K^-\pi^+\eta$ analysis, which are consistent with the observed values from the present analysis, multiplied by a factor 1.25 to account for the smaller helicity 1 contribution. The quoted values correspond only to the $K\pi$ decay.

were refit with the present BW resonance parametrization with $R = 10 \text{ GeV}^{-1}$ to allow a direct comparison with the $\bar{K}_0^*\pi^-$ data. The results of the fits are shown in Figure XII.9, where the $K\eta$ data have been reduced to account for the difference between the upper limit of the t' sampling interval between the two channels. The fit parameters under these conditions are listed in Table XII.7. The t' slope in the $K\eta$ case was determined to be $9.2 \pm 0.8 (\text{GeV}/c)^2$, appropriate for the intensity distribution of that work. Here, the t' slope from the amplitude was $4 \pm 0.5 \text{ GeV}^2$, as given earlier, or $8 \pm 1 \text{ GeV}^2$ in the intensity. The numbers are consistent within error. A common value of 8.5 GeV^2 is assumed. With these values, the mass

dependence as given by the BW is integrated from threshold to $2 \text{ GeV}/c^2$. Both integrals must be corrected for the unseen decays, a factor of 0.237 in the $K\eta$ case, and 0.67 in the $\bar{K}_0\pi$ case. Letting I be the integration, the branching ratio is found to be

$$\frac{\sigma(K^{*-}(1780)) \cdot BF(K_3^{*-}(1780) \rightarrow K^-\eta)}{\sigma(K^{*-}(1780)) \cdot BF(K_3^{*-}(1780) \rightarrow K^-\pi)} = \frac{I_{K\eta}}{I_{\bar{K}_0\pi}} \frac{0.67}{0.237} = 0.41 \pm 0.053. \quad (\text{XII.18})$$

This should be compared with the value 0.50 ± 0.13 obtained previously^[68]. That number was based on a simple parametrization of the beam momentum dependence of the total cross-section, as opposed to a direct comparison of like production amplitudes. A simple ratio from the sum of the points shown in Figure XII.9 gives 0.455 ± 0.063 , in agreement with the result obtained with the fit.

Figure XII.9: $|F_+|^2$ in the $K\eta$ and $\bar{K}_0\pi$ channels. The radius parameter is fixed to 10. The data have been corrected for the unseen decay modes.

Parameter	$K\eta$ Value	$\bar{K}_0\pi$ Value
Peak Value	67 ± 6	92 ± 8
Mass (GeV/c^2)	1.74	1.74
Width (GeV/c^2)	0.195	0.195

Table XII.7: Fit results for $K\eta$ to $K\pi$ branching ratio calculation. The radius parameter is fixed at 10.

XII.9 Evidence for Higher Mass Production

Recent results from studies in $s\bar{s}$ channels have shown that the effect of diffractive N^* production can interfere with meson production to enhance the observation of meson resonance structure in the region of the N^* overlap^{[87][88]}. In principle, the same effect could be present in the $\bar{K}_0\pi^-p$ final state, even though the N^* system is not produced diffractively.

In the region of $\bar{K}_0\pi^-$ mass above $2 \text{ GeV}/c^2$, where spin 4 and 5 K^* 's would be expected^{[73][77][89]}, there is no evidence in the unnormalized moments for the production of such states. This implies that the corresponding production amplitudes are small. However, in the region of N^* overlap, the larger N^* amplitudes might produce effects in the corresponding $\bar{K}_0\pi^-$ mass spectrum, through interference effects with the small $K\pi$ amplitudes. In particular, because the average phase of a BW amplitude is 90° , an interference effect might take the form of the imaginary part of the $\bar{K}_0\pi^-$ BW amplitude, producing a bump close to the mass of a resonance in the $\bar{K}_0\pi^-$ mass spectrum.

In Figure XII.10b, the $p\pi^-$ mass spectrum for $M_{p\pi^-} \leq 1.8 \text{ GeV}/c^2$ exhibits significant Δ and N^* production. The corresponding $\bar{K}_0\pi^-$ mass spectrum in Figure XII.10a shows that in addition to the clear $K_2^*(1430)$ and $K_3^*(1780)$ peaks, that there is an excess of ~ 200 events in the $2.02\text{--}2.14 \text{ GeV}/c^2$ mass range. Although the spin-parity associated with the enhancement is not determined here, it is plausible that the events are due to the production of the $K_4^*(2050)$ via the $K^* \rightarrow N^*$ interference mechanism.

Figure XII.10: $m_{K^0\pi^-}$ in the N^* overlap. The bump near $1.95 \text{ GeV}/c^2$ is indicative of spin 4 production.

If the bump is due to G wave production, it would follow that there could be interference effects between the G wave and other smaller J amplitudes in the region below $2 \text{ GeV}/c^2$. In fact, the final point in Figure XII.2 is near zero in the $|F_+|$ solution, due to the t_{62} moment being positive in that mass bin. The moment would be expected to be positive if there were a significant contribution from $D_+ - G_+$ interference.

XIII. Conclusions

The analysis of the $\bar{K}_S^0\pi^-$ meson system produced in the reaction $K^-p \rightarrow \bar{K}_S^0\pi^-p$ at $11 \text{ GeV}/c$ incident momentum has been described. The data sample is the largest of its kind, to date, and the first to combine the large statistics typical of a spectrometer experiment with the uniform acceptance not unlike that of a bubble chamber experiment.

The simple event topology, consisting of a single V^0 and no more than 2 additional tracks, the large acceptance, and the use of a sophisticated kinematic fitting program result in a very pure data sample. That this is the case is indicated by the lack of any clear need for background terms in the subsequent amplitude fits. These amplitudes show clear $K^*(890)$, $K_2^*(1430)$, and $K_2^*(1780)$ production, predominantly by means of natural parity isoscalar exchange in the t -channel. Additional vector states at $\approx 1.4 \text{ GeV}/c^2$ and $\approx 1.7 \text{ GeV}/c^2$ are observed, and evidence for the production of a higher lying 2^+ amplitude near $\approx 1.95 \text{ GeV}/c^2$ is obtained.

The 1^- partial waves show resonant structure at a mass of $1.73 \text{ GeV}/c^2$, which is identified as resulting from the production and decay of the 1^- member of the $L = 2$ multiplet. A second 1^- resonant state is observed at a mass of ≈ 1.42 . This resonance, strongly produced in $\bar{K}_S^0\pi^-\pi^+n$ reactions, is produced only weakly, consistent with results from a recent analysis in the $K^-\pi^+n$ channel^[71]. The properties of this state are consistent with its assignment as the first radial excitation of the $1^- K^*(892)$. These data confirm the results of a prior analysis of the $\bar{K}_S^0\pi^-\pi^+n$ final state,^[72] where these two states were observed and called the $K_1^*(1790)$ and the $K_1^*(1410)$. A lower mass and larger width for the $K_1^*(1790)$ is inferred from this current experiment, although the partial wave data appear to be consistent with the prior one^[72].

In addition to the well-known $K_2^*(1430)$ resonance, the D_+ amplitude exhibits resonant amplitude and phase behavior at $\approx 1.95 \text{ GeV}/c^2$. The resonance parameters are in agreement with those obtained previously in this experiment from the partial wave analysis of the $\bar{K}_S^0\pi^+\pi^-$ system recoiling against a neutron. This state is probably the first radial recurrence of the $K_2^*(1430)$, although it could also be

the 2^+ partner of the $K_4^*(2050)$ in the 3F $q\bar{q}$ ground state triplet.

The 3^- $K_3^*(1780)$ resonance, one of the $L = 2$ triplet states, is seen, at a mass of ≈ 1.72 GeV/c², below the Particle Data Group value, but not inconsistent with other recent measurements.

A $\bar{K}_0\pi^-$ enhancement around 2.1 GeV/c² is observed. Although it is not possible to show that this due to resonance production, the interpretation that it results from an interference effect between the N^* amplitudes and that describing the production and decay of the 4^+ $K_4^*(2060)$ is consistent with the data.

Appendix A. Moment-Amplitude Relations

The set of equations relating the observed moment structure to the underlying amplitude behavior is listed here. Spectroscopic notation is used to define the amplitudes, with an additional subscript to indicate naturality. A factor of $\sqrt{2L+1}$ is implicit in each amplitude. Note that the physics of interest here is that of strong interactions, so that parity is a good quantum number. As such, only bilinear combinations where both terms have natural parity, or both unnatural parity, can occur. Other relevant definitions are given in the first section of Chapter XII.

$$\begin{aligned}
 t_{00} &= |S|^2 + |P_0|^2 + |P_-|^2 + |P_+|^2 + |D_0|^2 + |D_-|^2 + |D_+|^2 + |F_0|^2 + |F_-|^2 + |F_+|^2 \\
 t_{10} &= 2|S||P_0|\cos(\varphi_S - \varphi_{P_0}) + \frac{1}{\sqrt{2}}|P_0||D_0|\cos(\varphi_{P_0} - \varphi_{D_0}) \\
 &\quad + \frac{2\sqrt{2}}{3}(|P_+||D_+|\cos(\varphi_{P_+} - \varphi_{D_+}) + |P_-||D_-|\cos(\varphi_{P_-} - \varphi_{D_-})) \\
 &\quad + \frac{2\sqrt{2}}{3}|D_0||F_0|\cos(\varphi_{D_0} - \varphi_{F_0}) + \frac{2\sqrt{2}}{3}(|D_+||F_+|\cos(\varphi_{D_+} - \varphi_{F_+}) + |D_-||F_-|\cos(\varphi_{D_-} - \varphi_{F_-})) \\
 t_{11} &= \sqrt{2}|S||P_-|\cos(\varphi_S - \varphi_{P_-}) + \sqrt{\frac{2}{3}}|P_0||D_-|\cos(\varphi_{P_0} - \varphi_{D_-}) - \sqrt{\frac{2}{3}}|P_-||D_0|\cos(\varphi_{P_-} - \varphi_{D_0}) \\
 &\quad + \frac{1}{\sqrt{3}}|D_0||F_-|\cos(\varphi_{D_0} - \varphi_{F_-}) - \frac{2\sqrt{2}}{3}|D_-||F_0|\cos(\varphi_{D_-} - \varphi_{F_0}) \\
 t_{20} &= 2|S||D_0|\cos(\varphi_S - \varphi_{D_0}) + \frac{3}{2}|P_0|^2 - \frac{1}{2}(|P_+|^2 + |P_-|^2) + \frac{2\sqrt{2}}{3}|D_0|^2 + \frac{\sqrt{2}}{3}(|D_+|^2 + |D_-|^2) \\
 &\quad + \frac{1}{\sqrt{2}}(|F_+|^2 + |F_-|^2) + \frac{1}{3}|P_0||F_0|^2 + 6\sqrt{\frac{2}{3}}|P_0||F_0|\cos(\varphi_{P_0} - \varphi_{F_0}) \\
 &\quad + 6\sqrt{\frac{2}{3}}(|P_+||F_+|\cos(\varphi_{P_+} - \varphi_{F_+}) + |P_-||F_-|\cos(\varphi_{P_-} - \varphi_{F_-})) \\
 t_{21} &= \frac{\sqrt{2}}{3}|P_0||P_-|\cos(\varphi_{P_0} - \varphi_{P_-}) + \frac{\sqrt{18}}{3}|D_0||D_-|\cos(\varphi_{D_0} - \varphi_{D_-}) + \sqrt{2}|S||D_-|\cos(\varphi_S - \varphi_{D_-}) \\
 &\quad + \frac{1}{\sqrt{3}}|P_0||F_-|\cos(\varphi_{P_0} - \varphi_{F_-}) + \frac{2\sqrt{2}}{3}|F_0||F_-|\cos(\varphi_{F_0} - \varphi_{F_-}) - \frac{2\sqrt{2}}{3}|P_-||F_0|\cos(\varphi_{P_-} - \varphi_{F_0}) \\
 t_{22} &= -\frac{1}{\sqrt{18}}(|P_+|^2 - |P_-|^2) - \frac{\sqrt{18}}{3}(|D_+|^2 - |D_-|^2) \\
 &\quad + \sqrt{\frac{2}{3}}(|P_0||F_+|\cos(\varphi_{P_0} - \varphi_{F_+}) - |P_-||F_-|\cos(\varphi_{P_-} - \varphi_{F_-})) - \sqrt{\frac{2}{15}}(|F_+|^2 - |F_-|^2) \\
 t_{30} &= \frac{5\sqrt{2}}{3}|P_0||D_0|\cos(\varphi_{P_0} - \varphi_{D_0}) - \frac{6}{\sqrt{38}}(|P_+||D_+|\cos(\varphi_{P_+} - \varphi_{D_+}) + |P_-||D_-|\cos(\varphi_{P_-} - \varphi_{D_-})) \\
 &\quad + 2|S||F_0|\cos(\varphi_S - \varphi_{F_0}) + \frac{5\sqrt{2}}{12}|D_0||F_0|\cos(\varphi_{D_0} - \varphi_{F_0}) \\
 &\quad + \frac{2\sqrt{18}}{3}(|D_+||F_+|\cos(\varphi_{D_+} - \varphi_{F_+}) + |D_-||F_-|\cos(\varphi_{D_-} - \varphi_{F_-})) \\
 t_{31} &= \frac{5\sqrt{2}}{3}|P_0||D_-|\cos(\varphi_{P_0} - \varphi_{D_-}) + \frac{1}{\sqrt{38}}|P_-||D_0|\cos(\varphi_{P_-} - \varphi_{D_0}) \\
 &\quad + \sqrt{2}|S||F_-|\cos(\varphi_S - \varphi_{F_-}) + \frac{2\sqrt{2}}{15}|D_-||F_0|\cos(\varphi_{D_-} - \varphi_{F_0}) + \sqrt{\frac{2}{3}}|D_0||F_-|\cos(\varphi_{D_0} - \varphi_{F_-}) \\
 t_{32} &= -\sqrt{\frac{2}{3}}(|P_+||D_+|\cos(\varphi_{P_+} - \varphi_{D_+}) - |P_-||D_-|\cos(\varphi_{P_-} - \varphi_{D_-})) \\
 &\quad - \frac{1}{\sqrt{3}}(|D_+||F_+|\cos(\varphi_{D_+} - \varphi_{F_+}) - |D_-||F_-|\cos(\varphi_{D_-} - \varphi_{F_-})) \\
 t_{40} &= -\frac{4}{3}(|D_+|^2 + |D_-|^2) + \frac{2}{3}|D_0|^2 + \frac{1}{11}(|F_+|^2 + |F_-|^2) + \frac{6}{11}|F_0|^2
 \end{aligned}$$

$$\begin{aligned}
& + \frac{A}{\sqrt{3}} |F_0| |F_0| \cos(\varphi_{D_+} - \varphi_{D_-}) - \frac{\sqrt{2}}{\sqrt{3}} (|P_+| |F_+| \cos(\varphi_{D_+} - \varphi_{F_+}) + |P_-| |F_-| \cos(\varphi_{D_-} - \varphi_{F_-})) \\
I_{41} &= \frac{\sqrt{2}}{\sqrt{3}} |D_0| |D_-| \cos(\varphi_{D_0} - \varphi_{D_-}) + \frac{\sqrt{2}}{\sqrt{3}} |F_0| |F_-| \cos(\varphi_{D_0} - \varphi_{F_-}) + \frac{\sqrt{2}}{\sqrt{3}} |F_0| |F_+| \cos(\varphi_{D_0} - \varphi_{F_+}) \\
I_{42} &= -\frac{\sqrt{2}}{\sqrt{3}} (|D_+|^2 - |D_-|^2) - \sqrt{\frac{2}{3}} (|P_+| |F_+| \cos(\varphi_{D_+} - \varphi_{F_+}) - |P_-| |F_-| \cos(\varphi_{D_-} - \varphi_{F_-})) \\
& \quad - \frac{\sqrt{2}}{\sqrt{3}} (|F_+|^2 - |F_-|^2) \\
I_{50} &= \frac{10\sqrt{3}}{11} |D_0| |F_0| \cos(\varphi_{D_0} - \varphi_{F_0}) - \frac{10\sqrt{3}}{11} (|D_+| |F_+| \cos(\varphi_{D_+} - \varphi_{F_+}) + |D_-| |F_-| \cos(\varphi_{D_-} - \varphi_{F_-})) \\
I_{51} &= \frac{10}{11} |D_0| |F_-| \cos(\varphi_{D_0} - \varphi_{F_-}) + \frac{10\sqrt{3}}{11} |D_-| |F_0| \cos(\varphi_{D_-} - \varphi_{F_0}) \\
I_{52} &= -\frac{10}{11} (|D_+| |F_+| \cos(\varphi_{D_+} - \varphi_{F_+}) - |D_-| |F_-| \cos(\varphi_{D_-} - \varphi_{F_-})) \\
I_{60} &= -\frac{23}{11\sqrt{13}} (|F_+|^2 + |F_-|^2) + \frac{100}{33\sqrt{13}} |F_0|^2 \\
I_{61} &= \frac{10\sqrt{13}}{33} |F_0| |F_-| \cos(\varphi_{F_0} - \varphi_{F_-}) \\
I_{62} &= -\frac{10\sqrt{13}}{11\sqrt{39}} (|F_+|^2 - |F_-|^2)
\end{aligned}$$

It is often interesting to use make linear combinations of these moments in order to project out some particular bilinear combination of amplitudes. These have been referred to on occasion throughout the body of the text of this thesis and are defined here,

$$\begin{aligned}
\sigma_0^D &= |P_0|^2 + \frac{1}{3} |S|^2 = \frac{1}{3} I_{00} + \frac{\sqrt{3}}{3} I_{20} \\
\sigma_+^D &= |P_+|^2 + \frac{1}{3} |S|^2 = \frac{1}{3} I_{00} - \frac{\sqrt{3}}{3} I_{20} - \sqrt{\frac{2}{3}} I_{22} \\
\sigma_-^D &= |P_-|^2 + \frac{1}{3} |S|^2 = \frac{1}{3} I_{00} + \frac{\sqrt{3}}{3} I_{20} + \sqrt{\frac{2}{3}} I_{22} \\
\sigma_0^D &= |D_0|^2 + \frac{1}{10} (|S|^2 + \sum |P|^2) = \frac{2}{15} I_{00} + \frac{1}{15} I_{20} \\
\sigma_+^D &= |D_+|^2 + \frac{1}{10} (|S|^2 + \sum |P|^2) = \frac{2}{15} I_{00} - \frac{1}{15} I_{20} - \frac{1}{\sqrt{10}} I_{21} \\
\sigma_-^D &= |D_-|^2 + \frac{1}{10} (|S|^2 + \sum |P|^2) = \frac{2}{15} I_{00} - \frac{1}{15} I_{20} + \frac{1}{\sqrt{10}} I_{21} \\
\sigma_0^F &= |F_0|^2 + \frac{1}{3} (|S|^2 + \sum |P|^2 + \sum |D|^2) = \frac{2}{3} I_{00} + \frac{2\sqrt{13}}{33} I_{20} \\
\sigma_+^F &= |F_+|^2 + \frac{1}{3} (|S|^2 + \sum |P|^2 + \sum |D|^2) = \frac{2}{3} I_{00} - \frac{2\sqrt{13}}{33} I_{20} - \frac{11\sqrt{39}}{10\sqrt{13}} I_{42} \\
\sigma_-^F &= |F_-|^2 + \frac{1}{3} (|S|^2 + \sum |P|^2 + \sum |D|^2) = \frac{2}{3} I_{00} - \frac{2\sqrt{13}}{33} I_{20} + \frac{11\sqrt{39}}{10\sqrt{13}} I_{42}
\end{aligned}$$

with $\sum |X|^2 \equiv |X_0|^2 + |X_+|^2 + |X_-|^2$.

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