Photoproduction of Λ and Σ^0 Hyperons off Protons in the Nucleon Resonance Region using CLAS at Jefferson Lab

by

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Abstract

The differential cross section and hyperon recoil polarizations of the photoproduction of the ground state hyperons, $\gamma p \rightarrow K^+\Lambda$ and $\gamma p \rightarrow K^+\Sigma^0$, have been measured with the CLAS at Jefferson Lab up to a photon energy in the lab of 2.325 GeV. The results for both channels show significantly larger cross section in the middle to forward angles than have been observed previously by the SAPHIR Collaboration.

Both reactions show significantly more backward peaking in the angular distributions than has previously been possible to observe. The backward peaking hints that hyperon resonances in the u-channel play a significant role in the production mechanism.

In addition, in the $\gamma p \rightarrow K^+\Lambda$ reaction, a previously unobserved bump in the cross section was observed at forward angles, centered on a W of 1.95 GeV with a width of approximately $\Gamma = 100$ MeV.

In both $\gamma p \rightarrow K^+Y$ reactions the recoil polarization in the forward direction seems reasonably well reproduced by t-channel interferences in a Regge model calculation as well as hadrodynamic models that include kaon resonances in the t-channel. The recoil polarization for $\gamma p \rightarrow K^+\Lambda$ shows a significant enhancement around a W of 1.9 GeV in the backward angles, which is a sign of resonance activity in this vicinity. The polarization of $\gamma p \rightarrow K^+\Sigma^0$ at backward angles is, in contrast, less pronounced and mostly consistent with zero.

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Chapter 1

Introduction

Ernest Rutherford fathered experimental sub-atomic physics when he and his co-workers first measured alpha particles scattering off of atoms within a thin gold foil and noted that the number of back-scattered alpha particles implied that the gold atoms contained a very small, dense, charged core surrounded by practically empty space. Much of particle physics can, like Rutherford's experiment, be described as "smashing things together and seeing what comes out where." This description, although accurate, does not convey either the difficulties involved in performing these experiments nor the knowledge that may be obtained from these observations. In this dissertation I present the results of analyzing the colliding of two particular types of particles, photons striking protons, and observing the distribution of the resultant production of two other types of particles, kaons and hyperons.

1.1 The Standard Model

Since Rutherford's experiment, a vast quantity of data has been accumulated about the interaction of particles. These data have contributed to the creation of a theory of particle interaction, the Standard Model, which explains the experimental observations in terms of so called elementary particles or fields. What follows is a brief description of the Standard Model with an emphasis on those ideas that are most relevant to kaons, hyperons, and their production when photons and protons interact.

1.1.1 Fundamental Particles and Fields

The universe as we know it is mostly empty space filled here and there with pockets of matter. The Standard Model classifies the constituents of this matter into three categories: quarks, leptons, and force-carrying gauge bosons. One often thinks of these three types of matter as being particles. It is, however, more accurate to say that the Standard Model represents these types of matter as interacting fields and the concept of "particles" as being excitations of these fields. For the purpose of brevity

I will not make a distinction between the standard use of the term particle and this more technical version.

These particles and their quantum numbers are listed in Tables 1.1 and 1.2.

1.1.1.1 Quarks

The nucleus of the atom, which Rutherford first inferred must exist from the results of his experiment, is now known to consist of a group of protons and neutrons. Just as it was shown that the atom was not a fundamental particle by observing sub-atomic structure, protons and neutrons have been shown to have structure and, therefore, are not considered fundamental. The observed structure has been explained by the existence of particles called quarks. Quarks are spin $\frac{1}{2}$ particles with fractional electric charge. They also have a “color charge,” which is related to the so-called strong force and which has one of three values (red, blue, and green for quarks, anti-red, anti-blue, and anti-green for anti-quarks). In addition, there are six flavors of quarks: down, up, strange, charm, bottom, and top. Protons and neutrons are composed of the two lightest quarks, up and down. The three most massive quarks, charm, bottom, and top, are each more massive than both the proton and neutron. The strange quark, while lighter than the proton, is still significantly more massive than either the up or down quarks. The present dissertation is particularly concerned with a measurement involving the strange quark.

1.1.1.2 Leptons

Outside the nucleus of the atom resides the electrons. Electrons along with their anti-particle, the positron, are the lightest of the charged leptons. The leptons, like quarks, come in six varieties, electrons, muons, and tau particles, along with their neutrino counterparts, the electron, muon, and tau neutrinos (ν_e, ν_μ, ν_τ). The electron, muon, and tau particles and antiparticles each have electric charge ± 1 , and do not carry a color charge. Neutrinos have no electric charge and are, in the standard model, massless, although recent experiments have shown evidence that at least some must have a (very small) mass.

1.1.1.3 Gauge Bosons

In the Standard Model there are four forces, electromagnetism, the weak force, the strong force, and gravity. These forces are carried by integer spin particles (bosons). The photon is the carrier of the electro-magnetic force, while the W and Z particles carry the weak force. At sufficiently high energies, the electro-magnetic and weak forces become unified into an “electro-weak” force. Interestingly, the weak force does not conserve quark flavor and so, for example, it can convert an s-quark into a d-quark. The strong force is carried by gluons which also carry the strong force’s charge, color. This means that gluons can interact directly with other gluons. Photons, in contrast,

Table 1.1: The fundamental fermions and their approximate masses and charges

Leptons, spin=1/2			
flavor		Mass($\frac{GeV}{c^2}$)	Electric Charge
ν_e	electron neutrino	$< 1 \times 10^{-8}$	0
e	electron	0.000511	-1
ν_μ	muon neutrino	< 0.0002	0
μ	muon	0.106	-1
ν_τ	tau neutrino	< 0.02	0
τ	tau	1.7771	-1

Quarks, spin=1/2, color= (red,green,blue)			
flavor		Approximate Mass($\frac{GeV}{c^2}$)	Electric Charge
u	up	0.003	2/3
d	down	0.006	-1/3
c	charm	1.3	2/3
s	strange	0.1	-1/3
t	top	175	2/3
b	bottom	4.3	-1/3

Table 1.2: Force carrying Bosons and their properties.

spin=1 force carrying Bosons			
Force	Name	Mass($\frac{GeV}{c^2}$)	Electric Charge
Electro-Magnetism	γ (photon)	0	0
Weak	W-	80.4	-1
	W+	80.4	+1
	Z^0	91.2	0
Strong (color)	g (gluon)	0	0

Table 1.3: A few select ground Mesons and Baryons.

Symbol	Name	Valence Quarks	Electric Charge	Mass($\frac{GeV}{c^2}$)	spin	isospin
Mesons						
π^+	pion	$u\bar{d}$	+1	0.140	0	1
K^+	kaon	$u\bar{s}$	+1	0.494	0	1/2
Baryons						
p	proton	uud	1	0.938	1/2	1/2
n	neutron	udd	0	0.940	1/2	1/2
Λ	lambda	uds	0	1.116	1/2	0
Σ^0	sigma 0	uds	0	1.193	1/2	1

do not have electric charge and can not, therefore, interact directly with each other. The result of this interaction of gluons is a phenomena called confinement in which particles with color are so tightly bound together by these interacting gluons that one does not see them as free objects at large distances from one another. Here large can mean approximately as small as the size of a proton, ~ 1 fm.

1.1.2 Baryons, Mesons, and their Resonances

Confinement leads to quarks grouping together into “colorless” groups of three (one red, one green, one blue, or the three anti-colors) or two (a color and its anti-color). This association of quarks into colorless objects results in two distinct categories, the baryons made up of groups of three quarks, and the mesons made up of groups of a quark and an anti-quark. In particular, protons are made of two up and one down quark, while neutrons consist of two down quarks and one up. These are the lightest baryons. The lightest meson, the pion, also consists of up and down quarks and comes in three varieties based upon its net electric charge. Table 1.3 lists a few of the mesons and baryons of interest in this analysis. The quark model and the “8-fold way” was first introduced by Gell-Mann in 1964 to explain the spectrum of mesons and baryons that had been observed by that time.

Of particular interest in this analysis is what happens when we consider the effects of the strange quark. Replacing one up or down quark with a strange quark leads to a new set of particles, the “strange particles”, which are more massive than their up and down counterparts. The lightest of these strange particles for the mesons are the kaon and the Λ and Σ^0 for the baryons. The difference between lambdas and neutral sigmas has to do with how the wave functions of these two particles change under the interchange of the up and down quarks. The near symmetry between up and down quarks is referred to as isospin, since the symmetry is mathematically identical to the normal spin associated with the intrinsic angular momentum of particles. A baryon with one strange quark is often referred to as a hyperon. With six flavors of quarks, it is easy to see that the number of combinations of three or two is, though finite, quite large.

By combining the three lightest quarks and taking advantage of the fact that quarks are fermions (they have spin $\frac{1}{2}$) and therefore the wave functions must be anti-symmetric it is possible to derive the list of possible states that correspond to the ground state mesons and baryons. Each of these mesons and baryons can be excited into higher energy states. These energy states are often called resonances because they decay strongly and therefore only exist for such a short period of time that they are not observed directly. These resonances can, however, be observed through their decay products.

There are many resonances that have been observed in previous experiments. Most of these observations were done by producing the resonances using the strong interaction between nucleons and pions. Although this was perhaps the easiest method

to observe resonances, it is possible that some resonances may be missed when only using this method since other nearby resonances that are produced with significantly more strength may mask the weaker resonance. In fact there are several resonances that have been predicted by simple quark model calculations that have not yet been observed, possibly for this reason. It is worth noting that non-strange resonances can decay strongly into a kaon and hyperon and thereby play an important role in hyperon production mechanisms. It has been shown using a quark model calculation [1] that some of the “missing resonances” might be observed in hyperon photoproduction more easily than in either pion photoproduction or nucleon resonance production from pion beams, because of the relatively clean electromagnetic production mechanism.

1.2 The Experiment

In this experiment, we begin with a proton at rest and hit it with a photon of known energy. We then observe what comes out, looking for evidence of a kaon and a ground state hyperon. At low energies, one of the production mechanisms for this reaction involves exciting resonances that then decay into the observed final state. These resonances are quite broad ($\Gamma \sim 150 \text{ MeV}$) and in regions in which they overlap the interference between them can have very large effects, sometimes dominating all other features. At higher energies, the resonances can not be formed since too much excitation of the quarks will separate them. This implies that there is a limited energy region, “the resonance region,” in which one will see so called “resonance physics.” At very high energies, the interaction of the photon and proton is best modeled as an interaction of the photon with one of the quarks of the proton. Such models are based on perturbative models of Quantum Chromo Dynamics (QCD), and so this region of energies is often called the “perturbative region.” In between these two regions in energy is a region where it is difficult to see resonances due to the large number of resonances that can contribute but where the assumptions of QCD are not valid.

One of the goals of the Thomas Jefferson National Accelerator Facility is to probe the proton with photons at energies both in and above the resonance region to observe the transition from resonance physics into perturbative physics.

1.2.1 Measured Observables of $\gamma p \rightarrow K^+ Y$

For the reactions $\gamma p \rightarrow K^+ \Lambda$ and $\gamma p \rightarrow K^+ \Sigma^0$, there are several observables available that can, in principle, be measured with the correct type of experiments. Many of these observables require that the incoming photon be polarized, either linearly or circularly. This analysis considers only the observables that can be measured with an unpolarized photon beam, namely the differential cross-section and the recoil polarization. Figure 1.1 shows a diagram of the reaction $\gamma p \rightarrow K^+ \Lambda$, including the decay of the hyperon into a proton and a π^- . Θ_{cm}^K is the angle of the kaon relative to the direction of the incoming photon in the center of mass reference frame. The angle

Figure 1.1: Diagram showing Θ_{cm}^K and $\Theta_{\Lambda_{R.F.}}^P$ for $\gamma p \rightarrow K^+ \Lambda$.

$\Theta_{\Lambda_{R.F.}}^P$ is the angle of the decay proton relative to the normal of the $\gamma p \rightarrow K^+ Y$ reaction plane in the hyperon rest frame.

The simplest observable to measure is the unpolarized differential cross-section which is roughly speaking the probability that a photon and proton under a given set of initial conditions will produce a kaon and hyperon with specific outgoing kinematics. Measuring the differential cross-section, in addition to identifying the kaon, requires measuring the energy of the incoming photon and the kinematics of the outgoing kaon. No particles in the final state other than the kaon need to be measured.

Measuring the recoil polarization, on the other hand, requires the measurement of the polarization of the recoiling hyperon. For most particles in other reactions, this is a difficult task which could not be accomplished in a large acceptance spectrometer such as the CLAS. Fortunately, the weak decay of the Λ provides an opportunity to

measure its recoil polarization.

The Λ decays 63.9% of the time into a proton and a π^- . The only other significant decay mode is into a neutron and π^0 . In either case the decay involves the conversion of the strange quark into a down quark and therefore violates quark flavor conservation. This violation requires that the decay be mediated by the weak force and in turn allows for the direction of the decay proton or neutron to be correlated with the Λ 's polarization. For the case where the lambda decays into a proton and a pi-minus, the decay distribution is: $I(\Theta_{\Lambda_{R.F.}}^P) = \frac{1}{4\pi}(1 - \alpha_- P \cos(\Theta_{\Lambda_{R.F.}}^P))$ where P is the polarization of the Λ , $\Theta_{\Lambda_{R.F.}}^P$ is the angle of the decay proton's momentum relative to the normal of the production plane in the lambda reference frame, and α_- is the self-analyzing power of the lambda which has been experimentally determined to be 0.642 ± 0.013 . It is therefore possible to infer the polarization of the lambda by observing the decay distribution of protons. A more in depth discussion of this phenomena may be found in Si McAleer's dissertation [2].

The Σ^0 , on the other hand, decays 99.9% of the time into a Λ and a photon. It can be shown that the lambda from this decay has $-\frac{1}{3}$ the polarization of the Σ^0 . By measuring the polarization of the Λ from the decay, we can in turn infer the polarization of the Σ^0 . This measurement relies on the assumption that the Λ reference frame is identical to the Σ^0 reference frame which, is on average true, and for a particular event is not a bad assumption given that the decay lambda has a momentum of only 74 MeV in the sigma's reference frame (compared to a mass of 1116 MeV).

The measurement of the recoil polarization of the hyperon is quite important because, beginning with an unpolarized initial state, this polarization can only result from the quantum mechanical interference of different production mechanisms. This interference enables much more sensitive discrimination between theoretical hyperon photo-production models than the differential cross-section alone.

1.3 A Brief Summary of My Participation in the CLAS

I began work with the CLAS collaboration in the fall of 1997. For this first year I spent approximately every other month at Jefferson Lab. The construction of the Region One drift chamber had recently been completed when I started, and I participated in the pre-installation checks and the installation. For the next several years I acted as one of the drift chamber experts, which meant that I was on call for a week at a time to deal with problems in the drift chambers during CLAS data taking runs.

After the first full running period in 1997, I begin looking at the tracking package's ability to distinguish detached vertexes from the primary vertex, a potentially valuable tool for identifying strange particles. This work led me naturally into working on the tracking code after finding a charge dependence in the vertexes. My work on

debugging the tracking code and improving the track detection efficiency continued intermittently for several years. Eventually this led me back to the development of the CLAS standard vertex-finding package, MVRT, to provide the best estimate of a primary vertex based on the results of the tracking package.

Eventually, after a coin flip between Rob Feuerbach and myself, it was determined that I would analyze the $\gamma p \rightarrow K^+ \Lambda$ and $\gamma p \rightarrow K^+ \Sigma^0$ reactions for my dissertation work, while Rob analyzed the reactions $e p \rightarrow e' K^+ \Lambda$ and $e p \rightarrow e' K^+ \Sigma^0$ for his dissertation. Originally my analysis was to be based on the first real photon data run at CLAS, g1a.

While the third real photon data run with a proton target, g1c, took place in 1999, I spent three months at Jefferson Lab acting as the offline “chef” responsible for “cooking” some of the raw data as quickly as possible to catch problems while the run was ongoing. I had considerable help from several people who did the initial calibrations of the various detectors in CLAS. For the next eight months after data taking for g1c finished, I continued to coordinate the efforts at calibrating the g1c data from Pittsburgh and worked on improving the cooking scripts used to process the data. In August of 2000 I turned the coordination of the calibration effort over to Luminita Todor who oversaw the final calibration and cooking of the g1c data.

As time went by it became obvious that there were potentially serious problems with the normalization of the g1a data set. Once the g1c data was available, support in the collaboration for the analysis of the g1a data set began to fade fast. Eventually I too abandoned the g1a data run and begin analyzing the portion of the g1c data set most similar to the g1a data. The g1c data set had the additional advantage that it was not only better calibrated than the g1a data set, but it also had more than 5 times more statistics at basically the same electron endpoint energy. Questions about the photon normalization remained, even in the g1c data set, and this led me to begin analyzing two other reactions, $\gamma p \rightarrow \pi^+ n$ and $\gamma p \rightarrow p \pi^0$, that have been well studied in previous experiments. These studies led me to the discovery of inefficiencies for single charged particles in certain time-of-flight paddles.

After taking into account these inefficiencies and others in the time-of-flight drift chambers, and tagger, both of the two single pion production channels were in reasonable agreement with the existing world data except for an overall scale factor. Eventually I stopped analyzing the $\gamma p \rightarrow p \pi^0$ channel because the resultant scale factor was very sensitive to the parameterization of the background and the cuts on the π^0 identification. To reduce the uncertainties caused by the background to a level on par with the $\gamma p \rightarrow \pi^+ n$ channel would have required considerable work. The overall scale factor from the $\gamma p \rightarrow \pi^+ n$ of 1.087 is believed to be the result of an overall tagger trigger inefficiency that is not accounted for in the photon normalization. Work by other individuals analyzing other channels should also observe this phenomenon.

1.4 Summary

This dissertation reports on the measurements of the production from a photon and a proton of a kaon and a hyperon. The energies of the observed photon-proton system were mostly in, but did extend above, what is normally considered the resonance region. It will be seen that there is evidence of resonance-like structure at energies higher than has previously been observed in this type of reaction.

Chapter 2

Theory and Previous Experiments

In this chapter I will describe most current world data for the observables of the reactions $\gamma p \rightarrow K^+\Lambda$ and $\gamma p \rightarrow K^+\Sigma^0$, and the present theoretical understanding of that data. Comparisons between the current data and theory with the results of this dissertation will be discussed in Chapter 8.

2.1 Existing Data

Prior to 1990, the reactions $\gamma p \rightarrow K^+Y$ was last studied experimentally in results published in 1969 [3] and 1972 [4]. The experiments were done with low luminosity which, combined with the relatively low cross section, resulted in quite low statistics. Recently, however, advances in accelerator technology have enabled experiments, including this one, to revisit hyperon photoproduction with sufficiently high luminosity to measure the cross sections with much greater precision in a reasonable amount of time. The first experiment to publish results for $\gamma p \rightarrow K^+Y$ was done with the SAPHIR detector using the electron stretcher ring ELSA located in Bonn, Germany.

The most recently published data from SAPHIR [5] measured differential cross sections and hyperon polarizations from threshold up to a photon energy of 2 GeV in the lab. The short path-length of tracks within the detector ($\sim 2\text{m}$) prevented the direct identification of kaons by time-of-flight. However, hyperon production events could be identified from the decay of the lambda into a proton and a π^- in coincidence with a kaon candidate track. This process resulted in 7591 $\gamma p \rightarrow K^+\Lambda$ events and 5869 $\gamma p \rightarrow K^+\Sigma^0$ events. These events were split among nine E_γ bins and ten $\cos(\Theta_{cm}^K)$ bins for the differential cross section measurement, and three energy bins with four $\cos(\Theta_{cm}^K)$ bins for the polarization measurements. The results may be seen in Figures 2.1–2.3.

At the time of this writing there are other experiments working on measurements of these channels at the GRAAL and SPRING8 accelerators, as well as some additional data from the SAPHIR Collaboration. Data from these analyses have not yet been published.

Figure 2.1: Differential cross section measurements of $\gamma p \rightarrow K^+ \Lambda$ from SAPHIR [5]. The curves are fits to Legendre polynomials.

2.2 Current Theory

There are several theoretical models that apply to the $\gamma p \rightarrow K^+ Y$ reactions. In this section I summarize several of the more recent models and their results. The models discussed below are representative of three basic types: Quark Models, Hadrodynamic Models, and Regge Models. Each has its advantages, disadvantages, and different regions of applicability.

Capstick and Roberts in their April 1998 paper [1] reported theoretical predictions of the decay of non-strange baryon resonances into strangeness-containing final states. The theory uses relativized quark wave functions for mesons and baryons and uses a 3P_0 pair-creation model to predict the strong decays of baryons. The 3P_0 model is based on the assumption that a quark anti-quark pair will be created with the quantum numbers of the vacuum. The results of the Capstick and Roberts paper predicted that several of the so called “missing resonances” should have large strange decay amplitudes and electromagnetic couplings and therefore could be observed with the $\gamma p \rightarrow K^+ Y$ channels. Although this quark model is useful for predicting what resonances might be observed, it does not actually make any direct predictions about

Figure 2.2: Differential cross section measurements of $\gamma p \rightarrow K^+ \Sigma^0$ from SAPHIR [5]. The curves are fits to Legendre polynomials.

the cross sections and hyperon polarization. For such predictions we turn to a few hadrodynamical models and a Regge model.

Hadrodynamical models, also known as isobar models, are based on the approximation that the hadronic degrees of freedom are the relevant quantities to consider in constructing an effective field theory. These models begin with the Born terms and then add resonances in the s, t, and u channels to fit the experimental data. Different models use different prescriptions for the coupling of the hadrons and include different resonances. The result is that each different model has a different number and selection of free parameters (typically a few per resonance) but since most have been fit to the same data from SAPHIR, their results for the differential cross section and polarization are generally consistent within the errors of the SAPHIR data set. Thus at this time it is not possible to distinguish resonant effects from other so called background effects.

In 1992, prior to the SAPHIR results, Williams, Ji, and Cotanch produced a hadrodynamical model [6], hereafter referred to as the WJC model, for hyperon photoproduction. Their model was based in part on a duality constraint which is a phenomenological recognition that the inclusion of all s, t, and u channel resonances results in an over-complete description of the reaction. They therefore limited the diagrams and resonances used in their model to those in Figure 2.4.

In July of 2001, Chiang, Tabakin, Lee, and Saghai [7] used the WJC model along with results from pion production to investigate the effects of final state interactions

Figure 2.3: Polarization results for $\gamma p \rightarrow K^+\Lambda$ (top) and $\gamma p \rightarrow K^+\Sigma^0$ (bottom) from SAPHIR [5]. The curves are fits to the data that are constrained to be 0 at $\cos(\Theta_{cm}^K) = \pm 1$

on Λ photoproduction. Although the rescattering of particles in final state interactions are often ignored, for $\gamma p \rightarrow K^+Y$, the effects can be quite important. Since the cross section for the photoproduction on protons of pions are two orders of magnitude larger than the kaon cross sections, even a small rescattering of these channels into hyperon-kaon channels will still be significant. the coupled-channels approach of Chiang, *et. al* showed an approximately 20% effect in the total cross section which demonstrated the importance of including these effects in future models. The coupled-channel approach, however, suffers from difficulty in implementing a large number of resonances into any such model. Figure 2.5 show the total cross section results when using the final state interactions with the WJC model directly, and Figure 2.6 shows the same curves with the WJC model modified by refitting the coupling parameters.

The George Washington University group of T. Mart, C. Bennhold, and collaborators begin with a coupled-channels analysis that predicts which resonances are the important ones in the two different reactions $\gamma p \rightarrow K^+Y$. They then incorporate these resonances in a single channel effective model [8] that is fit to the SAPHIR data. In their model they include Born terms plus three “well-established” nucleon resonances: $S_{11}(1650)$, $P_{11}(1710)$, and $P_{13}(1720)$. In addition, they allow K^* exchange in the t-channel. In order to account for the finite size of the hadrons, form factors were added along with contact terms to restore gauge invariance. Using the Capstick and Roberts paper as a guide, they added a D_{13} resonance to account for the bump around 1.9 GeV. Their results indicated that the D_{13} resonance contributes strongly to the lambda photoproduction. This D_{13} , if confirmed, would be one of the “missing resonances.” Figures 2.7 and 2.8 show their results for $\gamma p \rightarrow K^+\Lambda$ total and differential cross sections.

Bijan Saghai [9] responded to the claims of a D_{13} by releasing results that showed that the SAPHIR data could be fit just as well without adding a D_{13} , but instead adding two hyperon resonances in the u channel. This was followed by papers from Janssen *et. al* [10–12], on the effects of the cutoff masses (Λ_c) used in the form factors used at the hadronic vertexes. The paper criticized the GW groups calculations by

Figure 2.4: Feynman diagrams for the WJC model [6]. Note that in the s-channel diagrams, the Δ resonances will not decay into Λ and a kaon because the Λ is an isospin 0 particle and the kaon is an isospin $\frac{1}{2}$ particle. The resonances used in the WJC model were: $\{N^*\} = \{N(1650), N(1710)\}$, $\{\Lambda^*\} = \{\Lambda(1405)\}$, $\{\Delta^*\} = \{\Delta(1620), \Delta(1900), \Delta(1910)\}$, $\{K^*\} = \{K^*(892), K_1(1270)\}$.

pointing out that including the D_{13} term alone required an unreasonably low cut off mass of 0.4 GeV. By including the u-channel hyperon resonances, equally good fits were obtained with the cutoff mass greater than 1.6 GeV. The inclusion of the hyperon resonances also reduced the sensitivity to the form factors associated with the contact term used to restore gauge invariance. Figures 2.9 and 2.10 show the Janssen fits to the SAPHIR $\gamma p \rightarrow K^+ \Lambda$ data with and without the hyperon resonances included.

In a different approach to kaon photoproduction, Guidal, Laget, and Vanderhaeghen have used a model [13] based on Regge trajectories to make predictions. The idea of Regge trajectories is that in the t-channel, one can replace all the diagrams of exchange of mesonic resonances of the same quantum numbers with a single diagram. This method has several advantages, not the least of which is that it has very few free parameters, namely the coupling constants of the trajectories with other particles. These free parameters were fit to data above 4 GeV. Gauge invariance was maintained by including a Regge-ized s-channel or u-channel term. Other parameters of the Regge models are determined by the mass spectra of the particles. Regge models approach perturbative results at very high energies, but the use of them as one nears threshold becomes less legitimate. They do, however, have the advantage that they represent well the effects of t-channel processes, so one can view discrepancies between the data and the Regge model as being caused by a non-t-channel produc-

Figure 2.5: Total cross section for coupled channels analysis [7] of $\gamma p \rightarrow K^+ \Lambda$ using the WJC model.

Figure 2.6: Total cross section for coupled channels analysis [7] of $\gamma p \rightarrow K^+ \Lambda$ using the WJC model with adjusted coupling parameters.

Figure 2.7: Total cross section measurement of $\gamma p \rightarrow K^+ \Lambda$ from SAPHIR [5] (black dots) and theoretical fit [8] that do not (dashed) and do (solid) include a D_{13} resonance in the s-channel. The open circles with large errors are from the “ABBHHM” Collaboration [3], published in 1969.

Figure 2.8: Differential cross section measurements of $\gamma p \rightarrow K^+ \Lambda$ from SAPHIR [5] and theoretical fits [8] that do not (dashed) and do (solid) include a D_{13} resonance in the s-channel. Total center of mass energy (W) is shown in each panel.

tion mechanism. Since this model is not intended to be applicable at the energies of this experiment, the model was not originally compared with the SAPHIR data. Predictions using the model are presented in Chapter 8.

2.3 Summary

The 1998 data from SAPHIR resulted in the production of several models that all could fit the data reasonably well. The differences between the models are not trivial, and therefore the interpretation of the data depends on which model is used. In particular, the existence of a D_{13} resonance at around 1.9 GeV remains questionable, as does the importance of the inclusion of hyperon resonances in the u-channel.

Figure 2.9: Differential cross section measurements of $\gamma p \rightarrow K^+ \Lambda$ from SAPHIR [5] and theoretical fits [10] that do not (dashed) and do (solid) include hyperon resonances in the u-channel.

Figure 2.10: Λ polarization measurements from SAPHIR [5] and theoretical fits [10] that do not (dashed) and do (solid) include hyperon resonances in the u-channel. ω_{lab} is the photon energy in the lab reference frame. Of particular interest is the third bin ($1400 \text{ MeV} < \omega_{lab} < 1600 \text{ MeV}$) as will be seen in Chapter 8.

Chapter 3

The Experimental Detectors and Electronics

In this chapter I will describe the experimental apparatus and conditions used for this measurement of the $\gamma p \rightarrow K^+ Y$ reaction. This experiment was performed in Hall B at the Thomas Jefferson National Accelerator Facility (TJNAF), located in Newport News, Virginia. A photon beam was produced using the bremsstrahlung radiation of an electron that had been accelerated to a fixed energy, 2.4 GeV, by the Continuous Electron Beam Accelerator Facility (CEBAF). The photon produced in a thin radiator foil was then “tagged” by measuring the energy loss of the electron. The photon beam was incident on a liquid hydrogen target at the center of the CEBAF Large Acceptance Spectrometer (CLAS). Particles produced in reactions of the photon with this target were detected and, for charged particles, their momentum was reconstructed from measurements made with CLAS. The reconstruction and analysis of these particles will be the topic of chapters 4 and 5.

3.1 The Continuous Electron Beam Accelerator Facility

The Continuous Electron Beam Accelerator Facility is so called because it accelerates electrons in a “Continuous Wave” mode, meaning that the electrons do not arrive in widely separated bursts, but rather come at regular intervals in approximately 2 ns wide “beam buckets.” When first produced, the electrons are accelerated up to an initial energy between 40-80 MeV in the injector, from whence they enter the main accelerator. This race-track shaped structure is comprised of three main parts, the linear accelerators, the recirculating magnets, and the beam separator/extractor. The last of these extracts the accelerated electrons from the main accelerator and sends them to one or more of the three experimental halls.

I will not discuss in detail the polarization of the electrons produced in the injector and delivered to the experimental hall, as the polarization information was not part

Figure 3.1: The key elements of The Continuous Electron Beam Accelerator Facility.

of this analysis. The beam polarization did, however, have an effect in the choice of endpoint energy of the electrons delivered to Hall B. The polarization orientation of the electron beam at the injector may be adjusted with a precision spin manipulator using electrostatic deflectors [14]. This orientation is the same for each electron, regardless of which hall the electron eventually ends up in. In order to deliver the maximum polarization of the electron beam to two halls simultaneously, it is necessary that the precession experienced by the electrons along both paths to the different halls be carefully controlled. It turns out that the precession is determined mainly by the number of recirculation passes around the accelerator. There are several energy combinations where the simultaneous polarization can be accomplished [15]. For the glc running period, the endpoint energies used were 3.115, 2.897, and 2.445 GeV. This analysis only covers the last of these energies.

3.1.1 The Injector

At the time of this experiment, the injector was made up of a polarized electron gun and part of a cryomodule which will be described in more detail in Sect. 3.1.2.

The electron gun used a photo-cathode made of gallium arsenide (GaAs) and an anode set to give an initial acceleration of 100 kV. The time and energy spread of the electrons was controlled with both the laser pulse width and a pre-buncher. Current was controlled by controlling the laser power and with the use of a beam chopper [16].

Polarized electrons were produced by illuminating the cathode with 499 MHz pulses from three different lasers, one for each hall. The pulses were timed to be 120° out of phase so that the net signal had a 1497 MHz structure [17]. The pulse widths of the lasers were between 60-80 ps FWHM. By independently setting the power of each laser, it was possible to control the amount of current sent to each hall without large losses in the chopper [16]. This control was important for this experiment since hall-B required a current of a few nA, whereas the other halls required currents ranging up to several hundred μA .

3.1.2 The Superconducting Linear Accelerators

After leaving the electron gun, the electrons are accelerated to an initial energy of 40-80 MeV before entering the first linear accelerator (linac). This acceleration is accomplished using the same type of RF cavities as are used in the linacs. Each linac was nominally capable of accelerating the electrons by 400 MeV per pass [14]. In practice it has been shown that the linacs are capable of significantly better performance than this design value.

The linear accelerators are made up of a series of RF cavities powered by independent klystrons that are adjusted so that the phase and energy gradient of the cavity maximizes performance. The RF cavities are made of niobium and cooled to about 2K using a liquid helium bath. At this temperature, niobium becomes superconducting, allowing current to flow with no resistance. The klystrons provide a high-powered electromagnetic wave inside the RF cavities which induces a charge distribution in the cavities similar to the one shown in Fig 3.2. The electron in a cavity, represented by a sphere in the diagram, is accelerated in the direction of the positive charges along the cavity walls. By careful selection of the right frequency ($\sim 1.5\text{GHz}$), it is possible to have the standing electro-magnetic wave in the cavity to be such that the electron is always accelerated in the same direction. Pairs of these cavities are grouped together and four pairs are then placed in a cryomodule for cooling to the necessary 2K [18].

3.1.3 The Recirculation Arcs

In order to enable the beam to pass through each linac multiple times, it is necessary to turn the beam 180° up to 9 times. This is accomplished by the use of a series of dipole magnets in the recirculation arcs, giving CEBAF a characteristic race-track appearance. Since the deflection of an electron in a magnetic field is proportional to the inverse of its momentum, electrons of different energies require different magnetic

Figure 3.2: Charge distribution in the superconducting RF cavities at one instant [18].

field strengths to have the same bending radius. This is accomplished by first splitting the beam into as many as five separate beams of different energies. Each of these beams pass through different dipole magnets on their path through the arcs. At the end of the arcs, the separate beams are recombined before entering the next linac. The arcs are designed to prevent the different beams from losing their synchronization upon recombination [14].

3.1.4 Beam Extraction and The Experimental Halls

After the beam has passed through the second linac and before entering the second recirculation arc, it passes through the beam switchyard. Here the beam is extracted for use in the three halls. The beam is first split according to its energy and then a kick is provided by 499 MHz RF separator cavities which can be placed in synchronization with one of the three electron gun lasers. Thus the beam for a particular hall is extracted at the specified energy. Due to this design, only one hall at a time can receive electrons at a given energy, with the exception of the highest energy, which all three halls can receive simultaneously [14].

The three experimental halls are labeled alphabetically: A, B, and C. Halls A and C both house two-arm spectrometers capable of making high precision measurements, relative to CLAS, but only over a limited phase space at any one time. Hall B is primarily the home of CLAS, which is a large acceptance spectrometer designed to have nearly 4π solid-angle coverage and therefore to be able to detect multiple particles simultaneously over a large region of phase space. CLAS will be discussed in more detail in Sect. 3.3.

Figure 3.3: Layout of the tagger elements in the alcove of Hall-B. The length of the tagger magnet is approximately 6m [19].

3.2 The Bremsstrahlung Photon Tagging System

It is necessary for this, and many other, experiments to use real photons, not electrons, as the beam incident on the target. The method used in this experiment to create the photons is to place a thin radiator in the electron beam from the accelerator. The electrons passing through the radiator give off bremsstrahlung radiation, and for a sufficiently thin radiator, this will be almost entirely made up of a single photon.

To a good approximation, the atom in the radiator acts as a spectator, and the photon and the electron share the energy of the incident electron. For high momentum incident electrons, it can be shown that both the produced photon and the recoil electron will, in the lab frame, continue basically in the same direction as the incident electron. The beam line then passes through a dipole magnet which separates out the electrons from the photons. Electrons which have not lost any energy are directed by the dipole into an underground beam dump, while electrons that have lost energy experience a greater change in direction. Electrons that have lost more than 20% of their incident energy pass through two layers of scintillation paddles to determine the electrons' energy and timing. The photons associated with the tagged electrons, meanwhile, pass through a collimator, and continue on toward the CLAS target [19]. The layout of the tagger elements may be seen in Fig. 3.3.

Figure 3.4: Electron beam profile from a harp scan just before run 21895. The fitted sigma in both directions is approximately 0.12 mm.

3.2.1 The Radiator

The radiator assembly for the tagger actually consists of several different foils and a “harp” consisting of two wires at 90° to each other. These are all mounted on an arm or “ladder” that is moved by a stepping motor, allowing different radiators to be moved into the beam line. The thicker foils were made of gold, while the thinner foils were made of gold deposited on carbon for support. The different foils have radiation lengths that vary from 10^{-6} to 10^{-4} [19]. For the running period used in this analysis, called the “glc” running period, the 1×10^{-4} radiation length foil was used most often, while a 2×10^{-5} radiation length foil was occasionally used.

The “harp” was used for measuring the electron beam profile and position at the radiator. The harp’s two wires were perpendicular to each other and at 45° to the direction of motion of the wires across the beamline. As these wires moved through the beam, scintillators downstream were used as an indirect measure of the intensity of the beam at the wire position [19]. An example of these profiles is shown in Fig. 3.4. Harp scans were done after any significant accelerator beam tuning to ensure the quality of the beam.

Figure 3.5: Geometry of the tagging system including “typical” electron trajectories labeled with the fraction of the incident electron’s energy that was radiated [19].

3.2.2 The Dipole Magnet

The dipole magnet used in the tagger was designed to enable electron beams up to an energy of 6.1 GeV to be used to produce real photons and still bend the non-radiating electrons through the 30° necessary to steer them into the beam dump. After construction, extensive tests were done of the uniformity of the magnetic field. These studies indicated that the energy resolution for operating below 4GeV should be accurate to better than 0.2%. In addition, ray-tracing calculations used to calculate the coincidence between the two scintillating planes (discussed below) were in good agreement with observed coincidences in the data [19].

3.2.3 The Hodoscope

The purpose of the hodoscope is three-fold. It is meant to measure the energy of electrons that radiated, and hence the energy of the radiated photon, to determine the timing of the event with sufficient precision to identify which 2 ns beam bucket the event occurred in, and to act as an important element in the trigger. To perform these functions, there are two rows of scintillating paddles, called the E and T planes for “Energy” and “Timing,” and all the associated electronics. Here I will discuss the two scintillating planes and save the discussion of the electronics for Sec. 3.5.

Figure 3.6: Scale drawing of a portion of the hodoscope showing the orientation of the numbered scintillating paddles in the E and T focal planes relative to some “typical” electron trajectories. The E counters are nearer the vacuum window than the T counters [19].

3.2.3.1 The E-plane

The E-plane consists of 384 scintillating paddles with a photo-tube on each end. These “E-paddles” are designed to cover areas that represent approximately equal sized energy bins. Their placement was chosen so that neighboring paddles would overlap by approximately $\frac{1}{3}$ of a paddle, thus allowing even finer energy resolution by using the coincidence of two paddles. The individual paddles were aligned perpendicular to the expected trajectory of electrons as they passed through the focal plane. This was done to maximize the utility of individual paddles while minimizing signals from back-scattered particles. [19].

In order to minimize electron energy loss prior to striking an E-paddle, the tagger was constructed with a large vacuum chamber that extended almost all the way to the e-counter focal plane. For ease of repairs and operations, the e-counters remained in a normal atmospheric environment, which was separated from the vacuum by a very strong yet thin vacuum window. Due to the size that was necessary, 20 cm by 9.6 m, this window was constructed out of a woven kevlar mesh that supports a thin layer of mylar. [20]

3.2.3.2 The T-plane

As can be seen in Fig. 3.5, the 61 “T-paddles” are placed farther away from the dipole magnet than the E-paddles in order that the measurement of the timing might be more accurate. Like the E-paddles, the T-paddles were aligned perpendicular to the trajectories of the electrons as they passed through the focal plane. In addition to being wider than individual E-paddles, the T-paddles were also thicker, again to improve the timing resolution. Again, individual T-paddles were placed so that they overlapped with neighboring paddles, however, the overlap for the T-paddles was approximately 10%. The purpose of the overlap in the case of the T-paddles was

Figure 3.7: Plot of T counter hits vs. coincident E counter. The correlation between the E and T counters is quite obvious [19]. k/E_0 is the ratio of the photon energy to the initial electron energy.

primarily to ensure that there were no gaps. The relative alignment of a sample of E and T paddles can be seen in Fig. 3.6.

Assuming no scattering off internal structures, an electron that has radiated a photon will have a trajectory that carries it through both detection planes. It is therefore possible to predict which E-paddles and T-paddles should fire in coincidence for “good” events. The results of comparing the coincidence seen in data runs (Fig. 3.7) and predictions made using the dipole field map were in quite good agreement. This coincidence is most useful for the high-energy electrons (low energy photons) since there was a significant background from nearly full energy electrons striking part of the beamline structure on the way to the beam dump [19].

3.2.4 The Collimator Assembly

In order to ensure that the photon beam is aimed at the target, a collimator assembly was placed downstream of the tagger. This assembly consisted of two collimators with sweeping magnets in between. The sweeping magnets were used to redirect any charged particles created by the photon beam hitting the first collimator into the second collimator. The collimators did add an additional level of complexity in that it became necessary to effectively measure any loss of photons on route to the target. This Beam-Line-Efficiency has often been mislabeled the “tagging efficiency.” The process of measuring it, discussed in Chapter. 5, used the total absorption shower counter which is discussed in Sect. 3.4.2. The diameter of the Primary Collimator used for this experiment was 8.61 mm at a distance of 5.80 m from the radiator which defined the photon beam to have a spread of 0.7 mrad. The Secondary Collimator, downstream of the Primary Collimator was sufficiently large, 17.3mm at 8.84m from the radiator, so as not to intercept this photon beam.

3.3 The CEBAF Large Acceptance Spectrometer (CLAS)

Once a real photon has made it past the collimator, it travels to the target at the center of the CLAS. Each photon has a small chance that it will interact with the target and produce one or more charged particles. The CLAS is designed to not only detect these particles, but to measure their momentum, time of flight, and path length with good accuracy. From these measurements and the direction of the curvature of the track, the mass and charge of each particle can be inferred, thus identifying them. A computer rendition of the CLAS can be seen in Fig. 3.8

In order to make these precision measurements, several sophisticated instruments have been brought together in the CLAS. The target contains liquid hydrogen which must be kept below 20 K so that it does not vaporize. Surrounding the target is the start counter which is used to initially detect any charged particles leaving the target, and ultimately is important for determining which tagged photon caused the interaction in the target. The drift chambers are used to track the particles and reconstruct their trajectories, while the toroidal magnet’s field causes these trajectories to curve so that the momenta can be determined. Finally, the time-of-flight scintillation paddles, the outer skin of the CLAS, measure the time that it took the particles to reach them. Other detectors in CLAS, such as the Cerenkov detector and the electromagnetic calorimeter, were not used in this analysis.

3.3.1 The Target

In order to look at the reactions, $\gamma p \rightarrow K^+ Y$, it is necessary to have a target with as high a density of atomic protons as reasonably possible. To do this for the g1 running

Figure 3.8: Computer rendering of the CLAS.

Figure 3.9: The g1 target cell. In operation the cell was oriented sideways, with the photon beam entering at the flanged end.

Figure 3.10: Schematic of the g1 target cell.

periods, a liquid hydrogen target was used. The target cell was made of kapton designed to be 180 mm long. When the length was measured and the end-cap shape taken into account the effective length was estimated at 178.5 ± 0.5 mm. Figures 3.9 and 3.10 show, respectively, the g1 target cell and the schematic illustrating the dimensions as measured after construction.

The density of liquid hydrogen can be determined from the temperatures and pressures which were recorded at the beginning of each run. Within each of the three periods of running with different end point energies, the temperature and pressure remained relatively constant, and so taking the average density as representative of the conditions for any individual run is a reasonable approximation. For the 2.4GeV running period of g1c, the average density was $0.0711 \pm 0.0001 \text{ g} \cdot \text{cm}^{-3}$ [22].

Figure 3.11: One start counter segment in place relative to the beam pipe. The target is not visible in this drawing but would be inside the beam pipe by the Scintillator [23].

3.3.2 The Start Counter

The start counter had several important requirements for its construction and operation. As part of the trigger, the start counter needs to detect, with a very high efficiency, charged particles created in the target that are not in the “shadow” of the torus coils. The timing resolution of the start counter needs to be good enough to determine in which 2 ns beam bucket the photon interacted with the target. In addition, the start counter needed to minimize the energy loss of particles passing through it and to fit in a very tight space, close to the target. These constraints led to an interesting design which kept everything but the scintillator paddles upstream of the target and out of the way of particles heading into the CLAS.

The basic layout of the start counter, seen in Fig. 3.11, is made up of three custom shaped scintillating paddles, each one covering two sectors. The special shape enables two phototubes to be attached upstream of the target to each paddle, increasing both the efficiency and timing resolution. The price one pays, however, is that due to the path light takes around the bend in the nose of the paddle, the calculation of the precise timing and location of one or more particles passing through a paddle in the start counter is complicated. This can be corrected for with information obtained from the tracking of the charged particles as they pass through the drift chambers, discussed in Sect. 3.3.4. A flattened diagram of an individual paddle may be seen in Fig. 3.12. The achieved resolution after software corrections were applied was 260 ps [23].

Figure 3.12: Diagram of an individual start counter paddle. The dimensions of each scintillator are in mm unless otherwise stated [23].

3.3.3 The Superconducting Toroidal Magnet

The advantage of a toroidal field over other, more simple designs, stems directly from the nature of the interaction between electrically charged particles and magnetism. When passing through a magnetic field, a charged particle experiences a force perpendicular to both its direction of motion and the magnetic field. The radius of curvature of the particle is simply $r = p_{\perp}(qB)^{-1}$, where $p_{\perp}$ is the component of the momentum perpendicular to the magnetic field. In order to measure the total momentum of a particle, then, it is desirable that the magnetic field be perpendicular to the total momentum. Since all the particles being analyzed are, to a good approximation, coming from along the beam-line, they will not have any momentum around the beam-line. From this it is easy to see that the best way to measure the total momentum of the particles is to have a magnetic field that circles the beam-line, a toroidal field. The one caveat to the above argument is that the toroidal magnetic field can not provide any coverage at 0° .

The magnet itself is made up of six superconducting coils. The kidney-shaped coils were designed so that particles moving in the forward direction would experience a stronger magnetic field than those going in a backward direction. This is desirable since forward going particles typically will be of higher momentum than backward going particles, and hence in order to measure them with the same accuracy one either needs a stronger field in the forward direction or considerably better spatial resolution. In order to operate properly, the superconducting coils are cooled to 4.5K by liquid helium. The magnet produces a peak field of 3.5 T at an operating current of 3790 A; however, during the g1c running period the current was set to 1920 A.

Much of the geometry of the CLAS detector elements are determined by the Toroidal Magnet. The six superconducting coils necessary to create the field provide shadow regions in which no particles can be detected and so are the natural places to put as much of the electronics and structural components of individual detectors. The result is a six-fold symmetric detector with the space between torus coils called “sectors.”

3.3.4 The Drift Chambers

The drift chambers are the heart of the charged particle tracking system. Effective tracking depends upon the ability to measure the location of a particular particle at several points along its trajectory away from the target while minimizing its multiple scattering. For the typical momentum of the particles that the CLAS is designed to measure, 0.2 – 2.0 GeV, drift chambers are a good choice since they enable position measurements to be made with a precision of a few hundred microns, while not requiring that the particle pass through as much material as in other methods.

Figure 3.13: The CLAS Toroidal Magnet after installation but before installation of any other systems. Note the standard-size human at the right.

3.3.4.1 How Drift Chambers Work

The effectiveness of track reconstruction is directly related to the accuracy and precision with which the drift chambers measure a track's position. Drift chambers work by measuring a signal caused by electrons produced when a charged particle passes through gas. To measure this signal, the drift chamber has a series of wires strung across the length of the chamber.

There are two types of wires used, sense wires and field wires. The sense wires are maintained at a high positive voltage while the field wires are maintained at a high negative voltage. One sense wire and the surrounding field wires is called a “cell.” The voltage on the wires creates an electric field in the cell which causes negatively charged particles to accelerate toward the sense wire. Thus, when a particle from the target ionizes the gas in the cell, electrons along the particle's path begin moving toward the sense wire. Due to collisions with the other molecules in the gas, the electrons move on average at a fairly constant “drift” velocity. This is where drift chambers get their name.

As the electron gets close to the sense wire, the electric field increases to the point where, between collisions with the gas molecules, the electron is accelerated to a sufficient energy to ionize the next atom it hits. This leads to two electrons accelerating toward the sense wire, then four, and so on The result of all this ionization near the wire is that what was once one electron, a very hard signal to

Figure 3.14: Horizontal cut through the CLAS detector centered at the beam line. The dotted lines represent a projection of the Torus coils. The dashed box is enlarged in Fig 3.17 [25].

Figure 3.15: Vertical cut through the CLAS drift chambers [25].

detect, has become a very large number of electrons, a much easier signal to detect. The time of arrival of these electrons at the sense wire is then recorded and later used to calculate the distance of closest approach (DOCA) of the track to the sense wire using a parameterized time-to-distance function. The docas were then used to fit the trajectory of the track through the detectors. This will be discussed in more detail in Chapter. 4.

3.3.4.2 Design of The CLAS Drift Chambers

For the CLAS, there were three “regions” of drift chambers built, called, aptly enough, regions 1, 2, and 3. The relative positioning of these regions can be seen in Figures 3.14 and 3.15. The first region, was close to the target (~ 1 m) and in approximately zero magnetic field. The second region was constructed between the torus coils discussed in Sect. 3.3.3, and therefore in a very strong field. The largest and final region was well outside the torus, and so was again in an area of almost no magnetic field. By measuring the position of a track in these three regions, it is possible to reconstruct the track’s trajectory with a precision of well under 0.5 mm. The three different regions had very different spacial and structural constraints, and therefore there are features that will be discussed below that were different for each. Before discussing these features, I will touch upon the design elements of the three regions that were similar.

Figure 3.16: Schematic of the region 3 drift chamber. Hardware common to all regions are labeled [25].

Figure 3.17: Hits in the region 3 drift chamber. Individual cells are represented by hexagons with field wires located at the corners and sense wires in the center. In the upper right the edge of a Cerenkov detector is shown [25].

Ideally, for a drift chamber, one would like there to be a one-to-one relationship between the time measured and the doca of the track to the wire. This is achieved when the electric field is cylindrically symmetric about the sense wire. To create such a field requires a cylinder at high voltage around the sense wires. Such “straw tubes” are indeed used in experiments tracking particles with very high momenta. In the CLAS, such tubes would introduce quite a lot of material in the path of the particle that was being tracked, and therefore a large amount of multiple-scattering could be expected.

To minimize the effect of multiple-scattering, while still approximating the idealized cylindrical symmetry, a quasi-hexagonal geometry was chosen. Each of the 35,148 sense wires was therefore surrounded by six field wires spaced so that at larger radial distances from the target the cell size increased. An added benefit of this layout is that the electrostatic forces on the wires are significantly decreased, enabling the tension within the wires to be lower than otherwise would be necessary. Features common to all regions are shown in Figure 3.16 which illustrates one of the region 3 drift chambers.

Each region had several layers of cells that a particle would pass through. There were two types of layers in each region, arranged in “super-layers” of, usually, six layers. The first type, the “axial” layers, were oriented so that the sense wires were perpendicular to the mid-plane of each sector. The “stereo” layers were tilted at 6° to

the axial wires to enable measurement of the position in the direction along the axial wires. In addition to the field wires associated with the stereo and axial layers, there was an added layer of field wires, so called “guard” wires, placed at the edges of each superlayer and held at a lower voltage, to simulate the electric field that would have come from an infinite number of layers of cells, thus making the electric field in each cell approximately independent of the cell’s position within the chamber. Figure 3.17 shows a computer rendering of the cross-section of one of the region 3 drift chambers illustrating the hexagonal configuration of the drift chamber cells.

Matching the geometry of the torus magnet resulted in six identical, wedge-shaped chambers for each region supported by two end-plates at an angle of 60° relative to each other. Circuit boards with preamps were mounted on these end-plates to enable the signals from the sense wires to reach the data acquisition electronics located outside the detector. The end-plates and associated electronics were constructed to minimize the area in which particles would not be detected.

The sense wires in all regions were made of $20\text{ }\mu\text{m}$ diameter gold-plated tungsten whereas the field wires were made of $140\text{ }\mu\text{m}$ diameter gold-plated aluminum. The gold-plating assured chemical inertness, while the choice of tungsten was based on durability consideration. Aluminum was chosen because of its low density which both decreases gravitational sag, thus lowering the required tensions, and minimizes multiple scattering. The diameter of the sense wires was decided upon in order that the electric field near the surface would be $\sim 280\frac{\text{kV}}{\text{cm}}$ to ensure that there was sufficient gain in the signal from the ionized trail. In contrast, the field wires were made larger to keep the electric field at the surface below $\sim 20\frac{\text{kV}}{\text{cm}}$ to minimize cathode deposits, the plating of the wires by junk from the gas.

The chamber gas was made up of a 90% argon and 10% CO_2 mixture. The decision to use this gas was based upon safety (it’s non-flammable), expense, the minimization of multiple scattering while still enabling reasonable signals and short enough collection time to prevent large backgrounds. The exact mixing fraction was chosen because it gave a reasonably large operating voltage plateau, had a reasonable saturated drift velocity, and fulfilled all of the above requirements.

The design of the region 1 chamber was strongly constrained by the spacing available between the main torus magnet’s coils and a “mini” torus that was used during electron runs. As a result there were only 4 stereo layers and there was a reduction in the most backward region of the axial superlayers from 6 to 3 layers. The cells ranged from 15 to 17 mm in diameter, and the sense wires were shifted from the center by $\pm 300\text{ }\mu\text{m}$ in order to provide a mini-stagger useful in resolving the left-right ambiguity of not knowing on which side of the sense wire a track passed. Unlike the other regions, region 1 was constructed as one unit, with the tensions in each sector carefully balanced against the others. The geometric constraints can be seen in Figures 3.18 and 3.19.

The region 2 drift cells ranged in diameter from 26 to 29mm in diameter. As the region 2 drift chambers were attached between the magnetic coil cryostats, the

Figure 3.18: Cross Section view of the region 1 drift chambers [25].

Figure 3.19: The assembly and installation process for region 1 [25].

Figure 3.20: The placement of the regions 2 and 3 drift chambers relative to the torus coils [25].

end-plates and associated electronics had to be thin to provide maximum azimuthal coverage. In addition the chambers were operating in a strong magnetic field and needed to be able to withstand sudden changes in the magnetic field as well as movements of the cryostat of up to 2 mm, if the cryostat was brought up to ambient temperature. To prevent damages to the chamber under such a circumstance, one end-plate was attached to the cryostat with springs and double hinges giving it considerable freedom of movement.

In comparison to regions 1 and 2, region 3 had ample space available which made it possible to build a very rigid structure. The 7 m long sectors are filled with cells ranging in diameter from 40 to 45 mm. The outer shell of the region 3 chambers was made of a Nomex-paper honeycomb sandwiched between two 300- μ m thick carbon-fiber skins [25]. The positioning of the region 2 and 3 drift chambers relative to the torus coils is illustrated in Figure 3.20.

3.3.5 The Time-of-Flight (ToF) Paddles

The time-of-flight paddles were the last detector element, used in this analysis, through which particles created in the target passed. The typical path lengths of these particles from the target to the ToF paddles were around 5 m. The ToF paddles covered an area outside the shadow of the torus coils in each sector from 8° to 142° . All paddles were made of 2 in thick scintillating material with photo-multiplier tubes (PMTs) attached at both ends.

The paddles were arranged in four panels in each sector, as seen in Fig 3.21. The first panel, holding paddles 1-23, spanned scattering angles up to 45° . These forward angle paddles were 15 cm wide and were connected to 2 in PMTs. The backward angle paddles were made of 22 cm wide paddles coupled to 3 in PMTs via bent and twisted light guides seen in Fig 3.22. A total of 57 scintillating paddles were used in each sector, with the last 18 being coupled together to form 48 logical counters [26]. The timing resolution of the ToF system, crucial for kaon identification, was approximately 150-200 ps.

3.4 Downstream Devices

There are three detectors along the beamline downstream of the target, the pair spectrometer (PS), the pair counter (PC), and the total absorption shower counter. The pair spectrometer was not used in this analysis, therefore I will not describe it.

3.4.1 The Pair Counter

The pair counter sits downstream of the target, just before the total absorption shower counter as shown in Figure 3.23. It consists of a veto scintillator paddle, a converter

Figure 3.21: View of one sector of the time-of-flight paddles showing the different panels. The forward angle paddles are on the right [26].

Figure 3.22: Light guide used to connect the backward angle ToF scintillation paddles to the 3 in phototubes [26].

Figure 3.23: General layout (not to scale) of the downstream devices. The photon beam enters from the left [19].

Figure 3.24: Beam profile from a pair-counter scan just before run 21765. The fitted width in both directions is approximately 0.81 in.

which produces $e^+ e^-$ pairs, another scintillator paddle, and four overlapping scintillator paddles. Like the harp scan, the pair counter can be used to measure the profile of the beam, but instead of measuring the electron beam's profile upstream of CLAS, it measures the photon beam's profile downstream of CLAS. This measurement was performed regularly during the course of the running period to ensure the quality and location of the photon beam. An example of a profile scan can be seen in Fig. 3.24.

3.4.2 The Total Absorption Shower Counter (TASC)

The total absorption shower counter (TASC) is used primarily as a means of determining the beamline-efficiency for the purpose of calculating an accurate estimate of the number of tagged photons that hit the target. The TASC sits downstream of CLAS and consists of four lead-glass blocks and phototubes. Each block measures $10 \times 10 \times 40 \text{ cm}^3$ and has an effective thickness of nearly 17 radiation lengths, thus providing nearly 100% efficiency [19].

Unfortunately, the TASC can only operate effectively with a beam current of a few 100 pA, so the TASC can not be used during production running. Special “normalization” runs were therefore necessary to determine the beam-line transmission efficiency. These runs were performed approximately once a day to ensure that

changes in beam conditions could correctly be accounted for. A description of the method used to determine this efficiency can be found in Chapter. 5.

3.5 The Trigger and Data Acquisition (DAQ) Systems

The trigger for any experiment is very important to the success of that experiment, but when attempting to measure cross-sections, it is absolutely critical. Inefficiencies in the trigger if not detected and understood will lead to systematic underestimation of the cross-section, and even angular distributions can be affected. On the other hand, a trigger that is too loose will cause the needless recording of too many background and accidental events that take up space and computing time. Below I will discuss the trigger and the associated data acquisition (DAQ) system for CLAS. Known trigger inefficiencies will be discussed in Chapter 5.

The “level 1” trigger for the glc running period required the coincidence between hits in the tagger, start counter, and time-of-flight paddles. Each of these systems is made up of multiple detection elements, and so a pre-trigger OR of the discriminated signal in each system was used to create one signal from each that could be used in the trigger module to check for coincidences. Before these pre-trigger processes, the signal from the detectors was split and then split again to go to analog-to-digital-converter (ADC) and time-to-digital-converter (TDC) boards. If a trigger was made, these boards were read out and the data assembled into a single “event.” Each event contained approximately 6kb of data. These events were then both combined into files of approximately 2 Gb in length and sent through a series of on-line monitoring programs which allowed problems to be detected during the run.

In the tagger, the pre-trigger module first took the signals from the E and T counters and passed them through a coincidence module that roughly matched E and T counters which could have come from a legitimate trajectory. The signal from this coincidence was then sent on to a chain of OR gates, called the master-OR, which then sent a signal on to be AND’d with the start counter signal before being sent to the level 1 trigger [19]. A schematic of the tagger electronics can be seen in Fig. 3.25.

Individual signals from phototubes in the coupled paddles in the start counter were first summed before being passed through a discriminator to allow a higher threshold than could have been set without first summing. As can be seen in Fig. 3.26, signals from the discriminators for each of the three coupled paddles were then OR’d and this signal was AND’d with the signal from the tagger [23].

The time-of-flight paddles had the signal from the two phototubes on each end summed, and these signals from adjacent paddles were then also summed to account for cases where particles passed from one paddle to the next. The signal was then discriminated and sent on to the level 1 trigger. A schematic of the ToF electronics and logic can be seen in Fig. 3.27 and 3.28.

Figure 3.25: The tagger electronics and pre-trigger logic [19]. The multi-hit TDCs shown here were, at the time of this experiment, actually single hit TDCs.

Figure 3.26: The start counter electronics and trigger logic [23].

Figure 3.27: The time-of-flight electronics [26].

Figure 3.28: The time-of-flight pre-trigger board logic [26].

Figure 3.29: The drift chamber electronics [25].

The signal from a sense wire was first enhanced in an signal translation board (STB). Each board processed signals from 16 sense wires before transmitting the signals to the amplification and discriminator boards (ADBs), which, as the name implies, first amplified the signal and if the signal was above a threshold set to minimize electronic noise, a signal was sent on to the TDCs. Figure 3.29 shows a schematic of the drift chamber electronics. At the time of the glc running period, the drift chambers were not part of the trigger.

3.6 Summary

This chapter has described the Continuous Electron Beam Accelerator Facility and the CEBAF Large Acceptance Spectrometer at The Thomas Jefferson National Accelerator Facility, along with the other detectors used in this analysis that are not central to CLAS, and a brief description of the trigger and data acquisition system.

Chapter 4

Raw Data Processing and Event Reconstruction

Data from the detectors were recorded on tapes in the tape silo at Jefferson Lab for every event that was accepted by the trigger described in Section. 3.5. Several calibration passes were done on a small fraction ($\sim 10\%$) of the data spanning the whole running period. A preliminary processing pass was then done which output so called “cooked” data, also in BOS format. The purpose of this pass was to perform those computations necessary for all analyses at once so as not to waste time and to have a standard approach to the common issues involved. This “cooked” data was then separately analyzed to extract the various quantities which were to be measured. During the preliminary pass, all raw data from TDCs were converted into times (in ns) and the time-consuming track reconstruction was performed. Also, periods in the run were identified to be excluded either because of inconsistencies in the raw data or momentary excess beam current which would make the photon normalization more difficult. This chapter briefly describes the event reconstruction, or cooking, that took place in the preliminary pass with the program called `a1c`. Chapter 5 describes in detail how these events were then analyzed to extract measurements of physical quantities.

4.1 Tagger Hit Processing

As mentioned in Chapter 3, the raw data output of the tagger is a collection of TDC values associated with each individual photo-multiplier tube (PMT) from various E and T paddles. The cooking process converts this data into a set of “good” hits that have an associated time and “tagged” photon energy.

The first step in this process was to convert the TDC value for a given PMT into a time. During the calibration stage, base peak position and slopes (in ps/channel) for each TDC were measured and stored in the map manager. These values were then used to convert the TDC channels into times.

After the times were calculated, a geometrical matching was done in which E-counters with hits were matched with T-counter hits if the particular E-T combination represented a possible combination from an electron that did not scatter after passing through the tagger radiator foil, thus reducing background associated with low energy electrons that scattered off of various beam-line components. This geometrical matching was combined with a time coincidence requirement that the E and T counters fired within 10 ns of each other. The geometrical matching was shown previously in Figure 3.7.

The E-T matching was followed by a re-binning from the 61 T-counters and 384 E-counters into 121 T-bins and 767 E-bins, where the additional bins were filled with events where adjacent counters had hits within a short time of each other (10 ns for E-counters and 5 ns for T-counters.) The advantage for re-binning the E-counters is that one effectively halves the energy resolution that can be achieved. For the T-counters it was found that the timing of individual T-counters shifted if the neighboring T-counter also fired, so the re-binning improves the ability to calculate an accurate time [27]. By design the E counters overlap by approximately 1/3 [19] and the T counters overlap by about 10%.

The final timing was determined by using the hits in the 121 individual T-bins, corrected for individual offsets, to identify the 2 ns beam bucket it was associated with, and using the RF machine time to correct the sub 2 ns timing. An overall offset between the tagger and the CLAS was also applied. The photon energy was determined from the electron beam endpoint energy and the individual E-bin hit.

The status of the tagger hit was reported by the tagger reconstruction software along with the time and energy. For the purpose of this analysis we only considered those hits that had a status value of 7 or 15, meaning that there were hits in both the left and right T-counter PMTs and one, in the former case, or more, in the latter case, unambiguous hits in the tagger hodoscope [27].

4.2 Start Counter Processing

The conversion of the TDC values associated with hits in the start counter into a start time is a process made more complicated by the geometry of the counters. After an initial conversion of each PMT TDC value into a time, it was necessary to know the location and number of tracks in a given paddle in order to calculate the most accurate time for the particles passage through the start counter. Since the propagation velocity for light from the passage of an incident track strongly depends on whether or not the light passed through the nose of a paddle, hits with more than one track in a given paddle (covering two adjacent sectors) were treated differently than those that had only one track. For the latter case the time was determined as the pulse-height (from the ADCs) weighted average of the two times. For the former case, the results of tracking (see Sect. 4.4) were used to determine where on the paddle the individual tracks hit, and use a position-corrected time from the closest PMT. [28]

In all cases, tracking information ($\beta \equiv v/c$ and the path-length from track vertex to the start counter) was then used to calculate the particle start time at the individual track vertex.

4.3 Time-of-Flight Processing

The time-of-flight (ToF) paddles were carefully calibrated to get the best particle identification possible. For each TDC, the channel, T , was converted into a time, t , using:

$$t = c_0 + c_1 T + c_2 T^2 + t_w \quad (4.1)$$

where c_0 is an overall offset, c_1 is the linear component of the TDC channel to time conversion, c_2 is a small correction for the non-linearity of the TDC, and t_w is a correction for time walk due to the pulse height. The times calculated for the left and right PMTs of the same ToF paddle, T_L and T_R , were then averaged to give a single time for that paddle:

$$T = \frac{T_L + T_R}{2} \quad (4.2)$$

Although the software in principle can take into account different propagation velocities toward the left and right PMTs, by adding an extra term, $-y \frac{v_R - v_L}{2v_L v_R}$, in practice the two velocities, v_L and v_R were set equal to each other and the extra term drops out. If only one TDC was present, say the left, the time was calculated as:

$$T = T_L - \frac{y}{v_L} \quad (4.3)$$

where y is the distance along the paddle to the PMT from the point where the particle entered the scintillator. This distance was determined from tracking. The energy measured in each PMT was calculated as:

$$E_{L,R} = \frac{10(A - P)}{A_{NMIP}} \text{MeV}, \quad (4.4)$$

where A is the ADC channel, P is the pedestal, and A_{NMIP} is the pedestal-subtracted channel corresponding to a minimum ionizing pulse (which would deposit 10 MeV in the paddle.) If there was only the left ADC, then the energy for the paddle was determined by:

$$E = E_L \exp(-y/\lambda_L), \quad (4.5)$$

where λ_L (λ_R) is the attenuation length of light propagating toward the left (right) tube. If ADC values for both tubes were present, the energy is given by:

$$E = \sqrt{E_L E_R \exp(y \frac{\lambda_R - \lambda_L}{\lambda_L \lambda_R})}. \quad (4.6)$$

If there were two hits in adjacent paddle that were both close in time and space, they were combined into a single hit where the new time was the energy weighted average of the times from the individual paddles. [29]

4.4 Drift Chamber Processing and Track Reconstruction

The raw drift chamber data is a collection of TDC values from hits in individual drift cells. In order to make use of these data, hits associated with individual tracks must be identified and each TDC value must be converted into a distance of closest approach of a track to the sense wire. These distances can then be used to extract the relevant kinematic parameters of the tracks' trajectories.

The tracking process begins by identifying "segments" of adjacent hits in individual superlayers of the drift chambers. Each segment required that hits be adjacent to one another and that there be at most one layer missing from within a superlayer. Segments were then compared sector by sector against a previously generated list of linked segments associated with simulated tracks, requiring segments in at least five of the six superlayers, and thereby an initial guess of the track's parameters was made. At this point, multiple tracks can share the same segments. Using the initial guess, so called Hit-Based-Tracking (HBT) was performed where, for each track, the particle's trajectory was fit by iteratively "swimming" [30] the particle through the various detector elements and the magnetic field, from a plane that passes through the nominal beam line. A χ^2 fit was performed with:

$$\chi^2 = \sum_i \frac{DOCA_i^2}{\sigma_{HBT}^2} \quad (4.7)$$

where $DOCA_i$ is the distance of closest approach of the "swum" track to the closest sense-wire in the i^{th} layer of the drift chamber and σ_{HBT} is set to the size of the drift cell. The results of the hit-based-tracking were then passed on to the time-based-tracking (TBT) routine which performed a similar fit to the tracks, but the χ^2 is instead defined as:

$$\chi^2 = \sum_i \frac{(DOCA_i - DIST)^2}{\sigma_{HBT}^2} \quad (4.8)$$

where DIST is the distance of closest approach calculated from the timing information of the sense wire hit. This distance was estimated using a drift time function (often referred to as the x -vs.- t function) using the time it took from when the track passed through the drift chamber (estimated initially from the β from HBT and the path length from the track's origin to the drift cell) to when the sense-wire fired, calculated from the TDC. The parameters of the drift velocity function were determined during the calibration pass. [31]

At the end of time-based-tracking, tracks that share segments were eliminated in favor of the track with the best χ^2 . The results were reported in the TBER bank and the TBTR bank. The TBER bank recorded the somewhat non-intuitive parameters of tracking for each track along with the elements of the (symmetric) error matrix and a few status bytes. The TBTR bank is the conversion of the TBER results into a slightly more intuitive form with the results reported in the lab coordinate system.

4.5 Vertex Determination

The determination of a best estimate for the event vertex was a project that I spent quite some time working on, although it plays only a small role within this analysis. The culmination of this work was the development of the MVRT bank which is now used as the standard event vertex within CLAS software. The following section describes this work.

4.5.1 Coordinate Systems and Parameters of the Tracking code

In order to make the explanation of the vertex determination procedures described in Section 4.5.3 easier, it is necessary to first explain the various coordinate systems used in tracking, and how these correspond to the parameters of the track fitting procedures. There were three main coordinate systems in tracking: the lab coordinate system, the sector coordinate system, and the “tracking” coordinate system. In all of these coordinate systems, the origin was the nominal center of the target, although for some running periods, the target may be quite a distance from this position. Also, in both the lab and sector coordinate systems, the z axis was defined as the nominal beam line with $\hat{z}_{lab}$ in the direction of the beam. The lab coordinate system is the more intuitive of the two coordinate systems, with the x_{lab} axis passing through the center of sector 1 and hence the y_{lab} axis points vertically upward. There were really 6 sector coordinate systems, one for each sector. They differ only in that the x_{sector} axis passes through the center of the specific sector, and the y_{sector} axis changes accordingly ($\hat{y} = \hat{z} \times \hat{x}$). It can easily be seen that the lab coordinate system, and the sector coordinate system for sector 1 are identical. The lab coordinates are related to the sector coordinates simply by:

$$\begin{pmatrix} x_{lab} \\ y_{lab} \\ z_{lab} \end{pmatrix} = \begin{pmatrix} x_{sector} \cos(\alpha) - y_{sector} \sin(\alpha) \\ x_{sector} \sin(\alpha) + y_{sector} \cos(\alpha) \\ z_{sector} \end{pmatrix} \quad (4.9)$$

where α is the angle of the center plane of the sector to the x axis, $\alpha = 60^\circ$ (Sector # – 1). The tracking coordinate system is defined from the sector coordinate system as:

$$\begin{pmatrix} x_{track} \\ y_{track} \\ z_{track} \end{pmatrix} = \begin{pmatrix} z_{sector} \\ x_{sector} \\ y_{sector} \end{pmatrix} \quad (4.10)$$

and is chosen so that the z_{sector} axis is aligned with the average magnetic field direction in the sector, i.e. parallel to the axial wires.

The parameters of the i^{th} track from the TBER bank are: $(q/p, \lambda, \phi, D_0, Z_0)$ where q is the charge of the particle, p is the momentum in the lab frame, λ is the dip angle relative to the sector's (x_{track}, y_{track}) plane, and ϕ is the angle of the track in the (x_{track}, y_{track}) plane from the tracking x axis. These parameters are all relative to the tracking coordinate system, so in order to convert them to the lab coordinate system, the sector number is needed, although it is not a free parameter. From these parameters, one gets the momentum of the particle at the track vertex in the lab coordinate system as:

$$\vec{p}_{lab} = \begin{pmatrix} p_{x_{lab}} \\ p_{y_{lab}} \\ p_{z_{lab}} \end{pmatrix} = \begin{pmatrix} p(\cos(\lambda) \sin(\phi) \cos(\alpha) - \sin(\lambda) \sin(\alpha)) \\ p(\cos(\lambda) \sin(\phi) \sin(\alpha) + \sin(\lambda) \cos(\alpha)) \\ p \cos(\lambda) \cos(\phi) \end{pmatrix} \quad (4.11)$$

The two additional parameters, D_0 and Z_0 define the point on the (y_{track}, z_{track}) plane that the track passes through with Z_0 equal to the z_{track} position of the intersection of the track with that plane, and D_0 equal to the minimum distance from the z_{track} axis line ($x_{track} = 0, y_{track} = 0$) that the track passes.

4.5.2 Track Vertexes

The results of the tracking code in the TBTR bank include a fit vertex for each track. An artifact of the creation process of the TBTR bank causes these vertexes to be located on one of the six planes corresponding to the (y_n, z_n) planes of the six sector coordinate systems. The particular plane was determined by the sector the track went through. For an individual track, this vertex was a reasonable first approximation for the vertex location, since the choice of position along the track to call the “vertex”, without additional information is somewhat arbitrary. This vertex was used when calculating the vertex corrected start counter time since the path length along the trajectory of the track from any point to this vertex is well defined. Using information beyond that of the individual track can, however, yield a slightly better event vertex.

4.5.3 MVRT: the Event Vertex

The event vertex, stored in the MVRT bank, was an attempt to use additional information from either the beam's measured location, if there was only one track in

Figure 4.1: The tracking parameter λ for a track in sector 1

Figure 4.2: The tracking parameter ϕ for a track in sector 1

Figure 4.3: The tracking parameter D_0 for a track in sector 1

Figure 4.4: The tracking parameter Z_0 for a track in sector 1

the event, or information from all the tracks together, to define a best fit for a single vertex for the event. In all cases, the trajectory of the track near the target was assumed to be linear, so the MVRT vertex at present is best when used only for unpolarized targets. This also means that the track parameter $(\frac{q}{p})_i$ was not used in the construction of the MVRT bank.

For events with only one charged track, that track's parameters at the track TBTR vertex were used to define a line. This was compared with a line representing the measured beam line, and the track's point of closest approach to that beam line was used as the event vertex. The beam's position was determined by examining the three-track MVRT vertex distribution during the calibration pass.

For events with more than one charged track, the parameters and errors from the fit results of each track were combined to create a single vertex. No additional fitting of the drift chamber hits was done. Also no modification of the track parameters to force them onto this new vertex was done. The MVRT routines were modeled after similar routines, for helical tracks, from the Crystal Barrel experiment [32], and used the method of least squares fitting algorithm for indirect measurements from Siegmund Brandt's Statistical and Computation Methods in Data Analysis [33].

The process of fitting the vertex begins with the creation of a vector, $\vec{\eta}$, made up of the relevant parameters of all the tracks:

$$\vec{\eta} = (\vec{x}_1, \vec{x}_2, \dots, \vec{x}_n), \quad (4.12)$$

where

$$\vec{x}_i = (\lambda_i, \phi_i, D_{0i}, Z_{0i}) \quad (4.13)$$

The associated error matrix, $C_{\vec{\eta}}$ is a block diagonal matrix composed of the 4 by 4 error matrices of each of the tracks, $C_{\vec{x}}^i$:

$$C_{\vec{\eta}} = \begin{pmatrix} C_{\vec{x}}^{i=1} & . & . & . & 0 \\ . & . & . & . & . \\ . & . & . & . & . \\ . & . & . & . & . \\ 0 & . & . & . & C_{\vec{x}}^{i=n} \end{pmatrix} \quad (4.14)$$

The aim of the MVRT routines is then to find the vertex, $\vec{U}_V = (u_V, v_V, w_V)$ and the associated symmetric 3×3 error matrix. The variables u_V , v_V , and w_V correspond to x_{lab} , y_{lab} , and z_{lab} but were introduced to prevent confusion with the tracking coordinates. To find a best estimate for the vertex, first the single track vertexes, U_{0i} , and directions, $\vec{C}_i$, were calculated from the i^{th} track parameters and sector angle, α_i :

$$U_{0i} = \begin{pmatrix} u_i \\ v_i \\ w_i \end{pmatrix} = \begin{pmatrix} -Z_{0i} \sin(\alpha_i) \\ Z_{0i} \cos(\alpha_i) \\ -D_{0i}/\sin(\phi_i) \end{pmatrix} \quad (4.15)$$

and

$$\vec{C}_i = \begin{pmatrix} c_u \\ c_v \\ c_w \end{pmatrix} = \begin{pmatrix} \cos(\lambda_i) \sin(\phi_i) \cos(\alpha_i) - \sin(\lambda_i) \sin(\alpha) \\ \cos(\lambda_i) \sin(\phi_i) \sin(\alpha_i) - \sin(\lambda_i) \cos(\alpha) \\ \cos(\lambda_i) \cos(\phi_i) \end{pmatrix} \quad (4.16)$$

From these we calculate the point of closest approach of each track to the trial vertex, $\vec{U}_V$, and the vector difference between the two points. This is done first by taking the difference between the track vertex and the trial vertex, $\vec{d}_i = \vec{U}_V - \vec{U}_{0i}$. The (signed) distance along the track to the point which is closest to the trial vertex is then calculated: $t_i = \vec{d}_i \cdot \vec{C}_i$. With this we get the difference vector:

$$\vec{F}_i = \vec{d}_i - t_i \cdot \vec{C}_i. \quad (4.17)$$

The Jacobian of the vectors $\vec{d}_i$ and $\vec{C}_i$ with respect to the track parameters can be calculated explicitly (using $c(x)$ for $\cos(x)$ and $s(x)$ for $\sin(x)$):

$$dd_i = \begin{pmatrix} \frac{\partial(\vec{d}_i)_u}{\partial\lambda_i} & \frac{\partial(\vec{d}_i)_v}{\partial\lambda_i} & \frac{\partial(\vec{d}_i)_w}{\partial\lambda_i} \\ \frac{\partial(\vec{d}_i)_u}{\partial\phi_i} & \frac{\partial(\vec{d}_i)_v}{\partial\phi_i} & \frac{\partial(\vec{d}_i)_w}{\partial\phi_i} \\ \frac{\partial(\vec{d}_i)_u}{\partial D_{0i}} & \frac{\partial(\vec{d}_i)_v}{\partial D_{0i}} & \frac{\partial(\vec{d}_i)_w}{\partial D_{0i}} \\ \frac{\partial(\vec{d}_i)_u}{\partial Z_{0i}} & \frac{\partial(\vec{d}_i)_v}{\partial Z_{0i}} & \frac{\partial(\vec{d}_i)_w}{\partial Z_{0i}} \end{pmatrix} = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & \frac{-D_{0i} c(\phi_i)}{s(\phi_i)^2} \\ 0 & 0 & \frac{1}{s(\phi_i)} \\ s(\alpha_i) & -c(\alpha_i) & 0 \end{pmatrix} \quad (4.18)$$

and,

$$dc_i = \begin{pmatrix} \frac{\partial(\vec{C}_i)_u}{\partial\lambda_i} & \frac{\partial(\vec{C}_i)_v}{\partial\lambda_i} & \frac{\partial(\vec{C}_i)_w}{\partial\lambda_i} \\ \frac{\partial(\vec{C}_i)_u}{\partial\phi_i} & \frac{\partial(\vec{C}_i)_v}{\partial\phi_i} & \frac{\partial(\vec{C}_i)_w}{\partial\phi_i} \\ \frac{\partial(\vec{C}_i)_u}{\partial D_{0i}} & \frac{\partial(\vec{C}_i)_v}{\partial D_{0i}} & \frac{\partial(\vec{C}_i)_w}{\partial D_{0i}} \\ \frac{\partial(\vec{C}_i)_u}{\partial Z_{0i}} & \frac{\partial(\vec{C}_i)_v}{\partial Z_{0i}} & \frac{\partial(\vec{C}_i)_w}{\partial Z_{0i}} \end{pmatrix} = \begin{pmatrix} -s(\lambda_i)s(\phi_i)c(\alpha_i) - c(\lambda_i)s(\alpha_i) & -s(\lambda_i)s(\phi_i)s(\alpha_i) + c(\lambda_i)c(\alpha_i) & -s(\lambda_i)c(\phi_i) \\ c(\lambda_i)c(\phi_i)c(\alpha_i) & c(\lambda_i)c(\phi_i)c(\alpha_i) & -c(\lambda_i)s(\phi_i) \\ 0 & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix} \quad (4.19)$$

From this we can calculate the Jacobian of $\vec{F}_i$ with respect to the track parameters,:

$$\begin{aligned}
B_i &= \begin{pmatrix} \frac{\partial(\vec{F}_i)_u}{\partial\lambda_i} & \frac{\partial(\vec{F}_i)_v}{\partial\lambda_i} & \frac{\partial(\vec{F}_i)_w}{\partial\lambda_i} \\ \frac{\partial(\vec{F}_i)_u}{\partial\phi_i} & \frac{\partial(\vec{F}_i)_v}{\partial\phi_i} & \frac{\partial(\vec{F}_i)_w}{\partial\phi_i} \\ \frac{\partial(\vec{F}_i)_u}{\partial D_{0i}} & \frac{\partial(\vec{F}_i)_v}{\partial D_{0i}} & \frac{\partial(\vec{F}_i)_w}{\partial D_{0i}} \\ \frac{\partial(\vec{F}_i)_u}{\partial Z_{0i}} & \frac{\partial(\vec{F}_i)_v}{\partial Z_{0i}} & \frac{\partial(\vec{F}_i)_w}{\partial Z_{0i}} \end{pmatrix} \\
&= dd_i - t_i \cdot dc_i - \vec{C}_i((\vec{C}_i)^T \cdot dd_i + (\vec{d}_i)^T \cdot dc_i), \tag{4.20}
\end{aligned}$$

and the Jacobian of $\vec{F}_i$ with respect to the trial vertex location:

$$\begin{aligned}
A_i &= \begin{pmatrix} \frac{\partial(\vec{F}_i)_u}{\partial u_V} & \frac{\partial(\vec{F}_i)_v}{\partial u_V} & \frac{\partial(\vec{F}_i)_w}{\partial u_V} \\ \frac{\partial(\vec{F}_i)_u}{\partial v_V} & \frac{\partial(\vec{F}_i)_v}{\partial v_V} & \frac{\partial(\vec{F}_i)_w}{\partial v_V} \\ \frac{\partial(\vec{F}_i)_u}{\partial w_V} & \frac{\partial(\vec{F}_i)_v}{\partial w_V} & \frac{\partial(\vec{F}_i)_w}{\partial w_V} \end{pmatrix} \\
&= \delta_{kl} - (\vec{C}_i)_k (\vec{C}_i)_l \\
&= \begin{pmatrix} 1 - c_u^2 & -c_u c_v & -c_u c_w \\ -c_u c_v & 1 - c_v^2 & -c_v c_w \\ -c_u c_w & -c_v c_w & 1 - c_w^2 \end{pmatrix} \tag{4.21}
\end{aligned}$$

where c_u , c_v , and c_w are, respectively, the u , v , and w components of $\vec{C}_i$. The individual i^{th} matrices were formed into block diagonal matrices analogous to $C_{\vec{\eta}}$ and then the error matrix was constructed by first constructing a matrix G_b :

$$G_b = (BC_{\vec{\eta}}B^T)^{-1} \tag{4.22}$$

$$\tag{4.23}$$

from which the error matrix and new trial vertex may be derived:

$$C_V = (A^T G_B A)^{-1} \tag{4.24}$$

$$\Delta U_V = -C_V A^T G_b F, \tag{4.25}$$

$$U_V^{new} = U_V + \Delta U_V. \tag{4.26}$$

Figure 4.5: MVRT vertex (in cm) for 3 track events in run 21785. Bins are 0.5 mm wide.

The new χ^2 from the change in tracks' parameters is:

$$\Delta\eta = C_{\bar{\eta}} B^T G_b (F - A\Delta U_V) \quad (4.27)$$

$$\chi^2 = (B\Delta\eta)^T G_b (B\Delta\eta) \quad (4.28)$$

We start with a first guess for the vertex as the average of all the TBTR vertexes. The χ^2 is then minimized in iterative steps to yield the final vertex. An example of the distribution of the vertexes for three track events is given in Figure 4.5. A Gaussian was fitted to the projection onto each axis of the three track vertexes to yield an estimate of the beam position. Figures 4.6 and 4.7 show the resultant beam position as a function of run number.

4.6 Summary

At the completion of the first pass of the analysis, each event of the raw data was processed into a “cooked” form from which anyone could perform their own analyses. The tagger TDCs and ADCs were converted into photon energy and time. The drift chamber hits were used to identify tracks and, for each track, measure the momentum and track length. From these tracks individual track vertexes and a single event vertex were calculated. In addition, the data from the start-counter and time-of-flight paddles were converted into times. Chapter 5 will explain how all this information was used to measure normalized yields for the various reactions.

Figure 4.6: 3-track MVRT x position (in cm) vs run number Figure 4.7: 3-track MVRT y position (in cm) vs run number

Chapter 5

Analysis of The Reconstructed Events

In this chapter I describe the analysis that was done to “cooked” data to produce the particle yields necessary to measure the cross section and polarization of the reaction $\gamma p \rightarrow K^+ Y$. The first step in this process was to identify kaons as cleanly as possible. To do this, several sources of background were eliminated including out-of-time events, tracks striking inefficient time-of-flight paddles, and particles detected near the edge of the detectors. After these cuts were applied, the kaons were identified by a momentum dependent mass cut and their time of flight. The hyperon yields were extracted in photon energy (E_γ) and $\cos(\Theta_{cm}^K)$ bins through a series of fits to the hyperon missing mass squared. The yields were normalized to the number of photons incident on the target. The correction for the acceptance of the detector is described in Chapter 6 and combined with these normalized yields to produce the measurements reported in Chapter 7. The related analysis $\gamma p \rightarrow \pi^+ n$ is discussed in Appendix A.

5.1 Binning

The choice of bins was made based on a balance between maximizing the statistics in each bin, and maximizing the information gained by reducing bin size. The final binning chosen was 25 MeV wide bins in E_γ and 20 bins in $\cos(\Theta_{cm}^K)$ for the differential cross section measurements. The bins at the edges of the acceptance were excluded due to the theoretical model dependence in calculating the acceptances. This eliminated all the bins with centers at $\cos(\Theta_{cm}^K)$ of -0.95 and 0.95, but was more restrictive near threshold.

For measuring the polarization, 50 MeV wide E_γ bins were chosen along with 10 bins in $\cos(\Theta_{cm}^K)$ and 10 bins in $\cos(\Theta_{\Lambda_{R.F.}}^P)$. The larger binning in energy and $\cos(\Theta_{cm}^K)$ was necessary because of the need to bin in an extra variable, and because of the reduced acceptance due to the need to require the detection of a proton.

5.2 Tagger and Start-Counter Coincidence

As described in Sect 3.5, the trigger for the g1 running period had a coincidence requirement between the tagger's master OR (MOR) and the start-counter. Any hit in the tagger that was not in coincidence with the start-counter was considered an out of time hit, and was rejected. To determine if a tagger hit was in coincidence with a start-counter hit, it was necessary to calculate the time that the particle which passed through the start-counter was created at the vertex. This resulted in a correction to the start-counter time based on the location of the vertex. For each track, the vertex correction was gotten from the TBID bank, which contains the time that the particle was at the "vertex" of that track based on the beta of the track and the start-counter hit associated with that track. The measured time for the tagger hit was propagated forward to when the photon would have reached the vertex as calculated from the MVRT bank. With an ideal detector and perfect calibration, we would expect that there would be no difference in the time.

The measured difference between the tagger and start-counter for all events can be seen in Figure 5.1. For events that had a well-identified π^+ with a missing mass consistent with a neutron, one can see in Figure 5.2 that most of the accidentals to either side of the central peak have disappeared. However, for well identified kaons that had a missing mass consistent with Λ (Figure 5.3) or Σ^0 (Figure 5.4), one can see that there was still a significant number of events outside of the central peak. These accidentals were made up primarily of mis-identified pions or protons, and in order to have a clean kaon signal it was necessary to cut these events out. Therefore, events were considered to be have a good tagger-start-counter coincidence only if they were in the same beam bucket as the peak (i.e., within ± 1.002 ns of the center of the central peak).

It is possible that these cuts exclude real pion, kaon, and proton events because of a poorly measured time in either the tagger or start-Counter. To estimate the magnitude of this effect a comparison was done, energy bin by energy bin. First, the missing mass for all positively charged particles (assuming they were pions) and with a $|startcounter - tagger| < 9ns$ was plotted. This was compared with the missing mass of well identified pions (see Appendix A) that were within the nominal start-counter - tagger cuts of $\pm 1ns$. Figure 5.5 shows these two plots for one energy along with the difference between the two plots. Fits of a Gaussian plus linear background to the neutron peak were then used to estimate the correction to the yield needed to compensate for the start-counter-tagger and pid cuts using the form: $correction = 1 + \frac{diff}{pid}$, where pid is the yield extracted from the well identified pions with the narrow timing cut, and $diff$ is the yield extracted from the histogram of the difference between all charged particles with a loose tagger-start counter cut and the well identified pions. The results may be seen in Figure 5.7.

A similar comparison was done for the proton and missing π^0 with results in Figures 5.6 and 5.8. A correction of exactly 1 is indicative of a failure to fit the peak

Figure 5.1: Time difference at the vertex between the start-counter and tagger for all events.

Figure 5.2: Time difference at the vertex between the start-counter and tagger for $\gamma p \rightarrow \pi^+ n$ events.

Figure 5.3: Time difference at the vertex between the start-counter and tagger for $\gamma p \rightarrow K^+ \Lambda$ events.

Figure 5.4: Time difference at the vertex between the start-counter and tagger for $\gamma p \rightarrow K^+ \Sigma^0$ events.

in the difference of the plots, primarily because there is not much of a peak to fit. Notice that in energies greater than the $\gamma p \rightarrow K^+Y$ threshold, there was effectively no correction to be made. There was no evidence for real $\gamma p \rightarrow \pi^+n$ or $\gamma p \rightarrow p\pi^0$ events in the accidentals region beyond 9ns. This type of comparison is not possible to do for kaons since the pion background is very large, thus a correction of half the $\gamma p \rightarrow \pi^+n$ correction was applied to the kaons. This correction was energy dependent and amounted to $\sim 2\%$.

5.3 Fiducial Cuts

The purpose of implementing fiducial cuts is to remove from the analysis data from regions of the detector that are difficult to model accurately, thus making the calculation of the acceptance and hence the cross section, more accurate. In this analysis there were two types of fiducial cuts applied. The first type of fiducial cut was a momentum and angle dependent cut that eliminated particles near the edge of the drift chambers. The second type of fiducial cut removed particles that struck time of flight paddles that were inefficient when compared to the paddles in neighboring sectors.

5.3.1 Momentum and Angle Cuts

Although the efficiency of CLAS was fairly constant near the mid-plane of each sector, several effects, including a more rapidly changing magnetic field and particles scattering from the torus structure back into CLAS, make it difficult to model the acceptance of CLAS as a particle approached the main torus coils. To avoid the pitfalls associated with these effects, we apply the standard fiducial cut, pseudo-spa, to all particles. Pseudo-spa is a function of lab momentum, angle, and vertex of the particle. The standard pseudo-spa cuts, established by other groups from the analysis of various reactions in electroproduction data, were used without modification. Figures 5.9 and 5.10 show the lab angular distributions (ϕ vs Θ_{cm}^K) for positive protons and kaons respectively. One can see that there was no significant ϕ dependence to the yields. The holes caused by the knocked out ToF paddles can be seen quite clearly in both the kaon and proton plots, especially the forward angle knockouts in sectors 2 and 3.

5.3.2 Time-of-Flight Knockout

The Time-of-Flight(ToF) system consists of 57 scintillation counters in each sector. The back 18 were paired together so that for each sector there were effectively 48 “paddles”. Initial calibration of the scintillator timing was done before cooking. In order to have consistency between the real data and simulated data, it was necessary

Figure 5.5: Missing mass for assumed $\gamma p \rightarrow \pi^+ n$ for all charged particles and ± 9 ns start-counter-tagger cut (nopid), well identified π^+ with standard ± 1.04 ns start-counter - tagger cut (pid), and the difference (diff.)

Figure 5.6: Missing mass for assumed $\gamma p \rightarrow p\pi^0$ for all charged particles and ± 9 ns start-counter-tagger cut (nopid), well identified p with standard ± 1.04 ns start-counter - tagger cut (pid), and the difference (diff.)

Figure 5.7: Correction to the $\gamma p \rightarrow \pi^+ n$ yield due to particle identification and start-counter-tagger timing.

Figure 5.8: Correction to the $\gamma p \rightarrow p\pi^0$ yield due to particle identification and start-counter-tagger timing. At energies above 1 GeV there is no visible peak, and the correction at very high energy are artifacts of the fitting procedure.

Figure 5.9: Angular distribution in lab coordinates (in degrees) of protons after fiducial cuts and ToF knockout.

to make sure that paddles that either were not working, or had major calibration problems, were identified and knocked out of both the simulated and real data.

To identify inefficient or poorly calibrated paddles, single charged track events were examined. The yield of these events was studied as a function of lab momentum and scintillator paddle hit. The yield for each sector in each bin was then divided by the maximum yield for that bin amongst all of the six sectors. These plots can be seen in Figures 5.11 through 5.16. These ratios were then examined for obvious systematic problems (on the order of 10%) that extended over large ranges in momentum. Several interesting types of features were found. First and foremost was a class of problems where, for a particular scintillator paddle, there was a swath of momentum bins which were inefficient. These were easily rejected (see for example paddle 35 of sector 1 in Figure 5.11.) The second class of phenomenon was an inefficiency that moved from one paddle at high momenta to more backward angle paddles at low momenta. These are believed to be associated with inefficiencies in the drift chamber. Although it might be possible to adequately model these inefficiencies by knocking out associated regions in the drift chambers during Monte-Carlo, it was decided that an easier and quicker

Figure 5.10: Angular distribution in lab coordinates (in degrees) of kaons after fiducial cuts.

method was to knock out all the affected time of flight paddles. Although this did reduce the statistics in some regions, it was much more straightforward to model the detector accurately with ToF knockouts. An example of these “DC” related problems can be seen in sector 4, paddles 43-48 of Figure 5.14.

There were a few prominent features for which we decided not to remove the associated scintillator paddles. One such case for which we did not remove any paddles was in sectors 1 and 2, at very low momenta, between paddles fifteen and twenty-five. This region was actually kinematically forbidden for two particle events (one charge, one neutral), but allowed in cases where there were at least three particles and only one was detected. The measured yield in this region was also considerably smaller than at higher momentum. The yield was changing rapidly across this region, so this discrepancy may also have been caused by slight misalignment between the sectors and did not represent either a real inefficiency in the detectors, or a problem which would affect the measurement of the reactions studied in this analysis.

The almost horizontal stripes visible in most sectors was caused by the knocking out of several tagger E-counters. Slight misalignments between the sectors then break

the symmetry of the detector and result in the yields in some bins being higher in one sector and the neighboring bins being lower. Plots similar to Figures 5.11 through 5.16 for the cases when the charged track was identified as either a pion or a proton were also examined. Several paddles were found for which minimum-ionizing pions did not always cause a trigger. Table 5.1 lists the paddles that were knocked out.

The inefficiencies represented in this list of knockouts represents problems associated with single track events. Since many of these Time-of-Flight related problems have not been observed in multiple charged particle events, it is believed that these inefficiencies are caused by problems with the signal going from the scintillators to the trigger logic after the signal to the TDC's had already been split off. Multiple charged-track events in which one track hits a "bad" paddle would still be seen because the second particle would trigger the event. The signal to the tdc from the first particle would still be seen, and so no inefficiency would be seen. Thus, the use of multiple charged particle states in CLAS to measure the single particle efficiency is limited when looking at single charged particle reactions by the need to also measure the trigger efficiency for each ToF paddle.

Figure 5.11: Bin-by-bin ratio of single charged particle yield in sector 1 to maximum yield of all sectors as a function of lab momentum and time-of-flight paddle hit.

Figure 5.12: Bin-by-bin ratio of single charged particle yield in sector 2 to maximum yield of all sectors as a function of lab momentum and time-of-flight paddle hit.

Figure 5.13: Bin-by-bin ratio of single charged particle yield in sector 3 to maximum yield of all sectors as a function of lab momentum and time-of-flight paddle hit.

Figure 5.14: Bin-by-bin ratio of single charged particle yield in sector 4 to maximum yield of all sectors as a function of lab momentum and time-of-flight paddle hit.

Figure 5.15: Bin-by-bin ratio of single charged particle yield in sector 5 to maximum yield of all sectors as a function of lab momentum and time-of-flight paddle hit.

Figure 5.16: Bin-by-bin ratio of single charged particle yield in sector 6 to maximum yield of all sectors as a function of lab momentum and time-of-flight paddle hit.

Sector	ToF id	comment
Sector 1:	24, 29 30...48	pions weak
Sector 2:	12, 13, 14, 15 18 33...48	DC problems DC problems
Sector 3:	17, 18 23 30, 33, 41 42 43 44 45 46, 48	pions weak dead pions weak pions weak dead
Sector 4:	28, 42 43...48	DC problems
Sector 5:	32 42 44, 45 46, 47 48	pions weak pions weak
Sector 6:	32 40...45, 48	dead

Table 5.1: List of knocked out time-of-flight paddles, and reasons for some cases. 'pions weak' means that problems were seen in the single pion yields. 'DC problems' represent paddles that have pathologies consistent with there having been inefficient regions in the drift chambers. The paddles associated with 'DC problems' would be different for negatively charged particles.

5.4 Charged Particle Identification

The problem of identifying the type of charged particle that is associated with a track in the CLAS detector is a task with many different requirements that strongly depend upon the reaction being studied and the observable being measured. When attempting to identify pions and protons, it is possible to make very wide cuts since there is very little, almost negligible, background. For identifying kaons, however, the task is much more difficult since the cross section for pion production is several orders of magnitude greater and therefore the tail of the pion measured mass extends beneath the measured kaon mass peak. As a result, this analysis used different methods for identifying pions, protons and kaons. Particles with $P_{lab} < 0.20 \text{ GeV}/c$ were uniformly cut for all reactions. No cut was made on the number of positively charged particles in each event. The identification of pions is discussed in Appendix A

5.4.1 Identifying Kaons

Detecting a K^+ in the CLAS involves the measurement of the time of flight (T), the path length (L), and the momentum (P) of the particle. The momentum and path length were determined from the reconstructed trajectory through the drift chambers and toroidal magnetic field. The time of flight measurement was based upon the ToF scintillator counters and the start time derived from the photon tagger. The mass was then determined by the simple formulas:

$$\beta = \frac{L}{cT} \quad (5.1)$$

$$\gamma = (1 - \beta^2)^{-1/2} \quad (5.2)$$

$$m = \frac{P}{\beta\gamma} \quad (5.3)$$

From this it can be shown that:

$$\frac{dm}{m} = \frac{dP}{P} + (1 + \frac{P}{m})(\frac{dL}{L} - \frac{dT}{T}) \quad (5.4)$$

and hence, since P, L, and T are approximately independent measurements,

$$\frac{\sigma_m}{m} = \sqrt{(\frac{\sigma_P^2}{P^2}) + (1 + \frac{P}{m})^2(\frac{\sigma_L^2}{L^2} + \frac{\sigma_T^2}{T^2})} \quad (5.5)$$

From this, it is obvious that there should be a momentum dependence to the width of the inferred mass. To take this into account, a mass cut that varies as a function of momentum was used. Figure 5.17 shows the momentum vs measured mass in a log

scale for a single run. Figure 5.18 shows the measured mass from a similar histogram sliced up into 100 MeV wide bins in lab momentum, with the additional constraint that the missing mass had to be consistent with $\gamma p \rightarrow K^+ \Lambda$. Each distribution was then fit with a Gaussian plus linear background and the results of the fit are displayed in Figure 5.19. For comparison, Figure 5.20 shows the fit results for $\gamma p \rightarrow K^+ \Sigma^0$. As one would expect, the results were nearly identical for the peak location(mean) and width(sigma). The slight rise in the mean as a function of P_{lab} represents a slight mis-calibration of $< 100ps$. This rise was taken into account when applying the mass cut for good kaons which was defined as being $\pm 2\sigma$ from the peak. The sidebands were defined in the regions 2–3 σ from the peak, and the “far” sidebands were defined as being 3–4 σ from the peak.

For purposes of estimating the systematic errors introduced by this method of identifying kaons, another method, based upon the time of flight of the particle was also used. This method used the time difference between the measured time of flight, and the time of flight if the particle were a kaon, assuming the measured momentum and track length were accurate. For an assumed particle mass, m_0 , this quantity is simply:

$$\Delta t = t_{meas} \left(1 - \sqrt{\frac{m_0^2 + p^2}{m_{meas}^2 + p^2}} \right) \quad (5.6)$$

For kaons two different timing cuts were used: the “Ktof” cut of $|\Delta t| < 1.0$ ns, and the “Ktof-tight” cut of $|\Delta t| < 0.6$ ns. Figures 5.21 and 5.22 show the time difference as a function of lab momentum of all particles, assuming the mass of kaons and protons respectively.

As will be seen, kaon selection based on these three different methods have varying systematic biases as a function of momentum and angle. The yield extraction procedure, discussed in Section 5.6, uses these three kaon identification cuts to estimate the systematic uncertainty in the extracted yields caused by these biases.

5.4.2 Identifying Protons

Detecting protons involves the same measurements as the kaon. However, due to the higher mass, the kinematics of the reaction cause protons with the same momentum to have a smaller β . This makes the measurement of the mass less sensitive to the timing calibrations. Also, since there was no significant source of background, there was less of a need to have as narrow a cut as possible. This allows us to make one wide cut in mass without either increasing the background or misidentifying real protons. To this end, we used a mass cut from 700 MeV to 1200 MeV.

Figure 5.17: Momentum vs measured mass summed over the whole of CLAS. The track density is plotted on a log scale. One can see the pion, kaon, and proton bands.

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Figure 5.18: Measured kaon mass distribution for different P_{lab} bins, showing yield fits used to define the widths of the momentum dependent mass cuts.

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Figure 5.19: Fit parameters and calculated signal-to-noise (stn) for the measured mass distribution of $\gamma p \rightarrow K^+ \Lambda$.

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Figure 5.20: Fit parameters and calculated signal-to-noise (stn) for the measured mass distribution of $\gamma p \rightarrow K^+ \Sigma^0$.

Figure 5.21: Time difference assuming kaon mass vs momentum. The horizontal white lines represent the Ktof and Ktof-tight cuts. The white curve represents the result of the mass cut ($\pm 3\sigma$ and $\pm 4\sigma$) for a particle with a path length of 5m. The log scale enhances the region between the Kaon and Pion peaks.

Figure 5.22: Time difference assuming proton mass vs momentum. The white curves represent the mass cuts used for a particle with a path length of 5m.

5.5 Kaon Missing Mass

In order to identify the hyperon that recoiled against from the charged kaon, the missing mass was calculated for the reaction. Before doing this, the CLAS standard energy-loss correction was applied to each particle. Figure 5.23 shows the calculated missing mass vs momentum for kaon candidates whose time-of-flight was within ± 1 ns of a kaon with the same measured momentum. This is the “Ktof” definition of a kaon candidate. Note that there was no noticeable momentum dependence to the Λ and Σ^0 peaks. A strict cut was used to identify the recoil particle in the extraction of polarization results. That cut was made between 1.08 GeV and 1.15 GeV in the missing mass for Λ particles, and cut from 1.16 GeV to 1.23 GeV was made for the Σ^0 particles. However, for extracting differential cross sections a more elaborate technique was used as described below.

5.6 Extracting The Yields

After the candidate kaon events were identified, the yields for the two reactions, $\gamma p \rightarrow K^+ \Lambda$ and $\gamma p \rightarrow K^+ \Sigma^0$ were measured by performing a log-likelihood fit of the missing mass squared in each of the E_γ and $\cos(\Theta_{cm}^K)$ bins. Figures 5.24 – 5.31 show six of the approximately 1200 sets of fits. For the momentum dependent mass cuts, all counts within 4σ of the peak were used. The function that was fit was a two Gaussian fit with a polynomial background, where the means of the two Gaussians were constrained so that the difference in the missing mass would be equal to the difference of the masses of the Λ and Σ^0 . The widths of the Gaussians were also constrained so that the width of the Σ^0 was not greater than the width of the Λ and no less than half that width.

For each type of particle identification, the background was fit as a constant, a linear, and a quadratic function, and the result with the minimum normalized χ^2 was taken. If this χ^2 value was less than 1.75 and the signal-to-noise ratio was greater than 2.5, the fit was accepted. The signal-to-noise was defined as the fitted area of the peak over the fitted area of the background within 2σ of the fitted mean of that peak. This was necessary to exclude bins where the background underneath the peaks due to misidentified pions (primarily in the Ktof method) were sufficiently large that the polynomial approximation for this pion background was very poor.

Assuming at least two of the three results passed the χ^2 and signal-to-noise cuts, the two results with the largest number of counts in each peak were averaged as our best estimate of the yield, and the variance between these two results was used as an estimate of the systematic error.

Figure 5.23: Missing mass vs. P_{lab} for + charged particles for which the time-of-flight was within $\pm 1\text{ns}$ of a kaon with the same momentum. Note that for $P_{\text{lab}} > 1.2 \text{ GeV}$ real pions will be identified as kaons with this cut.

Figure 5.24: Missing mass squared fits showing the resultant number of counts for the Λ and Σ^0 , the ratio of the widths, and the χ^2 for a constant(green), linear(purple), and quadratic(blue) background. The light blue points in the bottom plot represent the counts that were between 3 and 4 σ in the mass cut.

Figure 5.25: Missing mass squared fits showing the resultant number of counts for the Λ and Σ^0 , the ratio of the widths, and the χ^2 for a constant (green), linear (purple), and quadratic (blue) background. The light blue points in the bottom plot represent the counts that were between 3 and 4 σ in the mass cut.

Figure 5.26: Missing mass squared fits showing the resultant number of counts for the Λ and Σ^0 , the ratio of the widths, and the χ^2 for a constant (green), linear (purple), and quadratic (blue) background. The light blue points in the bottom plot represent the counts that were between 3 and 4 σ in the mass cut.

Figure 5.27: Missing mass squared fits showing the resultant number of counts for the Λ and Σ^0 , the ratio of the widths, and the χ^2 for a constant (green), linear (purple), and quadratic (blue) background. The light blue points in the bottom plot represent the counts that were between 3 and 4 σ in the mass cut.

Figure 5.28: Missing mass squared fits showing the resultant number of counts for the Λ and Σ^0 , the ratio of the widths, and the χ^2 for a constant(green), linear(purple), and quadratic(blue) background. The light blue points in the bottom plot represent the counts that were between 3 and 4 σ in the mass cut.

Figure 5.29: Missing mass squared fits showing the resultant number of counts for the Λ and Σ^0 , the ratio of the widths, and the χ^2 for a constant (green), linear (purple), and quadratic (blue) background. The light blue points in the bottom plot represent the counts that were between 3 and 4 σ in the mass cut. This plots shows how the background from misidentified pions can distort the fit

Figure 5.30: Missing mass squared fits showing the resultant number of counts for the Λ and Σ^0 , the ratio of the widths, and the χ^2 for a constant(green), linear(purple), and quadratic(blue) background. The light blue points in the bottom plot represent the counts that were between 3 and 4 σ in the mass cut.

Figure 5.31: Missing mass squared fits showing the resultant number of counts for the Λ and Σ^0 , the ratio of the widths, and the χ^2 for a constant(green), linear(purple), and quadratic(blue) background. The light blue points in the bottom plot represent the counts that were between 3 and 4 σ in the mass cut.

5.7 Measuring The Hyperon Recoil Polarization

As discussed in Section 1.2.1 the proton from the decay of the Λ is preferentially oriented in the direction of the Λ polarization, so that one expects the acceptance-corrected proton yield, $I(\Theta_{\Lambda_{R.F.}}^P)$ to have the form: $I(\Theta_{\Lambda_{R.F.}}^P) = 1 - \alpha P \cos(\Theta_{\Lambda_{R.F.}}^P)$ where α is the self analyzing power of the Λ , approximately (0.64). The fit to a line therefore enables one to extract the recoil polarization of the Λ . The Σ^0 s decay more than 99% of the time into a photon and a Λ . It can be shown that the resultant Λ polarization is $-\frac{1}{3}$ of the original polarization of the Σ^0 , and so it is also possible to extract the Σ^0 recoil polarization using this method, although the uncertainty in the measurement will be three times as large.

Two methods were used to obtain the proton yields necessary to extract the polarization information. First, kaons were selected with the momentum dependent mass cuts using the $\pm 2\sigma$ cut, to minimize the pion contamination. Separately, the Ktof cut was used. After identifying a kaon and proton, a strict cut on the missing mass was made (described in Section 5.5) and all events within that cut were counted. These yields were then fit to a linear function multiplied by the acceptance in that bin.

For each bin the measured polarization using the two different kaon identification methods were averaged and the variance was taken as the systematic uncertainty. In most cases this systematic uncertainty was much smaller than the statistical error. Figure 5.32 shows the fits of the proton distribution yields for one bin in E_γ and $\cos(\Theta_{cm}^K)$. The blue points are from kaons identified with the momentum dependent mass cut where as the green points are from the Ktof cut. In each plot the lower data points are the uncorrected yield and the upper plots are the acceptance corrected yields. The linear fits are also shown. The plot in the lower right shows the polarization for the two different methods of kaon identification and in magenta the average of the two with their variance used as an additional systematic error.

5.8 Photon Normalization

Photon normalization refers to the process by which the number of photons in a given energy bin which reached the target during the running period was determined. Obviously this is an important aspect of any cross-section measurement, as the cross-section is proportional to the acceptance corrected yield divided by this number of photons, N_γ . Throughout the lifetime of CLAS, there have been several methods developed, each an improvement upon its predecessors, to calculate these numbers. The method I will describe below was the one used in this analysis. It was originally developed by Eugene Pasyuk and Jim Ball and is known as the “G-flux” method.

Figure 5.32: Fits of the acceptance corrected proton yields' angular distributions in the Λ_{cm} relative to the normal to the reaction plane. The histogram in proceed from $\cos(\Theta_{cm}^K)$ of -0.9 in the upper left to a $\cos(\Theta_{cm}^K)$ of $+0.9$ at the bottom. The histogram in the bottom right is a plot of the resultant polarization. More detail may be found in the text.

5.8.1 The G-flux Method

The G-flux method uses random, out-of-time hits in the tagger which were not associated with the trigger as the starting point of calculating N_γ . The process was performed for each ten second interval between scaler events and calculated for each T-counter in the tagger. Finer energy binning was achieved by redistributing the results from a particular T-counter to individual E-bins according to the proportion of hits in the E-bin associated with the T-counter using the so called “E-T matrix”.

In order to calculate N_γ in a particular counter, the G-flux method takes advantage of the single-hit nature of the TDC’s. The events in a T-counter, i , were considered out-of-time if they occurred between 30 ns and 180 ns after the T-counter that triggered the event, as determined by coincidence with the start-counter. We shall refer to the measured number of these hits in a particular counter as N_{out}^i . Since the TDCs were single hit TDCs, any hit before the 30 ns cutoff would preclude there being any hits within the time window. Therefore, the number of hits in that T-counter that occur before the 30 ns, N_{early}^i , was subtracted from the total number of triggers in the whole tagger, N_{tot} , to give the number of random triggers, $N_{trig}^i = (N_{tot} - N_{early}^i)$, that could have a hit in the out-of-time window of the i^{th} counter.

To calculate the number of expected out-of-time hits, it is necessary to realize that if there was more than one hit in the out-of-time period of a single trigger, it would count only once. Therefore, the probability of there being a hit within this time period follows the Poisson distribution for there being at least one hit. This is

$$P_{\geq 1} = 1 - e^{-R^i \Delta T_{out}} \quad (5.7)$$

where R^i is the rate in a single T-counter and ΔT_{out} is the out-of-time window, 150ns. From this, we can get the expected number of out-of-time hits,

$$N_{out}^i = P_{\geq 1} \cdot N_{trig}^i = (1 - e^{-R^i \Delta T_{out}}) N_{trig}^i. \quad (5.8)$$

Solving for the rate gives:

$$R^i = \frac{-1}{\Delta T_{out}} \ln(1 - \frac{N_{out}^i}{N_{trig}^i}). \quad (5.9)$$

Before calculating N_γ from R , it was necessary to estimate the number of “tagged” photons that actually make it to the target, the “beam-line transmission function”. We measured this function in a series of “normalization” runs in which the luminosity of the beam was decreased so that a total absorption counter can be placed in the beam-line, downstream of the target. The beam-line efficiency, ϵ^i was then calculated for a particular counter by taking the ratio of hits in the tagger Master OR (MOR) in coincidence with the total absorption counter (TAC) to the hits in the tagger: $\epsilon^i = (\frac{MOR * TAC}{MOR})^i$. After this, there was still a correction that needs to be made for the loss of photons in the region between the target and the total absorption counter. This correction has been calculated to be approximately 4% ($\eta = 0.039 \pm 0.002$) [34].

The last information needed to calculate the number of photons that hit the target was the live-time corrected running time (ΔT_{run}). This was obtained by multiplying the run gated clock time (ΔT_{rg}), by the live time (L_t):

$$L_t = N_{tot}/T_{total} \quad (5.10)$$

$$\Delta T_{run} = \Delta T_{rg} L_t \quad (5.11)$$

where T_{total} is the total number of triggers measured by the scalars. With this it was possible to calculate the number of photons that hit the target at a particular energy:

$$N_{gflux}^i = R^i \Delta T_{run} (\epsilon^i)_{norm} \quad (5.12)$$

$$N_\gamma^i = N_{gflux}^i (1 + \eta) \quad (5.13)$$

5.8.2 Corrections to G-flux

The G-flux methodology is predicated on the assumption that the tagger trigger efficiency for a given energy bin is equivalent to the efficiency of the tagger recording a hit in the same energy bin when some other tagger element actually caused the trigger. This assumption breaks down if there is a corruption of the signal to either the trigger or the TDC's. In order to take such an effect into account, the normalized yield of all particles was measured with both the 2.445 GeV endpoint energy and with the 3.115 GeV endpoint energy. These yields can be seen in Figure 5.33. The notable feature of comparing these two normalized yields is that the CLAS acceptance for particles at a given E_γ will not change, but the tagger detector element associated with that photon energy does. The ratio of these two normalized yields was taken, and values less than 1 were used as a measure of the inefficiency of the tagger in the 2.445 GeV data set. Values greater than 1 were interpreted as inefficiencies in the 3.115 GeV data set. This ratio can be seen in Figure 5.34.

It is clear that ideally one would like to measure the trigger inefficiencies for the individual elements within the tagger, particularly for the 61 T-counters. Unfortunately, this method does not allow one to directly measure the inefficiency for any particular T-counter. However, one can see that the largest correction near 1.1, almost a factor of 2, is in the region of T-counters 35-37 which had several problems during the running period. From Figure 5.34 we can see that the average correction was $\sim 97\%$ which indicates that on average the 3.115 GeV data had a higher normalized yield. Since the cross-section doesn't change as a function of time, this indicates an overall trigger inefficiency in the 2.445 GeV data.

Figure 5.33: Normalized all particle yields for two endpoint energies. The open squares are from the 2.445 GeV data set and the filled triangles are from the 3.115 GeV data set.

Figure 5.34: Ratio of the 2.445 GeV normalized yield to the 3.115 GeV normalized yield. The red line is a fit to a constant value.

5.9 Summary

In this chapter I have described the cuts used to reduce the background as well as the method of identifying kaons and extracting the yields for $\gamma p \rightarrow K^+\Lambda$ and $\gamma p \rightarrow K^+\Sigma^0$. Using the G-flux method to calculate the number of photons in each energy bin allows us to obtain the normalized yields. In the next chapter I describe the methods used to calculate the acceptance, the final step in extracting the cross-sections.

Chapter 6

Detector Simulation

The purpose of the detector simulation is to compute the acceptance, as a function of the various kinematic variables of a given reaction for the CLAS detector to detect and correctly identify that reaction. For this analysis we used two different methods to estimate this acceptance. The first method uses a GEANT based simulation of the CLAS (GSIM) to create simulated events that are then processed just like the real data. This method represents the most accurate means of estimating the detector acceptance short of actually measuring it, but it suffers from being very slow. The second method that we used to estimate the acceptance of the detector was based upon a sophisticated parameterization of the acceptance. The advantage of the second method is that it is much quicker and therefore leads to a much smaller statistical uncertainty, although the price is a larger systematic uncertainty. Because of the speed of the parameterized calculation, it was also possible to investigate the size of effects that different sources of inefficiency have on the acceptance.

It is worth noting that what we have termed acceptance here is similar to what would have been called the product of the geometric acceptance and detector efficiency in other more simple experiments. Due to the complexity of the CLAS, separating these two quantities is not necessarily useful and so we combine the two to form what we call the acceptance.

6.1 GSIM

The GSIM [35] based acceptance calculation began with the generation of 1 million $K^+\Lambda$ and 1 million $K^+\Sigma^0$ events. It should be noted that this was a physics Monte Carlo as no background channels were simulated. The events were generated pseudo-randomly from a distribution that was flat in both the energy of the photon (E_γ) and $\cos(\Theta_{cm}^K)$. The primary vertexes of the generated events were distributed evenly along the nominal beam axis between ± 10 cm. Each generated event was then run through GSIM which “swam” the particles from the generated event through the simulated detector and creating a set of simulated hits in various detector elements. The geometry of the detectors in GSIM used the best estimates for the location of the

various detector elements [36]. The swimming of the particle through the detector includes calculations for multiple-scattering, particle decays, and interactions with scattering from the structural elements of the CLAS.

The output from GSIM is a set of simulated events produced with perfect timing resolution for each detector, so after an event is produced GSIM, it was processed through a program called gpp [37] which smeared the output of the time-of-flight paddles and the drift chambers, by adding a pseudo-random time from a Gaussian distribution to each detector. The widths of the smearing were functions of both the length of the paddle for the time-of-flight paddles, and, in the case of the drift chambers, the distance of closest approach, of the track to the wire (DOCA). In Figures 6.1-6.3 the smearing of the drift chamber hits can be seen as a function of the distance of closest approach and Figure 6.4 shows the time-of-flight smearing as a function of the length of the paddle for the left TDC value. The smearing of the right TDC was similar. The nominal values of the smearing widths were determined using the residuals measured during the calibration procedures. The overall scale relative to the nominal value for each detector was chosen so as to have the widths of the measured kaon mass from the simulated events be consistent with that of the real data. This scale factor was 1.0 for the time-of-flight paddles, 2.0 for region 1, 2.5 for region 2, and 3.0 for region 3. The program gpp was also capable of knocking out “dead” detector elements, including individual drift cells. This feature, however, was not used in this analysis since the time-of-flight knockout procedure used in the analysis was deemed to sufficient in its ability to take into account the dead and inefficient detectors.

After the output of GSIM was smeared, the simulated data, like the real data, was processed through alc (see Chapter 4) to produce “cooked” events. These cooked events were then analyzed with the same analysis as the real data. The same fiducial cuts and time of flight knockouts were applied and the same cuts were used in identifying the particles. Figure 6.5 shows a time difference plot for the simulated data to be compared with Figure 5.21 from the real data. As with the real data, the missing mass squared off of the kaon was fit in each bin. A sample fit may be seen in Figure 6.7. The resultant yield for each reaction in each bin was divided by the number of events of that type that were generated in that bin to calculate the acceptance. The resultant acceptance for $\gamma p \rightarrow K^+ \Lambda$ and $\gamma p \rightarrow K^+ \Sigma^0$ as a function of $\cos(\Theta_{cm}^K)$ and E_γ may be seen in Figures 6.8 and 6.9 respectively.

6.2 Parameterized Monte Carlo

The parametric Monte Carlo (PMC) acceptance calculation is based upon the idea that the acceptance can, to first order, be broken down into a geometric acceptance and a correction for the probability of particles decaying in flight. After an event is generated, the PMC first applies the same momentum and angle dependent fiducial cuts from pseudo-spa and then uses the E-loss package to calculate the average energy

Figure 6.1: DOCA smearing from GPP versus DOCA value in region 1 of the drift chambers.

Figure 6.2: DOCA smearing from GPP versus DOCA value in region 2 of the drift chambers.

Figure 6.3: DOCA smearing from GPP versus DOCA value in region 3 of the drift chambers.

Figure 6.4: time-of-flight smearing of the left TDC value as a function of the length of the paddle.

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Figure 6.5: Time difference assuming kaon mass vs. momentum for the GSIM generated data. The horizontal white lines represent the Ktof and Ktof-tight cuts. The white curve represents the result of the mass cut ($\pm 3\sigma$ and $\pm 4\sigma$) for a particle with a path length of 5m. The log scale enhances the region in the tail of the kaon peak.

Figure 6.6: Fit parameters and calculated signal-to-noise (stn) as a function of lab momentum for the measured mass distribution of simulated GSIM data for the $\gamma p \rightarrow K^+ \Lambda$ reaction. This figure compares to Figure 5.19 in the real data.

Figure 6.7: Missing mass squared fits to GSiM generated data showing the resultant number of counts for the Λ and Σ^0 , the ratio of the widths, and the χ^2 for a constant (green), linear (purple), and quadratic (blue) background. The light blue points in the bottom plot represent the counts that were between 3 and 4 σ in the mass cut.

Figure 6.8: Calculated GSIM Acceptance for $\gamma p \rightarrow K^+ \Lambda$. The 1 million events were thrown evenly in E_γ and $\cos(\Theta_{cm}^K)$.

Figure 6.9: Calculated GSIM Acceptance for $\gamma p \rightarrow K^+ \Sigma^0$.

loss that each particle would experience when passing through the target and start counter. Multiple scattering of the particles in the target and start counter were accounted for by smearing the momentum vector with a Gaussian distribution. The trajectory of the particle was then calculated by “swimming” the particle through the magnetic field and recording which detector elements were hit. The swimming of the tracks used the full knowledge of the CLAS geometry as also implemented in GSIM. Two very small corrections ($< 1\%$) were made. The first for drift chamber knockouts was based on the standard glc list of weak cells and the requirements of the tracking code to have five of six layers hit per segment in the drift chamber super layers and five of six segments needed per track. The second correction was for events where two charged particles were produced with similar kinematics such that the tracking algorithms might have had difficulty distinguishing between them.

The probability of a particle’s decay was then used to calculate a weight for each event representing the chance that the particle was detected by CLAS. If the particle hit a ToF paddle which was not in the knockout list (see Table 5.1) the probability that the particle survived without decaying before reaching the ToF paddles, $P_{ToF} = e^{-L/(\gamma\beta c\tau)}$, was used as the initial weight. If this survival probability was less than 99.8%, then an additional weight was added proportional to the probability that the particle survived without decaying before leaving the toroidal magnetic field. This probability was represented as the probability of decaying ($1 - P_{ToF}$) multiplied by a weighting factor for the decay products hitting the original paddle or its immediate neighbors:

$$W_i = \frac{\int_{\frac{3}{4}L}^L e^{-\frac{x}{\gamma\beta c\tau}} g_i(x) dx}{\int_0^L e^{-\frac{x}{\gamma\beta c\tau}} dx}, \quad (6.1)$$

where L is the length of the track, and $g_i(x)$ is the fraction of decays at a distance x along the path that would hit the i^{th} ToF paddle. $g_i(x)$ was calculated using the average probability for the charged decay product, from isotropic two body decays, hitting either the paddle in question. This probability was set to 0 for knocked out paddles, otherwise it was calculated numerically by first calculating the angle in the lab that a decay product at 90° in the rest frame of the original particle would be boosted to, and then integrating the resultant cone’s intersection with the paddle. This calculation of the intersection assumed that the track was, in the lab frame, almost perpendicular to the face of each of the three paddles, and that the neighboring paddles were the same size as the original paddle. The resultant final weight was therefore:

$$W = K_i P_{ToF} + (1 - P_{ToF})(K_i W_i + K_{i+1} W_{i+1} + K_{i-1} W_{i-1}) \quad (6.2)$$

where K_i is 0 if the paddle was knocked out and 1 otherwise. The acceptance was then calculated bin-by-bin as the sum over all events in each bin of this weight divided by the number of thrown events in that bin. The acceptance was calculated in this

manner using 40 million events each for $\gamma p \rightarrow K^+\Lambda$ and $\gamma p \rightarrow K^+\Sigma^0$ and the results are shown in Figures 6.10 and 6.11, respectively. The average statistical error in each bin was $< 1\%$.

6.3 Comparison

The primary reason for having used these two quite different methods for calculating the acceptance of CLAS was to be able to simultaneously have a statistically very precise estimate of the acceptance while also being able to estimate remaining systematic uncertainties. The ratio of the efficiencies calculated by the two different methods for both $\gamma p \rightarrow K^+Y$ reactions can be seen in Figures 6.12 and 6.13. The bin-by-bin ratio with errors can be seen in the “linearized” version of the plots in Figures 6.14 and 6.15. The two methods were initially compared in a “no decay” mode (not shown) in which the decays of particles in GSIM were turned off and P_{ToF} was set equal to 1 in the PMC calculation. After this initial comparison was found to give statistically consistent results, the decays were turned back on. For the kaon events in this analysis, the decay effects dominate the acceptance and good agreement between the two methods was only achieved after the modeling of the inclusion of the $P_{\frac{3}{4}}$ decay weighting was included. This additional weighting beyond just the P_{ToF} survival had an overall effect of approximately 12%, but also significantly increased the agreement in areas where the time-of-flight knockouts have a large effect.

Despite these efforts, there remained some discrepancies in the acceptance, most notably at the forward-most and backward-most angles. These bins were excluded from the data. There also remains some discrepancies, caused by the time-of-flight knockouts, visible in the center bins. To quantify the remaining discrepancies between the two methods, the ratio for bins with a non-zero acceptance and inside the fiducial borders, discussed in Section 5.1, was fit to a constant value using a χ^2 fit, which is equivalent to taking the weighted mean of the values with the weighting factor as: $\frac{1}{\sigma^2}$. To estimate the systematic error, for each bin an additional $\sigma_{systematic}$ was added in quadrature to the statistical error, and the ratio was iteratively refit as the systematic error was increased until the normalized χ^2 of the fit was less than 1. The systematic error was then taken to be the average of the last two iterations. The resultant fit indicated that the ratio between the GSIM and PMC efficiencies was essentially 1 for $\gamma p \rightarrow K^+\Lambda$ and $\gamma p \rightarrow K^+\Sigma^0$ with a systematic error of $\sim 6.5\%$.

Comparison of these two methods was also done for the acceptance calculations for the $\gamma p \rightarrow \pi^+n$ reaction, discussed more in Appendix A, and gave an overall ratio within 3% of 1 and with a systematic error of approximately 6.5%. The remaining differences are believed to be caused by the more accurate modeling by GSIM of the tracking efficiency, the start-counter-tagger coincidence cuts, and the effects of decays. The overall scale factors were therefore applied to the parametric calculations before they were used to calculate the cross-sections.

Figure 6.10: parametrically calculated acceptance for $\gamma p \rightarrow K^+ \Lambda$.

Figure 6.11: parametrically calculated acceptance for $\gamma p \rightarrow K^+ \Sigma^0$.

Figure 6.12: Ratio of GSIM acceptance to parametrically calculated acceptance for $\gamma p \rightarrow K^+ \Lambda$. The black lines represent cuts that were applied to the data to exclude bins with very high uncertainties in the acceptance calculation. Bins with boxes in them are bins in which the real data was rejected for other reasons.

$$\gamma + p \rightarrow K^+ + \Sigma^0$$

Figure 6.13: Ratio of GSIM acceptance to parametrically calculated acceptance for $\gamma p \rightarrow K^+ \Sigma^0$. The black lines represent cuts that were applied to the data to exclude bins with very high uncertainties in the acceptance calculation. Bins with boxes in them are bins in which the real data was rejected for other reasons.

Figure 6.14: Ratio of GSiM acceptance to parametrically calculated acceptance for $\gamma p \rightarrow K^+ \Lambda$. The top plot shows the bins grouped by energy and has very little energy dependence. The bottom plot groups the bins by angle and shows that the angular dependence is mostly constrained to the forward-most and backward-most bins which are removed from the data.

Figure 6.15: Ratio of GSiM acceptance to parametrically calculated acceptance for $\gamma p \rightarrow K^+ \Sigma^0$. The top plot shows the bins grouped by energy and has very little energy dependence. The bottom plot groups the bins by angle and shows that the angular dependence is mostly constrained to the forward-most and backward-most bins which are removed from the data.

6.4 Summary

The acceptance of the CLAS detector was modeled using two different methods which agreed to within a 1% in overall scale with a systematic uncertainty of approximately 6.5%. Combined with the yields extracted in Chapter 5, these acceptance calculation produce the measured differential cross sections reported in Chapter 7.

Chapter 7

Results

In the previous chapters, I have gathered all the factors necessary to assemble the measured cross section and recoil polarization for the two reactions: $\gamma p \rightarrow K^+\Lambda$ and $\gamma p \rightarrow K^+\Sigma^0$. In this chapter, I first discuss the estimate of the systematic uncertainties, both those that were implemented in each bin individually and the global factors that remain. I then present the results of this analysis, the differential cross sections and the hyperon recoil polarizations.

7.1 Estimates of The Systematic Uncertainties

In any experiment, the estimation of the systematic uncertainty is an important, if sometimes tricky, part of the process. In this section, I discuss the systematic uncertainties of this measurement by referencing the three major pieces of the analysis: the yield extraction, the acceptance calculation, and the photon normalization. In the case of the yield extraction, most of the systematic uncertainty was estimated on a bin-by-bin basis and integrated into the results. The rest of the uncertainty in the yield, as with the uncertainty for the acceptance and normalization, has been left as a global estimate, and therefore was not integrated into the individual bins. Table 7.1 summarizes these systematic uncertainties including the average total uncertainty (systematic from kaon identification and statistical added in quadrature) from the bin-by-bin yield extractions.

7.1.1 Uncertainties in The Yield Extraction

As mentioned in Section 5.6, in each kinematic bin the missing mass (of the hyperons) squared was fit using two different methods of identifying the kaons and the variance between the resultant yields was considered an estimate of the systematic uncertainty and added in quadrature to the average statistical uncertainty reported by the fits. The fractional systematic uncertainty as a function of bin can be seen in Figure 7.1.1. The bins are sorted by energy and angle, such that the quasi-periodic spikes below bin 500 are associated with the backward-most angle bins. One can see that the

Source Effect	Correction	Where Applied	Uncertainty
Yield			
Kaon ID		in bin-by-bin results in E_γ and $\cos(\Theta_{cm}^K)$	5.8%(Λ), 5.3%(Σ^0) AVERAGE TOTAL already in results
Start-Counter-Tagger cuts	1.0-1.02	E_γ bin by E_γ bin	1.0% point to point
$\frac{d\sigma}{d\Omega}$ of $\gamma p \rightarrow \pi^+ n$	1.087	global scale	3.0% global uncertainty
Normalization			
photon counting(G-flux)		E_γ bin-by- E_γ bin	3.0% point to point
Photon Attenuation	1.039	globally	0.2% accounted for in $\gamma p \rightarrow \pi^+ n$
Pt-to-Pt Trigger (In)Efficiency Comparison of 2.4GeV and 3.1GeV endpoint Data	1.00-1.62	E_γ bin by E_γ bin	1.0% point to point
Target			
Density= $0.0711(\frac{g}{cm^3})$ [22]		globally	0.14% accounted for in $\gamma p \rightarrow \pi^+ n$
Length=178.5mm		globally	0.28% accounted for in $\gamma p \rightarrow \pi^+ n$
Acceptance			
comparison of GSIM to MC_Generator		global scale	6.9%(Λ) 6.3%(Σ^0) point to point and global
Overall Systematic (excluding Kaon ID effects on yield):			8.2%(Λ), 7.7%(Σ^0)

Table 7.1: Sources of systematic uncertainties and their estimated values.

Figure 7.1: The estimated fractional total uncertainty from the yield extraction for the reactions $\gamma p \rightarrow K^+ \Lambda$ and $\gamma p \rightarrow K^+ \Sigma^0$. This uncertainty is estimated from the variance in the yield results from different methods of identifying kaons added in quadrature to the statistical error. In almost all cases the systematic error dominates.

average of approximately 5.5% would be a significant overestimate for many bins. As discussed in Section 5.2, the yields were corrected for losses due to the start-counter - tagger coincidence cuts. This correction, though energy dependent, was generally less than 2%, and we estimate that this has an uncertainty of 1%.

7.1.2 Uncertainties in The Photon Normalization

The G-flux method presented in Chapter 5 is one in a chain of progressively improved algorithms for calculating the number of photons that hit the target for each energy bin. The estimated systematic uncertainty for the G-flux method, based on comparisons to the results of previous normalization methods, is approximately 2%. The correction to G-flux based upon the normalized yield of all particles at different end-point energies showed an average overall correction of approximately 3%. Although the correction in each individual bin is statistically very well determined, in light of this overall scale, we estimate the uncertainty due to the photon normalization to actually be 3%. The beamline photon attenuation calculation that is part of the G-flux calculation introduces an extra negligible 0.2% uncertainty.

7.1.3 Uncertainties in The Acceptance Calculations

As discussed in Chapter 6, the detector acceptance was calculated using two different methods that gave very good agreement on average ($< 1\%$) within the bins used in this analysis. However, the variance between the two results on a bin-by-bin basis required an additional systematic uncertainty of $6.5 - 7.0\%$. This uncertainty is believed to be a result of the decay parameterization and tracking efficiency issues. We use this number as our estimate of the systematic uncertainty of the acceptance. Although we quote a global uncertainty for this value, it is actually a point-to-point uncertainty, unlike, for example, the uncertainty on target length. We believe that any systematic biases in the results vary fairly smoothly across several bins, however, and that with rare exception 7% jumps between neighboring bins are unlikely.

7.1.4 Other Sources of Uncertainty

As seen in Appendix A, the measured cross sections for $\gamma p \rightarrow \pi^+ n$ are a factor of 0.9198 lower than the SAID parameterization. This has been attributed to an overall trigger inefficiency, therefore we rescaled the Λ and Σ^0 yields accordingly. This introduces the potential for an overall uncertainty caused by this rescaling. Although the estimated uncertainty on the 0.9198 from the fit is quite small, $\sim 0.2\%$, the difference in results as a function of angle leads us to estimate the systematic uncertainty from this correction as being 3% .

Since the target length and density were needed to calculate the cross section, uncertainty in their measurement contribute to the overall uncertainty. The estimated uncertainty for both these quantities turn out to be quite small. The results of both are, however, subsumed in the normalization to the pion data and therefore they do not add any additional uncertainty. They have been included in Table 7.1 for completeness.

7.1.5 Summary of The Systematic Uncertainty

The estimated systematic uncertainties listed in Table 7.1 are in two categories, the uncertainty from the yield extraction which is accounted for on a bin-by-bin basis and included in the reported results, and those that represent an overall uncertainty in the measurement. The reported average uncertainty for the bin-by-bin correction of the yield extraction is not included in the estimate of the overall systematic uncertainty since it has already been included in the individual data points.

7.2 Measured Cross Sections

The differential cross section was obtained by measuring the reaction yield (Y) as a function of photon energy and kaon production angle, and then correcting for the

acceptance of the detector(A), the number of photons(N_γ) that passed through the target, and the density ($\rho = 0.0711(\frac{g}{cm^3})$) [22] and length ($L = 178.5 \text{ mm}$) of the target:

$$\frac{d\sigma(E_\gamma, \Theta^K)}{d\cos(\Theta_{cm}^K)} = \frac{Y(E_\gamma, \Theta_{cm}^K)}{\delta\cos(\Theta_{cm}^K)A(E_\gamma, \Theta_{cm}^K)N_\gamma(E_\gamma)} \frac{H_{gpm}}{N_A\rho L} \quad (7.1)$$

where H_{gpm} is the number of grams per mole of the target ($1.00794\frac{g}{mole}$), N_A is Avogadro's number, and $\delta\cos(\Theta_{cm}^K)$ is the size of the $\cos(\Theta_{cm}^K)$ bin.

The resultant differential cross sections for $\gamma p \rightarrow K^+\Lambda$ can be seen in Figure 7.2, which displays the cross section as a function of E_γ for fixed $\cos(\Theta_{cm}^K)$ bins, and Figure 7.3 presents the cross section as a function of $\cos(\Theta_{cm}^K)$ for fixed E_γ bins. The corresponding plots for $\gamma p \rightarrow K^+\Sigma^0$ are in Figures 7.4 and 7.5. A table of the results can be found in Appendix B for $\gamma p \rightarrow K^+\Lambda$ and Appendix C for $\gamma p \rightarrow K^+\Sigma^0$. The error bars include statistical errors and estimated systematic uncertainty from the yield extraction, but do not include the overall estimated uncertainty from Table 7.1. These results will be discussed in detail in Chapter 8.

7.3 Recoil Polarization

As discussed in Section 5.7, the recoil polarization for both the Λ and Σ^0 were extracted by fitting the decay proton's angular distribution. Figures 7.6 and 7.7 show the recoil polarization averaged over all photon energies. The results plotted as a function of W , the total center of mass energy, for the Λ can be seen in Figure 7.8 and for the Σ^0 in Figure 7.10. The results are also shown as a function of $\cos(\Theta_{cm}^K)$ for fixed E_γ bins in Figures 7.9 and 7.11. The results include both statistical error and estimated systematic uncertainties. Because of the somewhat limited statistics, the data points were binned such that their fractional statistical uncertainty was less than 0.25. The horizontal error bars, when present, represent the total width of the original bins used in the re-binned data point.

Figure 7.2: Differential cross section for $\gamma p \rightarrow K^+ \Lambda$ as a function of E_γ in GeV.

Figure 7.3: Differential cross section for $\gamma p \rightarrow K^+ \Lambda$ as a function of $\cos(\Theta_{cm}^K)$. All energies, E_γ , are in GeV.

Figure 7.4: Differential cross section for $\gamma p \rightarrow K^+ \Sigma^0$ as a function of E_γ in GeV.

Figure 7.5: Differential cross section for $\gamma p \rightarrow K^+ \Sigma^0$ as a function of $\cos(\Theta_{cm}^K)$. All energies, E_γ , are in GeV.

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Figure 7.6: Recoil polarization in $\gamma p \rightarrow K^+ \Lambda$ averaged over all E_γ . The green curves are fits to the data constrained to be 0 at $\cos(\Theta_{cm}^K) = \pm 1$.

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Figure 7.7: Recoil polarization in $\gamma p \rightarrow K^+ \Sigma^0$ averaged over all E_γ . The green curves are fits to the data constrained to be 0 at $\cos(\theta_{cm}^K) = \pm 1$.

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Figure 7.8: Recoil polarization in $\gamma p \rightarrow K^+ \Lambda$ as a function of W .

Figure 7.9: Recoil polarization in $\gamma p \rightarrow K^+ \Lambda$ as a function of $\cos(\Theta_{cm}^K)$. All photon energies, E_γ , are in GeV.

Jefferson Lab / CLAS, Carnegie Mellon University, 10-02

Figure 7.10: Recoil polarization in $\gamma p \rightarrow K^+ \Sigma^0$ as a function of W .

Figure 7.11: Recoil polarization in $\gamma p \rightarrow K^+ \Sigma^0$ as a function of $\cos(\Theta_{cm}^K)$. All photon energies, E_γ , are in GeV.

Chapter 8

Discussion of Results

In Chapter 7 I presented the results of this analysis: the differential cross sections and recoil polarizations for the reactions $\gamma p \rightarrow K^+ \Lambda$ and $\gamma p \rightarrow K^+ \Sigma^0$. In this chapter I discuss those results and compare them to the existing data and theory. I also include a simple Legendre moment analysis to extract a total cross section and shape parameters. I conclude the chapter with suggestions for future work in analyzing these reaction channels.

In all histograms, the solid blue circles are the CLAS data and the open red triangles are the previous data from SAPHIR [5]. It should be noted that the binning used in this analysis and that of SAPHIR differ, with the centers of the SAPHIR bins typically between two of the CLAS bins; we therefore plot the SAPHIR data only once in the closest bin to our data.

The theory curves are calculated at the centers of the CLAS bins. The solid red curves are results calculated from the KAON-MAID web site:

<http://www.kph.uni-mainz.de/MAID/kaon/kaonmaid.html>

This represents the George Washington group's analysis [8] of the published SAPHIR data. The KAON-MAID calculation is a hadrodynamic model that in addition to the born terms includes the $K^*(893)$ and $K_1(1270)$ in the t-channel. For both $K \Lambda$ and $K \Sigma^0$ production the $S_{11}(1650)$, $P_{11}(1710)$, and the $P_{13}(1720)$ were included. In addition, in the $K \Lambda$ reaction, a $D_{13}(1900)$ was included in the s-channel. For the $K \Sigma^0$ case the $S_{31}(1900)$ and $P_{31}(1900)$ were also included.

The magenta curves are from another hadrodynamic model fit to the published SAPHIR data. The curves were provided by Stijn Janssen [10, 12] from his model "B" for calculating the $\gamma p \rightarrow K^+ \Lambda$ and $\gamma p \rightarrow K^+ \Sigma^0$ observables. The model uses the same resonances as KAON-MAID, but a different prescription for the hadronic form factors. In addition, it includes hyperon resonances in the u-channel.

The blue dash-dotted curves are results from a Regge model calculation [13], which emphasizes the effects of the t-channel. The model includes trajectories for both the K and K^* mesons as well as a Regge-ized s-channel nucleon exchange. The results were obtained from a private copy of the authors' program.

8.1 The Differential Cross Sections

Figure 8.1 shows the measured differential cross sections plotted against the previously mentioned theories and the SAPHIR data. The $\gamma p \rightarrow K^+\Lambda$ reaction contains 921 new data points while there are 783 new $\gamma p \rightarrow K^+\Sigma^0$ points.

There are several features apparent in comparing the new CLAS results to the SAPHIR data that are common to both the $\gamma p \rightarrow K^+\Lambda$ and $\gamma p \rightarrow K^+\Sigma^0$ reactions. Although at low energies in the backward angles, the data is in good agreement with the SAPHIR results, as one moves forward in angle (for example at $\cos(\Theta_{cm}^K) = -0.05$) and to higher photon energies, a discrepancy arises in that SAPHIR observed a significantly lower cross section in both the Λ and Σ^0 channels. This discrepancy can not be explained by the different binning of the data. We have looked for sources of this discrepancy in our extraction of the yields, photon normalization, and acceptance calculations. There are no known sources of uncertainty in our analysis that could explain these discrepancies.

Notably, the discrepancies for both reaction channels are not the same in their energy or angular dependencies. It is also worth noting that the discrepancies with the SAPHIR data are in regions where the $K^+\Lambda$ and $K^+\Sigma^0$ yields are very clean and so it is extremely unlikely that the CLAS measurements are overestimating the yield. See Fig 5.27 for an example. The comparison of the $\gamma p \rightarrow \pi^+n$ channel cross sections with SAID indicates that it is not caused by a problem with the CLAS photon normalization procedure and make it very difficult to explain the angular dependencies. The magnitude of angular and energy dependencies would rule out all of the known possible sources of problems in the acceptance calculation. It is therefore our conclusion that the overall difference in cross section measurements between SAPHIR and CLAS can not be attributed to any known or suspected peculiarities in CLAS or this analysis.

8.1.1 Comparisons of $\gamma p \rightarrow K^+\Lambda$ Results

In addition to the discrepancies observed in the forward angles in comparing the $\gamma p \rightarrow K^+\Lambda$ data to the SAPHIR results, significantly more backward peaking can also be seen in Figure 8.4. The lack of backward peaking in the SAPHIR data is an artifact of the lack of statistics necessary to resolve this feature.

This backward peaking is all the more striking when compared in Figure 8.1 against the KAON-MAID model. The difference in results at backward angles seems most prominent in the region around W of 1.9 GeV, where the KAON-MAID model predicts the largest effect for the D_{13} resonance (see $\cos(\Theta_{cm}^K) = -0.65$). In fact, our data has even greater peaking than the model does in this region, although it falls off with increasing W much more quickly. This should not, however, be taken as evidence for the existence of the D_{13} , as backward peaking is also indicative of the involvement of a u -channel process. It is clear that a refit to only the new differential

cross section data will be just as ambiguous.

The Janssen model calculations fall very close to the SAPHIR data from which they were fit and does not show the backward peaking nearly as strongly as the KAON-MAID model. A refit of the new CLAS data will probably be able to reproduce such a feature without additional diagrams because of the u -channel hyperon resonances already included.

As mentioned previously, in the forward direction we generally see a much larger cross section at all energies. In addition, however, starting in the $\cos(\Theta_{cm}^K) = +0.45$ bin, a new bump emerges at a W of about 1.95 GeV which is much more narrow than the one at 1.9 GeV in the backward angles. There is no sign of the 1.9 GeV bump in the forward angles, although it might be lost beneath the higher cross section from t -channel production. This bump at 1.95 GeV, with an estimated width of $\Gamma = 100$ MeV, may indicate that there is a resonance that contributes significantly at this energy but which has not previously been seen in this channel. The previously published data from SAPHIR had insufficient statistical power to observe this feature.

It is important to note that given the extent in $\cos(\Theta_{cm}^K)$ that this 1.95 GeV bump covers, it is unlikely to be caused by any systematic problems with the CLAS detector. Any systematic problem with the detector should have also showed up in the $\gamma p \rightarrow K^+\Sigma^0$ channel since the kinematics for kaons in two reactions are very similar.

The Regge model calculation is, in contrast to the KAON-MAID model, quite smooth as a function of W and generally over-estimates the $\gamma p \rightarrow K^+\Lambda$ production cross-section. Interestingly, it is best at the backward-most angles, which is not what one would expect for t -channel production. By looking at Figure 8.4 at an E_γ of 1.7 GeV, we see that this is more a coincidence caused by the backward peaking of the cross-section. The Regge model can not be expected to reproduce the bumps and wiggles seen in W distributions since it does not explicitly include any nucleon resonances which would normally cause such features.

8.1.2 Comparisons of $\gamma p \rightarrow K^+\Sigma^0$ Results

The comparison of the $\gamma p \rightarrow K^+\Sigma^0$ reaction is much less dramatic. The W distributions in Figures 8.5 - 8.7 do not show significant structure beyond the large peak around 1.9 GeV which is believed to be caused by the large collection of N^* 's and Δ 's in this region. There is some hint to an overall backward peaking at values of W greater than 2 GeV, better seen in the bottom plots of Figure 8.8.

In the forward direction, there is perhaps a hint of a broad bump centered just above 2.0 GeV that is not reproduced in any of the models, however, the lack of data in the forward-most bins at the highest energies makes this bump more difficult to distinguish from an overall rise caused by t -channel processes. In this vein, the Regge model seems to do reasonably well at the highest energies for $\gamma p \rightarrow K^+\Sigma^0$, in that the data looks to be asymptotically approaching the Regge model prediction.

Figure 8.1: Comparison of the CLAS $\gamma p \rightarrow K^+ \Lambda$ differential cross section (blue circles) with SAPHIR [5] data (red triangles). The solid red line is calculated from KAON-MAID [8], the blue dash-dotted line is calculated from the Guidal-Laget Regge model [13], and the magenta dashed line is model B provided by Janssen [10].

Figure 8.2: Comparison of the CLAS $\gamma p \rightarrow K^+ \Lambda$ differential cross section (blue circles) with SAPHIR [5] data (red triangles). The solid red line is calculated from KAON-MAID [8], the blue dash-dotted line is calculated from the Guidal-Laget Regge model [13], and the magenta dashed line is model B provided by Janssen [10].

Figure 8.3: Comparison of the CLAS $\gamma p \rightarrow K^+ \Lambda$ differential cross section (blue circles) with SAPHIR [5] data (red triangles). The solid red line is calculated from KAON-MAID [8], the blue dash-dotted line is calculated from the Guidal-Laget Regge model [13], and the magenta dashed line is model B provided by Janssen [10].

Figure 8.4: Comparison of the CLAS $\gamma p \rightarrow K^* \Lambda$ differential cross section (blue points) for a few representative energy bins chosen so that all the SAPHIR data from 1998 [5] (red triangles) are included. The solid red line is calculated from KAON-MAID [8], the blue dash-dotted line is calculated from the Guidal-Laget Regge model [13], and the magenta dashed line is model B provided by Janssen [10].

Figure 8.5: Comparison of the CLAS $\gamma p \rightarrow K^+ \Sigma^0$ differential cross section (blue circles) with SAPHIR [5] data (red triangles). The solid red line is calculated from KAON-MAID [8], the blue dash-dotted line is calculated from the Guidal-Laget Regge model [13], and the magenta dashed line is model B provided by Janssen [12].

Figure 8.6: Comparison of the CLAS $\gamma p \rightarrow K^+ \Sigma^0$ differential cross section (blue circles) with SAPHIR [5] data (red triangles). The solid red line is calculated from KAON-MAID [8], the blue dash-dotted line is calculated from the Guidal-Laget Regge model [13], and the magenta dashed line is model B provided by Janssen [12].

Figure 8.7: Comparison of the CLAS $\gamma p \rightarrow K^+ \Sigma^0$ differential cross section (blue circles) with SAPHIR [5] data (red triangles). The solid red line is calculated from KAON-MAID [8], the blue dash-dotted line is calculated from the Guidal-Laget Regge model [13], and the magenta dashed line is model B provided by Janssen [12].

Figure 8.8: Comparison of the CLAS differential cross section (blue points) for a few representative energy bins chosen so that all the SAPHIR data from 1998 [5] (red triangles) is included. The solid red line is calculated from KAON-MAID [8], the blue dash-dotted line is calculated from the Guidal-Laget Regge model [13], and the magenta dashed line is model B provided by Janssen [12].

8.2 The Recoil Polarization

The CLAS measurements of the recoil polarization for $\gamma p \rightarrow K^+ Y$ are compared with the 12 previously published data points from SAPHIR and the theoretical models from the previous section in Figures 8.9- 8.12. There are several notable features in these polarization results. First, the results are consistent with the sparse (red) data from SAPHIR. The greater statistics in CLAS, however, allows one to observe interesting trends that are not as visible in the SAPHIR data.

At higher energies and forward directions, the Regge model, despite the previously noted discrepancies with the cross section measurements, does quite well for both the Λ and Σ^0 . This would be consistent with the polarization in the forward/high-energy region being caused by the interferences in the t -channel. The Regge model, however, is not able to reproduce the backward angle polarization measurements, in particular the change of sign in the $K\Lambda$ reaction and the rapid diminishing of the polarization in the $K\Sigma^0$ reaction. This implies that the polarization must be caused by the interference of non- t -channel portions of the production mechanism.

The KAON-MAID model does a reasonably good job with the $\gamma p \rightarrow K^+ \Lambda$ channel, although at photon energies between 1.4 GeV and 1.85 GeV (see Figure 8.10) the model significantly underestimates the polarization at backward angles. In Figure 8.9 this manifests itself as a bump in the two bins between -0.5 and -0.3 in $\cos(\Theta_{cm}^K)$. This should be compared with the polarization predicted by Janssen *et al.* [10] as seen in Figure 2.10. The Janssen model, which incorporates hyperon resonances in the u -channel, predicted the enhanced polarization seen in this region, however this enhancement is not due primarily to the u -channel resonances, but rather to the change in the hadronic form factors.

In contrast to the $\gamma p \rightarrow K^+ \Lambda$ channel, all the theoretical predictions for the $\gamma p \rightarrow K^+ \Sigma^0$ channel polarizations only agree at the forward most energy bins, where the Regge model does equally well. Since the Sigma is an isospin 1 particle which permits delta resonances to contribute, there are a large number of resonances that can contribute in this region. The limited number of resonances actually used in these models for Sigma photoproduction leaves wide open the possibility that by adding more resonances they will be able to reproduce the measured polarization in the backward angle by including a few more known resonances.

Unfortunately, it will always be the case, due to the nature of the experiments necessary to measure these observables, that the data from the cross section measurements will have much better statistical power than the data from the polarization results. Theoretical models that fit their free parameters to both the cross section and polarization will therefore favor the fitting of the cross section. This could be avoided by artificially increasing the weight of the polarization data in the fits. Although, ultimately, it would be best not to have to add such an ad hoc element to the fitting procedure, doing so could shed light upon what are the important resonances needed to fit the polarization results and thus enable the inclusion of these in the standard

fitting procedure to get a result that matches well both the differential cross section and the polarization.

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Figure 8.9: Recoil polarization in $\gamma p \rightarrow K^+ \Lambda$ as a function of W . The red curve represents the predictions from KAON-MAID [38], a dominately s-channel model, and the blue dot-dashed curve are the results of a t-channel dominated Regge-model [13]. The magenta dashed curve is the Janssen model B [10]. The red data points are the 12 data points from SAPHIR [5]. The blue data points are the CLAS measurements.

Figure 8.10: Recoil polarization in $\gamma p \rightarrow K^+ \Lambda$ as a function of $\cos(\Theta_{cm}^K)$. The red curve represents the predictions from KAON-MAID [38], a dominately s-channel model, and the blue dot-dashed curve are the results of a t-channel dominated Regge-model. [13]. The magenta dashed curve is the Janssen model B [10]. The red data points are the 12 data points from SAPHIR [5]. The blue data points are the CLAS measurements.

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Figure 8.11: Recoil polarization in $\gamma p \rightarrow K^+ \Sigma^0$ as a function of W . The red curve represents the predictions from KAON-MAID [38], a dominately s-channel model, and the blue dot-dashed curve are the results of a t-channel dominated Regge-model [13]. The magenta dashed curve is the Janssen model B [12]. The red data points are the 12 data points from SAPHIR [5]. The blue data points are the CLAS measurements.

Recoil Polarization for $\gamma + p \rightarrow K^+ + \Sigma^0$

Figure 8.12: Recoil polarization in $\gamma p \rightarrow K^+ \Sigma^0$ as a function of $\cos(\theta_K^{c.m.})$. The red curve represents the predictions from KAON-MAID, a dominantly s-channel model, and the blue dot-dashed curve are the results of a t-channel dominated Regge-model. The magenta dashed curve is the Janssen model B [12]. The red data points are the 12 data points from SAPHIR [5]. The blue data points are the CLAS measurements.

8.3 A Legendre Moment Analysis

In this section I describe a fit of the differential cross section with Legendre polynomials to extract a total cross section. This total cross section is not an independent measurement, and should not be used for extracting parameters of theoretical models.

In an ideal experiment with complete angular coverage, one can safely extract the total cross section by fitting the angular distribution with Legendre polynomials in $\cos(\Theta_{cm}^K)$. As long as one has the complete angular coverage, the Legendre polynomials are orthogonal functions and therefore the fitting of higher order polynomials will not effect the fit results of lower order polynomials. Figure 8.13 shows the measured differential cross section at $E_\gamma = 1.9$ GeV from CLAS, the KAON-MAID prediction, and a fit to:

$$\frac{d\sigma}{d\cos(\Theta_{cm}^K)} = \frac{\sigma_{tot}}{2} \left(1 + \sum_{i=1}^3 C_i P_i(\cos(\Theta_{cm}^K)) \right) \quad (8.1)$$

where σ_{tot} is the total cross section, the P_i are the Legendre polynomials, and the C_i are the energy dependent shape parameters. This figure demonstrates the problem with this simple type of analysis to extract a total cross section when one does not have complete coverage. Although the fit is a good fit across all bins that have data, one can see that in the forward most bin where there is no data to constrain the fit, the ability of the Legendre polynomial to accurately extract the total cross section is hampered if the cross section suddenly turns over. Of course, any method of extracting the total cross section under such circumstances would be model dependent.

Despite this difficulty and our lack of data in the forward-most and backward-most bins, we have fit the data to Legendre polynomials up to order four. The results may be seen in Figures 8.14 to 8.17 where the lowest order term has been converted into a cross section. In addition to the parameters of the fit the χ^2 has been plotted for each fit along with the results of an F-test. Points above the line pass the F-test and are therefore considered a significant improvement in the fit over the fit to one lower order. Alternatively, this can be thought of as the additional fit parameter for that order is significantly not zero when the point in the F-test plot is above the line.

In Figures 8.18 - 8.27 the best fit results for the total cross section and the Legendre moments are presented. In each plot, the best fit result is plotted in dark blue, over the range in which the fit is considered valid. The lowest order fit is also plotted in light blue over the whole range to give some sense as to how sensitive the fit results are. The order of the best fit result is chosen to be the highest order which itself passes, and all lower orders pass, the F-test. It is also required that the order continue to pass the F-test for the neighboring bins at higher energies.

Figure 8.13: A Legendre fit to the differential cross section and the prediction from KAON-MAID [8].

Figure 8.14: Fit results for $\gamma p \rightarrow K^+ \Lambda$: total cross section(top) χ^2 (middle two) and Ftest(bottom) in which points above the line would be considered a significant improvement to fit to that order of the next lower order.

Figure 8.15: Fit results from Legendre polynomials for $\gamma p \rightarrow K^+ \Lambda$

Figure 8.16: Fit results for $\gamma p \rightarrow K^+ \Sigma^0$: total cross section(top) χ^2 (middle two) and Ftest(bottom) in which points above the line would be considered a significant improvement to fit to that order of the next lower order.

Figure 8.17: Fit results from Legendre polynomials for $\gamma p \rightarrow K^+ \Sigma^0$

Figure 8.18: Total cross section from Legendre polynomial fit for $\gamma p \rightarrow K^+ \Lambda$. The “lowest” data are from a fit to 0^{th} order and the “best” data are from the highest order fit that was significantly not zero.

Figure 8.19: Fit result from Legendre polynomials for $\gamma p \rightarrow K^+ \Lambda$. The “lowest” data are from a fit to 1^{st} order and the “best” data are from the highest order fit that was significantly not zero.

Figure 8.20: Fit result from Legendre polynomials for $\gamma p \rightarrow K^+ \Lambda$. The “lowest” data are from a fit to 2nd order and the “best” data are from the highest order fit that was significantly not zero.

Figure 8.21: Fit result from Legendre polynomials for $\gamma p \rightarrow K^+ \Lambda$. The “lowest” data are from a fit to 3^{rd} order and the “best” data are from the highest order fit that was significantly not zero.

Figure 8.22: Fit result from Legendre polynomials for $\gamma p \rightarrow K^+ \Lambda$. The “lowest” data are from a fit to 4th order and the “best” data are from the highest order fit that was significantly not zero.

Figure 8.23: Total cross section from Legendre polynomial fit for $\gamma p \rightarrow K^+ \Sigma^0$. The “lowest” data are from a fit to 0^{th} order and the “best” data are from the highest order fit that was significantly not zero.

Figure 8.24: Fit result from Legendre polynomials for $\gamma p \rightarrow K^+ \Sigma^0$. The “lowest” data are from a fit to 1st order and the “best” data are from the highest order fit that was significantly not zero.

Figure 8.25: Fit result from Legendre polynomials for $\gamma p \rightarrow K^+ \Sigma^0$. The “lowest” data are from a fit to 2^{nd} order and the “best” data are from the highest order fit that was significantly not zero.

Figure 8.26: Fit result from Legendre polynomials for $\gamma p \rightarrow K^+ \Sigma^0$. The “lowest” data are from a fit to 3^{rd} order and the “best” data are from the highest order fit that was significantly not zero.

Figure 8.27: Fit result from Legendre polynomials for $\gamma p \rightarrow K^+ \Sigma^0$. The “lowest” data are from a fit to 4th order and the “best” data are from the highest order fit that was significantly not zero.

8.4 Suggestions for Future Work

Obviously, the more observables for hyperon production that are measured, the more constraints can be placed on the theoretical models. In that regard comparison of the results of this analysis with the results of the soon-to-be released electroproduction counterpart to this experiment, done by Robert Feuerbach, should reveal more about the underlying hyperon production mechanism. Also the work being done by Robert Bradford to extract the circularly polarized beam-recoil double polarization observables from the full glc data set should be quite enlightening. The g8 run in CLAS will enable a similar extraction for a linearly polarized beam. There is also a proposal to do a similar experiment with a polarize target. Together these should enable a “complete” measurement of the hyperon photo production.

To improve the analysis of the $\gamma p \rightarrow K^+ Y$ channel for this data set, one must improve on the understanding of the largest sources of systematic uncertainties. The largest uncertainty ($\sim 7\%$) comes from the acceptance/efficiency calculations, therefore any improvement would require a better understanding of this function. Much of the present uncertainty is the result of the need to knock out time-of-flight paddles to account for single particle trigger and drift chamber inefficiencies. A better understanding of these issues might enable one to model the detector with improved confidence.

The next major source of uncertainty is largely bin dependent. In bins where the pion background begins to seep into the kaon selection cuts, the yield extraction becomes the dominant source of uncertainty, to the point where we have excluded many bins in the forward angles at high energies. Improving the rejection of this pion background could significantly reduce the uncertainty in these regions, but only if the fraction of real kaons rejected by the tighter cuts can be well modeled.

In bins where the pion mis-identification is not a problem, the photon normalization becomes the next biggest source of uncertainty. In this regard, improvements in the $\gamma p \rightarrow \pi^+ n$ cross section measurement would be sure to help. The largest source of uncertainty in that measurement, not counting acceptance, comes from the extraction of the yield. At present we believe the errors we report put upper and lower bounds on this yield, but understanding the shape of the neutron peak and the background underneath it could significantly improve these results. This in turn would reduce the uncertainty in our understanding of the photon normalization.

The polarization measurements fortunately seem fairly insensitive to all of these problems. Improvements in the polarization measurement can therefore be made simply by analyzing the remainder of the glc data set to increase the statistics and extend the measurement to higher photon energies.

Chapter 9

Conclusions

The differential cross section and hyperon recoil polarizations of $\gamma p \rightarrow K^+\Lambda$ and $\gamma p \rightarrow K^+\Sigma^0$ have been measured with the CLAS at TJNAF up to a photon energy in the lab of 2.325 GeV. The results for both channels show significantly larger cross section in the middle to forward angles than have been observed previously by the SAPHIR Collaboration.

The W dependence of the $\gamma p \rightarrow K^+\Lambda$ reaction at backward angles looks similar to what is predicted by the KAON-MAID model which includes a D_{13} resonance at 1.9 GeV. The predictions from this model, however, do not match the polarization data in this energy region and therefore we must conclude that the hyperon photo-production mechanism model of KAON-MAID still needs improvements. It seems likely that the inclusion of hyperon resonances in the u-channel, as done by Janssen *et al.* may be able to replicate this phenomenon.

There is an as yet unexplained feature in the differential cross section at a W of about 1.95 GeV in the forward most bins which may indicate the importance of resonance physics at this higher energy. The resonance-like bump has a full-width, Γ , of approximately 100 MeV.

Unfortunately, there is not a unique explanation for these phenomena in the $\gamma p \rightarrow K^+\Lambda$ data. It is possible that the two bumps are both the results of resonances, but that the 1.9 GeV bump is masked in the forward channel by the rising t-channel processes and the 1.95 GeV bump is masked in the backward angles by the 1.9 GeV bump. It is also possible that these two bumps are actually the result of the interference of a single resonance with some other terms that causes the width and centroid of the bump to change as a function of angle. The fact that the bump at 1.95 GeV in the backward angle is in the same kinematic region as a very strong polarization may indicate that the bump is caused by the interference with u-channel resonances. It is also possible that the large polarization in this region may be accounted for by a change of the hadronic form factors used in calculating the theoretical models.

It seems plausible that the shape of the differential cross section in the $\gamma p \rightarrow K^+\Sigma^0$ channel will be easily fit by existing models with only minor modifications to the input

parameters. Alternatively, it certainly is possible to reproduce the observed differential cross section with a different set of s -channel N^* and Δ^* resonances, and there are many such resonances that reasonably might contribute but that have not been included in the current models. It is certain that the differential cross section alone is not sufficient to identify the relevant resonances. The polarization measurements from this analysis, however, clearly indicate that the present models are not adequate. Indeed, matching theoretical predictions to the polarization results should constrain these models much more tightly than previous experimental data.

The experimental results of this analysis show several new and interesting features in hyperon photoproduction off protons. Although there are several possible theoretical interpretations to these new results, when combined with the results of several ongoing or planned experiments, it should be possible to significantly constrain these models and identify the key components of the hyperon photoproduction mechanism.

Appendix A

Analysis of $\gamma p \rightarrow \pi^+ n$

This appendix describes the work done to measure the differential cross-section for $\gamma p \rightarrow \pi^+ n$ using the CLAS detector with the 2.445 GeV data set of the glc running period. The measurement was compared to the SAID-2001 parameterization of the world's data for this reaction, and an overall scale factor was extracted. The MAID-2000 model for this reaction was also compared with SAID and the present data. This was interpreted as a trigger inefficiency correction to the yields of $\gamma p \rightarrow K^+ Y$.

A.1 Binning

The binning chosen for the $\gamma p \rightarrow \pi^+ n$ reaction channel was identical to that chosen for the $\gamma p \rightarrow K^+ Y$ channels: 25 MeV wide bins in E_γ starting from the lowest bin centered at 500 MeV, and 20 bins in $\cos(\Theta_{cm}^{\pi^+})$. In energy, these bins start in the “2nd resonance” region. In angle, CLAS had only partial coverage in the extreme forward and backward bins, and so these were rejected, leaving 18 bins spanning $-0.9 < \cos(\theta) < +0.9$.

A.2 Identifying Pions

For identifying positive pions, we used the time difference between the measured time of flight, and the expected time of flight if the particles were pions. Figure A.1 shows this time difference as a function of lab momentum. For the pion identification, cuts of $-2 \text{ ns} < \Delta t < 1 \text{ ns}$ were used in order to include all well-reconstructed pions, while still having a good π -p separation.

A.3 Event Based Cuts and Global Corrections

The standard PSEUDO-SPA fiducial cut was first applied to all particles. Particles hitting paddles from the Time-of-Flight Knockout list (Table A.1) were excluded

Figure A.1: Measured ToF minus expected ToF assuming pion mass vs lab momentum. The white lines represent the cuts used in identifying the pions.

since these paddles had inefficiencies related to problems with the Drift Chamber, Calibrations, or the single charged particle trigger. The table of knockouts used for pion cross sections is the same as the table used for the kaon cross sections. Figure A.2 shows the angular distribution in the lab of identified pions. The effects of the Time-of-Flight knockouts are quite clear, especially in sectors 2 and 3. At large angles, only sectors 3 and 6 had well-functioning ToF counters and drift chambers. In addition to the standard fiducial cuts, particles with an energy-loss corrected momentum less than 200 MeV were rejected

A Start-Counter - Tagger cut of ± 1 ns was applied to all events. For the reaction $\gamma p \rightarrow \pi^+ n$ this cut combined with the pion identification procedure was found to eliminate approximately 4% of pions with a significant energy dependence. This was determined by comparing the missing mass off all charged particles with no Start-Counter-Tagger coincidence cut (NoPID) with the missing mass for accepted events (PID) for each energy bin. A sample comparison at one photon energy can be seen in Figure A.3. The difference between the PID and NoPID was taken (DIFF). The DIFF and PID spectrums were then fit to a Gaussian with a quadratic background yielding Y_{diff} and Y_{pid} . The correction was then simply: $correction = 1 + \frac{Y_{diff}}{Y_{pid}}$. This correction was applied to the extracted pion yields before calculating the cross-section. Figure A.4 shows the energy dependence of this correction.

A.4 Fitting the $\gamma p \rightarrow \pi^+ n$ missing mass spectra

After the π^+ particles were identified, the missing mass recoiling from the π^+ was fit in each bin to extract the yield. Sample fits can be seen in Figures A.5-A.10. Several estimates were made. First a single Gaussian was fit with three possible polynomial backgrounds. The Gaussian fit with the lowest normalized χ^2 was selected. Then a simple summing was done in the missing mass region between $\pm 4\sigma$ of the fit peak. The area of the peak from the single Gaussian fit was averaged with the summed counts to give the yield. The error was calculated by adding the variance of the two numbers in quadrature to the average statistical error. This averaging was done since for many bins, the single Gaussian fit did not represent the shape of the neutron peak very well, and it was believed that the fit value underestimated the actual yield by overestimating the background. The sum between $\pm 4\sigma$, on the other hand, had no background subtraction in it, and therefore was considered an overestimate. The double Gaussian fits to the peak were not used since for many bins considerable effort was needed to get good initial parameters such that the fit would converge. A better estimate of the yields can probably be obtained by doing this double Gaussian fit. In most cases where the fit did converge it agreed with the extracted yield, within the estimated error.

The G-flux method (see Section 5.8.1) was used to obtain the photon normalization, with corrections based on the comparison of the normalized all-particle yields from two different endpoint energies, 3.1 GeV and 2.4 GeV, as discussed in sec-

Table A.1: The list of knocked-out Time-of-Flight paddles, and causes for their rejection. “Pions Weak” means that problems were seen in the single pion yields. “DC problems” represent paddles that have pathologies consistent with there having been inefficient regions in the Drift Chambers. The paddles associated with “DC problems” would be different for negatively charged particles.

Sector	ToF id	comment
Sector 1:	24, 29 30...48	pions weak
Sector 2:	12, 13, 14, 15 18 33...48	DC problems DC problems
Sector 3:	17, 18 23 30, 33, 41 42 43 44 45 46, 48	pions weak dead pions weak pions weak pions weak dead
Sector 4:	28, 42 43...48	DC problems
Sector 5:	32 42 44, 45 46, 47 48	pions weak pions weak
Sector 6:	32 40...45, 48	dead

Figure A.2: Angular distribution in lab coordinates of pions after fiducial cuts and ToF knockout.

2002/08/15 11.13

Figure A.3: The effect of the Tagger - ST coincidence and π^+ identification cuts. Missing mass for assumed $\gamma p \rightarrow \pi^+ n$ for all charged particles with a ± 9 ns Start-Counter-Tagger cut (nopid), well identified π^+ with standard ± 1.04 ns Start-Counter - Tagger cut (pid), and the difference between these two cases(diff).

tion 5.8.1.

A.5 Acceptance Calculations for $\gamma p \rightarrow \pi^+ n$

The detection acceptance of CLAS for $\gamma p \rightarrow \pi^+ n$ was calculated using two different models. The first was a GEANT based simulation of the detector called GSIM, and the second was a parameterized acceptance calculation, called MC_Generator, based on a geometric acceptance and small a correction for decaying pions. In both cases the thrown events were generated using version SM01 of the SAID parameterization to calculate the differential cross-section. For any set of kinematics, if a pseudo-random number was less than the SAID cross section, or less than a fixed minimum value (to ensure that some events were generated in all bins) the event was generated and considered to be a “thrown” event.

2002/08/14 15.36

Figure A.4: Correction to the $\gamma p \rightarrow \pi^+ n$ yield due to particle identification and Start-Counter-Tagger timing.

A.5.1 GSIM

The GSIM calculation of the acceptance produced a simulated set of CLAS events that had the timing of the Time-of-Flight paddles and Drift Chambers smeared by the program gpp. The resultant data was “cooked” identically to the real data and then processed in the same fashion as the real data. The acceptance was then calculated by taking the resultant yield in each bin and dividing it by the number of events generated in that bin. The resultant acceptance for $\gamma p \rightarrow \pi^+ n$ can be seen in Figure A.11.

A.5.2 Parametric Monte-Carlo, MC_Generator

The parametric acceptance calculation was based on the approximation that the acceptance could be broken down into two parts, a geometric acceptance and a cor-

rection weight for the decay of the pion. This model is discussed in more detail in Section 6.2. The basic geometric acceptance is based on the PSEUDO-SPA fiducial function. Corrections for energy loss and multiple scattering in the Target and Start-Counter were also applied. Particle trajectories were calculated and “swum” out to the Time-of-Flight paddles using the full CLAS geometry and track generating software. The probability that the particle would decay in flight before hitting the ToF paddles was then calculated. A weight equal to the survival fraction was added to the accepted hits if the ToF paddle was not in the knockout list. If the decay probability was greater than 0.2%, the probability of the decay products hitting either the original paddle or the immediate neighbors was approximated numerically. The method is described in more detail in Section 6.2. The sum of the weight for all events in a particular bin was then divided by the number of thrown events in that bin to calculate the acceptance which can be seen in Figure A.12.

A.5.3 Comparison of the Two Acceptance Calculations

The ratio of the two acceptance models’ results can be seen in Figure A.13. The global ratio for all bins was then fit using a total error that for each bin consisted of that bin’s statistical error added in quadrature to an overall error: $\sigma_{total}^2 = \sigma_i^2 + \sigma_{sys}^2$. This fit can be seen in Figure A.14. The fit result for the ratio between the two acceptance calculations of 97.4% is attributed to GSIM’s more accurate modeling of tracking inefficiency, the Start-Counter - Tagger cut, and the decays of the pions. Thus, the parametric calculation was used for its statistical power, but scaled by this factor.

A.6 $\gamma p \rightarrow \pi^+ n$ Cross-Section and comparison to SAID

The differential cross-section was obtained by measuring the reaction yield(Y) as a function of photon energy and pion production angle and then correcting for the acceptance of the detector(A), the number of photons(N_γ) that passed through the target, and the density(ρ) and length(L) of the target:

$$\frac{d\sigma(E_\gamma, \Theta_{cm}^\pi)}{d\cos(\Theta_{cm}^\pi)} = \frac{Y(E_\gamma, \Theta_{cm}^\pi)}{\delta\cos(\Theta_{cm}^\pi)A(E_\gamma, \Theta_{cm}^\pi)N_\gamma(E_\gamma)} \frac{H_{gpm}}{N_A\rho L} \quad (\text{A.1})$$

where H_{gpm} is the number of grams per mole of the target, N_A is Avogadro’s number, and $\delta\cos(\Theta_{cm}^\pi)$ is the size of the $\cos(\Theta_{cm}^\pi)$ bin. A table of these values is in in Appendix D. The value we measured can be seen as a function of E_γ in Figures A.7 - A.7 along with the comparison to the SAID and MAID parameterizations of the results of previous experiments. The top panel shows the cross section on a log scale as a function of E_γ . The bottom panel shows the ratio of CLAS results divided by the

SAID parameterization. The green curve in the bottom panel is the MAID/SAID ratio. We consider the wiggles in this curve to reflect the size of wiggles that might be expected to be caused by uncertainties in the parameterization. Error bars include both the statistical error and the estimated systematic uncertainty from the yield extraction. The fit result to an overall scale factor for each $\cos(\Theta_{cm}^\pi)$ histogram is plotted in Figure A.15 from which we extract an overall scale factor of $\frac{CLAS}{SAID} = 0.9198$, or alternatively $\frac{SAID}{CLAS} = 1.087$. This discrepancy is attributed to an overall trigger inefficiency. The first and last $\cos(\Theta_{cm}^\pi)$ bins were excluded since the CLAS coverage in these bins is far from complete. An estimate of the systematic error in this measurement is included in Table A.2

A.7 Summary

The reaction $\gamma p \rightarrow \pi^+ n$ has been measured with the CLAS across a large number of energy bins and angles. The results are overall lower than the SAID parameterization of previous measurements by a scale factor of 92%. Despite this discrepancy the continuous coverage across such a large phase space should contribute greatly to the understanding of this reaction.

Source Effect	Correction	Where Applied	Uncertainty
Yield			
Pion Fit			in bin-by-bin
Start-Counter-Tagger cuts	1.0-1.04	E_γ bin by E_γ bin	2.0% point to point
Normalization			
photon counting(G-flux)		E_γ bin-by- E_γ bin	3.0% global
Photon Attenuation	1.039	globally	0.2% global
Pt-to-Pt Trigger (In)Efficiency Comparison of 2.4GeV and 3.1GeV endpoint Data	1.00-1.62	E_γ bin by E_γ bin	1.0% point to point
Target			
Density= $0.0711(\frac{g}{cm^3})$ [22]		globally	0.14% global
Length=178.5mm		globally	0.28% global
Acceptance			
comparison of GSIM to Parametric	0.974	global scale	6.38% point to point
Overall Systematic (excluding pion fit effect on yield):			7.40% global and point to point

Table A.2: Sources of systematic uncertainties and their estimated values.

Figure A.5: Fits to the π^+ missing mass squared. The top plot uses a single Gaussian to fit the peak with a polynomial background (top green=constant, middle magenta=linear, bottom blue=quadratic.) The fit values on the left represent the area of the Gaussian. The sum values on the right are the total number of counts between the lines, defined by $\pm 4\sigma$ from the fit with $\sqrt{(N)}$ as the error. The average in the lower left is the average of the two methods where the variance of the two results was added in quadrature to the average error. The lower plot is a two Gaussian fit to the same data.

Figure A.6: Fits to the π^+ missing mass squared. The top plot uses a single Gaussian to fit the peak with a polynomial background (top green=constant, middle magenta=linear, bottom blue=quadratic.) The fit values on the left represent the area of the Gaussian. The sum values on the right are the total number of counts between the lines, defined by $\pm 4\sigma$ from the fit with $\sqrt{(N)}$ as the error. The average in the lower left is the average of the two methods where the variance of the two results was added in quadrature to the average error. The lower plot is a two Gaussian fit to the same data.

Figure A.7: Fits to the π^+ missing mass squared. The top plot uses a single Gaussian to fit the peak with a polynomial background (top green=constant, middle magenta=linear, bottom blue=quadratic.) The fit values on the left represent the area of the Gaussian. The sum values on the right are the total number of counts between the lines, defined by $\pm 4\sigma$ from the fit with $\sqrt{(N)}$ as the error. The average in the lower left is the average of the two methods where the variance of the two results was added in quadrature to the average error. The lower plot is a two Gaussian fit to the same data.

Figure A.8: Fits to the π^+ missing mass squared. The top plot uses a single Gaussian to fit the peak with a polynomial background (top green=constant, middle magenta=linear, bottom blue=quadratic.) The fit values on the left represent the area of the Gaussian. The sum values on the right are the total number of counts between the lines, defined by $\pm 4\sigma$ from the fit with $\sqrt{(N)}$ as the error. The average in the lower left is the average of the two methods where the variance of the two results was added in quadrature to the average error. The lower plot is a two Gaussian fit to the same data.

Figure A.9: Fits to the π^+ missing mass squared. The top plot uses a single Gaussian to fit the peak with a polynomial background (top green=constant, middle magenta=linear, bottom blue=quadratic.) The fit values on the left represent the area of the Gaussian. The sum values on the right are the total number of counts between the lines, defined by $\pm 4\sigma$ from the fit with $\sqrt{(N)}$ as the error. The average in the lower left is the average of the two methods where the variance of the two results was added in quadrature to the average error. The lower plot is a two Gaussian fit to the same data.

Figure A.10: Fits to the π^+ missing mass squared. The top plot uses a single Gaussian to fit the peak with a polynomial background (top green=constant, middle magenta=linear, bottom blue=quadratic.) The fit values on the left represent the area of the Gaussian. The sum values on the right are the total number of counts between the lines, defined by $\pm 4\sigma$ from the fit with $\sqrt{(N)}$ as the error. The average in the lower left is the average of the two methods where the variance of the two results was added in quadrature to the average error. The lower plot is a two Gaussian fit to the same data.

Figure A.11: The results of 1 Million events from the GSIM calculation of the acceptance for $\gamma p \rightarrow \pi^+ n$ as a function of $\cos(\Theta_{cm}^{\pi^+})$ and E_γ .

Figure A.12: The results of 40 million events from MC_Generator of the acceptance for $\gamma p \rightarrow \pi^+ n$ as a function of $\cos(\Theta_{cm}^{\pi^+})$ and E_γ .

Figure A.13: Ratio of GSIM acceptance to MC_Generator acceptance for $\gamma p \rightarrow \pi^+ n$.

Figure A.14: Ratio of GSIM acceptance to MC_Generator acceptance for $\gamma p \rightarrow \pi^+ n$.

Figure A.15: Ratio of CLAS to SAID as a function of $\cos(\Theta_{cm}^{\pi})$. The values were obtained by fits at each angle to the $\frac{\text{CLAS}}{\text{SAID}}$ ratio across photon energies from 0.6 GeV to 1.6 GeV.

Figure A.16: Each plot contains the differential cross-section on a log scale for $\gamma p \rightarrow \pi^+ n$ (top) and the ratio to the SAID parameterization (bottom). The blue curve is the SAID value, the green curve is the MAID value, and the points are from CLAS.

Figure A.17: Each plot contains the differential cross-section on a log scale for $\gamma p \rightarrow \pi^+ n$ (top) and the ratio to the SAID parameterization (bottom). The blue curve is the SAID value, the green curve is the MAID value, and the points are from CLAS.

Figure A.20: Each plot contains the differential cross-section on a log scale for $\gamma p \rightarrow \pi^+ n$ (top) and the ratio to the SAID parameterization (bottom). The blue curve is the SAID value, the green curve is the MAID value, and the points are from CLAS.

Appendix G

List of Cooked Files Used

Below is the list of cooked files used in this analysis from the pass-1 cooking of the 2.445 GeV data set of the glc running period.

```
run21763_.pass1.A00.00 run21763_.pass1.A00.01
run21763_.pass1.A01.00 run21763_.pass1.A01.01
run21763_.pass1.A02.00 run21763_.pass1.A02.01
run21763_.pass1.A03.00 run21763_.pass1.A03.01
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run21979_.pass1.A05.00 run21979_.pass1.A05.01
run21979_.pass1.A06.00 run21979_.pass1.A06.01
run21979_.pass1.A07.00 run21979_.pass1.A07.01
run21979_.pass1.A08.00 run21979_.pass1.A08.01
run21979_.pass1.A09.00
run21981_.pass1.A00.00 run21981_.pass1.A00.01
run21981_.pass1.A01.00 run21981_.pass1.A01.01
run21981_.pass1.A02.00
run21982_.pass1.A00.00 run21982_.pass1.A00.01
run21982_.pass1.A01.00 run21982_.pass1.A01.01
run21982_.pass1.A02.00 run21982_.pass1.A02.01
run21982_.pass1.A03.00 run21982_.pass1.A03.01
run21982_.pass1.A04.00 run21982_.pass1.A04.01
run21982_.pass1.A05.00 run21982_.pass1.A05.01
run21982_.pass1.A06.00 run21982_.pass1.A06.01
run21982_.pass1.A07.00 run21982_.pass1.A07.01
run21982_.pass1.A08.00 run21982_.pass1.A08.01
run21982_.pass1.A09.00 run21982_.pass1.A09.01
run21982_.pass1.A10.00 run21982_.pass1.A10.01
run21982_.pass1.A11.00 run21982_.pass1.A11.01
run21982_.pass1.A12.00