A measurement of the differential cross section for
the reaction \( \gamma n \rightarrow \pi^- p \) from deuterium

by

Wei Chen (陈 伟)

Department of Physics
Duke University

Date: _______________________

Approved:

Prof. Haiyan Gao, Advisor

Steffen Bass

Albert Chang

Calvin Howell

Henry Weller

Dissertation submitted in partial fulfillment of the requirements for the degree of
Doctor of Philosophy in the Department of Physics
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Abstract

Strong interactions exhibit two distinct behaviors: confinement and asymptotic freedom. Although the quantum chromodynamics (QCD) in terms of quarks and gluons is the fundamental theory of strong interactions, at lower energies, the hadronic processes are usually described by meson-exchange models. With two distinct pictures of strong interaction: quark-gluon and nucleon-meson, it has been very challenging to understand the transition between them. Many exclusive processes such as pion photoproductions have been used to study the transition region. In recent experiment of charged pion photoproduction [1, 2] a broad enhancement has been seen around $\sqrt{s} = 2.1$ GeV in the scaled differential cross section followed by the entry into the scaling region at center-of-mass (c.m.) angle $\theta_{c.m.} = 90^\circ$. With very limited data points, further confirmation of these findings requires fine energy and angular scans of the differential cross section at photon energies around a few GeV. The JLab g10 data are ideal for this purpose.

We present a measurement of the differential cross section using the g10 data for the $\gamma n \rightarrow \pi^- p$ process from the CLAS detector at JLab in Hall B for photon energies between 1.0 and 3.5 GeV and $\theta_{c.m.}$ between 50° and 115°. We confirm a previous indication of a broad enhancement in the scaled differential cross section $s^7 d\sigma/dt$ and a rapid fall-off in the c.m. energy region of about 400 MeV following the enhancement. Our data show an angular dependence of this enhancement for the first time as the suggested scaling region is approached for $\theta_{c.m.}$ from 70° to 105°.
To Iris Chen (陈奕如)
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1

Introduction

Quantum chromodynamics (QCD) has been well accepted as the fundamental theory of strong interactions. Although it has been very successful in the perturbative region where the coupling constant is very small, little is known at the non-perturbative region where QCD is not solvable analytically. This is due to the confinement property of the strong interaction. In principle, the strong interactions are between quarks and gluons. However, at lower energies, the effective degrees of freedom are colorless nucleons and mesons. And dynamical models such as meson-exchange models are better suited to describe the interactions at this energy region. Thus the strong interactions show two distinct pictures: quark-gluon and nucleon-meson. There's a gap between the two regions where little is known either theoretically or experimentally. In the medium energy region such as a few GeV, low energy dynamical models usually fail while perturbative QCD (pQCD) is not applicable yet at such low energy. The constituent counting rule (CCR) for exclusive processes was proposed as a signature of the transition from the nucleon-meson to the quark-gluon degrees of freedom. The CCR predicts that [3, 4] the differential cross sections for exclusive two-body process at fixed center-of-mass (c.m.) angle $\theta_{\text{c.m.}}$ scales as $d\sigma/dt \propto s^{-(n-2)}$. 
with \( n \) the number of fields involved in the process and \( s, t \) the Mandelstam variables. For two body scattering process \( AB \rightarrow CD \), the Mandelstam variables are defined as

\[
\begin{align*}
s &= (P_A + P_B)^2 \\
t &= (P_A - P_C)^2 \\
u &= (P_A - P_D)^2,
\end{align*}
\]

where \( P_X \) is the 4-momentum of \( X \) particle, and \( X \) is \( A, B, C \) and \( D \). The CCR connects the energy dependence of the differential cross section at fixed \( \theta_{\text{c.m.}} \) to the fundamental degrees of freedom involved in both the initial and the final state of the exclusive two-body process (numbers of quarks, leptons and photons). The CCR has been confirmed in a number of exclusive reactions in specific kinematic regimes [5, 6, 7, 8, 9, 10, 11, 12, 13].

Recent developments in both experiments and theories have renewed the interest in the CCR and the transition from the nucleon-meson degrees of freedom to the quark-gluon degrees of freedom. The deuteron photodisintegration process has been studied extensively in various experiments at Thomas Jefferson Lab National Accelerator Facility (JLab). It has been shown that the onset of the scaling in the deuteron photodisintegration process is surprisingly low, corresponding to a transverse momentum of \( P_T = 1.1 \text{ GeV/c} \) [10, 11, 12, 13], where pQCD is not supposed to work at such low energies. Meanwhile, new developments in theory have suggested that the parton orbital angular momentum and helicity-flipping amplitudes seem to play important roles in exclusive two-body processes [14]. Although \( pp \) elastic scattering data have shown global scaling according to CCR at a wide range of angles, the differential cross sections seem to oscillate around the CCR prediction. It has been shown that the helicity-flipping amplitudes and their interference can explain the deviation of the \( pp \) scattering data from the CCR prediction [15]. Pion photo-
production processes such as $\gamma n \rightarrow \pi^- p$, $\gamma p \rightarrow \pi^+ n$ and $\gamma p \rightarrow \pi^0 p$ are fundamental processes that are ideal candidates for the study of strong interaction such as nucleon resonances and the transition from nucleon-meson to quark-gluon degrees of freedom above the resonance region. The differential cross sections of charged pion photoproduction on proton and deuteron [1, 2] have been measured at intermediate energies at JLab. These measurements confirmed the scaling at $\theta_{c.m.} = 90^\circ$ for both channels. Furthermore, just before the onset of scaling behavior, the scaled cross section drops by a factor of 4 in a very narrow c.m. energy range of a few hundred MeV around $\sqrt{s} = 2.5$ GeV. Is this a signature for the transition from the nucleon-meson degrees of freedom to the quark-gluon degrees? The very limited data did not allow for a detailed investigation of either the true nature of the apparent scaling behavior or the observed enhancement, or the drop in the differential cross section.

Many baryon resonances are predicted to be in this energy region by the constituent quark model [16, 17]. Some of them are 4-star resonances, but many others have not been seen experimentally. So the observed enhancement might be one or more of these baryon resonances. However for the pion photoproduction, there are very limited data available to allow for a partial wave analysis to determine the resonance couplings and amplitudes.

The major effort has been devoted to measurement of the differential cross sections of the $\pi^-$ photoproduction from the deuteron with unprecedented statistical precision. With photon energy up to 3.5 GeV and large angular coverage between $45^\circ$ and $145^\circ$, it provides a detailed scan of the energy region of interest. Chapter 2 gives the physics motivation of the experiment. Chapter 3 shows an overview of the experimental apparatus in Hall B at JLab. Chapters 4 - 7 describe in detail the procedure of the analysis of the data. In the end, Chapter 8 presents the results and discussion.
2.1 A Brief History of the Strong Interaction

The study of the strong interaction began with the discovery of the neutron by James Chadwick, the particle that does not carry charge while able to bind with protons to form nuclei. This binding force, which is much stronger than the electromagnetic and gravitational force, is the new force - strong force. Unlike the electromagnetic forces mediated by the massless photons which can extend to infinity, the nuclear force is very short-ranged (∼ 2 × 10^{-15} m). In 1935, Hideki Yukawa predicted that the nuclear force is mediated by massive particles (see Fig. 2.1). He calculated the mass from Heisenberg’s uncertainty principle and predicted that the mass of the meson responsible for the nuclear force is about 100 MeV [18]. ¹ His work didn’t catch much attention at the time because no known particle had these properties. In 1947, Powell et al. [19] discovered strongly interacting particles in cosmic rays with mass very close to the predicted Yukawa particle. The pion meson was discovered. The prediction and discovery of pion meson significantly influenced the particle physics research and

¹ Calculate pion meson mass from the uncertainty principle: \( \Delta X \Delta P \sim \hbar \) where \( \hbar \sim 6.5 \times 10^{-22} \) MeV·s and \( \Delta X \sim 2 \times 10^{-15} \) m. \( M \sim \frac{\Delta P}{c} = \frac{\hbar}{\Delta X c} = \frac{6.5 \times 10^{-22} \times 3 \times 10^8}{2 \times 10^{-15}} \text{MeV}/c^2 \sim 100 \text{MeV}/c^2. \)
advanced the understanding of the nuclear force and the strong interaction.

Figure 2.1: In 1935, Hideki Yukawa predicted that a particle of about 100 MeV mediates the interactions between protons and neutrons. The particle was later discovered experimentally, and it is the $\pi$ meson. The Feynman diagram shows the pion-exchange description of the strong interaction. Yukawa was awarded the Nobel Prize in 1949 (the picture is from Nobelprize.org).

The meson-exchange description of the nuclear force has been very successful in describing interactions at lower energies such as the nucleon-nucleon interaction [20], nuclear electroweak interactions [21], pseudoscalar meson-meson scattering [22], pion-nucleon scattering [23] and pion photoproduction [24].

In the meson-nucleon picture of nuclear interaction, mesons and nucleons are considered fundamental particles. It was until the establishment of the Standard Model and Quantum chromodynamics (QCD) that the meson-nucleon description of the strong interaction was no longer considered fundamental. With progress in experimental physics, more and more particles were discovered and people started thinking that there were more fundamental building blocks behind all the hadrons. In 1964, the quark model was proposed by Gell-Mann [25] and Zweig [26] independently.
In 1973, the three physicists discovered the asymptotic freedom, the intrinsic property of strong interaction and QCD. They shared the 2004 Nobel Prize (pictures are from Nobelprize.org).

Initially, they proposed three flavors of quarks, up ($u$), down ($d$), strange ($s$) and later three additional quarks were proposed, charm ($c$), top ($t$), and bottom ($b$). These six quarks are the minimum set of particles required to build all the known hadrons, and they all have been experimentally confirmed. Although the quark model has been very successful in describing the properties of hadrons, there were some difficulties such as the $\Delta^{++}$ state. With spin $\frac{3}{2}$, $\Delta^{++}$ is required to have three identical $u$ quarks, while according to the Pauli exclusion principle, there must exist an extra quantum number for them to obey the spin-statistics. The introduction of the color charge solved those problems. The theory behind the strong interaction remained unknown yet, and many field theories arose trying to implement the strong interactions. There are some special properties these theories have to address, such as confinement and antiscreening. Quarks always come out as a bound state such as baryons ($qqq$) and mesons ($q\bar{q}$) no matter how hard physicists tried to find free quarks. Although quarks are strongly bound and confined, they seem to be free particles inside the
hadrons suggested by the SLAC deep inelastic electron-proton scattering data [27]. Symanzik suggested that a negative beta function is necessary to explain the SLAC data, however almost all the proposed field theories gave positive beta functions. In 1973, Gross, Politzer and Wilczek (see Fig. 2.2) demonstrated that non-abelian gauge theories (Yang-Mills theories) gave negative beta functions, and asymptotic freedom was discovered [28, 29]. Quantum chromodynamics was born. The strong interactions are mediated by gluons carrying color charges. Similar to the quantum electrodynamics (QED), QCD is also built based on the simple concept of gauge principle. However the fundamental difference of the two gauge theories arise from the fundamental symmetry groups they are based on. As opposed to QED which is based on U(1) symmetry, QCD, based on SU(3), is far more complicated. Due to the non-abelian nature of the SU(3) group, the gluons carry color charges and as a result, the interactions are not only between quarks, but the gluons interact with themselves. This gives rise to the antiscreening as opposed to the screening in QED, and it explained both confinement and asymptotic freedom nicely. QCD has survived various experimental tests. One well-known test is of the coupling constant $\alpha_s$ which characterize the strength of the strong interaction. Asymptotic freedom suggests that quarks at high energy/momentum transfer or small distance behave like free particles while at low energy/momentum transfer or large distance they bind into hadrons or create new quarks out of vacuum.

The coupling constant can be calculated through the renormalization group equation as a function of momentum transfer $Q^2$, and to the first order, it is

$$\alpha_s(Q^2) = \frac{1}{\beta_0 \ln(Q^2/\Lambda^2)}$$  \hspace{1cm} (2.1)

\footnote{Story says that ’t Hooft first announced that the beta function of Yang-Mills theory was negative at Symanzik’s talk.}
where $\beta_0 = \frac{33-2N_f}{12\pi}$ is the first order $-\beta$ calculated by Gross, Politzer and Wilczek ($\beta$ is negative), $\Lambda$ is the energy scale and $N_f$ is the number of active flavors of quarks at this energy scale $Q$. The coupling decreases logarithmically. Although it doesn’t give us the absolute value of the running coupling constant, it is straightforward to predict $\alpha_s$ at any $Q^2$ given one pair of $\alpha_s(Q_0^2)$ and $Q_0^2$ (requires at least a next-to-leading order calculation to get accurate results). As can be seen from Eq. 2.1, $\alpha_s(Q^2)$ vanishes as $Q^2$ approaches infinity. Asymptotic freedom has been tested and demonstrated in various experiments at different energy scales such as $e^+e^-$ annihilation and deep inelastic lepton-nucleon scattering. Fig. 2.3 shows the QCD calculation of $\alpha_s(Q^2)$ with 4-loop approximation compared with experimental measurements over a wide range of energy scales. Experimental results are in excellent agreement with theory. It can be shown that QCD is the only one to describe the data, and any other functional forms of $Q$ such as $1/Q$, $1/Q^2$, or even leading order approximation of QCD fail to describe the data in the full energy range [30]. Since the coupling constant is running, it has been a convention to extrapolate the measured $\alpha_s$ to $Q = M_Z$ ($Z$ boson mass $M_Z = 91.2$ GeV). The extrapolated values of $\alpha_s(M_Z)$ from various processes are shown in Fig. 2.4, and the current world average is $\alpha_s(M_Z) = 0.1176(20)$ [31].

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3 Here we used the convention $c = 1$ so mass and momentum has the same unit.
Figure 2.3: Summary of measurements of $\alpha_s(Q)$ as a function of the respective energy scale $Q$ (used the nature unit $c = 1$). Figure is taken from Ref. [30]. The curves are the QCD predictions with a 4-loop approximation. The current world average is $\alpha_s(M_Z) = 0.1189 \pm 0.0010$ [31].
2.2 Nucleon-Meson or Quark-Gluon?

When the Standard Model was established, nucleons and mesons were no longer considered fundamental particles.\footnote{In fact, proton was “known” not to be a point-like particle much earlier before the birth of the Standard Model due to the knowledge of its anomalous magnetic moment first measured by O. Stern \textit{et al.} in 1933.} Instead, quarks and gluons are the fundamental particles and internal degrees of freedom of the strong interaction. Due to the asymptotic freedom property of QCD, perturbative methods can be used at high energy when $\alpha_s$ is small. Extensive theoretical and experimental research has been done, and QCD has proved to be very successful at very high energies where pQCD
is applicable. However at lower energy, where $\alpha_s$ becomes large and perturbative methods are not applicable, QCD can’t be solved exactly. Although the quark-gluon degrees of freedom are responsible for the strong interactions, in low energy regimes (so called confinement region), the effective degrees of freedom are actually nucleon-meson, since at this energy scale, quarks and gluons are strongly bound together as colorless objects such as nucleons and mesons. The “old” meson-exchange picture brought by Yukawa still gives a good approximation of the nuclear interactions such as nucleon-nucleon, nucleon-meson, meson-meson interactions. The natural question to ask is where is the transition region between these two types of degrees of freedom: meson-nucleon and quark-gluon. Perturbative QCD works at higher energies while meson-exchange models describe the interactions at lower energies. They both have limitations, and there’s a gap between the two methods where neither of them works. Study of the interactions in this medium energy region will be critical to move forward the understanding of the strong interaction and help the development of models in this energy regime.

2.3 Constituent Counting Rule

The Constituent Counting Rule (CCR) was proposed as a signature for the transition from the nucleon-meson to the quark-gluon picture. For an exclusive reaction $AB \to CD$ at high energy and large momentum transfer, CCR predicts that the differential cross sections start to scale as \[ d\sigma / dt (AB \to CD) \sim s^{2-n} f(t/s) \quad (2.2) \]
where $s, t$ are Mandelstam variables and $n$ is the total number of fundamental fields such as photons, quarks and leptons on both side of the reaction. For example, for nucleon-nucleon elastic scattering, $n = 12$ and $d\sigma/dt \sim s^{-10}$ and similarly for
nucleon-meson elastic scattering, \( d\sigma/dt \sim s^{-8} \). The differential cross sections scale at fixed \( t/s \) when \( s \to \infty \), which is \( t/s \to \frac{\cos(\theta_{\text{c.m.}}) - 1}{2} \). Thus the scaling law at fixed \( \theta_{\text{c.m.}} \), as \( s \to \infty \) gives a fundamental connection between the differential cross sections and the fundamental degrees of freedom involved in both the initial and final states of the two-body process. Scaling means that the reaction can be described by a reduced set of physical quantities. One well-known example is the Bjorken scaling which suggests that the structure function of electron-nucleon deep inelastic scattering (DIS) depends only on the Bjorken scaling variable \( x \) at large momentum transfers. Similarly, the CCR predicts that the differential cross sections at fixed c.m angles depend only on \( s \) at large momentum transfers.

2.3.1 Theoretical Background

The CCR was first derived in 1973 from dimensional analysis by Brodsky and Farrar [3], and Matveev et al. [4] independently. Here we show how to derive this scaling law from dimensional analysis. For processes such as the ones shown in Fig. 2.5, each external fermion line contributes a dimension of \( E^{1/2} \) to the invariant amplitude, assuming the spinor normalization as \( u^\dagger u = 2E \). Since each constituent of a hadron carries a fraction of the total hadron momentum, no matter how the integration over the momentum is performed, it can never introduce extra dependence on \( s \), i.e. the fermion external line must contribute \( E^{1/2} \sim s^{1/4} \) to the amplitude. Similarly, each boson propagator contributes a dimension of \( s^{-1} \) while each fermion propagator contributes \( s^{-1/2} \). For an irreducible and simply connected Feynman diagram without loops, if there are \( m \) external fermion lines (\( m \) must be even), there must be \( m/2 - 1 \) boson propagators and \( m/2 - 2 \) fermion propagators which is true topologically no matter how the fermions and bosons are connected. Thus the dimension of the amplitude of \( m \) fermions is

\[
\mathcal{M} \sim (s^{1/4})^m \cdot (s^{-1})^{m/2 - 1} \cdot (s^{-1/2})^{m/2 - 2} = s^{2 - m/2}.
\] (2.3)
This equation works not only for the hadron-hadron scattering, it also works for photon involved reactions such as pion photoproduction. For each external boson line attached to the diagram which increases the degrees of freedom by 1, it creates an extra fermion propagator. With itself dimensionless, each external boson line will contribute $s^{-1/2}$ to the amplitude due to the extra fermion propagator. Thus the amplitude of an interaction with $n$ total number of fermions and bosons is

$$M \sim s^{2-n/2}, \quad (2.4)$$

and the differential cross section follows immediately as

$$\frac{d\sigma}{dt} \sim s^{-2}|M|^2 \sim s^{2-n}. \quad (2.5)$$

In summary, Eq. 2.2 holds at large $s$ where the $n$ is the total number of degrees of freedom of the system, i.e. the total number of external bosons and fermions involved in the reaction.

In their 1975 paper [32], Brodsky and Farrar argued that the crucial component of the dimensional analysis of CCR relies on two assumptions: (a) the hadrons are the composite of constituents carrying a finite fraction of the hadronic momentum, and (b) no extra mass scale appears in the amplitude. They showed that the renormalizable field theories with assumptions such as the absence of infrared effects and the accumulation of logarithms fit the requirements naturally. In 1980, Lepage and Brodsky [33] proved in the framework of pQCD that the scaling law of the exclusive process at large momentum transfer is rigorous up to powers of the running constant $\alpha_s(Q^2)$. They showed that for sufficiently large transverse momentum $P_T$, the cross section is

$$\frac{d\sigma}{dt}(AB \rightarrow CD) \rightarrow \left(\frac{\alpha_s(P_T^2)}{s}\right)^{n-2} \left(\ln \frac{P_T^2}{\Lambda^2}\right)^{-2\sum_i \gamma_i} f(\theta_{c.m.}) \quad (2.6)$$
Figure 2.5: Feynman diagrams for meson-nucleon scattering and pion photoproduction from the nucleon. For $m$ external fermions, $m/2 - 1$ boson propagators are needed to connect them, and thus create $m/2 - 2$ fermion propagators. Each external boson line will create an extra fermion propagator.

where $P_T$ is the transverse momentum, and $\gamma_i$ is a helicity dependent constant. At high momentum transfer, $f(\theta_{\text{c.m.}})$ is just an overall constant while at lower transverse momentum, it varies logarithmically with $P_T^2$ slowly.

2.3.2 Experimental Evidence

Even before the CCR was established, there were other models and speculations of the behavior of the differential cross sections of exclusive hadronic processes at large $s$ and $t$. For example, for $pp$ elastic scattering, Gunion et al. [35] predicted that

$$\frac{d\sigma}{dt} \propto s^{-12}\sin(\theta_{\text{c.m.}})^{-12}.$$  \hspace{1cm} (2.7)
Figure 2.6: $d\sigma/dt$ versus $s$ of the $pp$ elastic scattering from Ref. [34]. $pp$ scattering shows excellent agreement with the CCR prediction at higher $s$ and $t$.

Landshoff and Polkinghorne [34] pointed out that however, if choosing the $pp$ data with the following cuts

$$s \geq 15 \text{ (GeV)}^2, \ |t| \geq 2.5 \text{ (GeV)}^2,$$

(2.8)

the fitted result is

$$\frac{d\sigma}{dt} \propto s^{-9.7\pm0.5},$$

(2.9)

which is in excellent agreement with the CCR prediction of $s^{-10}$ (see Fig. 2.6). Since then, the CCR has been tested by numerous experiments. In 1994, White et al. [6] reported the results from the BNL Alternate Gradient Synchrotron (AGS) experiments of 20 exclusive two-body hadron-hadron scatterings at large $t$. At the incident momentum of 5.9 GeV/c, 19 out of 20 channels studied are consistent with the CCR.
prediction which is surprising considering the relatively low incident momentum. In addition to the exclusive hadron-hadron scattering, the exclusive photoproduction processes at large $t$ have been studied as well. Anderson et al. [5] reported 6 channels of meson photoproduction cross section, and the data show that the cross sections approximately follow the CCR prediction.

2.4 Recent Development

2.4.1 Experiments at JLab

As the frontier of medium energy experimental nuclear physics research, JLab has recently conducted several exclusive measurements related to the CCR. Deuteron photodisintegration has been extensively studied theoretically. Besides the pQCD, various other models attempt to understand the behavior of this channel at medium energy. These models include the reduced nuclear amplitude model (RNA), the hard quark rescattering mechanism model (HRM), the quark-gluon string model (QGS) and the asymptotic meson exchange current model (AMEC). The cross sections have been measured at SLAC [7, 8, 9] up to $E_\gamma = 2.8 \text{ GeV}$ and the agreement with CCR was found at $\theta_{\text{c.m.}} = 90^\circ$. With much higher beam energy, JLab has carried out several experiments to study deuteron photodisintegration further. It was first measured with an electron beam up to 4.0 GeV in Hall C at JLab [10], and the differential cross sections at $\theta_{\text{c.m.}} = 89^\circ$ and $69^\circ$ show good agreement with CCR. Later on, it was studied with electron beam energies of 5.0 and 5.5 GeV [11], and with higher energies, it was found that $\theta_{\text{c.m.}} = 70^\circ, 53^\circ$ and $37^\circ$ all exhibit CCR above certain $E_\gamma$. It also suggested that the onset of the scaling is quite low corresponding to a transverse momentum of $P_T = 1.3 \pm 0.1 \text{ GeV/c}$. It has also been measured in Hall B at JLab [12, 13] with the CLAS detector and a tagged photon beam up to 3 GeV. With large angular coverage, it reported differential cross sections covering $\theta_{\text{c.m.}} = 35^\circ$ to $145^\circ$ and concluded that the CCR scaling is reached for
proton transverse momentum above 1.1 GeV/c which is consistent with the previous findings. This indicates that the quark-gluon regime may have been reached above this momentum.

Pion photoproductions also have been studied in JLab in several experiments. Experiment E94-104 \[^1\] was carried out in 2001 in Hall A on single charged pion photoproduction on liquid hydrogen (LH2) and liquid deuterium (LD2) targets. With the two High Resolution Spectrometers (HRS), and the untagged bremsstrahlung photon beam, both the cross sections of $\pi^+$ and $\pi^-$ channels were measured with
Figure 2.8: The scaled differential cross section $s^7 \frac{d\sigma}{dt}$ versus center-of-mass energy for the $\gamma n \rightarrow \pi^- p$ at $\theta_{c.m.} = 90^\circ, 70^\circ, 50^\circ$.

high accuracy up to an electron beam energy of 5.6 GeV at various fixed c.m. angles. As shown in Fig. 2.7 and Fig. 2.8, the results [2] show good agreement with the world data. The data show scaling behavior predicted by the CCR at $\theta_{c.m.} = 70^\circ and 90^\circ$ in both channels. The data at $\theta_{c.m.} = 50^\circ$ do not display scaling behavior. For the similar transverse momentum $P_T$ value to the $90^\circ$, it may require higher energies for the forward angles to reach the onset of the scaling behavior. In addition to the observed scaling behavior at certain angles, there is another very interesting feature from the data that shows an apparent enhancement in the scaled differential cross
section at $\theta_{\text{c.m.}} = 90^\circ$ below the scaling region, with $\sqrt{s}$ ranging approximately from 1.8 GeV to 2.5 GeV in both channels, as shown on the top panel of Fig. 2.7 and 2.8. This effect has been observed in existing neutral pion photoproduction [36] data. Without any conclusive statements at present, some speculations may be made. The observed enhancement around 2.2 GeV might relate to some baryon resonances, as some of the well known baryon resonances ($\Delta$, $N^*$’s around 1.5 GeV and 1.7 GeV) are clearly seen in the scaled cross section below 2.2 GeV. Around 2.2 GeV, there are three 4-star $N^*$ resonances ($G_{17}(2190)$, $H_{19}(2220)$ and $G_{19}(2250)$) and one 2-star resonance ($D_{15}(2200)$) [31]. In addition, there are many more baryon resonances predicted to be in this energy region while having not been seen experimentally, i.e. the so called “missing resonances” [16, 17]. All these resonances could contribute to this broad enhancement. Brodsky and Teramond [39] suggested that the resonance structure observed in $pp$ elastic scattering is associated with the threshold production of hadrons with $s$ and $c$ quarks. Similarly, in our case, the enhancement may be associated with the threshold production of strange hadrons and $s$ quark degrees of freedom. Therefore, more detailed studies in this photon energy region are important.

### 2.4.2 Theoretical Development

The pQCD models which predicted the CCR also suggest helicity conservation for the exclusive processes. Some processes such as $pp$ elastic scattering have shown the deviation and oscillation around the CCR predictions. This may be a result of the helicity-flipping amplitude and orbital angular momentum of the constituent quarks. The nucleon spin has long been considered as arising from the valence quark spins while all the quarks are in the $s$-orbit. The EMC data [40], first put an end to the simple description of proton spin in 1989. In this double polarization muon and proton scattering measurement, the spin dependent structure function for the proton and its integral were found much smaller than expected. This implies that the quark
spin only constitute a very small portion of the nucleon spin. The orbital angular
momenta of valence quarks and spin of gluons may contribute significantly to the
nucleon spin. Later on many follow-up experiments [41] confirmed the conclusion. Ji
et al. derived six wave-function amplitudes needed to describe the three-quark sys-
tem of the nucleon wave function using a systematic way to enumerate independent
amplitudes of a light-cone wave function [14]. A general term in the hadron wave
function appears as
\[
\int \prod_{i=1}^{n} d[i] \ (k_{1,\perp}^{+})^{l_{z1}}(k_{2,\perp}^{+})^{l_{z2}}\ldots(k_{(n-1),\perp}^{+})^{l_{z(n-1)}} | \psi_n(x_i, k_{\perp i}, \lambda_i, l_{zi}) a_1^{\dagger}a_2^{\dagger}\ldots a_n^{\dagger}|0 \rangle , \tag{2.10}
\]
where \(k_{i,\perp}^{\pm} = k_{ix} \pm k_{iy}\) and the + (−) sign applies when \(l_{zi}\) is positive (negative), and
\(d[i] = dx_i d^2 k_{\perp i}/(\sqrt{2x_i}(2\pi)^3)\) with the overall constraint on \(x_i\) and \(k_{\perp i}\) implicit. In
the light-cone frame, the longitudinal mass dimension is ignored, and only transverse
dimensions are left. It shows that the leading behavior of the wave function amplitude
\(\psi_n(x_i, k_{\perp i}, l_{zi})\) with \(n\) partons and orbital angular momentum projection \(l_z\) goes as
\[
\psi_n(x_i, k_{\perp i}, l_{zi}) \sim \frac{1}{(k_{\perp i}^{2})^{n+[l_z]+\min(n',+l_{z}')}]}^{2-1} . \quad (2.11)
\]
Based on the wave functions, they derived a generalized counting rule. They
concluded that for a process \(A + B \rightarrow C + D + \cdots\), the fixed-angle differential cross
section scales as [42]
\[
\Delta \sigma \sim s^{-1-\Sigma n_H(l_{zH}+l_{zH})-1} , \tag{2.12}
\]
where \(\Delta \sigma\) is the exclusive cross section integrated over any \(\theta_{c.m.}\), \(l_{zH}\) is the \(z\)-component of the orbital angular momentum, and \(n_H\) the number of fields summing
over all the particles in both initial and final state. This generalized counting rule
suggests that the orbital angular momentum of the quarks play roles in the ampli-
tude. If we assume that the orbital angular momentum \(l_{zH} = 0\), this scaling rule
reduced to the original CCR \cite{3}

\[
\Delta \sigma \sim s^{-1-N_M-2N_B} = s^{-1-\sum_{n_H}^{(n_H-1)}} = s^{-1-N_M-2N_B},
\]

where \(N_M, N_B\) are the number of mesons and baryons, respectively. For \(pp\) elastic scattering, both CCR and the generalized counting rule predict that the helicity-conserving amplitudes \(M(\++ \to \++), M(\+- \to \+-)\) and \(M(\-- \to ++)\) scale as \(1/s^4\), while the generalized counting rule suggests that helicity-flipping amplitudes \(M(\++ \to \+-) \sim 1/s^{4.5}\) and \(M(\--+ \to +++) \sim 1/s^5\). Thus, the helicity-flipping amplitudes and their interference are probably responsible for the deviation from the original CCR prediction. Dutta and Gao \cite{15}, studied the energy dependence of the \(pp\) elastic scattering as well as the pion photoproduction data using the generalized counting rule. As shown in Fig. 2.9 and Fig. 2.10, they found that by including the helicity-flipping amplitudes and their interference, they can describe better the \(pp\) and pion photoproduction oscillatory behavior, and concluded that the helicity-flipping amplitudes (related to the quark orbital angular momentum contribution) play an important role at GeV energies.

Another explanation of the deviation from CCR was proposed by Zhao and Close \cite{43} as the restricted locality of the quark-hadron duality. The quark-hadron duality was first discovered by Bloom and Gilman \cite{44, 45} about 40 years ago. As shown in Fig. 2.11, they found that the structure function \(\nu W_2\) of \(ep\) scattering shows resonance structures at lower \(Q^2\), and as \(Q^2\) increases, it converges to a smooth scaling-limit curve. The low energy resonance behaviors appear to be strongly correlated to the high energy scaling-limit. Zhao and Close explain that the sum over the resonances is correlated with the scaling-limit as a result of destructive interference between high density of overlapping resonances at high energies. At lower energy, due to fewer resonances, the destructive interference breaks down when the partial waves do not cancel locally. This “restricted locality” gives rise to the oscillation
Figure 2.9: The fit to the \( pp \) elastic scattering data without helicity-flipping amplitude. The fit failed to describe the data at lower energies. Figure is taken from Ref. [15].

about the smooth curve. They further showed that the restricted locality of duality can produce oscillations and significant deviation from the CCR.

For example, for \( \pi^+ \) photoproduction, by adopting effective Lagrangians which treats mesons as elementary particles, the transition amplitudes can be derived for
Figure 2.10: The $pp$ elastic scattering data were fitted with helicity-flipping amplitudes, and it shows much better agreement with data. Figure is taken from Ref. [15].

the $s$- and $u$-channels, i.e. the direct and the virtual resonance excitations as

$$M_{fi}^{s+u} = e^{-(k^2+q^2)/6\alpha^2}$$

$$\times \left\{ \sum_{n=0}^{\infty} (\mathcal{O}_{cc}^{d} + (-\frac{1}{2})^n \mathcal{O}_{cc}^{e}) \frac{1}{n!} \left(\frac{k \cdot q}{3\alpha^2}\right)^n \right\} + \sum_{n=1}^{\infty} (\mathcal{O}_{ci}^{d} + (-\frac{1}{2})^n \mathcal{O}_{ci}^{e}) \frac{1}{(n-1)!} \left(\frac{k \cdot q}{3\alpha^2}\right)^{n-1}$$

$$+ \sum_{n=2}^{\infty} (\mathcal{O}_{ii}^{d} + (-\frac{1}{2})^n \mathcal{O}_{ii}^{e}) \frac{1}{(n-2)!} \left(\frac{k \cdot q}{3\alpha^2}\right)^{n-2} \right\}, \quad (2.14)$$

where the multiplets are degenerate in $n$. The spin structures, charge and isospin operators have been subsumed in the symbol $\mathcal{O}$. For $n = 3$, the $L$-dependent multiplets
Figure 2.11: The structure function $\nu W_2$ of $ep$ scattering shows resonance structures at lower $Q^2$ while converges to a smooth curve at higher $Q^2$. The low energy resonance behaviors seem to be strongly correlated to the high energy limit. This is the so called “quark-hadron duality”. Figure is from Ref. [45].

vanish at $\theta_{c.m.} = 90^\circ$, while for $n = 4$, partial waves including $P$, $F$ and $H$ remains and give nonvanishing terms in the cross section. As can be seen in Fig. 2.12, the nondegeneracy cases for $n \leq 2$ including known resonances give a flat scaling of cross section, while $n \leq 4$ gives rise to the sizable oscillations through the breaking of the restricted locality.
Figure 2.12: The differential cross section data compared to the theories for $\pi^+$ photoproduction at $\theta_{\text{c.m.}} = 90^\circ$. The solid curve denotes degeneracy breaking for $n \leq 2$, and the dotted for $n \leq 4$. 
2.5 CAA g10 Analysis

As discussed in this chapter, the study of the GeV region of the $\pi^-$ photoproduction is very important. The CLAS g10 run period with the goal of a high statistics search for the pentaquark $\Theta^+$ state is ideal for the study of the differential cross section $s^7 \frac{d\sigma}{dt}$ as a function of $\sqrt{s}$ as well as a detailed study of the angular dependence of the scaling behavior for the $\pi^-$ photoproduction.

The data were taken on a 24-cm liquid deuterium target, with two settings of the torus magnet current (I=3375 A and I=2250 A) and with two charged particle triggers. The incident electron beam energy was 3.767 GeV, which corresponds to a maximum $\sqrt{s}$ of 2.8 GeV for $\gamma n \rightarrow \pi^- p$ process. Our CAA proposal [46] presented to the CLAS collaboration to analyze these data to study the $\gamma n \rightarrow \pi^- p$ process was approved.

For the process of interest, the $\gamma n \rightarrow \pi^- p$ quasifree process from the deuteron, only events with two charged particles were analyzed. Using two-body kinematics and the photon energy from the CLAS photon tagger, the spectator proton can be reconstructed by detecting the $\pi^-$ and the proton in coincidence. The approved analysis allows a detailed mapping of the dramatic transition region suggested by the Hall A E94-104 data [1], and also helps to elucidate the discrepancy between the Bonn data [47] and the Hall A data [1].
This experiment was carried out in Hall B at Thomas Jefferson Lab National Accelerator Facility (JLab) using the electron beam from the Continuous Electron Beam Accelerator Facility (CEBAF). The electron beam was incident on a gold radiator and bremsstrahlung photons were produced and tagged by the Hall B photon tagging spectrometer. The photon beam was incident on a liquid deuterium (LD2) target, and the produced charged particles were detected by the CEBAF Large Acceptance Spectrometer (CLAS).

A brief overview of the CEBAF and Hall B experimental apparatus is given in this chapter.

3.1 The Continuous Electron Beam Accelerator Facility

The CEBAF \cite{48} at JLab as shown in Fig. 3.1 and 3.2, produces continuous electron beam up to 6 GeV and is planned to be upgraded to 12 GeV. The unique racetrack-shaped accelerator design of CEBAF with two antiparallel linear accelerators (linacs) connected by nine recirculation beam arcs allows the linacs to be used several times during the acceleration of electrons. The electron injected with few eV energy will
Figure 3.1: An aerial view of CEBAF. The linacs are shown in dashed lines, and the three existing halls in circles.

be accelerated five times on each linac and reach up to 6 GeV. Each linac consists of 20 cryomodules, and there are 8 superconducting radio frequency (s.r.f) cavities in each cryomodule operating at very low temperature of 2 K. With s.r.f technology, CEBAF requires an average of 20 megawatts of power to operate. CEBAF can simultaneously delivers electron beam to three experimental halls, Hall A, B and C. A Hall D is under construction now. All three halls can run experiments independently with different beam currents. The core instruments of Hall A are a pair of High
Resolution Spectrometers (HRS) with momentum resolution of $2 \times 10^{-4}$ and angular resolution of 2 mrad at the momentum of about 4 GeV/$c$. Hall C houses the High Momentum Spectrometer (HMS) and Short Orbit Spectrometer (SOS). Both Hall A and C can handle high beam current (typically 100 $\mu$A) while Hall B with its CEBAF Large Acceptance Spectrometer (CLAS) usually operates at the current of 1 - 50 nA.

3.2 Hall B and CEBAF Large Acceptance Spectrometer

Hall B is the smallest one among the three experimental halls. Unlike the other two halls with spectrometers of high resolution and small acceptance, Hall B features the CEBAF Large Acceptance Spectrometer (CLAS). Due to its near $4\pi$ acceptance, it can only operate at a relatively low current ($\sim 1/1000$ of the current delivered to the two other halls). With large acceptance for both charged and neutral particles, it provides detection of multi-particle final states and is ideal for exclusive measure-
ments. As shown in Fig. 3.3, in addition to the CLAS, Hall B has one of the world’s largest photon tagging systems.

When the electron beam is incident on a radiator, under the influence of the electromagnetic field of the nuclei in the radiator, the electron emits one high energy bremsstrahlung photon, and loses energy. The Hall B photo tagging spectrometer [49] shown in Fig. 3.4 was used to measure the energy and time of the bremsstrahlung photon by tagging degraded electrons from the radiator.

The degraded electrons were bent by a uniform dipole magnetic field, and then detected by a scintillator hodoscope. The scintillator hodoscope consists of two planes of scintillators: the energy plane and the timing plane. There are 384 scintillators (E-counters) in energy plane and 61 scintillators (T-counters) in the timing plane. Each scintillator is overlapping with neighboring counters by one-third, and thus provides 767 energy bins, and 121 effective timing paddles. The photon tagger tags photons from 20% to 95% of electron beam energy \( E_e \), and with 767 energy channels, it provides 0.1% \( E_e \) resolutions.

The CLAS detector [50] with nearly \( 4\pi \) coverage, can detect charged particles with \( p > 0.2 \) GeV/c. As shown in Fig. 3.5, it is comprised of 6 azimuthal symmetric sectors. In each sector, there are three drift chamber regions [51], with two superlayers in each region. One of the superlayers is tilted \( 6^\circ \) to the other to provide azimuthal information. Region 2 was in a toroidal magnetic field generated by six superconducting torus coils. The magnetic field is always perpendicular to the momentum of the outgoing particles to maximize the curvature of the track. Outside of the drift chambers are the 48 time-of-flight (TOF) [52] paddles in each sector. When a charged particle with \( p > 0.2 \) GeV/c comes out of the target cell, the start counter which is just outside of the target will provide the timing of this event, and the 3 drift chamber regions will provide the tracking information. With the magnetic field in Region 2, the momentum, and charge of the charged particle
Figure 3.3: The Schematic layout of the Hall B and the CLAS detector in Hall B and other equipment.
can be obtained. However momentum alone is not enough to identify the particle, with TOF and track length, the velocity can be determined, and then mass can be calculated from momentum and velocity.
Figure 3.5: Top: A schematic view of the CLAS detector (from top and cut along the beam line); Bottom: Schematic view of the CLAS detector (cut perpendicular to the beam line).
The CLAS detector due to its large coverage and complicated detector subsystems usually generates large amount of data during experiments. The raw data collected were first calibrated and “cooked” (data reconstruction) which was computationally very time-consuming. It usually takes several months for one pass of calibration and cooking. This work is based on the well calibrated and cooked data. The data are about 40 terabytes in size and this chapter gives a detailed description of the procedure to select events of interest from the data.

In this analysis, we extracted the differential cross sections of the channel of interest. The differential cross section was given by

$$\frac{d\sigma}{d\Omega_{c.m.}} = N \frac{1}{t_G \epsilon} \frac{A}{N_\gamma \rho L N_A} \frac{1}{d\Omega_{c.m.}},$$

(4.1)

where $N$ is the total number of events detected, $\epsilon$ is the acceptance (see Chapter 5, $N$ is the number of events, $N_\gamma$ is the integrated photon flux, $\rho$ is the density of liquid deuterium (0.163 g/cm$^3$), $L$ is the length of the target (24 cm), $N_A$ is the Avogadro number ($6.022 \times 10^{23}$ mol$^{-1}$) and $d\Omega_{c.m.}$ is the angular bin size in the c.m frame of $\pi^- p$. $t_G$ is the transparency to account for the FSI effect (see Chapter 6, and the
transparecy \( (t_G) \) means the ratio of cross sections for exclusive processes from nuclei to those from nucleons. In this chapter, we describe how to extract the number of events.

4.1 Running conditions

During the g10 run period, the following running conditions were used:

- **Energy** of the incident electron beam (incident upon a gold radiator of \( 10^{-4} \) radiation length): \( E_0=3.767 \) GeV;

- **Beam Current**: 30 nA (production runs), 0.1 nA (normalization runs);

- Torus magnetic field: 2250 A (low field setting) and 3375 A (high field setting). Approximately equal beam time for each setting;

- **Target**: 24 cm long and 4 cm in diameter filled with liquid deuterium positioned at 25 cm upstream relative to the nominal CLAS center;

- **Tagged Photon Energy Range**: 0.75 - 3.58 GeV;

- **Trigger**: at least two charged particles in different sectors.

There were three different trigger conditions for low field setting: noEC (EC information was not used in the trigger, 20%), withEC (EC information was in the trigger, 60%), withEC_noASYNC (EC was used, but tagger was not in the trigger, 20%), with only one trigger condition for the high field data: noEC. In this analysis, all the runs with trigger condition noEC are analyzed, i.e. all the high field data, and 20% of the low field data. Dr. Mibe did a comparison of data with and without EC in trigger, and they agreed with each other within statistical error. Since our measurements are limited by systematic uncertainties instead of statistical, we decided to analyze only the trigger configuration of noEC. Bad runs, bad files, runs without
deuterium target, and normalization runs have been excluded. The data files used in this analysis are listed in Ref. [53].

The plots presented here if not specified are taken from the high field data. The high field data and the low field data were analyzed separately using the same criteria.

4.2 Initial Skimming

4.2.1 Photon Energy Corrections

The tagger was calibrated using the procedure described in Ref. [54] which gives the correction factor for the tagger E-counter relative to each other at the level of $10^{-3}$. The correction is a function of the E-counter ID (shown in Fig. 4.1) and of the run number. The run-dependent correction [55] was determined using the kinematically complete final states $\gamma d \rightarrow pp\pi^-$. As can be seen, the electron beam energy drifted

![Figure 4.1](image_url)

**Figure 4.1:** The photon energy correction factor determined using the pair spectrometer equipped with microstrip detectors (MS) for each E-counter.
slightly from run to run. Fig. 4.2 shows the run-dependent correction factor as a function of the run number, typically 0.1%.

![Graph](image)

**Figure 4.2:** The photon energy run-dependent correction factor from the $\gamma d \rightarrow pp\pi^-$ process.

The corrected final photon energy is then given by: $E_{\gamma}^{corr} = E_{\gamma} \cdot \text{Corr}_{ps} \cdot \text{Corr}_{pp\pi}$, and the overall correction is about 0.5%.

### 4.2.2 Particle Identification

The particle ID in the PART bank was used to identify the particle type. One proton and one $\pi^-$ track were required in this analysis. In order to ensure that each event was associated with at least one “good” photon, the time difference, $t_{diff}$, between a photon and a track was required to be within 2 ns. The time difference is defined as:

$$t_{diff} = t_{hadron} - t_{\gamma}$$ (4.2)
where

\[ t_{\text{hadron}} = t_{\text{sc}} - \frac{l_{\text{path}}}{\beta_{\text{hadron}}} c \]  

(4.3)

\[ t_\gamma = t_{\text{tagger}} - \frac{z}{c} \]  

(4.4)

where \( t_{\text{sc}} \) is the scintillator counter (SC) time for the track, \( t_{\text{tagger}} \) is the time of the photon reconstructed in the tagger, \( z \) is the z-vertex of the track, and \( l_{\text{path}} \) is the path length from the vertex to SC. The timing from the start counter is not used in calculating the vertex time in g10 analysis because it was saturated due to the high photon flux in g10 runs. In addition, the status of tagger events is required to be 7 or 15. Status 7 means one unambiguous hit was reconstructed in the tagger, and status 15 means more than one unambiguous hits were reconstructed.

Fig. 4.3 shows the number of good photons. For some events, more than one good photon were found. In this analysis, \( N_\gamma > 0 \) is required. Events with \( N_\gamma > 1 \) are not rejected even if there is an ambiguity regarding which photon is involved in the reaction. Instead, each photon is used to calculate the missing momentum and missing mass, and after applying strict cuts on missing mass, missing momentum, vertex timing, this ambiguity is removed. In order to check this, three data files (42923_00, 42924_00, 43220_01) were analyzed, and among all the survived events (about 22,000), only 1 event was double counted.

The spectator proton is identified from energy-momentum conservation:

\[ M_x^2 = (P_\gamma + P_d - P_p - P_{\pi^-})^2 \]  

(4.5)

where \( P_\gamma, P_d, P_p \) and \( P_{\pi^-} \) are the 4-momenta of the identified particles. As shown in Fig. 4.4, there is a prominent peak around the proton mass. The missing proton identification requires \( |M_x - M_p| < 0.1 \text{ GeV} \) in the data skimming process, and more accurate cuts are determined after the data skimming.
4.2.3 Charged Particle Energy Loss

Before being detected in the drift chambers, charged particles lose energy when interacting with the target cell, start counter and the air inside the CLAS detector. For this reason, the energy must be corrected using the ELOSS package [56] which was updated by Eugene Pasyuk for the g10 configuration. This correction had been tested by Nathan Baltzell, and is in agreement with GSIM simulation [57].

4.2.4 Charged Particle Momentum Corrections

There are 3 different subroutines for momentum corrections for g10 data set. The corrections have been determined from three independent analyses [58]:

1. 4-momentum conservation in $\gamma d \rightarrow pp\pi^-$ reaction;

2. Transverse momentum conservation in $\gamma d \rightarrow pp\pi^-$ reaction;
Figure 4.4: Missing mass distribution for high field data after PID cut only. There is a prominent peak around the proton mass. The $\Delta$ around 1.2 GeV and $N^*$ around 1.5 are also seen in this spectrum.

3. The $K_s^0$ mass in $Ks^0 \rightarrow \pi^+\pi^-$ decay mode.

In this analysis, the second method was used. The corrections depend on kinematics and the torus field and are typically less than 1%.

4.2.5 Rejection of Events During The Beam Trip

The beam was not always stable during each run period. The period of unstable beam was detected by the sync utility [59]. As shown in Fig. 4.5, the beam trip rates (the ratio of the number of triggers during beam trip and the total triggers) were about 2% to 10%, while in some runs, the trip rates were as high as 35%. The photon flux was measured by means of sampling the “out-of-time” hits in the T-counter [60], and the bad beam intervals were removed from the data. In this analysis, the events
during the bad beam intervals were removed as well.

![Figure 4.5: The beam trip rates for each run.](image)

### 4.2.6 Summary of Data Skimming

All the g10 high field and 20% of low field pass-2 data with the liquid deuterium target of the g10 running period were skimmed using bos2root package [61] with the corrections mentioned in this section. The event rate (selected events/total events) for each run was about 2.22%. Some runs and files with very low rates were rejected in the following analysis. About 3.2 billion events from high field data and 620 million events from low field data were analyzed, and 74 million events and 15 millions events survived the skimming, respectively.
4.3 Finer Cuts

After the initial skimming, the total file size reduced from 40 terabytes to 12 gigabytes, and what to do next is to determine the finer cuts.

4.3.1 Z-vertex Cuts

The target for g10 experiment is a 24-cm long cell filled with liquid-deuteron. It is mounted at \( z = -25 \) cm of the CLAS coordinate. As shown in Fig. 4.6, the common z-vertex lies between \(-37\) cm to \(-13\) cm, and to reject events from outside of the cell and the interaction of the photon beam with the target cell entrance and exit windows, the z-vertex cut employed in this analysis requires \(|z + 25| < 11\) cm, that is cutting 1 cm at each end of the target.

![Figure 4.6](image)

**Figure 4.6**: z-vertex of the reconstructed events.
4.3.2 Missing Mass Cuts

During the data skimming, the missing mass cut for the identification of the spectator proton is very loose (\(|M_x - M_p| < 0.1 \text{ GeV}\)). More strict cuts should be applied to identify the spectator. As shown in Fig. 4.7, the gaussian fit shows that the missing mass is peaked at 0.94 GeV, 0.2% higher than the actual proton mass and \(\sigma\) is about 0.0103 GeV. To further identify the spectator, the missing mass is required to be within \(3\sigma\) of the mean value. The \(\sigma\) value is actually not constant, and it depends on kinematics. Thus, the missing mass (squared) distributions were studied for each photon energy bin. As show in Fig. 4.8, \(\sigma\) increases with \(E_\gamma\) linearly. The angular dependence of the missing mass resolution was studied as well. As shown in Fig. 4.9, most part of the resolution is within 10% of the average. For the highest energies, the angular dependence shows the same behavior. For a gaussian distribution, varying \(3\sigma\) cut by 10%, the events vary only by about 0.5%. Thus, it is sufficient to ignore the angular dependence and apply only the energy-dependent \(3\sigma\) cuts. The mean value and cuts of the missing mass squared are shown in Fig. 4.10 and were compared with simulation. While both the data and the MC show similar energy-dependent trend in the missing mass resolution, the MC has better missing mass resolution.

This channel under study is a very clean process, and there is almost no background in the missing mass distribution. There are two peaks in the missing mass spectrum as shown in Fig. 4.11. The other one is centered at 1.14 GeV with \(\sigma = 39\) MeV and it is more than \(5\sigma\) away from the proton peak, and at least \(3\sigma\) away from the proton missing mass cut, so the contamination if there is any, is negligible. The missing mass distribution is fitted with a gaussian plus a linear background, and the background is about 2% (Fig. 4.12 top panel), and the same analysis was done to the simulated data, and the background is about 1% (Fig. 4.12 bottom panel). The missing mass resolution is actually dependent on the photon energy, and a detailed
background study for each energy bin was done. As shown in Fig. 4.8, the missing mass resolution becomes worse at higher photon energies, and the background shown in Fig. 4.13 increases with energy from about 2% to 6%. The same procedure was applied to the simulated data, and the background is about 0.7% to 1.3%. The in-flight $\pi^-$ decay, which affects the mass of this track, makes the tail of the missing mass spectrum and contributes to the background in the simulation. The net background is the difference between that from the data and the simulation.

### 4.3.3 Photon Tagger And Photon Energy Cuts

The CLAS photon-tagging system can tag photon with energy between 20% and 95% of the energy of incident electron beam. In g10 experiment with electron beam energy of 3.767 GeV, the photon energy ranges from 0.75 GeV to 3.58 GeV. The 384 E-counters in the tagger together with 61 T-counters, provide 767 photon energy
Figure 4.8: Sigma of the missing mass (squared) distribution of the skimmed high field data.

bins, with energy resolution of 0.1% \( E_e \) [49]. The photon flux is estimated from the out-of-time tagger hits together with the tagging efficiency measured by total absorption shower counter. The photon flux versus E-counter ID for high field data is shown in Fig. 4.14. In this analysis, we study the cross sections with \( E_\gamma \) between 1 GeV and 3.6 GeV.

4.3.4 Vertex Time Cuts

The proton and pion vertex time is derived from the time-of-flight (TOF) counters, and the photon vertex time is determined by the tagged electron. Combined with the RF signal, the vertex time can be determined accurately. The time difference between hadron and photon defined in Eq. 4.2 is shown in Fig. 4.15 and Fig. 4.16, and the resolution is about 150 ps for pions and 170 ps for protons. Times of both protons and pions are required within 1 ns of that of photon which ensures that they
Figure 4.9: Angular dependence of the missing mass (squared) resolution of the skimmed high field data. They are within 10% of the average values for all the photon energies.

are associated with the same beam bunch.

4.3.5 Missing Momentum Cuts

After applying the aforementioned cuts, the process selected is $\gamma d \rightarrow \pi^- pp$, which includes the quasifree process, three-body breakup, and final state interaction of the quasifree process. For a quasifree process, the spectator proton is not involved in the reaction, the missing momentum distribution should follow the Fermi distribution. As shown in Fig. 4.17, the missing momentum distribution agrees well with the Monte Carlo simulation below 0.2 GeV/c while there are much more events at higher momentum then expected. In this analysis, these events are rejected by requiring $P_x < 0.2$ GeV/c. Fig. 4.18 shows the neutron mass distribution calculated from
Figure 4.10: Mean value (middle) of the missing mass squared, and cuts (top and bottom) for data (black) and MC (red). They both show the same trend, but MC has better missing mass resolution.

\[ M_n^2 = (P_p + P_{\pi^-} - P_\gamma)^2 \]  

before and after the missing momentum cuts. After applying the cuts, the neutron mass distribution agrees with the simulation.

4.3.6 Fiducial Cuts

In this analysis, fiducial cuts are used to cut out the regions where the CLAS acceptance is not well simulated. This package is developed by Dr. Mibe [62] for both high and low field torus settings. It defines cuts on scattering angle (\(\theta\)) to cut out the very forward and very backward tracks, and cuts on the azimuthal angle (\(\phi\)) to reject tracks too close to the torus coils or in the gaps between sectors.
4.3.7 Data Quality Check

In the skimming process, test runs, and junk runs are ignored, and only production runs are analyzed. However, even in the production runs, there could be some problems, one way to check the quality of the data is to calculate the yield of each run. Some runs with very low yield were rejected. As shown in Fig. 4.20 and 4.21, the yields for runs between 43034 and 43194 are smaller than others. This is caused by a new hole in DC Sector 3. As can be seen in Fig. 4.22, sector 3 has one hole in the beginning, and then a new one appeared around run 43034, and this new hole was not fixed until run 43194, but the first hole had never been fixed. However, since the DC efficiency used in MC simulation is the average of all runs, so the effect of both holes were taken into account in the simulation and all runs between 43034 and
Figure 4.12: Top: Missing mass distribution of high field data. Bottom: Missing mass distribution from MC simulation. The curves are the fits to the data with gaussian plus linear background.
43194 have been used in this analysis.

4.3.8 Study of Background Using Hydrogen Runs

There might be some contamination from the events scattered off the target cell or even cosmic rays. There are runs with empty target which can be used to see the effect of these contamination. One better way to estimate the contamination is to analyze the hydrogen runs using the same procedure. Since for a hydrogen target, if there is no contamination, no $p\pi^-$ final state should be found.

Another advantage of analyzing hydrogen runs is that we can examine the possible contribution from the false events from misidentified particles. As shown in Fig. 4.19, the vertex distribution of events with one proton and one $\pi^-$ (no other cuts) clearly shows the events from the target cell window and from hydrogen.
There are many events with $p\pi^-$ in the final state from the hydrogen target. However, from the missing mass and missing momentum distributions, we can see that most of the events are not really quasifree $p\pi^-$ events. After applying the same cuts described before, most of these events did not survive. For hydrogen runs (43238 - 43241) analyzed, the survived events compared to D2 runs are list in Table 4.1.

Table 4.1: List of survived $p\pi^-$ events from H2 and D2 target.

| H2 RUN | Flux  | Events |
|--------|-------|--------|
| 43238  | 1.3e11| 258    |
| 43239  | 2.65e11| 525    |
| 43240  | 2.24e11| 417    |
| 43241  | 2.35e11| 441    |

| D2 RUN | Flux  | Events |
|--------|-------|--------|
| 42924  | 8.44e10| 94777  |
| 42925  | 7.7e10  | 87197  |
Figure 4.15: vertex time of the reconstructed proton relative to photon (high field data), the curve is the gaussian fit of the data.

The survived events are very few for H2 runs compared with D2 runs. As we can see from the Table 4.1, the number of events normalized by total flux for H2 runs is only about 0.17% of the D2 runs. This means that the possible contaminations from the cell or other sources is very small and can be neglected.

4.3.9 Summary

All the g10 high field data were analyzed first, and then the same criteria were applied to analyze the low field data. After all the cuts, 26 million high field events and 5 million low field events survived. The cuts and the number of the events that survived are listed in Table 4.2 and 4.3.
Figure 4.16: vertex time of the reconstructed pion relative to photon (high field data), the curve is the gaussian fit of the data.

Table 4.2: Summary of the cuts used in this analysis

| Name               | Initial skimming | fine cuts                                      |
|--------------------|------------------|-----------------------------------------------|
| Good beam          | √                | √                                             |
| $N_p, N_\pi$       | $N_p, N_\pi = 1$ | $N_p, N_\pi = 1$                              |
| $N_\gamma$         | $N_\gamma > 0$   | $N_\gamma > 0$                                |
| Timing             | $|t_\pi - t_\gamma| < 2$ ns or $|t_p - t_\gamma| < 2$ ns | $|t_\pi - t_\gamma| < 1$ ns and $|t_p - t_\gamma| < 1$ ns |
| Missing mass       | $|M_x - M_p| < 0.1$ GeV | $|M_x^2 - M_p^2| < 3\sigma$                   |
| Missing momentum   | no               | $P_x < 0.2$ GeV/$c$                           |
| Vertex             | no               | $|Z_{vertex} + 25| < 11$ cm                   |
| Fiducial           | no               | yes                                           |
Figure 4.17: Missing momentum distribution (after all cuts except missing momentum cut) of the data (black) and Monte Carlo simulation (red)

Table 4.3: Summary of the events survived after skimming and all cuts.

|                                | High field | Low field |
|--------------------------------|------------|-----------|
| Total Events Analyzed          | 3,208 M    | 620 M     |
| Events Survived Skimming       | 74 M       | 15 M      |
| Events Survived All Cuts       | 26 M       | 5 M       |
Figure 4.18: Comparison of neutron mass distribution between the data (black) and the Monte Carlo simulation (red) before (top) and after (bottom) the missing momentum cuts.
**Figure 4.19:** The distributions of $p\pi^-$ events from LH2 runs 43238 - 43241. The CUT 1 in the plot includes: missing momentum cut, z-vertex cut, and timing cuts. The CUT 2 includes: missing mass cuts.

**Figure 4.20:** Events rate of skimming (upper) and after all cuts (lower)
Figure 4.21: Yield (N(selected)/flux) for each run

Figure 4.22: Comparison of momentum versus $\theta$ distribution of DC sector 3, before (top) and after (bottom) the new hole appeared. An extra hole can be seen on the bottom plot.
Extracting the differential cross section requires an understanding of the acceptance of the CLAS detector. It needs not only the geometry of the detector but also the efficiency of the detector subsystems.

5.1 Monte Carlo Simulation

The CLAS acceptance is determined from Monte Carlo simulation with efficiency extracted from the data and is described below in detail.

5.1.1 Event Generator

The event generator used in this analysis is the fsgen [63] based on PYTHIA package, which generates the $\gamma n \rightarrow \pi^- p$ events uniformly in the center-of-mass frame. The Fermi motion of the initial neutron was modeled using the deuteron momentum distribution obtained from the Bonn potential. No information on the cross section is included in fsgen. A new generator was developed which generates events according to the yield $N/\epsilon$, where $\epsilon$ is the acceptance extracted using generator fsgen. The acceptance of a two-particle final state is almost independent of the event distribution
from the generator.

5.1.2 Event Processing

The generated data have been processed by GSIM to simulate the CLAS response, considering CLAS as an ideal detector where all the subsystems are working properly as designed. Then the output of GSIM was processed by GSIM Post Processor (GPP), which smears the timing and spatial resolutions of the time-of-flight and drift chambers subsystems. The dead regions and inefficiency of the subsystems were also taken into account by GPP. After that, the data were processed by the reconstruction program user_ana which was used to cook the g10 data.

The output data of the simulation were analyzed by the same event selection program used to analyze the g10 data. The acceptance of the CLAS detector can be extracted by comparing the generated events and the ones that survived all the cuts.

5.1.3 CLAS Detector Subsystem Inefficiency

Since the CLAS detector used in the g10 experiment is divided into six regions by superconducting magnets, and each acts as an independent detector covering $\pi/3$ of the azimuthal angle acceptance. This feature provides a good internal check by comparing the differential cross sections from each sector.

The yield for the $\gamma n \rightarrow p\pi^-$ channel is studied in detail at $E_\gamma \sim 1$ GeV, where the flux is maximized. The yield is defined as

$$Y(\theta_{c.m.}) = \frac{N(\theta_{c.m.})}{\epsilon(\theta_{c.m.})}$$

(5.1)

where $N$ is the number of events, and $\epsilon$ is the acceptance calculated using Monte Carlo (MC) simulation. $N$ is sector dependent due to different inefficiencies in different sectors such as drift chamber holes caused by dead wires. The GSIM (MC simulation for CLAS detector) simulates the CLAS geometry and GPP takes into
account the drift chamber (DC) wire efficiency and resolutions. Ideally, if all the inefficiencies were properly simulated, the acceptance $\epsilon$ would compensate the loss of events due to the detector inefficiency, and the yield should be the same for each sector. As shown in Fig. 5.1, the yields at $E_\gamma \sim 1.15$ GeV show large discrepancies among sectors, which indicate poor understanding of the CLAS detector. Several other studies also show similar discrepancies.

![Figure 5.1](image)

**Figure 5.1:** Normalized yield versus $\theta_{\text{c.m.}}$ for $\gamma n \rightarrow p \pi^-$ at $E_\gamma \sim 1.15$ GeV. Large discrepancies were found amongst sectors.

In order to detect a charged particle, one needs the tracking information in the DC, and timing information from time-of-flight (TOF). TOF signals form level one trigger, and level two trigger requires “likely tracks” in DC. The DC wire efficiency is used in the MC simulation, while TOF is generally considered to be 100% efficient in the simulation. The CLAS detector subsystem including DC and SC have been studied in detail in the following sections. The understanding of the detector efficiency
has been greatly improved, and the results show that the DC and SC efficiency used in the MC has greatly improved the agreements of cross sections among six CLAS sectors.

5.2 DC Efficiency

The CLAS drift chamber system is designed to achieve a 0.5% momentum resolution and a 2 mrad angle resolution. To provide such high resolution, 18 drift chambers were built, and 3 in each sector. As shown in Fig. 5.2, these 3 drift chambers are referred to as Region 1 (R1), 2, and 3, respectively. Each drift chamber consists of two “superlayers”: axial and stereo layers. The stereo superlayer is tilted 6° relative to the axial layer to determine the azimuthal angle. Each superlayer consists of 6 layers of drift cells except for the stereo layers of R1 which have only 4 layers due to the space constraint. There are about 120 to 192 sense wires in each layer, and in total there are more than 35,000 sense wires (total 130,000 wires including field wires). The aging effect of the wires affects the efficiency of the wires. Aside from the aging effect, electric failures such as ADB boards are responsible for most of the event loss.

5.2.1 Maurizio-Li Algorithm

Maurizio and Li developed an algorithm [64] (Maurizio-Li Algorithm) to extract the wire efficiency using the DC0 bank in raw data which contains the DC hits. The basic idea of this algorithm is that CLAS contains six identical sectors, and the hit occupancy of each layer should follow the same distribution for all sectors. If the occupancy of some wires is much lower than most of the others, these wires are inefficient. The DC hit occupancy is extracted using PDU [65]. As shown in Fig. 5.3, the hit occupancy distribution of layer 7 for all sectors, each layer does follow the same distribution, but there are some fluctuations due to noisy hits. The algorithm
then tries to find the expected occupancy for each wire by averaging over a group of neighboring wires excluding the very hot wires and dead wires. The bold red curve in Fig. 5.3 represents the expected occupancy distribution and the bold blue curve represents 90% of the expected values. The hit distribution is then normalized by the expected distribution. As can be seen, the expectation lies between most of the wires, and most of the wires have “efficiencies” between 0.9 and 1.1. This algorithm, focusing on identifying dead wires and very low efficiency wires, assigns 100% efficiencies to all the wires above the 0.9 line. Since most of the wires are in good condition, and the event loss is mainly due to the bad ADB boards (each failed ADB board causes loss of $6 \times 8$ wires), the Maurizio-Li algorithm can easily identify these bad regions, and the MC simulation using efficiency table from this algorithm has been quite accurate.
5.2.2 Causes of Inaccuracy in DC Efficiency

The DC efficiency is considered to be accurate, but for high statistics differential cross section study, it may still need improvement. Three major causes of inaccuracy in DC efficiency have been found: cable swapping, dead TOF paddles, and contaminations.

Swapped Cables

There are over 35,000 sense wires. The EG3 group found that groups of wires from same connectors have very low time-based-track (TBT) hits. This later was confirmed [66] by studying the DC “efficiency” using the TBLA bank. Due to these cable swappings, the geometry of these regions is mixed up and thus fewer tracks can be reconstructed from these regions. Even worse is some region (16×6) in S5R1 where low TBT hits can not be explained by simple cable swapping. The previous
Figure 5.4: DC efficiency of Sector 3 for high field data using the Maurizio-Li algorithm.

The CLAS level one trigger system requires two TOF hits. Thus, the inefficiency of TOF paddles would cause loss of events and low occupancy of DC wires leading to these paddles. There are some known dead paddles, such as paddle 9 and 11 in sector 3. As can be seen in Fig. 5.4, the DC efficiency of sector 3 from the Maurizio-Li algorithm shows very low efficiency for a big region (wires between 15 and 40, layers above 15). This low DC occupancy is believed to be caused by the two dead TOF paddles. There are similar low efficiency regions in other sectors. To avoid over-correcting for these regions, we have to find a new algorithm to solve this problem.

Contaminations

Since the DC efficiency is determined by DC wire hits, the contamination (hits caused by particles not from the target) distribution from upstream devices needs to be studied. The wire occupancy distributions for g10 and g2 are studied for all
layers. As can be seen in Fig. 5.5, sector 5 and 6 have much more hits than other sectors in the backward regions. Especially in g2, the occupancy in sector 5 and 6 is 100 - 200% more than other sectors. With better shielding in g10, the occupancy in sector 5 and 6 is still much higher (about 30 - 80%) than other sectors. This huge discrepancy is only observed in R3 which is the outermost region of the drift chambers and indicates that the huge number of hits is caused by particles coming from outside of the CLAS detector (upstream devices such as tagger dump). As a result, the algorithm may not clearly identify the proper average value, and thus favor sector 5 and 6, and underestimates efficiencies for other sectors.

5.2.3 New Algorithm

A new algorithm is proposed to solve all the aforementioned problems. The essential of this algorithm is the so-called “excluded-layer method”.

---

**Figure 5.5**: DC0 occupancy of Layer 36 for g10 (upper), and g2 (bottom) runs.
There are 34 layers of drifting cells. It does not require all the layers to fire to reconstruct a track. Reconstructing a track requires at least 4 track segments in six superlayers allowing 2 superlayers without segments. Also, the track segment in each superlayer allows 2 layers out of six to be missing. Once a track is reconstructed, one can determine the wires it passed through in the missing layers according to the trajectory. With this information, we can define the wire efficiency \( \epsilon \) as

\[
\epsilon(i) = \frac{N_h(i)}{N_h(i) + N_r(i)}
\]

(5.2)

where \( N_h(i) \) is the number of tracks passing wire \( i \) with hit in this wire, and \( N_r(i) \) is the number of tracks passing wire \( i \) without firing this wire. When counting \( N_h(i) \), one should only count the track which remains a good track (with at least 4 good track segments) excluding the layer \( i \). In other words, the efficiency \( \epsilon(i) \) is the ratio of good hits of wire \( i \) and total good tracks (determined without counting this wire) passing this wire.

This method takes only the TBT hits which exclude almost all the noise hits and contaminations which hardly forms good tracks. Since it is a ratio of number of tracks, the effect of a bad TOF paddle cancels. In addition, regions with swapped cables will have much fewer reconstructed tracks, thus are correctly assigned low efficiency.

5.2.4 DC Efficiency Using New Algorithm

In this new method, the TBLA bank is required which contains the TBT hits information. However in g10 pass-2 cooked data, there is no TBLA bank available. 31 files of raw data were recooked for each field to get the additional TBLA bank. It provides enough statistics and most of the wires are quite stable during the whole run period. However, some ADB boards were dead for a while and fixed later. Eq. 5.2 can not take into account the dead time of these ADB boards. For a track passing
through the superlayer associated with a dead ADB board, the whole track segment in this superlayer would be missing. As a result, the number of tracks and number of hits for this region are both zero during the dead time, thus, the result from Eq. 5.2 for this region is only the efficiency when the ADB board is alive. For these regions, unlike Maurizio-Li algorithm which can take into account the dead time for these regions, special treatment needs to be done in this new algorithm. First, since they were completely dead when the ADB board failed, and close to 100% efficient when the board was fixed, a weight factor needs to be assigned as the efficiencies for these wires which is proportional to the alive time of these wires. All dead ADB boards were identified by checking each PDU occupancy for each run and the results are consistent with the log books. These ADB boards are listed in Table 5.1, and the weight \( w \) is defined as

\[
w = \frac{N_{\text{alive}}}{N_{\text{total}}} \tag{5.3}
\]

where \( N_{\text{alive}} \) is the number of triggers when this board is alive, and \( N_{\text{total}} \) is the total triggers during the whole run period of the same field setting. The wire efficiencies are weighted by the factors listed in Table 5.2, and in addition, each group of wires from these ADB boards is assigned a group number which is used in GPP such that all the wires die at the same time. For some reason, in TBLA bank there is no information for wires \( > 120 \) in superlayers 2. Since these very backward small regions were not found any problem in the Maurizio-Li algorithm, the wires there are assigned efficiencies of 100%.
Table 5.1: Dead ADB boards and their dead time

| Field | Sector | Layers | Wires | Time                                      |
|-------|--------|--------|-------|-------------------------------------------|
| 3375  | 2      | 31 - 36| 57 - 64| Dead since Run 43008                      |
| 3375  | 3      | 13 - 18| 81 - 88| Dead between Run 43042 - 43194           |
| 3375  | 3      | 19 - 24| 85 - 92| Dead between Run 43042 - 43194           |
| 3375  | 5      | 25 - 30| 9 - 16 | Dead between Run 42987 - 43188 and after Run 43209 |
| 2250  | 1      | 13 - 18| 153 - 160| Dead since Run 42762                   |
| 2250  | 1      | 13 - 18| 57 - 64 | Fixed since Run 42772                   |
| 2250  | 1      | 19 - 24| 61 - 68 | Fixed since Run 42772                   |
| 2250  | 1      | 7 - 12 | 115 - 120| Fixed since Run 42571                   |

Table 5.2: Weight factors for the partially dead ADB boards.

| Field | Sector | Layers | Wires | Weight |
|-------|--------|--------|-------|--------|
| 3375  | 2      | 31 - 36| 57 - 64| 0.2606 |
| 3375  | 3      | 13 - 18| 81 - 88| 0.5635 |
| 3375  | 3      | 19 - 24| 85 - 92| 0.5635 |
| 3375  | 5      | 25 - 30| 9 - 16 | 0.2515 |
| 2250  | 1      | 13 - 18| 153 - 160| 0.8614 |
| 2250  | 1      | 13 - 18| 57 - 64 | 0.1287 |
| 2250  | 1      | 19 - 24| 61 - 68 | 0.1287 |
| 2225  | 1      | 7 - 12 | 115 - 120| 0.5335 |

The final DC efficiencies are shown in Fig. 5.6 and 5.8, and the old efficiencies are shown in Fig. 5.7 and 5.9. A recent study on single proton efficiency using the new efficiency map has shown better agreement between data and simulation [67].

The efficiencies are put into CLAS calibration database:

CLAS_CALDB_RUNINDEX=calib_user.RunIndexWeig10a,
and a summary of the efficiencies for both high field and low field data are shown in Table 5.3. The efficiency of more than 75% of the wires are above 95%. The tracks passing too close to the sense wires which results in the signals with low pulse height and long duration may contribute to part of the small inefficiency [51].
Figure 5.6: DC efficiency for high field data using new algorithm
Figure 5.7: DC efficiency for high field data using old algorithm.
Figure 5.8: DC efficiency for low field data using new algorithm.
Figure 5.9: DC efficiency for low field data using old algorithm.
Table 5.3: Summary of DC efficiencies for both high field and low field data.

| Efficiency range | # of wires | %     |
|------------------|------------|-------|
| 0.00 - 0.05      | 2525       | 7.16% |
| 0.05 - 0.10      | 198        | 0.56% |
| 0.10 - 0.15      | 148        | 0.42% |
| 0.15 - 0.20      | 62         | 0.18% |
| 0.20 - 0.25      | 82         | 0.23% |
| 0.25 - 0.30      | 74         | 0.21% |
| 0.30 - 0.35      | 69         | 0.20% |
| 0.35 - 0.40      | 82         | 0.23% |
| 0.40 - 0.45      | 83         | 0.24% |
| 0.45 - 0.50      | 88         | 0.25% |
| 0.50 - 0.55      | 133        | 0.38% |
| 0.55 - 0.60      | 129        | 0.37% |
| 0.60 - 0.65      | 162        | 0.46% |
| 0.65 - 0.70      | 187        | 0.53% |
| 0.70 - 0.75      | 202        | 0.57% |
| 0.75 - 0.80      | 330        | 0.94% |
| 0.80 - 0.85      | 429        | 1.22% |
| 0.85 - 0.90      | 783        | 2.22% |
| 0.90 - 0.95      | 2392       | 6.78% |
| 0.95 - 1.00      | 26448      | 74.98%|
| 1.00             | 668        | 1.89% |

| Efficiency range | # of wires | %     |
|------------------|------------|-------|
| 0.00 - 0.05      | 2510       | 7.12% |
| 0.05 - 0.10      | 174        | 0.49% |
| 0.10 - 0.15      | 120        | 0.34% |
| 0.15 - 0.20      | 71         | 0.20% |
| 0.20 - 0.25      | 67         | 0.19% |
| 0.25 - 0.30      | 85         | 0.24% |
| 0.30 - 0.35      | 75         | 0.21% |
| 0.35 - 0.40      | 75         | 0.21% |
| 0.40 - 0.45      | 95         | 0.27% |
| 0.45 - 0.50      | 89         | 0.25% |
| 0.50 - 0.55      | 95         | 0.27% |
| 0.55 - 0.60      | 128        | 0.36% |
| 0.60 - 0.65      | 158        | 0.45% |
| 0.65 - 0.70      | 182        | 0.52% |
| 0.70 - 0.75      | 215        | 0.61% |
| 0.75 - 0.80      | 349        | 0.99% |
| 0.80 - 0.85      | 441        | 1.25% |
| 0.85 - 0.90      | 702        | 1.99% |
| 0.90 - 0.95      | 2313       | 6.56% |
| 0.95 - 1.00      | 26679      | 75.63%|
| 1.00             | 651        | 1.85% |

Contaminations and noise hits may affect the tracking efficiency as well.

5.2.5 Summary

The Maurizio-Li algorithm did very good job in understanding the overall wire efficiency, however could not give accurate efficiency for some particular regions. This new algorithm solved these problems. For the DC inefficiency caused by swapped cable, it is hard to simulate in GPP, the easiest solution is to put fiducial cuts to remove the tracks which went through these regions.
5.3 TOF Efficiency

The scintillation counters are crucial in the experiment not only in the sense they provide TOF for particle identification but also provide level one trigger.

There’s no TOF efficiency in the MC simulation (GSIM and GPP). In some analyses, only several very bad paddles were rejected.

5.3.1 Direct Evidences of TOF Inefficiency

The TOF efficiency has not been evaluated in the g10 data. However, a detailed differential cross section [68] study suggests that TOF efficiency needs to be taken into account. Several direct evidences have been found supporting the TOF efficiency correction.

Timing information is crucial to identify charged particles. As can be seen in Fig. 3.5, CLAS is covered by 4 panels of TOF paddles in each sector to provide timing information. There are extra detectors such as Gas Cherenkov counters (CC) and Electromagnetic calorimeters (EC) placed in the forward direction overlapping with the first panel of TOF paddles (22 paddles). For particle identification, TOF hits are first searched for matching with the reconstructed tracks. If no matching TOF hits were found, the timing from CC or EC is used. Among all the reconstructed $\pi^-$ in the first panel, 3% were found using EC or CC timing. If considering only the events with at least 3 charged particles, about 5.6% of the $\pi^-$ events were identified without TOF hits. Some of these tracks are caused by the geometry of the TOF paddles which is already simulated in GSIM. However the rates of this kind of tracks varies among sectors. For example, there are 50% more tracks using EC or CC in sector 3 than sector 1 which can not be explained solely by geometry and indicates a non-negligible inefficiency of TOF.

Each TOF paddle consists of two photo-multiplier tubes (PMTs), so for each
paddle there are four signals, two TDCs and two ADCs. Only one dynode signal is required to form the Level 1 trigger to get a prompt response, and as a result, some hits may not have all the 4 signals. The TOF hit status was plotted in Fig. 5.10, where status 15 means 4 good signals (2 ADCs and 2 TDCs), and all the status below 15 means missing at least one of the signals. As can be seen, most of the hits are good, but still a lot of hits missed some signals which may result in poor timing resolution for these hits. Among all the reconstructed charged particles, about 3% are found with status < 15.

For a clean selection of events, all the particles with TOF status < 15 are rejected in this analysis.
5.3.2 Method to Extract TOF Efficiency

The following analysis is based on the cooked high field data, and one file in each run was used (about 4% of the high field data).

TOF signals form Level 1 trigger, and Level 2 trigger from DC will select the event if a “likely track” was found. Level 2 trigger allows some of the DC inefficiency because it requires only 3 segments out of 5 superlayers (first superlayer is not used) and 4 layers out of 6 in each superlayer. If most of the DC wires are in good condition, the TOF hit distribution depends only on the TOF efficiency (assuming background is small and symmetric). The SCRC bank in the cooked data is used to study the TOF hit distribution. Since some particles hit two neighbouring TOF paddles, in the SCRC bank, this kind of hits are combined into one hit. A single hit has a status between 1 to 15, while a double hit is assigned \( status = status(1) \times 100 + status(2) \). Only hits with \( status \geq 15 \) are selected which includes all good single hits, and any double hits. Ideally, the TOF occupancy should be the same for all sectors if there’s no inefficiency.

However, as shown in Fig. 5.11, many paddles were found very low occupancy compared to the paddles in other sectors. Generally, the paddles in the first panel (first 22 paddles) are relatively consistent except for sector 3 where there are several big holes in DC and two dead paddles in SC. The paddles beyond the first panel show quite large discrepancies. The two runs shown in Fig. 5.11 represent two different torus field settings, and two different triggers: 42731 - low field with EC in the trigger, 42923 high field no EC in the trigger.

If there is no background TOF hits left after the status cut, the occupancy should be proportional to its efficiency for each paddle,

\[
N(s, p) = N_{real}(s, p) \times \epsilon(s, p)
\]

(5.4)

where \( s \) refers to sector, and \( p \) paddle and \( N(s, p) \) is the number of TOF hits, and
Figure 5.11: TOF Occupancy for run 42731 (top) and 42923 (bottom).
$N_{\text{real}}$ is the number of charged particles passing through. However, there’s no easy way to get $N_{\text{real}}$. We know that most of the paddles are working nearly 100% efficient. For a really good paddle $N \simeq N_{\text{real}}$ because the efficiency is almost 1.

A naive way to define the TOF efficiency $\epsilon$ is

$$\epsilon(s,p) = \frac{N(s,p)}{\max_i N(i,p)}$$  \hspace{1cm} (5.5)

That is to pick one best paddle among 6 symmetric paddles around the beam line as the normalization. In this case we assume this paddle is 100% efficient. Since $\max_i N(i,p) \leq N_{\text{real}}$, we can get an upper limit of the efficiency for all the TOF paddles.

The calculated efficiency is shown in Fig. 5.12 for 4 runs, 42731 (v2, withEC), 42733 (v1, noEC), 42922 (v2, noEC) and 42923 (v1, noEC). As can be seen, the efficiencies remain the same for most of the paddles, even if the torus settings, and triggers changed. Since DC was in the Level 2 trigger, the only concern about this method is that the occupancy difference between sectors may be caused by the inefficiency in DC. However, the chance for a Level 2 trigger to reject a track is almost zero even if two superlayers the track passing through are missing as long as other three superlayers don’t have holes. The almost constant efficiency calculated indicated that the DC inefficiency contribution is negligible. If the DC inefficiency caused a large loss of the TOF occupancy, the affected regions would shift when the torus field changed. To further confirm this, we studied several pairs of runs: 42771 and 42772, 42521 and 42522. Run 42771 has two extra consecutive holes in sector 1, region 2 compared to 42772, while 42522 has missed half of the whole superlayer 3 in sector 2. A comparison of the TOF efficiency was performed [69], and the efficiency remained the same. We can safely conclude that the TOF efficiency does not depend on the DC inefficiency for the g10 run period except for the forward region of sector 3. In sector 3, the DC has more holes than any other sectors. The occupancy for
the first panel of sector 3 is about 20% lower than other sectors when there’s no EC in the trigger. The exact cause remains unknown. Nevertheless, we can still get reasonable efficiency for these paddles using runs with EC in the trigger.

This efficiency map then is used in MC simulation in which each particle has a probability equal to the efficiency to fire the TOF.

The $\gamma n \rightarrow p\pi^-$ channel was simulated for $E_\gamma \sim 1.15$ and 1.45 GeV. A comparison of yields (defined in Eq. 5.1) from the simulation with and without TOF efficiency is shown in Fig. 5.13 and 5.14. As can be seen, using such a simple way to correct for the TOF, much of the discrepancy was removed for both $E_\gamma \sim 1.15$ and 1.45 GeV. Due to the different kinematics, overall correction is dependent on the energy, angle and torus current. Fig. 5.15 shows the ratio between the MC acceptance with TOF
efficiency and without TOF efficiency at different photon energies as a function of \(\theta_{\text{c.m.}}\). As we can see, the overall correction varies, and is around 20%.

Later, we extracted all the cross sections to 3.56 GeV using this single efficiency map, and they are in very good agreement with world data to highest energy. This gives us confidence that the naive method gives a very good estimation of the upper limit of the TOF efficiency.

5.3.3 Summary

In this section, we implemented a method to correct for the TOF inefficiency. The TOF efficiency for each paddle is proportional to the total good TOF hits in each paddle. This method provides the TOF efficiency completely independent of the channel under study. The differential cross sections for \(\gamma n \rightarrow p\pi^-\) are studied using this method at \(E_\gamma \sim 1.15\) and 1.45 GeV. They are compared with world data [47, 70, 71, 72]. As shown in Fig. 5.16 and 5.17, the differential cross sections using TOF corrections agree with the world data very well. With these corrections, the internal consistency (among sectors) and external consistency (with world data) are both achieved. Another indirect way to measure the TOF efficiency is using inclusive proton and \(\pi^-\) together with MC simulation [73]. The TOF efficiency from that method is consistent with the method described in this section. The final cross sections from both methods are in very good agreement with each other. This TOF inefficiency may not be caused by the scintillation counters themselves, since they are known with very good efficiency. The electronics, light guides, PMTs may contribute to the inefficiency. The inefficiency caused trigger inefficiency. Although it is about 10%, it actually affects the channels with more than two charged particles much less. For example, for channel with two charged particles, assuming every TOF paddle has 10% chance of not firing the trigger, both tracks have to give good trigger or this event will be lost. The total probability of losing this event is \(1 - 0.9^2 = 19\%\).
Figure 5.13: Top: yield at $E_\gamma \sim 1.15$ GeV before applying the TOF efficiency; Bottom: yield after applying the TOF efficiency.
Figure 5.14: Top: yield at $E_\gamma \sim 1.45$ GeV before applying the TOF efficiency; Bottom: yield after applying the TOF efficiency.
which is significant, and the TOF inefficiency has to be taken into account. While for channel with three charged particles, as long as two of them provide good trigger, the event can be collected. Thus, the trigger inefficiency is $0.1^3 + 3 \cdot 0.9 \cdot 0.1^2 = 2.8\%$ which is quite small. It is also consistent with the findings from the studies of other channels such as $\phi$-production which didn’t see the effect of the TOF inefficiency.
Figure 5.16: Differential cross sections for $\gamma n \rightarrow p\pi^-$ using different TOF efficiencies at $E_\gamma \sim 1.15$ GeV.

Figure 5.17: Differential cross sections for $\gamma n \rightarrow p\pi^-$ using different TOF efficiencies at $E_\gamma \sim 1.45$ GeV.
5.4 Acceptance

The acceptance is defined as the ratio of all the survived events and the total generated events. About 54 million (high field) and 67 million (low field) events were simulated for $E_\gamma$ between 1 GeV and 3.5 GeV and after applying all the cuts defined in the real data analysis, about 13 million and 18 million events survived. The acceptance is calculated as a function of $\theta_{c.m.}$ for each $E_\gamma$ bin. Both DC and SC efficiencies were used in the simulation to better describe the detector response.
Correction For Final State Interaction

Since deuterium is used as the effective neutron target in this experiment, the nuclear transparency effect, i.e., the final state interaction, must be considered to account for the loss of particles on the exit channel of the reaction of interest. According to the Glauber formulation, the nuclear transparency of a particle at position $r$ with momentum $p$ is related to the total cross section of particle with nucleons inside the nucleus as \[ t(r, p) = \exp\left\{-\int_0^\infty ds \rho(r + \hat{p}s)\sigma\right\}, \tag{6.1} \]
where $\hat{p}$ is the unit vector of the momentum $p$, $s$ is the measure of distance in the direction of $\hat{p}$, $\sigma$ is the effective final state interaction cross section, and $\rho(r + \hat{p}s)$ is the density of the nuclear medium. The transparency for $d(e, e'p)$ has been well measured and is in good agreement with the Glauber prediction [75, 76]. Based on the Glauber formulation, the transparency for $d(\gamma, \pi^- p)p$ can be deduced from that of $d(e, e'p)$ by comparing the cross section of the final state interaction following the procedure used in the Hall A experiment on the same channel [1].

In $d(e, e'p)$ process, the loss of the events is due to $pn$ interaction in the final
state. While in $d(\gamma, \pi^-)p$, there are three particles (two protons and one $\pi^-$) in the final state. Both the $pp$ interaction and $\pi^-p$ interaction contribute to the event loss. So the transparency can be calculated by

$$\frac{\ln (t_{\pi^-p}(\vec{p}))}{\ln (t_{pn}(\vec{p}))} = \frac{\sigma_{\pi^-p}(\vec{p})}{\sigma_{pn}(\vec{p})}$$

(6.2)

$$\frac{\ln (t_{pp}(\vec{p}))}{\ln (t_{pn}(\vec{p}))} = \frac{\sigma_{pp}(\vec{p})}{\sigma_{pn}(\vec{p})}$$

(6.3)

where $\sigma_{pp}$, $\sigma_{pn}$ and $\sigma_{\pi^-p}$ are the total interaction cross sections, $t_{\pi^-p}$, $t_{pp}$ are the transparencies for $pp$ interaction and $\pi^-p$ interaction, respectively, and $t_{pn}$ is the transparency for $d(e,e'p)$ and a fitted value of $0.904 \pm 0.013$ [75] was used in the calculation. The total transparency $t_{d(\gamma, \pi^-)p}$ then is given by

$$t_{d(\gamma, \pi^-)p} = t_{\pi^-p}(\vec{p}_{\pi^-}) \cdot t_{pp}(\vec{p}_p).$$

(6.4)

All the cross sections used in the calculation were taken from the 2006 Review of Particle Physics [31]. Here we followed the same procedure developed in Ref. [1]. The transparency is calculated for each photon energy bin. As shown in Fig. 6.1, the transparencies at different photon energies, the typical correction is around 20%. Recently, we worked with Laget to calculate the FSI effect, and the preliminary results from Laget’s approach [77] agree with that from the method described above.

The $\pi^-p$ total cross section used in this calculation is shown in Fig. 6.2. The fine structure in Fig. 6.1 is actually a result of the resonance in the $\pi^-p$ total cross section.
Figure 6.1: The calculated nuclear transparencies for different photon energies. The red points are the results from the aforementioned method with 10% error. The pits in the plots are caused by the resonance structures of the world data used in the calculation. The black curve is the preliminary results from Laget’s approach.
Figure 6.2: Cross sections used in the FSI calculation. Figure is taken from [31]
The systematic uncertainties of this analysis are discussed in this section which includes systematic errors from the following sources:

- Luminosity, including target thickness and photon flux;
- “Background” subtraction;
- Acceptance, including generators and TOF efficiency correction;
- Cuts related;
- FSI correction.

7.1 Luminosity

The uncertainty of the luminosity has been studied in another g10 analysis [78], and was estimated to be less than 10% by comparing the $\gamma p \rightarrow \phi p$ differential cross section with CLAS g6 results and other world data. This estimation was thought of as being too conservative and the actual uncertainty is likely much smaller than 10%.
The uncertainty of the luminosity includes two parts: one is the uncertainty of the target thickness and the other is the photon flux. The target temperature and pressure were constantly monitored and the uncertainty of the density is less than 1% [80]. The photon flux is estimated using the tagging ratios obtained during the normalization runs. There are 3 normalization runs shown in Fig. 7.1 and the tagging ratios are within 1.5%. Run 42582 and Run 42735 agree very well within 0.5% while Run 42734 is 1.5% higher. However, this is a very short run (only 1/10 of the data of others) and was terminated due to some problems and thus this normalization run should be discarded. The tagging ratio related uncertainty then is about 0.5%.

Run-by-run fluctuation also contributes to the systematic uncertainty of the total flux. We studied the normalized yield for each run, and group all the runs according to different experimental conditions, and the normalized yields are shown in Fig. 7.2 - 7.5. Then the distributions were fitted using gaussian distribution, and the standard deviation of the distribution indicates the run-by-run fluctuation under the same experimental conditions, and the difference of the mean value indicates the beam current related uncertainties. As summarized in Table 7.1, all the standard deviations are less than 2% and all the mean values are within 2% of the average. Similar studies have been done for the low field data, and the standard deviations are around 2% and the mean values are about 2.8%.

The total uncertainty of the luminosity is calculated and is always less than 4%. In this analysis, we give a upper limit of the uncertainty to 5% which is consistent with the recent CLAS $\pi^0$ study [37].
Figure 7.1: The tagging ratio for g10 run period. The uncertainty is within 1.5%.

Table 7.1: Summary of the run-by-run fluctuation of 4 groups of runs

| id | Runs            | Mean Value | Sigma (Sigma/Mean)     | Beam Current |
|----|-----------------|------------|------------------------|--------------|
| 1  | 42923 - 42949   | 1.44813    | 0.0279055 (1.9%)       | 25 nA        |
| 2  | 42950 - 43055   | 1.40134    | 0.0252473 (1.8%)       | 30 nA        |
| 3  | 43056 - 43198   | 1.3579     | 0.0207091 (1.5%)       | 28 nA        |
| 4  | 43198 - 43228   | 1.42617    | 0.0181932 (1.3%)       | 28 nA        |
Figure 7.2: Normalized yield distribution of the channel under study for Runs 42923 - 42949 (25 nA).

Figure 7.3: Normalized yield distribution of the channel under study for Runs 42950 - 43055 (30 nA).
**Figure 7.4**: Normalized yield distribution of the channel under study for Runs 43056 - 43198 (28 nA) with extra holes in DC.

**Figure 7.5**: Normalized yield distribution of the channel under study for Runs 43198 - 43228 (28 nA).
7.2 “Background” Subtraction

It has been studied using the H2 runs that the contamination from the target cell can be neglected. As for the spectator proton, it can be easily identified by the missing mass cut. As shown in Fig. 4.13, the fitted background in the missing mass squared is about 2% to 7% depending on the photon energy for the data, while for simulation, it is smaller using the default DC smearing parameters. However, when tuning the DC smearing parameters, the missing mass resolution gets worse, and the fitted background becomes larger. When the DC smearing parameters are doubled, the simulation can reproduce the same missing mass resolution and background. This means some of the background in the real data could come from the badly reconstructed real events. Thus, in this analysis, we do not subtract the background, instead we assign a systematic uncertainty to account for this, which is the difference of the fitted “background” between the data and simulation using the default DC smearing parameters. As shown in Fig. 7.6, this uncertainty is about 1% to 6%. 
Figure 7.6: The systematic uncertainties assigned according to the “background” in the missing mass squared distribution.

7.3 Acceptance

The Monte Carlo simulation has been done to extract the acceptance. See Chapter 5 for details. The $p\pi^-$ events were generated using phase space generator which works well for two body final state. Another generator was used to generate events according to the normalized yield. Both generators gave almost identical acceptance and the difference in acceptance is less than 1%. The TOF inefficiency used in the simulation contributes to a correction around 20%. The systematic uncertainty for the TOF correction is also very important. Since the TOF efficiency is defined as the ratio of the occupancy of the paddles, it is very important to check if the ratio is stable or not. In fact, the TOF efficiencies calculated from different runs throughout the whole run period are quite stable. We calculated the standard deviation of the TOF
occupancy normalized by the sum of the six sectors. The fluctuation, as can be seen in Fig. 7.7, is less than 1.5% for most of the paddles. Thus we assign a systematic uncertainty of 1.5% to the TOF efficiency and for this channel, we have two particles, so the overall uncertainty due to the TOF efficiency is about 2.1% (1.5% × √2). The overall uncertainty of the acceptance is about 2.3% (2.1% + 1.5%).

![Figure 7.7](image)

**Figure 7.7**: The fluctuation of the normalized TOF occupancy for each paddle. The dashed line is 1.5%, and the uncertainties of most of the paddles are well below 1.5%.

### 7.4 Cuts Related

We applied the energy dependent missing mass cut to identify the spectator. The missing mass resolution shows small angular dependence and it is within 10% of the average resolution. As shown in Fig. 7.8, we varied the missing mass cuts by ±10% and found out that it changed the final cross section by only about 0.5%. Thus, this
angular dependence is small. and 1% systematic uncertainty is assigned to it. The missing momentum cut is $P_x < 0.2 \text{ GeV/c}$, the timing cut is 1 ns, and the z-vertex cut is $|z + 25| < 11 \text{ cm}$. We tightened these cuts by 10% in both data and MC simulation. As can be seen in Fig. 7.9, the overall change in final normalized yield is less than 1%. Fig. 7.10 shows the change of normalized yield due to varying the cuts as a function of photon energy. We assign 1% uncertainty to the missing mass cuts, and 0 - 4% uncertainty to the vertex and missing momentum cuts depending on energy. The overall uncertainties due the cuts are about 1 - 4.1%.
Figure 7.8: Vary the missing mass squared cuts by ±10% for both data and MC simulation. The normalized yield changes within 0.5%.
Figure 7.9: Cuts 1: $P_x < 0.2$ GeV/$c$ and $|z + 25| < 11$ cm; Cuts 2: $P_x < 0.18$ GeV/$c$ and $|z + 25| < 10$ cm; Between the two cuts, both events and acceptance change about 10%, and the normalized yield changes within 1%.
Figure 7.10: The change of normalized yield due to varying missing mass cuts is well below 1% for all the energies (top panel). The change due to the vertex cuts and missing momentum cuts is about 0 - 4%.
7.5 Final State Interaction

The final state interaction (FSI) correction is estimated in Chapter 6. It depends on the transparency measurement for \( d(e, e'p) \) and the total cross sections of \( pn, pp \) and \( \pi^-p \) scattering. The systematic uncertainty of the FSI correction was assigned to be 5% in the approach [81] to account for the uncertainties of all quantities used in the calculation. More recent study found out the FSI correction obtained in this approach is consistent at the level of 10% with that from Laget’s calculation [77]. In this analysis we assign 10% systematic uncertainty to the FSI correction.

7.6 Summary

The systematic uncertainties have been studies in detail for each possible source. The details are listed in Table 7.2 and Table 7.3. The overall systematics are quadrature sum of all the systematics. The final systematic uncertainties are energy dependent and are calculated for each data point. The angular dependence of the systematics was found very small, thus only energy dependent systematic errors are applied.

| Systematic error sources | Errors       |
|--------------------------|--------------|
| Luminosity               | 5%           |
| FSI Correction           | 10%          |
| Acceptance               | 2.3%         |
| Background               | 1 - 6%       |
| Cuts related             | 1 - 4.1%     |
| Overall                  | 11.5 - 13.7% |
Table 7.3: List of the overall systematic errors

| Photon Energy (MeV) | Errors (%) |
|---------------------|------------|
| 1020                | 11.54      |
| 1300                | 11.59      |
| 1500                | 11.65      |
| 1700                | 11.73      |
| 1900                | 11.82      |
| 2100                | 12.03      |
| 2300                | 12.23      |
| 2500                | 12.49      |
| 2700                | 12.67      |
| 2900                | 12.84      |
| 3100                | 13.09      |
| 3300                | 13.28      |
| 3500                | 13.70      |
8

Results

8.1 Differential Cross Sections

The differential cross section was given by

\[ \frac{d\sigma}{d\Omega_{\text{c.m.}}} = \frac{N}{t_G\epsilon} \frac{1}{N_\gamma} \frac{A}{\rho LN_A} \frac{1}{d\Omega_{\text{c.m.}}}, \]  

(8.1)

where \( t_G \) is the transparency to account for the FSI effect (see Chapter 6), \( \epsilon \) is the acceptance, \( N \) is the number of events, \( N_\gamma \) is the integrated photon flux, \( \rho \) is the density of liquid deuterium (0.163 g/cm\(^3\)), \( L \) is the length of the target (24 cm), \( N_A \) is the Avogadro number (6.022 \( \times \) 10\(^{23}\) mol\(^{-1}\)) and \( d\Omega_{\text{c.m.}} \) is the angle bin size in the c.m frame of \( \pi^- p \). Fig. 8.1 shows the typical \( N, t_G \) and \( \epsilon \) distributions at \( E_\gamma = 1.15 \) GeV.

The differential cross sections can be calculated up to \( E_\gamma = 3.56 \) GeV. For photon energy below 2 GeV, the cross sections are extracted every 50 MeV, and since the cross sections drop quickly as photon energy increases, we used larger bins (100 MeV) for photon energy above 2 GeV. The angle bin size is always 5\(^\circ\). The differential cross sections are extracted independently from both high field and low field data, and they
Figure 8.1: The event distribution (top panel), nuclear transparency (middle panel), and acceptance (bottom panel) for $E_\gamma = 1.15$ GeV are in excellent agreement with each other (one example is shown in Fig. 8.2). The ratios of the differential cross sections from low field and high field data at the same kinematic bins are calculated for lower photon energies (below 2 GeV) and higher photon energies (above 2 GeV). As we can see in Fig. 8.3, the mean values of the distributions are 1.009 and 1.002, respectively, which means the average difference of differential cross sections between low field and high field data is below 1%. The
$\gamma n \rightarrow \pi^- p$ at $\sqrt{s} = 1.72$ GeV

Figure 8.2: Differential cross section at $E_\gamma = 1.10$ GeV from high field and low field data compared with world data and theories. The differential cross sections from both high and low field data are in excellent agreement with each other as well as the world data.

standard deviations are around 5% and 7%, respectively, and they are consistent with the overall systematic uncertainties. In the end, the two cross sections are combined together by statistical weighting.

For $E_\gamma$ between 1.05 GeV and 1.75 GeV, the differential cross sections are compared to world data [1, 47, 70, 71, 72, 82], shown in Fig. 8.4 - 8.18, and the cross sections are in good agreement with world data. In addition, several theoretical predictions are plotted as well. As for higher energies, there are few available world data, and the SAID and MAID are not well constrained in this region, thus only the g10 results are shown in Fig. 8.19 - 8.23.

There are three theoretical curves shown in the figures. They are SAID, MAID
Figure 8.3: The distributions of ratios between differential cross sections from low field and high field data for photon energies below 2 GeV (top panel) and above 2 GeV (bottom panel).
and Regge approach. SAID (the Scattering Analysis Interactive Dial-in) [83] is a program providing the most up-to-date partial wave analysis of the scattering data such as nucleon-nucleon and pion-nucleon scattering. The SAID with only minimum model-dependence (such as Born and vector-meson exchange terms and some phenomenological terms) fits all the available scattering data in the database, and provides information such as resonance properties and couplings. The most recent SAID fit for the pion photoproduction is the FA09 version [38] valid up to $E_\gamma = 2.7$ GeV ($\sqrt{s} = 2.4$ GeV). MAID [84] is based on the unitary isobar model with unitarized background and thirteen 4-star resonances below 2 GeV. The most recent version of MAID is MAID07 [85] which fits the world data up to $E_\gamma = 1.6$ GeV ($\sqrt{s} = 2$ GeV). A Regge approach is a general method of continuing the scattering $S$-matrix to the complex energy and momentum plane [86, 87, 88]. It creates singularities in the amplitude (so called Regge poles) which are related to the resonances. So a Regge approach establishes an important connection between the high energy scattering amplitude and the bound states and resonances. The presented Regge model [96] analyzed all the world data on $\pi^+$ and $\pi^-$ photoproduction for $3 < E_\gamma < 8$ GeV to determine the non-resonant amplitude assuming that resonance contributions are negligible in this region. Then the amplitude is used to determine the cross sections and other observables for $E_\gamma < 3$ GeV.

As we can see, the SAID analyses (SP09) [38] agree with our data and world data better than other predictions, for example the MAID07 [85] shown as the red dotted line. The FA09 shows excellent agreement with our data especially in forward angles. The backward agreement is not as good as the forward angles. For the $\pi^-$ channel, FA09 seems to be valid up to $\sqrt{s} = 2$ GeV. Our data will definitely help to extend the SAID fit to much higher energy. We are now working with GWU group to include our data in the partial wave analysis. The very preliminary fit with our data is shown in Fig. 8.24. The Regge model shows good agreement at higher
energies $E_\gamma > 2$ GeV and very forward angles. However, it predicts significantly smaller cross sections than the world data at lower energies. Considering the model assumes that the resonance contribution is negligible, this discrepancy indicates the possible resonance contributions to the cross sections such as $N(2190)$, $N(2220)$ and $N(2250)$ which are the 4-star resonances in this energy region.

Figure 8.4: Differential cross section at $E_\gamma = 1.05$ GeV compared with world data.
\[ \gamma n \rightarrow \pi^- p \text{ at } \sqrt{s} = 1.72 \text{ GeV} \]

**Figure 8.5:** Differential cross section at \( E_\gamma = 1.10 \text{ GeV} \) compared with world data.

\[ \gamma n \rightarrow \pi^- p \text{ at } \sqrt{s} = 1.74 \text{ GeV} \]

**Figure 8.6:** Differential cross section at \( E_\gamma = 1.15 \text{ GeV} \) compared with world data.
\[ \gamma n \rightarrow \pi^- p \text{ at } \sqrt{s} = 1.77 \text{ GeV} \]

**Figure 8.7:** Differential cross section at \( E_\gamma = 1.20 \) GeV compared with world data.

\[ \gamma n \rightarrow \pi^- p \text{ at } \sqrt{s} = 1.80 \text{ GeV} \]

**Figure 8.8:** Differential cross section at \( E_\gamma = 1.25 \) GeV compared with world data.
Figure 8.9: Differential cross section at $E_\gamma = 1.30$ GeV compared with world data.

Figure 8.10: Differential cross section at $E_\gamma = 1.35$ GeV compared with world data.
Figure 8.11: Differential cross section at $E_\gamma = 1.40$ GeV compared with world data.

Figure 8.12: Differential cross section at $E_\gamma = 1.45$ GeV compared with world data.
\[ \gamma n \rightarrow \pi^- p \text{ at } \sqrt{s} = 1.92 \text{ GeV} \]

\begin{figure}
\centering
\includegraphics[width=\textwidth]{figure8_13}
\caption{Differential cross section at \( E_\gamma = 1.50 \) GeV compared with world data.}
\end{figure}

\begin{figure}
\centering
\includegraphics[width=\textwidth]{figure8_14}
\caption{Differential cross section at \( E_\gamma = 1.55 \) GeV compared with world data.}
\end{figure}
Figure 8.15: Differential cross section at $E_\gamma = 1.60$ GeV compared with world data.

Figure 8.16: Differential cross section at $E_\gamma = 1.65$ GeV compared with world data.
Figure 8.17: Differential cross section at $E_\gamma = 1.70$ GeV compared with world data.

Figure 8.18: Differential cross section at $E_\gamma = 1.75$ GeV compared with world data.
Figure 8.19: Differential cross section at $E_\gamma = 2.00$ GeV.
\( \gamma n \rightarrow \pi^- p \) at \( \sqrt{s} = 2.28 \text{ GeV} \)

\( \gamma n \rightarrow \pi^- p \) at \( \sqrt{s} = 2.44 \text{ GeV} \)

**Figure 8.20:** Differential cross section at \( E_\gamma = 2.30 \text{ GeV} \).

\( \gamma n \rightarrow \pi^- p \) at \( \sqrt{s} = 2.44 \text{ GeV} \)

**Figure 8.21:** Differential cross section at \( E_\gamma = 2.70 \text{ GeV} \).
Figure 8.22: Differential cross section at $E_\gamma = 3.00$ GeV.

Figure 8.23: Differential cross section at $E_\gamma = 3.40$ GeV.
Figure 8.24: Differential cross sections compared with various theoretical predictions. The blue solid circles are the g10 data, and the blue solid lines are the latest SAID solution which included the g10 data and g1c $\pi^+$ data [38] and Spring-8 $\pi^0$ data [89](Figure provided by I. Strakovsky).
8.2 Scaled Differential Cross Sections

The scaled differential cross section is defined as

\[ s^7 \frac{d\sigma}{dt} = s^7 \frac{d\sigma}{d\Omega_{\text{c.m.}}} \frac{d\Omega_{\text{c.m.}}}{dt} = s^7 \frac{d\sigma}{d\Omega_{\text{c.m.}}} \frac{\pi}{E_{\gamma_{\text{c.m.}}} P_{\pi^-_{\text{c.m.}}}}, \]  

(8.2)

where \( E_{\gamma_{\text{c.m.}}} \) and \( P_{\pi^-_{\text{c.m.}}} \) are photon energy and \( \pi^- \) momentum, respectively, in the c.m frame of \( \pi^-p \), \( s \) and \( t \) (Mandelstam variables) are invariant mass squared and momentum transfer squared, respectively.

The scaled differential cross sections, \( s^7 \frac{d\sigma}{dt} \), are studied for different angles in center-of-mass frame for each 5° bin from 45° to 110° as a function of \( \sqrt{s} \) and compared with world data \([1, 71, 72, 82, 90, 91, 92, 93, 94, 95]\). As shown in Fig. 8.25 - 8.38, the g10 results are compared with world data and Regge model. The error bar on g10 data includes both statistical and systematic uncertainties. At lower energies, where there are overlap regions with world data, our data have excellent agreement with previous world data. There are some significant differences between our data and Bonn data \([47]\). These differences were also observed in Hall A E94-104 experiment \([1]\). We note that the Bonn data \([47]\) were not included in the recent partial wave analysis (FA06) \([37]\). In the angles such as \( \theta_{\text{c.m.}} = 50°, 70°, 90° \) which the Hall A experiment has focused on, we have excellent agreement in all the overlapped angles and energies. As can been seen, the scaling does not start yet at the very forward angles where higher energy is required to reach the onset. The cross sections do seem to approach the scaling for \( \theta_{\text{c.m.}} \geq 70° \). Though, we do need higher energy data to further confirm the scaling at angles other than 70° and 90°. The Regge model seen in many angles such as in Fig. 8.31 and 8.32, shows agreement at both extremely low energies and extremely high energies while completely missing the enhancement. The Regge model didn’t include resonances and this indicates that the resonances may contribute to the observed enhancement.
Figure 8.25: The scaled differential cross section, $s^7 \frac{d\sigma}{dt}$ versus center-of-mass energy for the $\gamma n \rightarrow \pi^- p$ at $\theta_{c.m.} = 45^\circ$.

Figure 8.26: The scaled differential cross section, $s^7 \frac{d\sigma}{dt}$ versus center-of-mass energy for the $\gamma n \rightarrow \pi^- p$ at $\theta_{c.m.} = 50^\circ$. 
\( \gamma n \rightarrow \pi^- p \) with \( \theta_{\text{cm}} \) around 55°.

**Figure 8.27:** The scaled differential cross section, \( s \frac{d\sigma}{dt} \) versus center-of-mass energy for the \( \gamma n \rightarrow \pi^- p \) at \( \theta_{\text{c.m.}} = 55° \).

\( \gamma n \rightarrow \pi^- p \) with \( \theta_{\text{cm}} \) around 60°.

**Figure 8.28:** The scaled differential cross section, \( s \frac{d\sigma}{dt} \) versus center-of-mass energy for the \( \gamma n \rightarrow \pi^- p \) at \( \theta_{\text{c.m.}} = 60° \).
Figure 8.29: The scaled differential cross section, $s^7 \frac{d\sigma}{dt}$ versus center-of-mass energy for the $\gamma n \rightarrow \pi^- p$ at $\theta_{\text{c.m.}} = 65^\circ$.

Figure 8.30: The scaled differential cross section, $s^7 \frac{d\sigma}{dt}$ versus center-of-mass energy for the $\gamma n \rightarrow \pi^- p$ at $\theta_{\text{c.m.}} = 70^\circ$. 

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\[ \gamma n \rightarrow \pi^- p \text{ with } \theta_{\text{c.m.}} \text{ around } 75^\circ \]

**Figure 8.31:** The scaled differential cross section, \( s \frac{d\sigma}{dt} \) versus center-of-mass energy for the \( \gamma n \rightarrow \pi^- p \) at \( \theta_{\text{c.m.}} = 75^\circ \).

\[ \gamma n \rightarrow \pi^- p \text{ with } \theta_{\text{c.m.}} \text{ around } 80^\circ \]

**Figure 8.32:** The scaled differential cross section, \( s \frac{d\sigma}{dt} \) versus center-of-mass energy for the \( \gamma n \rightarrow \pi^- p \) at \( \theta_{\text{c.m.}} = 80^\circ \).
$\gamma n \rightarrow \pi^- p$ with $\theta_{\text{cm}}$ around 85°

**Figure 8.33:** The scaled differential cross section, $s^7 \frac{d\sigma}{dt}$ versus center-of-mass energy for the $\gamma n \rightarrow \pi^- p$ at $\theta_{\text{c.m.}} = 85^\circ$.

$\gamma n \rightarrow \pi^- p$ with $\theta_{\text{cm}}$ around 90°

**Figure 8.34:** The scaled differential cross section, $s^7 \frac{d\sigma}{dt}$ versus center-of-mass energy for the $\gamma n \rightarrow \pi^- p$ at $\theta_{\text{c.m.}} = 90^\circ$.  

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\( \gamma \, n \rightarrow \pi^- \, p \) with \( \theta_{\text{cm}} \approx 95^\circ \)

**Figure 8.35:** The scaled differential cross section, \( s \frac{d\sigma}{dt} \) versus center-of-mass energy for the \( \gamma n \rightarrow \pi^- p \) at \( \theta_{\text{c.m.}} = 95^\circ \).

\( \gamma \, n \rightarrow \pi^- \, p \) with \( \theta_{\text{cm}} \approx 100^\circ \)

**Figure 8.36:** The scaled differential cross section, \( s \frac{d\sigma}{dt} \) versus center-of-mass energy for the \( \gamma n \rightarrow \pi^- p \) at \( \theta_{\text{c.m.}} = 100^\circ \).
Figure 8.37: The scaled differential cross section, $s^7 \frac{d\sigma}{dt}$ versus center-of-mass energy for the $\gamma n \rightarrow \pi^- p$ at $\theta_{\text{cm}} = 105^\circ$.

Figure 8.38: The scaled differential cross section, $s^7 \frac{d\sigma}{dt}$ versus center-of-mass energy for the $\gamma n \rightarrow \pi^- p$ at $\theta_{\text{cm}} = 110^\circ$.  

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8.3 Discussion

Fig. 8.39 shows the scaled differential cross section $s^7 \frac{d\sigma}{dt}$ as a function of $\sqrt{s}$ for $\theta_{c.m.} = 90^\circ$ for three different channels. The results from this experiment are shown in the middle panel as red solid circles with statistical uncertainties, and the systematic uncertainty is shown as a band. The error bars for E94-104 [1] include both the statistical and systematic uncertainties, while only statistical uncertainties are shown for the $\pi^0$ data [37] and the $\pi^+$ data [38]. All other world data are collected from Refs. [5, 97]. Our data are consistent with the E94-104 results [1] within experimental uncertainties. With fine photon energy bins and high statistical precision, our data confirm a broad enhancement around $\sqrt{s}$ of 2.1 GeV in the scaled differential cross section. Our data also confirm a marked fall-off of the differential cross section in a narrow energy window of about 400 MeV above this enhancement and the onset of the CCR scaling for $\sqrt{s}$ around 2.8 GeV as suggested by an earlier Jefferson Lab experiment [1] (shown as green solid squares). Similar behavior has been seen in the recent CLAS g1c $\pi^+$ photoproduction data [38] (magenta open squares in Fig. 8.39).

While this fall-off may be taken as a signature for the transition from nucleon-meson degrees of freedom to quark-gluon degrees of freedom, theoretical studies in this region are needed to confirm this speculation. Also shown are the results of the SAID SP09 partial wave analysis [38] (blue), the MAID07 model [85] (cyan), and the prediction from a Regge approach [96] (black).

In the Regge calculation, no baryon resonances in this energy region were included. And the results didn’t predict the enhancement seen in our data. Thus the deviation is speculated to be due to baryon resonances [96].

While the SAID SP09 fit has been greatly improved by the CLAS $\pi^0$ [37], the $\pi^+$ data [38] and the Hall A $\pi^-$ data [1, 2], it does not give as good a description of our data near the peak of the enhancement. Further, it fails to constrain the $\pi^-$
Figure 8.39: Scaled differential cross section $s^7 d\sigma/dt$ as a function of $\sqrt{s}$ for $\theta_{c.m.} = 90^\circ$ for three different channels. The upper panel is for the $\gamma p \rightarrow \pi^+ n$ process, the middle panel is for the $\gamma n \rightarrow \pi^- p$ process, and the lower panel is for the $\gamma p \rightarrow \pi^0 p$ process. The green solid squares are results from Ref. [1] and the results from this experiment are shown as red solid circles. Results from Dugger et al. [37] on neutral pion production are shown as magenta solid squares. The magenta open squares are recent CLAS data on $\pi^+$ production [38]. The SAID SP09 results [38] are shown as the blue curves in all three panels. The prediction from a Regge approach [96] is shown in the top and middle panels by black curves. The black open circles are the world data collected from Refs. [5, 97].
channel and does not describe our data well above $\sqrt{s}$ of 2.3 GeV. The precision data presented here will help to further constrain the SAID fit and will allow for a determination of the corresponding neutron electromagnetic parameters for resonances classified as 4-star by the PDG [31]. These studies will be reported in a future publication.

Fig. 8.40 shows the scaled differential cross section $s^7 \frac{d\sigma}{dt}$ as a function of $\sqrt{s}$ for $\theta_{c.m.} = 50^\circ$ to $115^\circ$ with an angular bin size of $5^\circ$ for the $\gamma n \rightarrow \pi^- p$ process. As in Fig. 8.39, the systematic uncertainties are shown as bands in Fig. 8.40. The blue arrows indicate the location of $\sqrt{s}$ corresponding to a pion transverse momentum ($P_T$) of 1.1 GeV/c. This $P_T$ value was suggested to govern the scaling onset by Refs. [11, 13]. We note the large discrepancy between our results and those from Ref. [47] at $\theta_{c.m.} = 75^\circ$ and $95^\circ$. We also note that the SAID fits [37, 38] did not include data from Ref. [47]. An angular-dependent feature in the scaled differential cross section is clearly seen in our data. The aforementioned broad enhancement around a $\sqrt{s}$ value of 2.1 GeV at $\theta_{c.m.} = 90^\circ$ seems to shift as a function of $\theta_{c.m.}$ from $\sqrt{s}$ of 1.80 GeV at $50^\circ$ to 2.45 GeV at $105^\circ$ as shown by the red dotted lines. It is not clear whether this enhancement dies off for $\theta_{c.m.} > 105^\circ$ or whether it shifts to further higher energies. We’ve studied very carefully to see if the angular dependence of the enhancement is introduced by the FSI correction or the $s^7$ factor. We found the same angular dependence with or without the FSI correction. We’ve also done a simple modeling of the enhancement with Breit-Wigner distributions and the $s^7$ factor seems to shift the peak by a few MeV, but compared to the observed angular dependence, this shift due to $s^7$ factor is negligible. The blue dashed lines indicate the locations of the nucleon resonances around 1.2 GeV and 1.5 GeV which, as expected, do not change with $\theta_{c.m.}$. However, such an angular dependent enhancement is not seen in the $\pi^+$ (see Fig. 8.41) and $\pi^0$ channels from the proton. The SAID FA09 prediction is also shown in Fig. 8.40 and Fig. 8.41 (blue curve), and it does show
Figure 8.40: Scaled differential cross section $s^7 \frac{d \sigma}{dt}$ for $\gamma n \rightarrow \pi^- p$ as a function of $\sqrt{s}$ for $\theta_{c.m.} = 50^\circ$ to $115^\circ$. The arrows indicate the location of $\sqrt{s}$ corresponding to a transverse momentum value of 1.1 GeV/c. The green solid squares are results from Ref. [1]. The results from this experiment are shown as red solid circles. The black open circles and open squares are the world data collected from Refs. [5, 97] and [47], respectively. Errors on the data from CLAS are the quadratic sum of the statistical and systematic uncertainties. The SAID SP09 results [38] are shown as the blue curves. The blue dashed lines indicate the known resonances, and the red dotted lines illustrate the angular dependent feature of the broad enhancement structure discussed in the text.
\( \gamma p \rightarrow \pi^+ n \)

Figure 8.41: Scaled differential cross section \( s^7 \frac{d\sigma}{dt}(10^7 \text{GeV}^2) \) for \( \gamma p \rightarrow \pi^+ n \) as a function of \( \sqrt{s} \) for \( \theta_{\text{c.m.}} = 50^\circ \) to 115\(^\circ\). The arrows indicate the location of \( \sqrt{s} \) corresponding to a transverse momentum value of 1.1 GeV/c. The green solid squares are results from Ref. [1]. The magenta open squares are recent CLAS data on \( \pi^+ \) production [38]. The black open circles and open squares are the world data collected from Refs. [5, 97] and [47], respectively. Errors on the data from CLAS are the quadratic sum of the statistical and systematic uncertainties. The SAID SP09 results [38] are shown as the blue curves.
an enhancement around $\sqrt{s}$ of 2.2 GeV which is not angular dependent. Thus, the angular dependence of the enhancement is a unique feature of the $\pi^-$ data.

The observed angular dependent enhancement structure in the $\pi^-$ channel could be due to some unknown resonances which couple differently to the neutron channel than to the proton channel. Polarization data from all three channels and partial wave analysis are necessary in order to understand the nature of this enhancement and its angular dependence in the $\pi^-$ channel.

Although the nature of the enhancement and its angular dependence remains unknown, there has been new development in theory trying to understand it. Recently, Laget [98] proposed that introducing the coupling to the $\rho^0N$ to the pion photoproduction may explain the data well at large $P_T$ for the energy range of $3 < E_\gamma < 10$ GeV. At low energies, the amplitude of the exclusive reaction includes all the possible intermediate states, while at high energies, the contribution of $\rho^0N$ dominates. In addition, the $\pi N$ scattering, although 10 times smaller than the $\rho^0N$, plays an important role through destructive interference with the pole amplitude. The two dominant amplitudes are shown in Fig. 8.42. The calculation with the two amplitudes were compared with experimental data. As shown in Fig. 8.43, the model reproduces the magnitude of the enhancement very well below $\sqrt{s} = 3$ GeV as well as the $s^{-7}$ scaling behavior following the enhancement.
Figure 8.43: The scaled cross sections of the $p(\gamma, \pi^+)n$ (top) and the $n(\gamma, \pi^-)p$ (bottom) reactions at $\theta_\pi = 90^\circ$. Basic Regge pole model: black dashed lines. Elastic pion rescattering cut included: blue dotted lines. Inelastic $\rho^0 N$ also included: red full lines. Filled circles: Ref. [1]. Empty circle: this experiment. Filled triangle: Ref. [5]. Figure is from Ref. [98].
8.4 Summary and Outlook

In this analysis, we have extracted the differential cross section for $\gamma n \rightarrow \pi^- p$, and studied the scaled differential cross sections at different c.m. angles with high statistics. The data can be downloaded from [99]. There is good agreement between high field and low field data, which gives confidence that the data analysis procedures are correct.

There are plenty of published data at very low incident photon energies below 1 GeV for this channel. However, there are very limited published data at low energy between 1 GeV and 2 GeV, and almost no data available between 2 GeV and 3 GeV. The CLAS g10 data give, for the first time, the opportunity to study the detailed energy dependence and angular dependence of the $\pi^-$ photoproduction differential cross section on a deuteron target over the wide range of c.m. angles at the incident photon energy between 1 GeV and 3.58 GeV.

The $\pi^-$ photoproduction cross sections decrease rapidly as the incident photon energies increase. The angular dependence changes as well, but the asymmetry of the cross sections between the forward and backward angles are observed throughout the whole photon energy range. The scaled differential cross sections show signs of scaling behavior at 70° and 90°, but not at 50°. The behavior of the scaled differential cross sections change dramatically between 50° and 70°, and it is interesting to understand such big a change within this small angular range.

Our data suggest a possible signature for the transition to the CCR scaling region, in the form of a fall-off of the scaled cross section over a narrow energy range. An angular dependent enhancement in the scaled differential cross section has been seen for the first time in our data, which is different from that of the $\gamma p \rightarrow \pi^+ n$ process, and it is also different from the latest SAID prediction.

Inspired by the $\pi^-$ results, a new experiment E08-003 [100] was proposed and
approved to measure the $\gamma p \rightarrow \pi^+ n$ differential cross sections with fine energy and angular bins up to $E_\gamma = 5.6$ GeV. This experiment E08-003 was carried out during the g12 run period with a beam energy of $E_e = 5.715$ GeV and a 40 cm liquid hydrogen target placed 90 cm upstream. A single-charged trigger has to be used for this channel with a beam current of 24 nA. During the g12 run period, we completed about 30 production runs with about $10^9$ triggers during about 50 hours of run time (about 40 hours live time). The data are now under calibration and reconstruction.

The DOE approved JLab 12 GeV upgrade will provide an excellent test platform in the near future for the single pion photoproduction. With $\sqrt{s}$ up to 4.6 GeV, one can further study in detail the possible substructure of the scaling, and oscillation around the CCR prediction.
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Biography

Wei Chen (陈伟) was born in Jurong (江苏句容), China in April, 1979.

Education:

• 2004/08 - 2010/03 Ph.D in Nuclear Physics, Duke University, USA.
• 2001/08 - 2004/07 M.S. in Theoretical Physics, Nanjing University, China.
• 1997/09 - 2001/07 B.A. in Theoretical Physics, Sichuan University, China.

Publication:

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