Determination of the Azimuthal Asymmetry of Deuteron Photodisintegration in the Energy Region $E_\gamma = 1.1 - 2.3$ GeV

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Dedication

This dissertation is dedicated to the memory of Dr. Barry L. Berman

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Abstract

Determination of the Azimuthal Asymmetry of Deuteron Photodisintegration in the Energy Region $E_\gamma = 1.1 - 2.3$ GeV

Deuteron photodisintegration is a benchmark process for the investigation of the role of quarks and gluons in nuclei. Existing theoretical models of this process describe the available cross sections with the same degree of success. Therefore, spin-dependent observables are crucial for a better understanding of the underlying dynamical mechanisms. However, data on the induced polarization (P_y), along with the polarization transfers ($C_{x'}$ and $C_{z'}$), have been shown to be insensitive to differences between theoretical models. On the other hand, the beam-spin asymmetry Σ is predicted to have a large sensitivity and is expected to help in identifying the energy at which the transition from the hadronic to the quark-gluon picture of the deuteron takes place.

Here, the work done to determine the experimental values of the beam-spin asymmetry in deuteron photodisintegration for photon energies between 1.1 – 2.3 GeV is presented. The data were taken with the CLAS at the Thomas Jefferson National Accelerator Facility during the g13 experiment. Photons with linear polarization of $\sim 80\%$ were produced using the coherent bremsstrahlung facility in Hall B. The work done by the author to calibrate a specific detector system, select deuteron photodisintegration events, study the degree of photon polarization, and finally determine the azimuthal asymmetry and any systematic uncertainties associated with it, is comprehensively explained. This work shows that the collected data provide the kinematic coverage and statistics to test the available QCD-based models. The results of this study show that the available theoretical models in their current state do not adequately predict the azimuthal asymmetry in the energy region 1.1 – 2.3 GeV.

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Chapter 1

Introduction

1.1 The Standard Model

Nuclear physics and particle physics aim to identify the fundamental particles and forces of the universe, and to determine the structure and characteristics of matter. The standard model of particle physics [1] is a theory concerning three of the four fundamental forces: the electromagnetic, the weak, and the strong force. The first step in the development of the standard model was done in 1960, when Glashow found a way to combine the electromagnetic and weak interactions by including a short-range neutral current, the Z^0 [2]. Quantum electrodynamics (QED) provides a theoretical basis for the interaction between electrically charged particles. The weak interaction is responsible for some nuclear phenomena such as beta decay. Finally, the strong interaction is responsible, on one hand, for holding the protons and neutrons together to form the nucleus of an atom, and on the other hand, for holding quarks together to form hadrons, such as protons and neutrons. The strong interaction is described in the framework of quantum chromodynamics (QCD).

The current formulation of the standard model was finalized in the mid-1970s upon experimental confirmation of the existence of quarks. Since then, the standard model has had great success in explaining a wide variety of experimental results. A brief historical overview that resulted in the development of the theories of fundamental interactions and the standard model is provided below.

1.1.1 The classical period

The modern era of elementary particle physics started with the discovery of the electron by Thomson in 1897 [3]. In his experiment, *cathode rays* were emitted by a hot filament and were then deflected in the magnetic field of permanent magnets. This experiment suggested that these *rays* were in fact a stream of particles that carried a negative electric charge. Later on, Thomson was able to measure the charge-to-mass ratio of these particles by passing them through electric and magnetic fields applied perpendicular to

each other, and by adjusting these until the deflection of the *rays* was zero. He called these particles *corpuscles*, and their charge the *electron*, with the latter name being applied to the particles themselves later on.

After Thomson's discovery of the electron and his correct prediction that electrons were constituents of atoms, physicists were puzzled by the neutral charge of atoms, as well as by the fact that atoms were much heavier than electrons. Physicists were also eager to explain how a compensating positive charge was distributed in an atom. Rutherford was able to answer these questions by scattering alpha particles on a thin sheet of gold (the famous Rutherford scattering experiment) [4]. His experiment determined that most of the atomic mass and the positive charge was concentrated in a tiny central part of the atom, called the nucleus. The nucleus of the hydrogen atom was then given the name *proton*.

In 1914, Niels Bohr, using his proposed model of the atom, was able to calculate the spectrum of hydrogen [5]. His model of an electron orbiting a proton proved inaccurate when he found out that the atoms weighed more than the charge of their nucleus suggested (the nuclei of hydrogen isotopes - protium, deuterium, tritium - carry the same charge, but not the same mass). This was resolved in 1932, when Chadwick discovered the neutron [6]. This discovery marked the end of the classical period of elementary particle physics.

1.1.2 Hadrons and leptons

The findings of the classical period raised the question of what holds the nucleus together. Yukawa suggested in 1934 that there was a force that overcomes the Coulomb repulsion between the protons, but since we could not experience it in the macroscopic scale, it must be acting only at very short distances [7]. The nearly constant density of nuclear matter allowed Yukawa to estimate the range of the nuclear force between pairs of nucleons to be about 1.4×10^{-15} m. Using the Heisenberg uncertainty principle, which states that $\Delta E \Delta t \geq \hbar$, one can calculate that the lifetime of a virtual particle is approximately $\hbar/mc^2$ and its range is $\hbar/mc$. From this, Yukawa was able to estimate that the mass of the mediator of the new force should be approximately $140 \text{ MeV}/c^2$, which is about 300 times the mass of the electron and about one sixth the mass of the proton. Yukawa's particle was then called a *meson*, which means middle-weight. Since the nuclear force between neutron and proton, proton and proton, and neutron and neutron, seems to be the same, there should

be three different kinds of mesons, a positively charged, a negatively charged, and a neutral, as shown in Fig. 1.1. Ten years after Yukawa's hypothesis, a group at Bristol University [8]

Figure 1.1. Nuclear force mediated by a neutral, a negative, and a positive pion respectively.

discovered, by studying cosmic rays, that there were actually two middle-weight particles, the pion (π) and the muon (μ), with the pion being the strong-force mediator as proposed by Yukawa. Later on it was found that the muon did not fall within the meson category, but is instead in the lepton category (light-weight), along with the electron.

During the same period as the discovery of antiparticles that arose from the union of special relativity and quantum mechanics came the discovery of neutrinos. The study of the beta decay $n \rightarrow p + e^- + \dots$, a process that results in a nucleon and the emission of an electron by the decay of another nucleon, suggested the existence of another particle. Physicists noticed that during this decay, energy was not conserved if one considers that p and e^- are the only decay products. This implied that the new particle would be responsible for the missing energy, and it had to be electrically neutral since the charge must be conserved. Later on, physicists discovered another neutrino, different from the one that was seen in beta decays. This led to the assignment of a quantum number to each of the known particles called the lepton number, which is conserved in decays. Neutrinos (and antineutrinos) were then categorized in the lepton category along with the electron, the muon, and their antiparticles (see Table 1.1) [9].

By the beginning of 1960, the number of particles observed had grown greatly, and hadron physics could be described as chaotic. The published cloud-chamber photograph by Rochester and Butler in 1947 [10] initiated the discoveries of new particles, with the

Table 1.1. The Lepton Family, 1962-1976

	Lepton number	Electron number	Muon number
Leptons			
e^-	1	1	0
ν_e	1	1	0
μ^-	1	0	1
ν_μ	1	0	1
Antileptons			
e^+	-1	-1	0
$\bar{\nu}_e$	-1	-1	0
μ^+	-1	0	-1
$\bar{\nu}_\mu$	-1	0	-1

first being the kaons (K^0 , K^+) and the Λ hyperon. The construction of the Brookhaven Cosmotron, the first modern particle accelerator, provided physicists with the opportunity to produce these new particles and study their properties. The experiments performed at Brookhaven showed that these particles are produced on a time scale of about 10^{-23} seconds, but they decay very slowly, with a lifetime of approximately 10^{-10} seconds. This fact, according to the Heisenberg uncertainty principle ($\Delta E \Delta t \geq \hbar$, $\Delta x \Delta p \geq \hbar$), suggested that they are produced by the strong interaction, but they decay through the weak interaction. These particles are referred to as *strange* because of this behavior.

The need for a method to classify these particles was pressing. The strongly interacting particles were divided into two families – mesons and baryons – and each one of these was distinguished by charge, mass, strangeness, and baryon quantum numbers. The baryon number assigned to all baryons is conserved in any physical process, whereas the strangeness number is conserved only in strong and electromagnetic interactions but not in weak interactions. Having this in mind, Gell-Mann and Ne’eman independently [11, 12, 13] introduced the so-called Eightfold Way to categorize elementary particles.

The Eightfold Way is a scheme of arranging the hadrons according to their symmetry properties, following specific geometrical patterns. For example, the eight lightest baryons with spin-1/2 were organized in a hexagonal array (called the baryon octet) shown in Fig. 1.2. Following the same principle, the eight lightest spin-0 mesons were organized in a different hexagonal array, called the pseudoscalar-meson octet. The vector-meson octet

corresponds to the hexagonal array of spin-1 mesons, and the baryon decuplet corresponds to a triangular array of spin-3/2 baryons, as shown in Figs. 1.3–1.5. Notice that particles of like charge lie along the downward-sloping diagonal lines, and the horizontal lines connect the particles that have the same strangeness quantum number.

Figure 1.2. The Eightfold Way classification scheme for spin-1/2 baryons.

When Gell-Mann was arranging the 3/2-spin particles into a decuplet, he realized that there were only nine particles known at that time that could be arranged in this array. The location at the very bottom of the decuplet suggested the existence of a particle with charge -1 and strangeness -3 . Gell-Mann was able to predict its mass and lifetime, providing experimentalists with the knowledge needed to search for this particle. Soon thereafter, a particle that carries these properties, the Ω^- , was discovered [14].

Figure 1.3. The Eightfold Way classification scheme for pseudoscalar mesons.

Figure 1.4. The Eightfold Way classification scheme for vector mesons.

Figure 1.5. The Eightfold Way classification scheme for spin-3/2 baryons.

1.1.3 The quark model

The reason for the success of the Eightfold Way was not obvious to physicists. In fact, the reason why one could separate the hadrons into such groups was puzzling. However, Gell-Mann and Zweig [15] independently proposed that all hadrons are composed of even more elementary constituents, the quarks. They proposed that quarks come in three different types (or flavors) – up, down, and strange – and have a charge that is a multiple of $\frac{1}{3}e$. They also proposed that each baryon is composed of three quarks, and each meson is composed of a quark and an antiquark.

The quark model of Gell-Mann and Zweig, however, suffered from two serious drawbacks: (1) no one had ever seen a free quark and (2) the model did not respect the Pauli Principle, with the Δ^{++} baryon, made up by three up quarks, being the best example for Pauli principle violation. It was expected that if quarks were the constituents of protons, one could easily produce free quarks in the laboratory. An attempt to make the quark model survive ended up with the notion of quark confinement. Quarks are confined within mesons and baryons such that it is not possible to free them. In addition to this, Greenberg [16] suggested that quarks not only come in three flavors but also in three “colors” (red, blue, and green).

Experiments carried out at the Stanford Linear Accelerator Center (SLAC) using high-energy positron and electron beams incident on a proton target showed that the charge of the proton is concentrated in three small regions within the proton. This gave strong support for the quark model, but to avoid any compromise of the quark model these “lumps”

were called partons. Finally, the discovery of a new meson, the J/Ψ meson [17], provided further support for the quark model. The large masses and long lifetime of this meson could not be explained using the then-common knowledge about mesons. This resulted in the proposal of a new quark, the charm quark, which motivated a search for even more quarks. The dominant hypothesis was that, since six leptons already had been discovered ($e, \nu_e, \mu, \nu_\mu, \tau, \nu_\tau$), six quark flavors must also exist. The reasoning behind this hypothesis was introduced many years earlier by Bjorken and Glashow [18], and even more compelling reasoning was given a couple of years later by Glashow, Iliopoulos, and Maiani [19]. The list of quarks was completed in 1995, when the top quark was discovered [20, 21, 22].

1.1.4 Bosons

In the standard model, particles interact through a mediator, referred to as gauge boson. The term boson refers to the fact that the mediator has integer spin and obeys Bose-Einstein statistics. Bosons can occupy the same quantum state at the same time. Each interaction has one or more gauge bosons associated with it. The massless photon is the gauge boson for the electromagnetic interaction, as described by QED. The gauge bosons mediating the weak interaction are the W^+, W^- , and Z , with charges +1, -1, and 0, respectively. The strong gauge bosons are responsible for the strong force that acts between quarks, and are called gluons. Quarks and gluons are the fundamental degrees of freedom in QCD. The nuclear force at low energies, on the other hand, can be very well understood by the exchange of light mesons, notably pions. Figure 1.6 shows the elementary particles of the standard model.

1.1.5 Quantum chromodynamics

Quantum chromodynamics is the accepted theory of the strong interaction in the standard model. The strong force occurs between particles with color charge, quarks and gluons. The color charge is an abstract property of these elementary particles that characterizes their role in strong interactions. As mentioned before, the color charge of the quarks can either be red, green, or blue, whereas an antiquark's anti-color charge can be either anti-red, anti-green, or anti-blue. Gluons have mixed combinations of color and anti-color

Fermions matter constituents spin = 1/2, 3/2, 5/2, ...

Leptons <small>spin = 1/2</small>			Quarks <small>spin = 1/2</small>		
Flavor	Mass GeV/c ²	Electric charge	Flavor	Approx. Mass GeV/c ²	Electric charge
ν_L lightest neutrino*	(0-0.13)×10 ⁻⁹	0	u up	0.002	2/3
e electron	0.000511	-1	d down	0.005	-1/3
ν_M middle neutrino*	(0.009-0.13)×10 ⁻⁹	0	c charm	1.3	2/3
μ muon	0.106	-1	s strange	0.1	-1/3
ν_H heaviest neutrino*	(0.04-0.14)×10 ⁻⁹	0	t top	173	2/3
τ tau	1.777	-1	b bottom	4.2	-1/3

Bosons force carriers spin = 0, 1, 2, ...

Unified Electroweak <small>spin = 1</small>			Strong (color) <small>spin = 1</small>		
Name	Mass GeV/c ²	Electric charge	Name	Mass GeV/c ²	Electric charge
γ photon	0	0	g gluon	0	0
W⁻	80.39	-1			
W⁺	80.39	+1			
W bosons					
Z⁰	91.188	0			
Z boson					

Figure 1.6. The elementary particles of the standard model and their characteristics [23].

charge, and are therefore self-interacting. This is illustrated in the reaction $q \rightarrow q + g$, where the color charge of the final-state quark may be different from the color charge of the quark in the initial state. Since the color charge is conserved at the vertex, the gluon must carry the color charge of the initial-state quark and the anti-color charge of the final-state quark.

The fact that the gluon has color charge results in the so called *quark confinement*. As two quarks separate, the gluon field between them forms a narrow tube of color charge. This tube has constant radius and energy density, and so the energy cost of separating color charges grows proportionally with the separation. At large enough distances, it is more energetically favorable for a new quark-antiquark pair to spontaneously appear, rather than to allow the tube to extend further. Quark confinement is exhibited in particle accelerator collisions by a process called hadronization [24] (see Fig. 1.7). For this reason, in QCD the combination of quarks that make up the baryons and mesons need to be colorless. For example, a quark can combine with an antiquark with the opposite color charge to form a meson of neutral color.

Figure 1.7. An example of hadronization due to quark confinement [25]. As energy is applied to the $c\bar{c}$ bound state, the quark-antiquark pair begins to separate. At a specific distance, it is more energetically favorable for a $d\bar{d}$ pair to be produced in the gluonic field between the $c\bar{c}$ pair. At all times the color charge of these combinations is conserved.

The strength of the strong interaction is governed in the QCD theory by the dimensionless coupling constant, α_s , also referred to as the running coupling constant. Unlike the coupling constant of QED, α_s depends strongly on the separation of the color charges. Specifically, α_s decreases with decreasing distance, such that at infinitely small distances the quarks behave as if they were free [26]. This property of QCD is referred to as *asymptotic freedom*, and as a result, it allows QCD predictions to be calculated perturbatively

at high energies. This is because the calculation of the infinite series of QCD interaction terms converges quickly due to the smallness of the coupling constant. On the other hand, the coupling strength at low energies is too large for perturbative calculations to converge, so non-perturbative methods must be employed to evaluate the strong interaction [24].

The study of QCD therefore splits naturally into the high-energy, very-short-distance regime, where perturbation theory is applicable, and the low-energy region of confinement in the bound states. The transition region, or the medium-energy region, poses a challenge for both theorists and experimentalists alike, while holding the key to our understanding of how QCD manifests itself in the non-perturbative regime.

The work presented in this document aims to study the medium-energy region, where the transition between the conventional mesonic picture of nuclear interactions to the picture where QCD degrees of freedom are dominant takes place. Chapter 2 introduces the framework used to study this transition region. Chapters 3 and 4 present the experimental facility and the work done to collect the data for the g13 experiment, including the work done by the author to calibrate specific detector system. Chapters 5 and 6 summarize the analysis steps followed by the author to select deuteron photodisintegration events and determine the degree of photon polarization, and Chap. 7 presents the method used by the author for determining the azimuthal asymmetry. Finally, Chap. 8 presents the studies performed by the author to determine the systematic uncertainties of the determined asymmetries and Chap. 9 presents the results obtained from this work.

Chapter 2

The Deuteron

2.1 Overview

This chapter introduces the formalism of deuteron photodisintegration. Specifically, an introduction is presented by encompassing theoretical calculations, previous measurements and a motivation for studying the azimuthal asymmetry of the reaction $\vec{\gamma} + d \rightarrow p + n$. This reaction was first investigated in the early 1930s, and it was used to better understand the structure of the neutron. Specifically, from a study of this reaction, Chadwick and Maurice Goldhaber [27] determined that the neutron mass was slightly larger than that of the hydrogen atom. Their measurement, showing that the neutron was heavier than the total mass of a proton plus an electron, led to two important outcomes: neutrons are not bound states of protons and electrons and should, therefore, be considered fundamental particles, removing any speculation that electrons are present in nuclei. Since then, this reaction has been well studied, both experimentally and theoretically. The deuteron provides us with a unique tool for studying the strong interaction since it is the only two-nucleon bound system, with its structure well understood. At lower beam energies, the deuteron has been used extensively to study properties of its constituents, as well as their bound state. On the other hand, at higher beam energies, the deuteron is especially important for the investigation of the role of quarks and gluons in nuclear reactions.

2.2 Formalism

In deuteron photodisintegration, a photon interacts with the deuteron, which then disintegrates into a proton and a neutron as shown in Fig. 2.1. This process can be described in terms of the three Lorentz-invariant Mandelstam variables [28], s , t , and u ,

$$\begin{aligned} s &= (\mathcal{P}_\gamma + \mathcal{P}_d)^2 = (\mathcal{P}_p + \mathcal{P}_n)^2, \\ t &= (\mathcal{P}_\gamma - \mathcal{P}_p)^2 = (\mathcal{P}_d - \mathcal{P}_n)^2, \\ u &= (\mathcal{P}_\gamma - \mathcal{P}_n)^2 = (\mathcal{P}_d - \mathcal{P}_p)^2, \end{aligned} \tag{2.1}$$

Figure 2.1. Diagram illustrating the kinematics of deuteron photodisintegration in the lab frame (above) and in the center-of-mass frame (below). The unprimed axes are defined so that the x axis is parallel to the Hall B floor and the z axis along the beam line.

where the variables $\mathcal{P}_\gamma$, $\mathcal{P}_d$, $\mathcal{P}_p$, and $\mathcal{P}_n$ in Eq. (2.1) denote the four-momenta of the photon, deuteron, proton, and neutron, respectively. The four-momentum of particle i is defined in terms of the particle's energy E , and its three momentum $\vec{p}$, as usual,

$$\mathcal{P}_i = (E, \quad \vec{p}_i). \quad (2.2)$$

Using four-momentum conservation

$$\mathcal{P}_\gamma + \mathcal{P}_d = \mathcal{P}_p + \mathcal{P}_n \quad (2.3)$$

and the definition of the Lorentz-invariant mass

$$\mathcal{P}_i \cdot \mathcal{P}_i = E_i^2 - \vec{p}_i \cdot \vec{p}_i = m_i^2 \quad (2.4)$$

the Mandelstam variables can be linearly combined to give the sum of the masses of the four particles involved in the process

$$s + t + u = \sum_{i=1}^4 m_i^2. \quad (2.5)$$

One can describe the full kinematics of this two-body to two-body reaction with two independent kinematical quantities only. Specifically, from four-momentum conservation the reaction $a + b \rightarrow c + d$ results in three equations and four unknowns (given the initial state $a + b$ is known). The unknowns correspond to the momenta of the outgoing particles as well as their scattering angles. A measurement of one of the unknowns, for example, the scattering angle of one of the outgoing particles, will result in the determination of the angle of the other outgoing particle along with the momenta of both particles. Typically, observables such as cross section and polarizations are evaluated as functions of E_γ and $\theta_{c.m.}$. The choice of θ in the center-of-mass frame is made due to the fact that in the two-body to two-body reaction, the distributions of the final-state particles, and thus the observables that depend on these distributions, are symmetric. From E_γ and $\theta_{c.m.}$ one can determine all other variables such as the momenta of both outgoing particles. For example, the outgoing proton momentum in deuteron photodisintegration as a function of the proton angle in the center-of-mass frame and the incident photon energy are calculated to be (assuming that $m_p \approx m_n$)

$$p_p = \frac{E_\gamma s \cos \theta_{cm} + (E_\gamma + m_d) \sqrt{s^2 - 4sm_p^2 - 4m_p^2 E_\gamma^2 \sin^2 \theta_{cm}}}{2(s + E_\gamma^2 \sin^2 \theta_{cm})}. \quad (2.6)$$

The allowed proton momenta for a given photon-energy bin can be determined by applying the condition $0^\circ \leq \theta_{cm} \leq 180^\circ$.

2.2.1 Cross section and observables

Studies of angular and energy dependencies of unpolarized and polarized physics observables give insights into the underlying dynamics of the reaction of interest. Certain characteristics of cross section and polarization observables, like their angular and energy dependence, are sensitive to the reaction dynamics. For this reason, a complete study in

terms of angular and energy dependence of all observables is necessary to form a complete description of the reaction mechanism. Here we describe the formalism, relating experimental observables such as polarizations and cross section to quantities which contain the physics information about hadron structure and reaction dynamics.

The most general decomposition of the deuteron photodisintegration cross section was first discussed by Donnelly and Raskin [29]. The cross section can be shown to have the form (from a latter work by Dmitrasinovic and Gross [30])

$$\begin{aligned}\frac{d^2\sigma}{d\Sigma_0} &= \frac{1}{2}\tilde{R}_T^{(I)} + \frac{1}{2}(|a_+|^2 - |a_-|^2)\tilde{R}_{T'}^{(II)} + \text{Re}(a_+a_-^*)\tilde{R}_{TT}^{(I)} + \text{Im}(a_+a_-^*)\tilde{R}_{TT}^{(II)} \\ &= \frac{1}{2}\left(\tilde{R}_T^{(I)} + p_\odot\tilde{R}_{T'}^{(II)} - p_{lin}\tilde{R}_{TT}^{(I)}\cos 2\phi - p_{lin}\tilde{R}_{TT}^{(II)}\sin 2\phi\right),\end{aligned}\quad (2.7)$$

where the $\tilde{R}$'s are the response functions, $a_\pm$ are the coefficients of the circular polarization states, ϕ is the angle between the photon polarization vector (assumed to be parallel to the x axis) and the reaction plane as indicated in Fig. 2.1, and $d\Sigma_0$ takes the appropriate form to transform the cross section to the reference frame of interest. These quantities are discussed in more detail below.

The cross section in Eq. (2.7) for deuteron photodisintegration is given as a function of the response functions, which are, in general, functions of $Q^2 = -q^2$ (with q being the four momentum of the photon), ν (the energy of the photon) and the scattering angle θ . This is commonly done for electrodisintegration or electron-deuteron scattering, since these response functions are used to differentiate between absorption of longitudinal and transverse photons [31] (the subscript T refers to the transverse component of the photon polarization). The cross-section result for photodisintegration is easily obtained from the cross-section formula for electrodisintegration (not shown here) by removing all longitudinal components of the response functions.

The response functions of Eq. (2.7) are related to sums over the squares of matrix elements of the deuteron current. In the helicity basis, with λ_γ the helicity of the photon, λ_1 and λ_2 the helicity of particles 1 and 2 in the final state, and λ_d the helicity of the initial deuteron, the current operator is written $\langle\lambda_1\lambda_2|J_{\lambda_\gamma}(p_\gamma)|\lambda_d\rangle$. Specifically, the response

functions, $\tilde{R}$'s, are related to the components of the deuteron's response tensor $R_{\lambda_\gamma \lambda'_\gamma}$,

$$R_{\lambda_\gamma \lambda'_\gamma} = \frac{m^2}{4\pi^2 W} \sum_{\lambda_1 \lambda'_1} \sum_{\lambda_2} \sum_{\lambda_d \lambda'_d} \rho_{\lambda_1 \lambda'_1}^{N(X)} \langle \lambda'_1 \lambda_2 | J_{\lambda_\gamma}(p_\gamma) | \lambda_d \rangle \rho_{\lambda_d \lambda'_d}^d \langle \lambda'_d | J_{\lambda'_\gamma}^\dagger(p_\gamma) | \lambda_1 \lambda_2 \rangle, \quad (2.8)$$

with $X = I$ or II , as given by the following relations:

$$\begin{aligned} \tilde{R}_T^{(I)} &= R_{++} + R_{--} \\ \tilde{R}_{TT}^{(I)} &= 2\text{Re}(R_{+-}) \\ \tilde{R}_{T'}^{(II)} &= R_{++} - R_{--} \\ \tilde{R}_{TT}^{(II)} &= -2\text{Im}(R_{+-}), \end{aligned} \quad (2.9)$$

and W being the invariant mass of the final state. The types $X = I$ or $X = II$ are used to distinguish between observables. Specifically, for unpolarized particles in the final state, the response functions of type II are zero.

In cases where the deuteron target is polarized only in the y direction, and where only the polarization of the outgoing particle 1 might be measured, the spin density matrix for particle 1 in the final state is given by $\rho^{N(X)}$ and that of the deuteron in the initial state by ρ^d . The spin density matrices are then given by

$$\begin{aligned} \rho_{\lambda\lambda'}^d &= \frac{1}{3} \begin{pmatrix} 1 & -ip_y^d & 0 \\ ip_y^d & 1 & -ip_y^d \\ 0 & ip_y^d & 1 \end{pmatrix} \\ \rho_{\lambda\lambda'}^{N(I)} &= \frac{1}{2} (1 + p_{y'}^N (\sigma_y)_{\lambda\lambda'}) \\ \rho_{\lambda\lambda'}^{N(II)} &= \frac{1}{2} (p_{x'}^N (\sigma_x)_{\lambda\lambda'} + p_{z'}^N (\sigma_z)_{\lambda\lambda'}), \end{aligned} \quad (2.10)$$

where p_y^d is the vector polarization of the deuteron target¹, p_i^N is polarization of outgoing particle 1 (measured with respect to the $x'y'z'$ coordinate system shown in Fig. 2.1), and σ_i 's are the Pauli matrices. It should be noted here that if all of the hadrons are unpolarized, type II response functions (denoted by the superscript) vanish and response functions of type I are nonzero. The determination of type II response functions requires measurement of the polarization of an outgoing nucleon [30].

¹For the purposes of this analysis an unpolarized deuterium target was used. The reference to polarized targets is only made for the sake of completeness.

The most general polarization of the incoming photon, ϵ , is a superposition of the circular polarization states $\epsilon_{\pm}$, which can be written with the expansion coefficients $a_{\pm}$ (see Eq. (2.7)) as $\epsilon = a_+ \epsilon_+ + a_- \epsilon_-$, where $|a_+|^2 + |a_-|^2 = 1$. The expansion coefficients can be written in terms of only three independent parameters, β , ϕ , and α ,

$$\begin{aligned} a_+ &= -\cos \beta e^{-i(\phi+\alpha)} \\ a_- &= \sin \beta e^{-i(\phi-\alpha)}, \end{aligned} \quad (2.11)$$

with $0 \leq \beta \leq \pi/2$. In Eq. (2.7), $p_{\odot} = \cos 2\beta$ and $p_{lin} = \sin 2\beta$ are the fractions of right circular and linear photon polarizations, respectively.

Equation (2.7) is used to describe the cross section in either the center-of-mass frame of the outgoing pn pair or the laboratory frame (using the appropriate form of $d\Sigma_0$), since the response functions $\tilde{R}$'s are covariant. To transform the cross-section equation to another frame of interest, one needs to use the appropriate form of $d\Sigma_0$

$$\begin{aligned} d\Sigma_0|_{c.m.} &= d\Omega_{1c.m.} \left(\frac{|\vec{p}_1|}{4p_{\gamma}W} \right)_{c.m.} \\ d\Sigma_0|_{lab} &= d\Omega_{1lab} \left(\frac{|\vec{p}_1|}{4p_{\gamma}m_d} \right)_{lab} \mathcal{R}, \end{aligned} \quad (2.12)$$

where $\mathcal{R}$ is a lab recoil factor

$$\mathcal{R} = \frac{W}{m_d} \left(1 + \frac{p_{\gamma}p_1 - E_1q \cos \theta}{m_d p_1} \right)_{lab}^{-1}, \quad (2.13)$$

and p_{γ} is the energy of the incident photon.

The cross section in Eq. (2.7) is the *exact* relativistic result, and all of the unknowns about strong dynamics are contained in the hadronic matrix element of the current, and therefore, in the response functions R (Eq. (2.9)) that are products of these currents. One can rewrite the cross-section expression to include the polarization observables defined by the density matrices of Eq. (2.10) (following the notation suggested in [32])

$$\begin{aligned} \frac{d\sigma}{d\Omega} &= \left(\frac{d\sigma}{d\Omega} \right)_0 [1 + p_y^N P_y + p_y^d T + p_{lin}(\Sigma + p_y^N T_1) \cos 2\phi + p_{\odot}(C_{x'} p_{x'} + C_{z'} p_{z'}) \\ &\quad + p_{lin}(O_{x'} p_{x'} + O_{z'} p_{z'}) \sin 2\phi], \end{aligned} \quad (2.14)$$

where $(\frac{d\sigma}{d\Omega})_0$ corresponds to the unpolarized cross section. In the center-of-mass system

$$\left(\frac{d\sigma}{d\Omega}\right)_0 = \frac{e^2 p_1}{8p_\gamma} R_T = \frac{\alpha p_1}{192\pi E_1^2 p_\gamma} f(\theta), \quad (2.15)$$

where

$$f(\theta) = \sum_1^6 (|F_{i+}|^2 + |F_{i-}|^2) \quad (2.16)$$

with the F 's being the helicity amplitudes of the reaction as defined in Table 2.1. In general, one can write the cross-section formula in terms of different scattering amplitudes. However, helicity amplitudes are considered to be the simplest and most useful amplitudes. The helicity amplitudes correspond to the elements of the scattering matrix with respect to the helicities of the initial and final particles involved, and are complex functions of E_γ and $\theta_{c.m.}$. The advantages of using helicity amplitudes is that the conditions imposed by symmetry transformations reduce the number of independent amplitudes, and therefore simplify the calculation. In the case of deuteron photodisintegration, there are only 12 independent amplitudes [30].

The explicit expressions for the observables in the cross section (Eq. (2.14)) in terms of the helicity amplitudes are shown in Table 2.2. The relation between the structure functions shown in Table 2.2 and the response function of Eq. 2.7 is given in Ref. [33].

Table 2.1. Helicity amplitudes of $\gamma + d \rightarrow p + n$ as defined in Ref [32].

$F_{1\pm}/2m$	$\langle \pm \frac{1}{2} \pm \frac{1}{2} J \cdot \epsilon_+ 1 \rangle$
$F_{2\pm}/2m$	$\langle \pm \frac{1}{2} \pm \frac{1}{2} J \cdot \epsilon_+ 0 \rangle$
$F_{3\pm}/2m$	$\langle \pm \frac{1}{2} \pm \frac{1}{2} J \cdot \epsilon_+ -1 \rangle$
$F_{4\pm}/2m$	$\langle \pm \frac{1}{2} \mp \frac{1}{2} J \cdot \epsilon_+ 1 \rangle$
$F_{5\pm}/2m$	$\langle \pm \frac{1}{2} \mp \frac{1}{2} J \cdot \epsilon_+ 0 \rangle$
$F_{6\pm}/2m$	$\langle \pm \frac{1}{2} \mp \frac{1}{2} J \cdot \epsilon_+ -1 \rangle$

The polarization observables can be divided into two groups: *single* and *double* polarization observables. The *single* polarization observables are the induced proton polarization P_y , the azimuthal asymmetry Σ , and the vector-polarized target asymmetry T . The induced polarization and azimuthal asymmetry can be extracted by taking the ratios of differential cross sections determined with unpolarized target and beam, and polarized

Table 2.2. Polarization observables of $\gamma + d \rightarrow p + n$ as defined in Ref [32].

Obse.	Structure function	Helicity amplitude combination
P_y	$R_T(y)$	$2Im \sum_{i=1}^3 \{F_{i+}^* F_{(i+3)-} + F_{i-} F_{(i+3)+}^*\} / f(\theta)$
T	$R_T(Im T_{11})$	$2Im \sum_{i=1}^2 \sum_{j=0}^1 \{F_{(i+3j)+} F_{(i+3j+1)+}^* + F_{(i+3j)-} F_{(i+3j+1)-}^*\} / f(\theta)$
Σ	R_{TT}	$2Re \sum_{i=1}^3 (-)^i \{-F_{i+} F_{(4-i)-}^* + F_{(3+i)+} F_{(7-i)-}^*\} / f(\theta)$
T_1	$R_{TT}(y)$	$2Im \sum_{i=1}^3 (-)^i \{-F_{i+} F_{(7-i)-}^* + F_{i-} F_{(7-i)-}^*\} / f(\theta)$
$C_{x'}$	$R_T(x')$	$2Re \sum_{i=1}^3 \{F_{i+}^* F_{(i+3)-} + F_{i-} F_{(i+3)+}^*\} / f(\theta)$
$C_{z'}$	$R_T(z')$	$\sum_{i=1}^6 \{ F_{i+} ^2 - F_{i-} ^2\} / f(\theta)$
$O_{x'}$	$R_{TT}(x')$	$2Im \sum_{i=1}^3 (-)^{i+1} \{F_{i+} F_{(7-i)+}^* + F_{i-} F_{(7-i)-}^*\} / f(\theta)$
$O_{z'}$	$R_{TT}(z')$	$2Im \sum_{i=1}^3 (-)^{i+1} \{F_{i+} F_{(4-i)-}^* + F_{(3+i)+} F_{(7-i)-}^*\} / f(\theta)$

beam, respectively.

$$P_y = \frac{1}{p_y^N} \left(\frac{d\sigma/d\Omega_+ - d\sigma/d\Omega_-}{d\sigma/d\Omega_+ + d\sigma/d\Omega_-} \right), \quad \Sigma = \frac{1}{p_{lin}} \left(\frac{d\sigma/d\Omega_{\perp} - d\sigma/d\Omega_{\parallel}}{d\sigma/d\Omega_{\perp} + d\sigma/d\Omega_{\parallel}} \right), \quad (2.17)$$

where $\pm$ refers to the proton spin being directed along (+) or against (-) the y' axis, and in the expression for Σ the photon is polarized either in the plane along the x axis ($\phi = 0$, denoted by $\parallel$) or out of plane, along the y axis ($\phi = \pi/2$ denoted by $\perp$). It should be noted here that for the analysis presented in this thesis, the opposite convention for $\Sigma(\sigma_{\parallel} - \sigma_{\perp})$ is used ($\Sigma \rightarrow -\Sigma$).

The *double* polarization observables are T_1 , $O_{x'}$ and $O_{z'}$, obtained with linearly polarized photons, and $C_{x'}$ and $C_{z'}$, obtained with circularly polarized photons.

2.3 Physics Motivation

The deuteron, the only $A = 2$ nucleus, provides the simplest microscopic test for investigating the role of quarks and gluons in nuclei. The process of deuteron photodisintegration,

$$\gamma + d \rightarrow n + p, \quad (2.18)$$

has been well studied experimentally. In the few-MeV energy region, cross-section and polarization measurements performed for deuteron photodisintegration are well described

by the “traditional” nuclear picture of meson exchanges. These “traditional” theoretical models include the Impulse Approximation [34], the Coupled Channel Approach [35], the Diagrammatic Approach [36], and the modified Current Conservation [37]. Until now, however, characteristics of the reaction in the few-GeV region challenge the validity of theoretical models that describe the reaction in the few-MeV region. Specifically, questions related to how and at which energies the transition takes place from the hadronic picture of the deuteron (established at energies below 1 GeV) to the quark-gluon picture (expected to hold well above 1 GeV) are still unresolved.

2.3.1 Theoretical models

The strong interaction is well described by QCD. QCD is widely used for high-energy calculations since the strength of the coupling constant allows the application of perturbative methods. At intermediate energies, the hadronic degrees of freedom fail to describe qualitative features of nuclear interactions and at the same time perturbative QCD calculations cannot be applied. Available theoretical models that attempt to describe reactions at these intermediate energies try to extend the applicability of pQCD to lower energies by factorization. The application of factorization requires the assumption that a large-momentum-transfer (hard) subprocess, which can be calculated in pQCD, is involved in the reaction. Non-perturbative contributions to the reaction are usually factored out from the calculation and replaced with measurable quantities. The kinematic conditions under which one can apply factorization are, however, still unclear. These hard processes are the origin of the Constituent Counting Rules (CCR). These rules predict a scaling of the differential cross section, and for this reason, such scaling has been used experimentally in studying the transition region from hadronic to partonic degrees of freedom.

Constituent Counting Rules

The prediction of the CCR is calculated by considering the carriers of the currents within the hadrons as structureless and describable in terms of a renormalizable field theory. The conditions of asymptotic scale invariance for the interaction between the constituents and finiteness at the origin of the hadronic wave function are considered for this

calculation [38]. Specifically, in the case where the strong interaction would be turned off, and thus the hadrons would be considered a collection of free quarks with equal momenta, the dimension of the connected invariant amplitude $M_{n_i \rightarrow n_f}$ with $n_i + n_f$ external lines is $(\text{length})^{n_i + n_f - 4}$. If all the invariants are large relative to the masses of the particles involved in the reaction and are proportional to s , then $M_{n_i \rightarrow n_f} \sim s^{-(n_i + n_f)/2 + 2}$. Under the assumption that any binding between the quarks in the hadrons does not modify this basic scaling result, then, by using $d\sigma/dt \sim s^{-2}|M|^2$, one obtains

$$\left(\frac{d\sigma}{dt}\right)_{AB \rightarrow CD} \sim s^{2-n} f(t/s) \quad (2.19)$$

($s \rightarrow \infty$, t/s fixed), where s , t (and u) are the usual Mandelstam variables and n is the number of elementary fields in the initial and final states.

In the deuteron photodisintegration process, $\gamma + d \rightarrow p + n$, $n = 13$, since the initial state has one photon plus 6 quark fields for the deuteron, while the final state has 6 quark fields (3 for the proton and 3 for the neutron). Therefore, at high incident photon energies ($E_\gamma \gg m_d$) and large angles (*i.e.* 90°), where the momentum transfer to each nucleon t_1 and t_2 is balanced, *i.e.*

$$\begin{aligned} t_1 \equiv t &= (\mathcal{P}_\gamma - \mathcal{P}_p)^2 = m_p^2 - m_d E_\gamma \left(1 - \sqrt{1 - \frac{4m_p^2}{s} \cos \theta_{c.m.}} \right) \rightarrow -s \sin^2 \frac{1}{2} \theta_{c.m.} \\ t_2 \equiv u &= (\mathcal{P}_\gamma - \mathcal{P}_n)^2 = m_n^2 - m_d E_\gamma \left(1 - \sqrt{1 - \frac{4m_n^2}{s} \cos \theta_{c.m.}} \right) \rightarrow -s \cos^2 \frac{1}{2} \theta_{c.m.} \end{aligned} \quad (2.20)$$

where $s = (\mathcal{P}_\gamma + \mathcal{P}_d)^2 = m_d^2 + 2m_d E_\gamma$,

the $\gamma + d \rightarrow p + n$ differential cross section is expected to scale as s^{-11} . The constituent counting rules have been initially calculated using pQCD methods, and for this reason, cross-section scaling had been considered as the onset of pQCD [38]. However, experimental data show evidence of this dimensional scaling at low energies, ~ 1 GeV, where pQCD is not applicable. Recently, these scaling laws have been calculated from the standpoint of non-perturbative approaches [39]. For this reason, Hadron Helicity Conservation (HHC) has been used as an independent test of pQCD.

Hadron Helicity Conservation

In the pQCD regime, the cross section and polarization observables should satisfy Hadron Helicity Conservation (HHC), in addition to the CCR. The HHC rule, which also stems from pQCD, requires the total helicity of the incoming and outgoing hadrons to be conserved [40, 41]

$$\sum_i \lambda_i = \sum_f \lambda_f, \quad (2.21)$$

where λ_i (λ_f) are the helicities of the initial (final) hadrons. This results from the fact that the dominant form factors are only those that conserve helicity. This is because the elementary vector interactions, such as photon–quark or gluon–quark couplings, conserve chirality, which leads to the conservation of the sums of the components of the hadronic spins along their respective momentum directions.

By applying the HHC rule, one can make predictions for many spin observables. Reference [42] predicts that the amplitudes of Table 2.1 will have the following asymptotic behavior at large t :

$$\begin{array}{ll} F_{1+}, F_{3-}, F_{5\pm} & \text{leading} \\ F_{2\pm}, F_{4\pm}, F_{6\pm} & \text{suppressed by } t^{-1} \\ F_{1-}, F_{3+}, & \text{suppressed by } t^{-2}, \end{array} \quad (2.22)$$

since the amplitude is suppressed by a power of t for each helicity flip. The implications of HHC for the observables is shown in Table 2.3. The limits at $\theta_{c.m.} = 90^\circ$ of some of the observables require additional assumptions about the relations between the helicity-conserving variables [43]. The HHC do not constrain the value of the azimuthal asymmetry in the pQCD regime¹. Specifically, in the pQCD regime the azimuthal asymmetry at $\theta_{c.m.} = 90^\circ$ is shown to be equal to -1 when the absorbed photon is isoscalar, and +1 when it is isovector [44].

The overall behavior of available data for deuteron photodisintegration (see Sec. 2.4) suggests a deviation from prediction of the CCR. In addition to this, several arguments exist against the applicability of pQCD in the energy region $E_\gamma = 1.1 - 2.3$ GeV, and ab-

¹A study that predicted the azimuthal asymmetry to be equal to -1 in the pQCD limits (using HHC) exists [33]. However, this was later shown that it was derived under incorrect assumptions [44].

Table 2.3. Implications of the HHC for some polarization observables in $\gamma + d \rightarrow p + n$. The 90° limits require additional assumptions [43].

Observable	HHC limit	Approach to HHC
$d\sigma/d\Omega$	$F_{1+}^2 + F_{3-}^2 + F_{5+}^2 + F_{5-}^2$	t^{-2}
P_y	0	t^{-1}
T	0	t^{-1}
T_1	0	t^{-1}
$C_{x'}$	0	t^{-1}
$C_{z'}$	$(F_{1+}^2 - F_{3-}^2 + F_{5+}^2 - F_{5-}^2)/(d\sigma/d\Omega) \rightarrow 0$ at 90°	t^{-2}
$O_{x'}$	0	t^{-1}
$O_{z'}$	$-2Im(F_{1+}F_{3-}^* + F_{5+}F_{5-}^*)/(d\sigma/d\Omega) \rightarrow 0$ at 90°	t^{-2}

solute pQCD calculations would not correctly predict the data at these energies. Because of this, several models have been proposed that approach the deuteron photodisintegration reaction using non-perturbative methods. Specifically, considerable efforts have been made on extending meson-exchange calculations to higher-energy scales [45] and QCD to lower-energy scales [38]. These efforts have shown that the energy region between a few hundred MeV and a few GeV provide the possibility of investigating the transition between the conventional model of nuclear interactions to the quark-gluon description based on QCD or QCD-inspired models.

Models for the transition region

The three main theoretical models that attempt to describe this reaction in the energy region $E_\gamma = 1.1 - 2.3$ GeV are the Reduced Nuclear Amplitudes model (RNA) [46], the Hard Rescattering Mechanism (HRM) model [47, 48], and the Quark-Gluon String Model (QGSM) [49, 50]. Other non-QCD-based models that attempt to describe the reaction by extrapolating the conventional mesonic picture of nuclear interactions to the higher-energy region also exist, but they are unsuccessful in describing the data on the differential cross section in this regime, and they will not be considered here.

In the RNA model, the binding of quarks inside the nucleons in the deuteron is described with empirical form factors. Specifically, the RNA is a pQCD extension to reactions involving nuclei, where the elementary scattering of the photon off a quark is considered to be perturbative and all non-perturbative dynamics of the hadron structure

are factored out by the form factors. This approach describes only the $\gamma + d \rightarrow p + n$ cross section with appropriate normalization factors for the data at $E_\gamma > 2$ GeV and $\theta_{c.m.} = 90^\circ$ and 70° . The HRM model investigates the mechanism in which the absorption of the incoming photon by a quark of one nucleon followed by a high-momentum transfer (hard) rescattering with a quark from the second nucleon produces the final two-nucleon state of large relative momentum (*i.e. large transverse momenta*). Finally, the QGSM assumes that the scattering amplitude at very high energies is dominated by the exchange of three valence quarks in the t -channel with any number of gluon exchanges between them.

The following sections provide more detailed descriptions of the models of interest. Since this work is focused on studying polarization observables (and specifically the azimuthal asymmetry, Σ , of deuteron photodisintegration), the model description is limited to the QGSM and HRM that provide predictions for the azimuthal asymmetry, Σ .

Quark-Gluon String Model

The QGSM is based on Regge phenomenology and has two basic ingredients: 1) a topological expansion in QCD, and 2) a space-time picture of the interactions between hadrons that takes into account the confinement of quarks [44]. This model is generally considered a microscopic (nonperturbative) model of Regge phenomenology. The amplitude for deuteron photodisintegration is described in first approximation by the exchange of three valence quarks in the t - (or u -) channels, as shown in Fig. 2.2. In this calculation, the exchange of three valence quarks corresponds to a nucleon Regge trajectory.

The Regge-pole analysis was formulated much before the development of QCD and is used to study the analytic properties of the scattering amplitude as a function of angular momentum. In this framework, the angular momentum is not restricted to be an integer, but is allowed to take any complex value. This phenomenological approach has been studied extensively and has been successful in describing data [51] that explore the strong interaction at high energies. Specifically, it has been shown that scattering amplitudes can be written in terms of moving poles in the complex angular-momentum plane and that all hadrons lie on trajectories which connect the particle mass with its total angular momentum. The QGSM has been successfully applied to many different problems of strong physics: hadronic mass spectrum [52], widths of resonances [53], relations between

total cross sections and residues of Regge poles [54], behavior of hadronic form factors, and lately, heavy-ion collisions at high energies [55].

Perturbative QCD calculations cannot be applied in the energy region $E_\gamma = 1.1 - 2.3$ GeV, since the value of the coupling constant is quite large (close to 1). However, it is possible to consider as a small parameter the quantity $1/N_c$, where N_c is the number of colors [56]. Another suitable parameter is $1/N_f$, where N_f is the number of flavors. The term “topological expansion” refers to the case where the ratio $N_f/N_c \rightarrow 1$ is fixed and the expansion of all amplitudes is carried out in terms of $1/N_f$ or $1/N_c$ [57].

Figure 2.2. Typical diagrams that are used to calculate the $\gamma + d \rightarrow p + n$ amplitude in the QGSM model. The exchange of three valence quarks in the t - (top) and u - (bottom) channels are shown.

The amplitude $T(\gamma + d \rightarrow p + n)$ can be written as [50]

$$\begin{aligned} \langle \mathcal{P}_p, \lambda_p; \mathcal{P}_n, \lambda_n | \hat{T}(s, t) | \mathcal{P}_d, \lambda_d; \mathcal{P}_\gamma, \lambda_\gamma \rangle \simeq \\ \bar{u}_{\lambda_p}(\mathcal{P}_p) \hat{e}_{\lambda_\gamma} [A(s, t) (\hat{\mathcal{P}}_p - \hat{\mathcal{P}}_\gamma) + B(s, t) m] \hat{e}_{\lambda_d} v_{\lambda_n}(\mathcal{P}_n), \end{aligned} \quad (2.23)$$

where m is the nucleon mass, $\mathcal{P}_\gamma$, $\mathcal{P}_d$, $\mathcal{P}_p$, and $\mathcal{P}_n$ are the four-momenta of the photon, deuteron, proton, and neutron, respectively, and λ_i denotes the helicity of the i^{th} particle. A and B are Lorentz-invariant amplitudes, and their explicit form can be found in Ref. [50].

These are explicitly used to describe the spin structure of the amplitude [50].

In terms of A and B , the differential cross section for the reaction $\gamma + d \rightarrow p + n$ is expressed as [44]

$$\frac{d\sigma_{\gamma+d \rightarrow p+n}}{dt} = \frac{1}{64\pi s} \frac{1}{(p_\gamma^{c.m.})^2} [S_t |B(s, t)|^2 + S_u |B(s, u)|^2 + (-1)^{I+1} 2S_{tu} \text{Re}(B(s, t)B(s, u))], \quad (2.24)$$

where I is the isospin of the interacting photon, *i.e.* $I = 1$ (or 0) for isovector (or isoscalar) photons. The kinematical functions S_t , S_u , and S_{tu} are defined as

$$\begin{aligned} S_t &= \frac{1}{6} \sum_{\lambda_\gamma, \lambda_d} \text{Sp}[\hat{\epsilon}_{\lambda_\gamma}(R(\hat{\mathcal{P}}_3 - \hat{\mathcal{P}}_1) + m)\hat{\epsilon}_{\lambda_d}(\hat{\mathcal{P}}_4 - m) \\ &\quad \times \hat{\epsilon}_{\lambda_d}^*(R(\hat{\mathcal{P}}_3 - \hat{\mathcal{P}}_1) + m)\hat{\epsilon}_{\lambda_\gamma}^*(\hat{\mathcal{P}}_3 + m)], \\ S_u &= \frac{1}{6} \sum_{\lambda_\gamma, \lambda_d} \text{Sp}[\hat{\epsilon}_{\lambda_d}(R(\hat{\mathcal{P}}_4 - \hat{\mathcal{P}}_1) - m)\hat{\epsilon}_{\lambda_\gamma}(\hat{\mathcal{P}}_4 - m) \\ &\quad \times \hat{\epsilon}_{\lambda_\gamma}^*(R(\hat{\mathcal{P}}_4 - \hat{\mathcal{P}}_1) - m)\hat{\epsilon}_{\lambda_d}^*(\hat{\mathcal{P}}_3 + m)], \\ S_{tu} &= \frac{1}{6} \sum_{\lambda_\gamma, \lambda_d} \text{Sp}[\hat{\epsilon}_{\lambda_d}(R(\hat{\mathcal{P}}_3 - \hat{\mathcal{P}}_1) + m)\hat{\epsilon}_{\lambda_d}(\hat{\mathcal{P}}_4 - m) \\ &\quad \times \hat{\epsilon}_{\lambda_\gamma}^*(R(\hat{\mathcal{P}}_4 - \hat{\mathcal{P}}_1) - m)\hat{\epsilon}_{\lambda_\gamma}^*(\hat{\mathcal{P}}_3 + m)], \end{aligned} \quad (2.25)$$

where $\hat{\epsilon}_d$ and $\hat{\epsilon}_\gamma$ denote the deuteron and photon polarization vectors, respectively, and R denotes the ratio of the invariant amplitudes, $R = A(s, t)/B(s, t)$. The authors treat this ratio R as a constant. Specifically, the value of R is considered a free parameter of the model and ranges between 1 and 2.5 [44].

To achieve consistency of the differential cross section, $d\sigma/dt$, with Regge asymptotics at large s and fixed t , the authors use the following parametrization of the amplitude $B(s, t)$

$$|B(s, t)|^2 = \frac{1}{C(p_\gamma^{c.m.})^2} |\mathcal{M}_{\text{Regge}}(s, t)|^2, \quad (2.26)$$

where $C = 36 \pm 3 \text{ GeV}^2$ is a constant that is determined from approximating the energy dependence of S_t , and $\mathcal{M}_{\text{Regge}}(s, t)$ is the Regge-pole exchange amplitude

$$\mathcal{M}_{\text{Regge}}(s, t) = F(t) \left(\frac{s}{s_0}\right)^{\alpha_N(t)} \exp\left[-i\frac{\pi}{2}\left(\alpha_N(t) - \frac{1}{2}\right)\right]. \quad (2.27)$$

Here $\alpha_N(t)$ is the trajectory of the nucleon Regge pole, and $s_0 = 4 \text{ GeV}^2 \approx m_d^2$. The

dependence of the residue $F(t)$ on t is such that it takes into account the nucleon pole and the contribution of non-nucleonic degrees of freedom in the deuteron [50]. There are several free parameters in this model and their values are determined from a comparison between this model and the experimental results for the differential cross section at $\theta_{c.m.} = 36^\circ$ and $E_\gamma > 1.5$ GeV [50].

As shown in Ref. [50], the energy behavior of the cross section crucially depends on the form of the Regge trajectory $\alpha_N(t)$. The form of $\alpha_N(t)$ that better describes the data is logarithmic with non-leading contributions:

$$\alpha_N(t) = \alpha_N(0) - \alpha'_N(0)T_B \ln(1 - t/T_B), \quad (2.28)$$

where $T_B = 1.7$ GeV².

The authors provide predictions for the azimuthal asymmetry, Σ , of deuteron photodisintegration, using its relation to the helicity amplitudes [44]. Figure 2.3 shows these predictions for the azimuthal asymmetry as a function of the proton angle in the center-of-mass frame for six different photon energies.

Hard Rescattering Mechanism

The Hard Rescattering Mechanism model considers that hard photodisintegration of the deuteron is a two-step process. The first step involves the interaction between the incoming photon and one quark from one of the nucleons. This interaction results in the knock-out of this quark, which then makes a hard rescattering with a quark from the second nucleon in the deuteron. Figure 2.4 shows schematically the reaction. This two-step process assumption allows the reaction amplitude to be expressed as a convolution of the deuteron wave function, the hard photon-quark interaction amplitude, and the amplitude of hard pn elastic scattering. The latter was estimated using existing data on $p + n \rightarrow p + n$, and therefore this amplitude is sensitive to the uncertainties of these data.

This model provides a convenient framework for calculating several observables. Specifically, the model allows the calculation of various helicity amplitudes that describe the polarization observables (as shown in Sec. 2.3). In the HRM, the differential cross section is calculated to be [47]

Figure 2.3. QGSM predictions for the azimuthal asymmetry Σ of deuteron photodisintegration as a function of the proton angle in the center-of-mass frame for six different incident photon energies. The predictions of QGSM show that the azimuthal asymmetry, Σ , is maximal at $\theta_{c.m.} = 90^\circ$ for all photon energies.

Figure 2.4. Quark rescattering diagram in the HRM.

$$\begin{aligned} \frac{d\sigma^{\gamma+p \rightarrow p+n}}{dt} &= \frac{4\alpha}{9} \pi^4 \frac{1}{s'} C\left(\frac{t_N}{s}\right) \frac{d\sigma^{p+n \rightarrow p+n}(s, t_N)}{dt} \\ &\times \left| \int \Psi_d^{NR}(p_z = 0, p_\perp) \sqrt{m_N} \frac{d^2 p_\perp}{(2\pi^2)} \right|^2, \end{aligned} \quad (2.29)$$

where $C\left(\frac{t_N}{s}\right) = -\frac{2t_N/s'}{1+2t_N/s'} \approx -\frac{t/s}{1+t/s}$ is a scaling factor that accounts for the difference between the hard propagators in this process and those occurring in wide-angle pn elastic scattering ($C\left(\frac{t_N}{s}\right) \approx 1$ at $\theta_{c.m.} = 90^\circ$ and varies slowly as a function of $\theta_{c.m.}$, with $t_N = (p_B - p_d/2)^2$). In Eq. (2.29), s is the usual Mandelstam variable, $s' = s - m_d^2$, $p_\perp$ is the relative transverse momentum of the nucleons in the deuteron, and Ψ_d^{NR} is the non-relativistic deuteron wave function. The differential cross section is calculated under the assumptions $p_\perp^2 \ll m_N^2 \ll s'$ and $\alpha \sim \frac{1}{2}$, where $\alpha \equiv \frac{p_{1+}}{p_{d+}}$ ($p_\pm$ corresponds to the light cone four-momentum defined as $(p_+, p_-, p_\perp)$, where $p_\pm = E \pm p_z$). The predictions of this model for the differential cross section fall within a band due to the uncertainty of the experimental data on $\frac{d\sigma^{p+n \rightarrow p+n}(s, t_N)}{dt}$ and due to the interpolation to the required (s, t) bins.

Within the HRM formalism, the helicity amplitudes of Table 2.1 can be related to the helicity amplitudes for pn scattering [58]. Thus, polarization observables for $\gamma + d \rightarrow p + n$ can be expressed in terms of helicity amplitudes for pn scattering as shown below

$$\begin{aligned}
P_y &= -\frac{2\text{Im}\{\phi_5^*[2(\phi_1 + \phi_2) + \phi_3 - \phi_4]\}}{2|\phi_1|^2 + 2|\phi_2|^2 + |\phi_3|^2 + |\phi_4|^2 + 6|\phi_5|^2}, \\
C_{x'} &= \frac{2\text{Re}\{\phi_5^*[2(\phi_1 - \phi_2) + \phi_3 + \phi_4]\}}{2|\phi_1|^2 + 2|\phi_2|^2 + |\phi_3|^2 + |\phi_4|^2 + 6|\phi_5|^2}, \\
C_{z'} &= \frac{2|\phi_1|^2 - 2|\phi_2|^2 + |\phi_3|^2 - |\phi_4|^2}{2|\phi_1|^2 + 2|\phi_2|^2 + |\phi_3|^2 + |\phi_4|^2 + 6|\phi_5|^2}, \\
\Sigma &= \frac{2\text{Re}\{|\phi_5|^2 - \phi_3^*\phi_4\}}{2|\phi_1|^2 + 2|\phi_2|^2 + |\phi_3|^2 + |\phi_4|^2 + 6|\phi_5|^2}.
\end{aligned} \tag{2.30}$$

where ϕ_i 's are the five independent helicity amplitudes of the pn elastic scattering

$$\begin{aligned}
\phi_1 &= \langle +, + | A_{pn \rightarrow pn} + A_{pn \rightarrow np} | +, + \rangle, \\
\phi_2 &= \langle +, + | A_{pn \rightarrow pn} + A_{pn \rightarrow np} | - , - \rangle, \\
\phi_3 &= \langle +, - | A_{pn \rightarrow pn} + A_{pn \rightarrow np} | +, - \rangle, \\
\phi_4 &= \langle +, - | A_{pn \rightarrow pn} + A_{pn \rightarrow np} | - , + \rangle, \\
\phi_5 &= \langle +, + | A_{pn \rightarrow pn} + A_{pn \rightarrow np} | +, - \rangle.
\end{aligned} \tag{2.31}$$

The following important conclusion can be made from the HRM predictions in Eq. (2.30):

1. P_y and $C_{x'}$ should be small at large $\theta_{c.m.}$, however, this prediction is not an indication of HHC.
2. The relative sign between P_y and $C_{x'}$ is an indication of the relative phase between ϕ_1 and ϕ_5 helicity amplitudes.
3. $C_{z'}$ is predicted to be positive with values $\approx 0.5 - 1$.
4. The relative sign between ϕ_3 and ϕ_4 determines the sign of the azimuthal asymmetry Σ , which is expected to be positive at $\theta_{c.m.} = 90^\circ$.

Figure 2.5 shows the predictions for the azimuthal asymmetry of deuteron photodisintegration, Σ , as a function of the proton angle in the center-of-mass frame for 6 different incident photon energies. Solid and dashed lines correspond to different approximations, referred to as “on-shell” and “off-shell” approximations, in the calculation of

Figure 2.5. HRM predictions for the azimuthal asymmetry, Σ , of deuteron photodisintegration as a function of the proton angle in the center-of-mass frame for six different incident photon energies. The thick and thin solid curves correspond to the “on-shell” version of the model and represent predictions based on large and small helicity non-conservation, respectively. The dashed and dotted lines represent two versions of the “off-shell” approximation [58].

the ϕ_5 helicity amplitude [58]. Thick and thin lines correspond to calculations under the assumption of large and small helicity non-conservation respectively [58]. The range of angular predictions varies for different photon energies to accommodate the assumption that the model is valid only in the hard regime, where $p_T = \sqrt{\frac{1}{2}E_\gamma m_d \sin^2 \theta_{c.m.}} > 1.0 \text{ GeV}/c$ (see Sec. 2.4.1).

2.4 Previous Measurements

During the last 20 years, a number of experiments have been performed that measured differential cross sections over a broad range in energy and angle [59, 60, 61, 62, 63, 64, 65, 66]. In addition, there are measurements of the induced proton polarization [67, 68] and the azimuthal asymmetry [69, 70]. Table 2.4 presents a list of published polarization

Table 2.4. List of published measurements of deuteron photodisintegration in the GeV photon energy region

Lab	Observable	E (GeV)	$\theta_{c.m.}$ ($^\circ$)	Reference
JLab	$d\sigma/dt$	0.5–3.0	10–160	[66]
JLab	$d\sigma/dt$	0.8, 1.5, 2.4, 3.2, 4.0	36, 52, 69, 89	[63]
JLab	$d\sigma/dt$	1.6, 1.9, 2.4	30, 36, 52, 70, 90, 110, 127, 142	[65]
JLab	$d\sigma/dt$	5.0, 5.5	37, 53, 70	[64]
SLAC	$d\sigma/dt$	0.8, 1.1, 1.3, 1.6	90	[59]
SLAC	$d\sigma/dt$	0.8, 1.0, 1.2, 1.6	52, 66, 78, 90, 113, 126, 142	[60]
SLAC	$d\sigma/dt$	1.4, 1.6, 1.8	90, 113, 142	
SLAC	$d\sigma/dt$	1.5, 1.9, 2.3, 2.7	37, 53, 89	[61]
Kharkov	P_y	0.5–1.125	90, 120	[71, 72, 73]
JLab	$P_y, C_{x'}, C_{z'}$	0.479–2.411	90	[67]
JLab	$P_y, C_{x'}, C_{z'}$	2.0	30–110	[68]
Bonn	Σ	0.233–0.818	114, 135	[74]
Yerevan	Σ	0.3–1.0	45, 60, 75, 95	[69]
Yerevan	Σ	0.8–1.6	90	[70]

data in the GeV energy region (for an extended list see Ref. [33]). The energies listed in Table 2.4 correspond to lab photon energies, and the angles correspond to proton center-of-mass angles.

There have been numerous experiments that were done in the few-MeV range

of photon energies. As mentioned before, theoretical models that use the conventional hadronic picture of nuclear interactions have been successful in describing the available cross-section and polarization data in this few-MeV energy region. On the other hand, experiments done in the high-energy region (of the order of GeV) have proven to be, theoretically, more challenging.

2.4.1 Cross sections

Figures 2.6 and 2.7 show published high-energy data on the differential cross section of deuteron photodisintegration as a function of E_γ and s , respectively. The data shown are from CLAS [66], SLAC [59, 60, 61], JLab Hall A [65], and Hall C [63, 64].

Figure 2.6 shows the differential cross section multiplied by the Mandelstam variable s^{11} as a function of E_γ compared to theoretical predictions. As shown in Table 2.4, the experiments carried out at JLab Hall C, SLAC, and Mainz determined the cross section for $\theta_{c.m.} \approx 36^\circ, 52^\circ, 69^\circ$ and 89° at energies from about 0.8 GeV to 5.5 GeV. Some data points at backward angles at 1.8 GeV also exist from NE8 (SLAC). The cross-section data above 1 GeV exhibit the s^{-11} fall-off of the cross section $d\sigma/dt$ at $\theta_{c.m.} = 90^\circ$ and 69° in agreement with the CCR. On the other hand, cross-section data at forward angles exhibit a slower fall-off (with approximate scaling of $\sim s^{-9}$) at lower energies until the onset of the s^{-11} scaling which is exhibited at about 4 GeV.

The cross-section data shown in Fig. 2.7 are multiplied by the factor s^{11} . The blue arrows on Fig. 2.7 indicate the value $p_T = \sqrt{\frac{1}{2}E_\gamma m_d \sin^2 \theta_{c.m.}} = 1.1 \text{ GeV}/c$, which corresponds to the center-of-mass proton transverse momentum. The scaling of the cross section was done to test the prediction of the CCR, that as mentioned before, suggests that point-like constituents, such as quarks and gluons (*i.e.* QCD degrees of freedom), are the correct description for the process.

Figure 2.8 shows the angular distribution of the differential cross section for various incident photon energies. Figures 2.6 and 2.8 show that the cross-section data agree with about the same degree of success with the QGSM and the HRM model. Specifically, the QGSM describes the forward-angle data reasonably well, even for values of $-t$ exceeding the nominal limits of the model (see Fig. 2.8). The most forward points support the presence

Figure 2.6. Photodisintegration cross section $s^{11}d\sigma/dt$ as a function of the incident lab photon energy. The calculations are from Kang, Erbs, Pfeil and Rollnik (solid line; 90° low-energy data only) [75], Lee (long dashed line; 90° low-energy data only) [45], Raydushkin (dot-dashed line), RNA of Brodsky and Hiller (dotted) [46], Quark-Gluon-String model (short dashed) [49, 50], and Frankfurt, Miller, Strikman, and Sargsian-HRM-(shaded region) [47, 48]. The data shown are from SLAC [59, 60, 61], Mainz [62], and Hall C [63, 64].

Figure 2.7. Photodisintegration cross section $s^{11}d\sigma/dt$ as a function of s for given proton scattering angles in the center-of-mass frame. Dashed lines correspond to the scaling prediction of s^{-11} for $p_T \geq 1.1$ GeV/ c . The vertical arrows indicate the value of s that corresponds to this proton transverse momentum p_T . Fits are not shown for $\theta_{c.m.} = 35^\circ, 45^\circ, 135^\circ$, and 145° where there are not enough data above 1.1 GeV. Data are from CLAS [66] (solid red circles), Mainz [62] (open black squares), SLAC [59, 60, 61] (solid down-pointing free triangles), JLab Hall A [65] (solid blue squares), and Hall C [63, 64] (solid up-pointing black triangles).

Figure 2.8. Photodisintegration cross section, $d\sigma/d\Omega$, as a function of the proton angle in the center-of-mass frame at various incident photon energies. Data are from Mainz [62] (open squares, average of the measured values in the given photon energy intervals), SLAC experiments [59, 60, 61] (full/green down-triangles), JLab Hall A [65] (full/blue squares), Hall C [63, 64] (full/black up-triangles), and Hall-B data [66] (full/red circles). The shaded black region corresponds to prediction of the HRM model, whereas the solid line corresponds to predictions of the QGS model.

of a local maximum at about 20° and this is successfully predicted by the QGSM. The predictions of the HRM (hatched band) were calculated using the best angular fit for fixed energy pp scattering data (due to lack of existing pn data), with the bands reflecting the poor accuracy of the data. The HRM more accurately predicts the data in the central angular region over the whole explored energy range, but underpredicts the cross-section data at forward and backward angles (apart from $E_\gamma = 1.8 - 2.5$ GeV).

2.4.2 Polarization observables

For energies above 1 GeV, there are only a few sets of polarization data for deuteron photodisintegration. As shown in Table 2.4, the induced polarization P_y was measured at Kharkov [71, 72, 73] and at JLab [67], the polarization transfer coefficients $C_{x'}$ and $C_{z'}$ at JLab [67], and the azimuthal asymmetry, Σ , at Yerevan [69]. Figures 2.9-2.12 show the results for these observables. There has been some debate whether the Kharkov data should be trusted due to the difficulties in calibrating their polarimeter, along with determining background contributions and false asymmetries, and thus the results of this experiment are not shown in the figures below [33].

Figure 2.9. Induced polarization, P_y , for deuteron photodisintegration at fixed $\theta_{c.m.} = 90^\circ$ (left panel) and fixed energy $E_\gamma \approx 2.0$ GeV (right panel). Data from Stanford [76], Tokyo [77], and JLab [67, 68] are shown in the left panel. Right panel shows only JLab results [67, 68]. The calculations shown in the left panel are from Kang, Erbs, Pfeil and Rollnik (dashed line) [75], HRM (dash dot line) [48] and from Schwamb and Arenhövel (solid line) [78]. Right panel shows calculations from HRM (dotted and dashed-dotted lines) [48].

Figure 2.10. Polarization transfer, $C_{x'}$, for deuteron photodisintegration at $\theta_{c.m.} = 90^\circ$ (top panel) and fixed energy $E_\gamma \approx 2.0$ GeV (bottom panel). The data are from Ref. [67] and Ref. [68]. The calculation is from Schwamb and Arenhöven (top panel solid line) [78]. The dashed line shows the prediction of the HHC. Calculations from HRM [48] are also shown in the bottom panel along with data from JLab.

Figure 2.11. Polarization transfer, $C_{z'}$, for deuteron photodisintegration at $\theta_{c.m.} = 90^\circ$ (top panel) and fixed energy $E_\gamma \approx 2.0$ GeV (bottom panel). See caption of Fig. 2.10 for description of data and calculations. The green dashed line in the lower panel corresponds to a QGS calculation [49].

The induced proton polarization, P_y is predicted to vanish at energies where HHC is applicable. The data for P_y show that the onset for HHC might occur close to $E_\gamma = 1$ GeV (see Fig. 2.9). The HRM makes predictions using small helicity non-conservation (HR-HC) and large helicity non-conservation (HR-nonHC) [48], with these calculations shown in the right panel of Fig. 2.9. The HR-HC predictions lead to small values of $C_{x'}$ and P_y , and $C_{z'}$ being consistent with 1. On the other hand, the HR-nonHC predictions also lead to small values for P_y (close to zero), but to large values for $C_{x'}$ (close to 0.5). $C_{z'}$ is also predicted to be large. It is worth mentioning that the calculations shown here correspond to the lower edge of the nominal validity range of the model. In addition, the predictions of the HRM model are based on the helicity amplitudes for pn scattering which were constrained using pp data. For this reason, there are large uncertainties in these predictions. This

Figure 2.12. Azimuthal asymmetry, Σ , of deuteron photodisintegration at $\theta_{c.m.} = 90^\circ$. The calculations are from HRM (left panel) [48] and QGSM (right panel) [49]. The data shown are from Yerevan [69, 70].

model in its current state is not very successful in describing the angular dependence of the induced polarization, P_y , at $E_\gamma \approx 2$ GeV. Predictions of the QGSM are not provided for the induced proton polarization due to its sensitivity to the relative phase between the helicity amplitudes.

Figures 2.10–2.11 show data for the polarization transfers $C_{x'}$ and $C_{z'}$ compared with theoretical calculations. HRM calculations are again not very successful in describing the data for $C_{x'}$, whereas both models (HRM and QGSM) describe the available data on $C_{z'}$ with about the same degree of success up to 90° . At larger angles, the QGSM is more consistent with the data, while the HRM calculations are above the experimental value.

For the azimuthal asymmetry, Σ , the only data available are from Yerevan [69, 70] at $\theta_{c.m.} = 90^\circ$ in the energy range $0.8 - 1.6$ GeV as shown in Fig. 2.12. These data are compared with the HRM calculations (left panel) and the QGSM (right panel). Although the Yerevan data have large uncertainties at large energies ($1.4 - 1.6$ GeV), they indicate that the azimuthal asymmetry at $\theta_{c.m.} = 90^\circ$ might be large (about 50%) at these energies and at $\theta_{c.m.} = 90^\circ$.

In general, the data for P_y at 90° are consistent with zero at energies above about 1 GeV, that is, at energies where the differential cross section-scaling begins. This would seem to be consistent with the predictions of pQCD. However, the polarization transfer observables $C_{x'}$ and $C_{z'}$ both peak near 1 GeV and do not decrease sufficiently fast to confirm the onset of HHC at 1 GeV. In fact, the comparison between these data and the theoretical calculations supports large helicity non-conservation. In addition, the measurements of the azimuthal asymmetry, Σ (for $\theta_{c.m.} = 90^\circ$), indicate that there may be a minimum near 1.2 GeV, with Σ increasing towards unity at higher energies.

2.5 Summary

It has been shown from several experiments that conventional meson-nucleon models describe the deuteron photodisintegration data well only in the few-MeV region. The scaling behavior of the differential cross section in the GeV energy region indicates that a new approach using quark degrees of freedom is needed to understand the reaction at these energies. Several theoretical models have been developed, and they describe the available data on the differential cross section with great success. Two of these models, the QGSM and HRM, make predictions also for polarization observables. Measurements of these observables can be used to test and constrain these models. The available data for the polarization transfers $C_{z'}$, are described with the same degree of success by both HRM and QGSM. The QGSM makes no predictions for the induced proton polarization, P_y , and polarization transfer, $C_{x'}$, due to the sensitivity of their calculation to the relative phase of the helicity amplitude. In general, the available data on the induced proton polarization and polarization transfers do not allow us to test the available theoretical calculations.

On the other hand, the situation is completely different for the case of the azimuthal asymmetry Σ . There is only one set of data for Σ in the energy region of interest, and only for $\theta_{c.m.} = 90^\circ$. Moreover, these data have large uncertainties, especially at higher energies, and do not provide an adequate test of the theoretical models at hand. The available theoretical models, QGSM and the HRM, are sensitive to the helicity amplitudes that define the azimuthal asymmetry, Σ , and the difference between their predictions of Σ (see Fig. 2.13) is quite large.

Figure 2.13. Azimuthal asymmetry, Σ , of deuteron photodisintegration at $\theta_{c.m.} = 90^\circ$ as a function of incident photon energy. The calculations are from QGSM (left panel) [49] and HRM (right panel) [48]. It is clear that there is a big difference in the predictions of the azimuthal asymmetry of deuteron photodisintegration between these two models.

Given the rich helicity structure of the reaction, and the observations that non-helicity-conserving mechanisms can mimic helicity-conserving behavior of a polarization observable made both in HRM studies and measurements of the proton recoil polarization, it is clear that measurements of additional polarization observables in the few-GeV range are needed to gain insight into the underlying non-perturbative quark-gluon dynamics. Each of the two theoretical models attempting to describe the reaction in the few-GeV range makes predictions for the energy and angular dependences of the azimuthal asymmetry, Σ . The shapes and the magnitudes of these dependences are very different (see Figs. 2.3, 2.5, and 2.13). This strong sensitivity of Σ to the kinematics and the dynamics of the reaction will give access to important aspects of the underlying physics. New data on Σ will allow us to test the available theoretical models and, in general, aid in understanding the underlying dynamical mechanisms.

Chapter 3

Experimental Facility

The work on this thesis was done using data from the g13 experiment [79] that was carried out at the Thomas Jefferson National Accelerator Facility (JLab) in Newport News, Virginia. JLab is focused on research at intermediate energies to study hadronic and nuclear physics, using a variety of targets, and both electron and photon beams. The current configuration of the accelerator allows electron-beam energies up to 6 GeV. An upgrade of the facility is currently underway and will allow electron-beam energies up to 12 GeV.

Facilities other than JLab are equipped with standard detectors that, on one hand, provide high momentum resolution ($dp/p \sim 10^{-4}$), but on the other hand, cover small solid angles (10^{-3} of 4π sr) and small momentum range ($\Delta p/p \sim 10\%$). These facilities are able to make precise measurements in a limited kinematical range. JLab, along with high precision spectrometers, is equipped with a detector that has a large angular coverage, and is thus best suited for measurements of asymmetries.

This chapter describes in general the experimental facility and the equipment used to collect the data for this analysis. Specifically, a description of the production of the beam, the detector systems used to identify reaction products, and an overview of the g13 experiment are presented.

3.1 CEBAF

The Continuous Electron Beam Accelerator Facility (CEBAF) [80] at JLab provides a high-duty-cycle electron beam up to 6 GeV in energy. A schematic diagram of the accelerator facility is shown in Fig. 3.1. Electrons are produced using radio-frequency-gain-switched lasers directed at a GaAs photocathode at 100 keV and are accelerated to 67 MeV by a cryounit and two cryomodules before being injected into the accelerator. The CEBAF electron guns can produce longitudinally polarized electrons beams with polarization of $\sim 85\%$ [79].

Figure 3.1. The Continuous Electron Beam Accelerator (CEBAF) at JLab.

After the electron beam is injected into the race-track-shaped accelerator, it can be circulated around up to five times. With each pass through one of the two linear accelerators (LINACs) shown in the Fig 3.1, the beam gains approximately 600 MeV in energy. The energy of the electron beam can vary from 1.2 GeV (one pass) to 6 GeV (five pass) by recirculating the beam through the accelerator. When the electron beam reaches the desired energy, it can be split into three independent beams and transported to the three end stations, called Halls, for simultaneous data collection for three different experiments.

The experimental halls are equipped with spectrometers for complementary experimental programs. Halls A and C operate at high electron-beam luminosities and are used in high-precision electron scattering experiments. The reaction products in both Halls are detected using high-resolution spectrometers. However, these high precision measurements come at the price of having a small acceptance (~ 1 sr for each spectrometer). Hall B is equipped with a large acceptance spectrometer, the CLAS, which was used in the g13

experiment. A detailed description of Hall B and the CLAS follows.

3.2 Hall B and the CLAS

Hall B houses the CEBAF Large Acceptance Spectrometer (CLAS), which provides efficient detection of particles over a large fraction of the full 4π sr solid angle. Figure 3.2 shows the experimental setup in Hall B [81]. Beamline detectors are used for measuring and monitoring the electron and photon beam characteristics. The photon beam is produced via the bremsstrahlung technique, and photons are tagged using the tagger spectrometer. Tagging refers to determining the photon energies based on the electrons timing and hit position information. Particles produced in a photoreaction in the target encounter first the start counter, which surrounds the target. The start counter (ST) [82] is used to determine the start time of a physics reaction. The torus magnet [83], which produces a uniform magnetic field primarily in the ϕ direction, bends outgoing charged particles which are then tracked in the drift chambers (DC) [84]. The time-of-flight system (TOF) [85] provides track timing information and it is used for particle identification. The Čerenkov counters (CC) [86], located at forward angles, can be used for electron-pion discrimination (not used for g13), and the electromagnetic calorimeters (EC) [87] are used to detect neutral particles.

Fast data acquisition and triggering rates allow experiments at Hall B to achieve luminosities of the order of $10^{34} \text{ cm}^{-2}\text{s}^{-1}$ [81]. The sophisticated equipment of Hall B, along with a variety of cryogenic and solid targets developed for the CLAS, makes Hall B an ideal environment for experimental research.

Figure 3.2. The experimental Hall B that is comprised of the CLAS, the photon tagger, and beamline detectors [81]

3.2.1 Photon beam

In general, there are two main mechanisms of producing beams for photoproduction experiments: the Compton back-scattering, and the bremsstrahlung techniques. The Compton back-scattering technique, which is mainly used at BNL (LEGS) and SPring8 (LEPS), refers to the process of photon emission by electrons stored in a storage ring. These highly-polarized, monochromatic photons are emitted when the electrons in the ring interact with a laser beam [88]. The bremsstrahlung technique [89], which is mainly used at JLAB (CLAS), ELSA, and MAMI, refers to the process of photon emission by a charged particle, such as an electron, when deflected by another charged particle, such as an atomic nucleus.

Hall B can operate in three different modes for photon beams: an unpolarized mode, a circularly polarized mode, and a linearly polarized mode. The unpolarized mode is produced by an unpolarized electron beam incident on an amorphous radiator such as carbon. This method produces a characteristic $1/E_\gamma$ photon-energy spectrum, with the photon energy being identified by measuring the initial and final energy of the electrons that produced the radiation. The g13 experiment took data with circularly and linearly polarized photon beams.

Circularly polarized photon beam

When the electron beam incident on an amorphous radiator is longitudinally polarized, the resulting bremsstrahlung photons are circularly polarized [90]. As mentioned before, CEBAF can produce longitudinally polarized electron beams with polarization up to 80-85% [79]. Their polarization is measured periodically using a Møller polarimeter (See Section 3.2.4). The degree of polarization transfer to the photon in the bremsstrahlung process is almost 100% when the photon carries off the maximum available energy. The polarization decreases with the fraction of the electron energy given to the photon. Figure 3.3 shows the degree of circular polarization of bremsstrahlung photons from longitudinally polarized electrons [90].

Figure 3.3. Degree of photon polarization of circular photons as a function of the momentum transfer to the bremsstrahlung photon k/E_0 , where k is the momentum of the photon and E_0 is the energy of the incident electron. Curves a and c correspond to an exact calculation and to an approximation, respectively [90].

Linearly polarized photon beam

The linearly polarized photon beam at Hall B is obtained via the coherent bremsstrahlung technique [91]. This technique makes use of the lattice structure of a crystal to create polarized photons in a direction which depends on the orientation of the crystal. Typically, a 20–50 μm thick diamond crystal is aligned and oriented at the desired angle to the beam line via a goniometer [92]. The goniometer is a device that allows an object to be rotated to a precise angular position. When the primary electron beam from CEBAF hits the thin diamond crystal a photons is produced, which is mixed with the incident electron beam. These two beams are separated using the tagging system described below. For a particular orientation of the crystal, an electron scattering from a reciprocal lattice vector will have a restricted momentum transfer that can be absorbed by the crystal. This will give rise to sharp structures (coherent peaks) in the photon energy spectrum. One way to visualize the spectrum of the polarized photons is by enhancement plots. The enhancement, or relative intensity plot, corresponds to the photon energy spectrum of the crystal radiator divided by the photon energy spectrum of the amorphous radiator, enhancing in this way the coherent bremsstrahlung spectrum. This is done to remove the incoherent bremsstrahlung effects from the photon energy spectrum produced in the diamond radiator. Figure 3.4 shows the enhancement as a function of photon energy, with the sharp structures corresponding to

Figure 3.4. Data and Monte Carlo simulation of the photon intensity spectrum produced by the process of coherent bremsstrahlung [92].

highly polarized photons. The series of peaks in the Fig. 3.4 correspond to the electrons scattering off a different reciprocal lattice vector, producing photons with different degrees of polarizations at different energies.

The degree of linear polarization of the photons depends on the position of the coherent peak relative to the incident electron-beam energy (E_0). Specifically, the degree of linear polarization increases with the ratio E_0/E_γ . For example, positioning the coherent peak at about 30% of the electron-beam energy the degree of photon polarization can reach up to $\sim 80\%$. Additionally, the linearly polarized photons are produced at very small aperture angles (for $E_\gamma = 6$ GeV, the characteristic angle is about $1/24$ rad), which requires the use of an active collimator located downstream of the tagging system in order to reduce the incoherent radiation, whose characteristic angle is twice that of the coherent peak [92].

For experiments aiming to extract polarization observables, the knowledge of the degree of photon polarization is essential. The photon polarization is obtained from fits to the experimental photon energy spectrum using a model that takes into account all relevant parameters, such as electron-beam energy, electron-beam spot size and divergence, coherent peak position, collimation, and multiple scattering in the crystal [93]. The extraction of the photon polarization is described in more detail in Chap. 6. Figure 3.5 shows the calculated

linear photon polarization and enhancement as a function of photon energy E_γ for 5 GeV (left) and 3 GeV (right) electron beams. As the edge of the coherent peak is moved to higher photon energies, the width of the peak increases. Using a 5 GeV electron beam, the polarization varies slightly over a 200 MeV region ($\sim 10\%$), but decreases rapidly in the tail of the distribution. Moreover, the intensity of the polarized photon beam decreases as the photon energy increases, and because of this g13 collected more triggers for the larger photon energies in an attempt to have equal uncertainties in all photon-energy bins.

Figure 3.5. Enhancement and degree of polarization for linearly polarized photons as a function of photon energy, E_γ , for 5 GeV (left) and 3 GeV (right) electron beams. The colors correspond to various setting of the goniometer, which is used to shift the edge of the coherent-photon peak [79].

3.2.2 Photon tagger

The photon tagging spectrometer is used to identify (tag) the energy of the photons incident on the physics target in the CLAS by measuring the energy of the electrons that underwent bremsstrahlung. This is done by bending the scattered electrons using a dipole magnet positioned downstream from the radiator and detecting them in the focal plane of the magnet using a hodoscope (thus determining the energy of the scattered electrons). Since the incident electron-beam energy is well determined by the accelerator (with uncertainties $<0.5\%$), one can determine the energy of the photon that initiated a

nuclear reaction in the physics target on an event-by-event basis. The photon energy is determined by the hit position of the scattered electron in the hodoscope. The hit timing information is used to identify the photon that initiated the nuclear reaction. Figure 3.6 shows a schematic diagram of the tagger.

Figure 3.6. Geometry of the Hall-B tagging system. The electron beam hits the radiator and produces bremsstrahlung photons. Then the electrons are bent by a dipole field with their radius of curvature defined by their energy. These electrons are then detected by two sets of scintillator detectors [94].

The Hall-B photon tagging system can tag photons with energies between 20% and 95% of the incident electron energy E_0 , and is capable of operating with electron-beam energies up to 6.1 GeV [94]. The magnetic field of the tagger is set so that electrons which did not interact with the radiator are directed into a beam dump located in the floor of the experimental hall. On the other hand, electrons which produced bremsstrahlung photons are dispersed in the tagger according to their momentum and detected by two detector planes.

The first detector plane (referred to as E-counters) is used to determine the energy of the scattered electrons by identifying their hit position in the hodoscope. The hit position is determined using 384 scintillator paddles which are 20-cm long, 4-mm thick, and range from 6 to 18 mm in width. The variable widths are used to ensure that each scintillator paddle covers a constant energy interval of $0.003 \times E_0$. However, the 384 paddles

are positioned to overlap each other by a third of their width on each side, creating 767 different tagger channels with a resolution of $0.001 \times E_0$ [94].

The second detector plane (referred to as T-counters) is used to determine the timing of the incident electrons. This plane is shifted by 20 cm downstream of the E-counters plane and consists of 61 scintillators, each 2-cm thick. The T-counters are organized in two groups. The first 19 counters correspond to photon energies of $\geq 0.75 \times E_0$, whereas the remaining scintillators span the energy range from $0.20 \times E_0$ to $0.75 \times E_0$. The width of the scintillators in each group varies to compensate for the $1/E_\gamma$ behavior of the bremsstrahlung cross section, in order to maintain a uniform counting rate between the different scintillators. The separation of the T-counters into two groups (with the first group considerably smaller in size) is done to allow experiments interested in the highest-energy photons to maintain adequate electron-photon matching. The T-counters overlap each other by a few millimeters to ensure that there are no gaps between the counters where electrons could escape undetected. There are two photomultiplier tubes attached to each end of each T-counter, resulting in a timing resolution of 110 ps [94]. This resolution of the tagger allows for determining unambiguously the photon that initiated the nuclear reaction in about 80% of all collected events.

3.2.3 CLAS

The CEBAF Large Acceptance Spectrometer (CLAS) provides particle identification and momentum determination for experiments using electron and photon beams for a big portion of the full solid angle ($\sim 2.8\pi$). Figure 3.7 shows the different detectors that comprise the CLAS. The torus magnet produces a magnetic field primarily in the ϕ direction and allows for momentum spectrometry. The momentum is determined from the curvature of the particle's trajectory in the CLAS. The latter is obtained from tracking the particle through three regions of drift chambers located within the area of the magnetic field. Fast information for event triggering in photon-beam experiments is provided by the start counter surrounding the target (not shown in the figure), and time-of-flight measurement is provided by a set of scintillators outside the drift chambers. The Čerenkov counters are used to distinguish between pions and electrons (not used in g13) and the electromagnetic calorimeters provide neutral particle identification [81]. A summary of the

CLAS characteristics is given in Table 3.1. A description of each system is provided in the sections below.

Figure 3.7. The CEBAF Large Acceptance Spectrometer (CLAS) in Hall B. It is composed of three regions of drift chambers (DC), electromagnetic calorimeters (EC), Čerenkov counters (CC), and time-of-flight scintillators (TOF). (A) Cross section of the CLAS along the beamline, with typical charged-particle tracks shown. (B) Cross section of the CLAS perpendicular to the beamline. The mini-torus shown in the middle is used only for electron runs [81].

Torus magnet

The torus magnet is composed of six superconducting coils (dashed lines in Fig. 3.7A) which separate the detector into six independent magnetic spectrometers (sectors). The magnetic field produced by the magnet bends charged particles, which are then tracked through the three regions of drift chambers for momentum determination. Figure 3.8 shows the magnetic field produced by the superconducting torus coils. The coils are approximately five meters in length and they are wound into a kidney shape using 216 layers of aluminum-stabilized NbTi/Cu wire. The shape of the magnet is such as to provide optimum curvature for the highest-momentum particles that are emitted at forward angles. Thus, at forward polar angles the magnetic field is the highest, whereas at backward polar angles the magnetic field is lower. When the magnet operates at the maximum design current of 3860 A, the integral magnetic field reaches 2.5 T·m in the forward region, and drops to 0.6 T·m at a 90° polar scattering angle. Since the center of the CLAS is relatively

Table 3.1. Summary of the CLAS characteristics [81].

Capability	Quantity	Range
Coverage	Charged-particle angle	$8^\circ \leq \theta \leq 140^\circ$
	Charged-particle momentum	$p \geq 0.2 \text{ GeV}/c$
	Photon angle (4-sectors)	$8^\circ \leq \theta \leq 45^\circ$
	Photon angle (2-sectors)	$8^\circ \leq \theta \leq 75^\circ$
	Photon Energy	$E_\gamma \geq 0.1 \text{ GeV}$
Resolution	Momentum ($\theta \lesssim 30^\circ$)	$\sigma_p/p \approx 0.5\%$
	Momentum ($\theta \gtrsim 30^\circ$)	$\sigma_p/p \approx (1 - 2)\%$
	Polar angle	$\sigma_\theta \approx 1 \text{ mrad}$
	Azimuthal angle	$\sigma_\phi \approx 4 \text{ mrad}$
	Time (charged particles)	$\sigma_t \approx (100 - 250) \text{ ps}$
	Photon energy (EC)	$\sigma_E/E \approx 10\%/\sqrt{E}$
Particle ID	π/K separation	$p \leq 2 \text{ GeV}/c$
	π/p separation	$p \leq 3.5 \text{ GeV}/c$
	π^- missidentified as e^-	$\leq 0.01\%$
Luminosity	Electron beam	$L \approx 10^{34} \text{ nucleon cm}^{-2} \text{ s}^{-1}$
	Photon beam	$L \approx 5 \times 10^{31} \text{ nucleon cm}^{-2} \text{ s}^{-1}$
Data Acquisition	Event rate	4 kHz
	Data rate	25 MB/s

Figure 3.8. (A) Magnetic field strength in the plane between two of the torus coils. The magenta lines correspond to the projection of one of the coils onto the midplane. (B) Magnetic field vectors in the cross section of the CLAS perpendicular to the beamline. The magenta blocks indicate the location of the six torus coils [81].

field-free, the torus magnet also allows for experiments with polarized targets [81].

Drift chambers

The drift chambers are used to map the trajectories of the charged particles passing through the magnetic field, and thus determine the particle's momentum. As mentioned before, the torus magnet separates the CLAS into six sectors. Each sector consists of three regions of drift chambers, located concentrically around the target, inside and outside the torus magnet (for a total of 18 different drift chambers). Region-1 drift chamber surrounds the start counter and is located in a low magnetic field. Region-2 drift chamber is located in the area of the highest magnetic field between the coils of the magnet, while Region-3 covers the area just outside of the coils (Fig. 3.7). The trapezoidal cross section of each chamber is designed to fill the sector volume between two coils, spanning 8° to 142° in polar angle. The endplates of each drift chamber are tilted by 60° so that they are parallel to the torus coils (and with adjacent drift chambers) resulting in a coverage of about 80% of the azimuthal-angle range (Fig. 3.7 (B)) [94].

The wire arrangement in each of the drift chambers forms hexagonal patterns, with the field wires (wires maintained at high negative potential) located at the vertices of the hexagons and a sense wire (maintained at high positive potential) located at the center of each hexagon (Fig 3.9). For each region of drift chambers, the wires are grouped into two

Figure 3.9. Representation of a Region-3 sector drift chamber showing the layout of its two superlayers. The field wires are located at the vertices of the hexagonal pattern, and the sense wire at the center. The highlighted cells indicate the passage of a charge particle that fired the corresponding drift cells [81].

superlayers, each consisting of six wire cells. The first superlayer of Region-1 drift chamber consists of only four layers of wire cells due to space constraints. The wires in the first superlayer in Region-2 and Region-3 drift chambers, referred to as axial superlayer, are aligned perpendicular to the magnetic field's bending plane, to best measure the track's curvature. The second superlayer in each of these two regions, referred to as the stereo superlayer, is tilted by 6° with respect to the first to provide azimuthal information. Region-1 drift chambers have the opposite configuration, with the stereo superlayers coming before the axial superlayers [81].

The detection of particles that pass through the drift chambers is done by ionizing the gas that fills the drift chambers, and accelerating these ions inside a strong electric field produced by the potential difference between the field and sense wires. When the electrons are produced and then accelerate through the electric field, they register a signal (current)

through the sense wires. If the electric field is strong enough, the electrons will gain enough energy to produce a cascade of ionization, amplifying the signal. The potential difference between the field and sense wires is kept at about 2000 V. The gas mixture used for the CLAS drift chambers is 90% argon and 10% carbon dioxide [81].

The geometry of the individual hexagonal cells allows a spacial resolution of 300-350 μm . The set of 10–12 wire layers per region increases the resolution to 100 μm in the bend plane of the torus field and 1 mm in azimuth. The resulting momentum uncertainty is $\delta p/p < 0.5\%$ and the angular resolutions $\delta\theta$, $\delta\phi$ are about 2 mrad [84].

Start counter

As mentioned before, the start counter is used to determine the time at which nuclear reactions occur in the target, for photon-beam experiments. The start counter surrounds the target and is the first detector encountered by a particle produced within the target through its path into the CLAS. The information from the start counter is used for event triggering and reconstruction. Specifically, it is used to synchronize the particles detected in the CLAS with the tagged photons that produced the reaction.

Figure 3.10. Representation of the start counter [95].

The start counter is composed of 24 scintillator paddles that are arranged az-

imutally in six sectors (four in each sector) around the target [82]. The paddles are used to measure the time at which tracks pass through them by collecting the radiated photons induced from ionization. The photons are collected by photomultiplier tubes that are attached on the upstream end of the scintillator paddle. Figure 3.10 shows the start counter, with the four different scintillator paddles in each sector and the photomultiplier tubes attached upstream. Figure 3.11 indicates the size of the start counter with the magenta line indicating a target cell of 40 cm in length.

The configuration of the start counter allows coverage of the entire azimuthal angle range and a polar angle range between $10^\circ \leq \theta \leq 140^\circ$, with these limits depending on the length and position of the target. The latter angular coverage is about the same as the one provided by the TOF scintillators (see below). Due to the fact that there is only one photomultiplier tube attached to each paddle, the timing resolution of the start counter is only ~ 250 ps. However, this resolution is accurate enough to allow the identification of the photon. The identification is done by comparing timing information from the tagger hits to the time the particle is detected in the start counter [95]. This way one can identify which beam bunch initiated the reaction, and then identify, using similar information, which photon initiated the reaction (see Analysis). In electron-beam experiments, the start counter is unnecessary, and thus removed since the start-time information can be obtained with the Čerenkov counters.

Čerenkov counter

The Čerenkov counter detection system is primarily used for electron-pion differentiation, and it covers the forward polar angles up to 45° . The detector is located outside of the drift chambers and before the time-of-flight detector system as shown in Figure 3.7. This system is a set of 36 Čerenkov counters for each sector (for a total of 216), splitting the polar angle of each sector into 18 regions, and the azimuthal in two regions, each with its own Čerenkov counter. They are filled with C_4F_{10} , which is a high-reflective-index gas ($n=1.00153$) that allows the emission of Čerenkov radiation [81]. Čerenkov radiation is the radiation emitted when a charged particle passes through a medium with a speed larger than the speed of light inside that medium. The C_4F_{10} gas mixture provides a high threshold for pions ($2.5 \text{ GeV}/c$), which aids in discriminating pions from electrons.

Figure 3.11. The start counter cross section indicating its dimensions and polar angle coverage with a 40 cm target cell [95].

Figure 3.12 shows a cross section of an individual Čerenkov counter. The elliptical and hyperbolic aluminum mirrors shown in this figure focus the emitted photons through a light collection cone onto a PMT. The over-all efficiency of the Čerenkov counters is above 99.5%, with the limiting factor being the imperfect reflectivity of the mirrors [86]. Electron–pion differentiation is not needed in the g13 experiment, since a photon beam was

Figure 3.12. A cross section of an individual Čerenkov counter showing an electron track (in orange) and Čerenkov radiation (blue dashed lines) [81].

used. For this reason the Čerenkov counters were disabled.

Time-of-flight system

The time-of-flight system provides timing and energy information for the particles detected in the CLAS. It consists of a series of scintillator paddles, positioned together to form a shell outside of the drift chambers (see Fig. 3.7) [85]. These plastic scintillators span the same solid angle as the drift chambers, between 8° to 142° in polar angle, and about 80% of the azimuthal angle. Just like all the other detectors comprising the CLAS, the time-of-flight detector is divided in 6 sectors, with sectors 1, 2, and 4 containing 57 scintillator paddles each, and sectors 3, 5, and 6 with 56 paddles each, due to space constraints. These paddles vary in length across the azimuth of the detector to cover as much space as possible. The 18 paddles positioned at backward angles are multiplexed into 9 logical paddles, resulting in a total of 48 different paddles in each sector [81].

Figure 3.13 shows the time-of-flight paddles within a sector. When a charged particle passes through the scintillator, it excites the material which in turn radiates photons. These photons are then transmitted towards the end of the paddle and collected by light guides directing them at the photomultiplier tubes that are attached at each end. The timing resolution of the paddles varies between 150 ps to 250 ps, depending on the length of the scintillator [81].

Electromagnetic calorimeters

Just outside the time-of-flight system, the particles encounter the electromagnetic calorimeter that is used to detect neutrons and photons. This system consists of two detectors, the forward electromagnetic calorimeter (EC), which covers the polar angle range from 8° to 45° , and the large-angle calorimeter (LAC), which covers the polar angle range from 45° up to 75° (Fig. 3.7). There are only two segments of the large-angle calorimeter covering only two sectors of the CLAS (total coverage of 120° in azimuthal angle) [96].

The forward electromagnetic calorimeter consists of 39 alternating layers of scintillators and lead sheets, with a total thickness of 16 radiation lengths (see Fig. 3.14) [87]. Each successive layer is rotated through 120° with respect to the previous one, thus, geo-

Figure 3.13. A set of 56 scintillator paddles that comprise one sector of the time-of-flight system [85].

Figure 3.14. A diagram of the forward electromagnetic calorimeter. The successive layers of alternating orientation are shown [87].

metrically forming an array of triangular cells. These cells are used to locate the area of energy deposition. Particles in the electromagnetic calorimeter interact with the lead sheets emitting radiation which is then amplified by the scintillators and converted into a current by a series of photomultiplier tubes connected to the end of the scintillator paddles. For PMT pulses corresponding to a photoelectron yield of $n_{pe} = 10^3$, which was the expected signal from a 1-GeV electron, the amplitude resolution was $\sigma_A/A \approx 4\%$, and the time resolution was $\sigma_t \approx 100 - 150$ ps (measured with a pulsed, variable intensity UV laser) [87]. The large-angle calorimeter has a very similar design and is used to detect electrons and neutral particles scattered at larger angles. Since in this analysis the final-state neutron is reconstructed using the missing-mass technique, information from the electromagnetic calorimeters was not used.

3.2.4 Beamline detectors

Additional detector systems are placed along the Hall-B beamline (see Fig. 3.2) and are used for monitoring various beam characteristics. Specifically, three beam-position monitors (BPM), which are composed of three resonant (RF) cavities each, are positioned at different distances upstream from the target and are used to monitor the position and intensity of the electron beam. Two of them are placed upstream from the tagger and one is between the tagger and the target but is not used in photon experiments. Electrons that pass through these cavities induce an electromagnetic field which depends on the position and intensity of the beam, identifying the beam position coordinates (x, y) . This information is used to center the electron beam on the CLAS target during electron-beam experiments, and on the tagger radiator during photon-beam experiments [81].

A Faraday cup is located downstream at the beam dump. This device is used to measure the electron-beam current, and it consists of a 4000-kg block of lead placed on ceramic supports inside vacuum [81]. The electrons are stopped and deposit all their energy in this device, determining the incident electron flux. The electron flux is used in electron-beam experiments for cross-section calculations.

The profile of the electron beam delivered to Hall B is determined by three devices referred to as harps. Each of these devices consists of a set of wires (20 μm and 50 μm tungsten, and 100 μm iron) oriented along the two orthogonal directions perpendicular to

the beam. During beam-profile monitoring, the harps are moved into the beamline with a direction of motion of 45° with respect to the horizontal axis, taking measurements at regular intervals. Harp scans are generally performed during the initial beam extraction at the beginning of each experiment, and after changes in the electron-beam configuration, since this monitoring is destructive to the beam quality [81].

In polarized-beam experiments, the longitudinal polarization of the electron beam is measured using the Hall-B Møller polarimeter. This measurement is usually done in experiments where the polarization of the electron is of interest, such as electron-beam experiments or photon experiments that use circularly polarized beams (see Sec. 3.2.1). The Møller polarimeter consists of a $25\text{-}\mu\text{m}$ thick cobalt-iron alloy foil, a set of quadrupole magnets downstream of the foil, and two detectors. For measurements, the foil is inserted into the beamline upstream of the photon tagger, tilted 20° from the beamline. The scattered electrons are bent by the two magnets into the detectors, and the asymmetry between the number of electrons scattered in each of the detectors yields the polarization of the electron beam [81].

In photon-beam experiments, the position and the size of the photon beam is measured by the photon-profile monitor [97]. This device consists of an array of crossed scintillator fibers, 2-mm wide, located downstream of the CLAS. These fibers respond to the passage of photons, identifying the photon beam profile (see Fig. 3.2) [81]. The photon flux is determined by two independent methods: a pair spectrometer positioned upstream of the target and a method called g-flux that uses the “out-of-time” electron hits in the tagger. The pair spectrometer is used to detect electron-positron pairs produced by a thin foil, identifying in this way the photon flux [81]. This procedure, removes approximately 1% of the photons from the beam, but provides crucial information, especially for cross-section analyses. Even though data from the pair spectrometer were collected during the g13 experiment, the photon flux is determined using the g-flux method. In this method, the photon flux is extracted by determining the transmission probability through the collimator using the good electron hits that do not trigger the CLAS (“out-of-time” hits), and the tagging ratio of each T-counter [98].

3.2.5 The g13 experiment

The g13 experiment [79] contains two running periods. The first period, called g13a that ran between October and December of 2006 with an additional week in March 2007, took data using a circularly-polarized photon beam, whereas the second period, called g13b that ran from mid-March through June 2007, used a linearly-polarized photon beam. Throughout the experiment, the target used was a 40-cm-long unpolarized liquid-deuterium target. A current of -1500 A in the torus magnet produced a magnetic field that bent negatively-charged particles away from the beamline. The current was chosen such as to maximize the CLAS acceptance for low-momentum π^- . A total of 50 billion events were collected during both running periods. All the data occupy about 130 TB of disk space [99].

Two different electron-beam energies of 1.990 GeV and 2.655 GeV were used during g13a, collecting about the same number of events. The electron-beam polarization was $\sim 84\%$ for the 1.990 GeV setting and $\sim 78\%$ for the 2.655 GeV setting. In order to maximize the photon polarization, photons with energies 1.990 GeV and 2.655 GeV were chosen for the various analyses (see Sec. 3.2.1). The tagged photon rate at the target was $\sim 10^7$ Hz [99].

For g13b, eight different electron-beam energies were used (between 3.3 GeV and 5.16 GeV) and data were collected for six different positions of the coherent photon peak at steps of 200 MeV, namely 1.1 GeV, 1.3 GeV, 1.5 GeV, 1.7 GeV, 1.9 GeV, 2.1 GeV, and 2.3 GeV [99]. This was done to maximize the photon polarization at each photon energy (see Sec. 3.2.1). At each coherent-peak position, the photon-polarization vector was rotated between two orthogonal orientations; one parallel to the Hall-B floor referred to as *Para* and one perpendicular to the floor, referred to as *Perp*. The position of coherent peak and the polarization of the photon beam was adjusted by accurate changes in the orientation of the diamond crystal that was mounted on the goniometer. At the time of the experiment, the photon polarization was estimated to vary between 70% and 90%. Due to the method used to produce the linearly polarized photon beam, the incident photon flux calculation is very difficult, if not impossible for g13b.

Target

The liquid-deuterium target used in the g13 experiment is shown in Fig. 3.15. The 40-cm long target had a slightly conical shape with a diameter at its widest point of 4 cm and a density of 0.1625 g/cm^3 [79]. The dimensions of the target yielded high luminosities while maintaining an acceptable rate of electrons in the tagger. In addition, the diameter of the target allowed for sufficient coverage of the $\sim 1\text{-cm}$ spread of the photon beam, as well as for any slight beam-position misalignments. The center of the target was positioned 20 cm upstream from the center of the CLAS to maximize acceptance for forward-going particles.

Figure 3.15. Conical 40-cm-long liquid-deuterium target, with a 4 cm diameter used in the g13 experiment [100].

Trigger and data acquisition

Event triggering is essential for every experiment using the CLAS to select events of interest for recording. During the g13 experiment, events were recorded at a rate of about 10 kHz, with a dead time of $\sim 15\%$. This time indicates the percentage of the time when the data-acquisition system (DAQ) is busy recording a physics event and does not accept any other events. There are two stages at which data-acquisition triggering can occur with the CLAS. The Level-1 trigger acts on the prompt signals from the various detector systems

and is used to set the timing of the gate signals for the analog-to-digital converters (ADCs) as well as the start signal for the time-to-digital converters (TDCs). The Level-2 trigger is based on pattern recognition of hits from the drift chambers and sends a fail signal if these tracks do not correspond to particle information from Level 1. For the purposes of the g13 experiment, only Level-1 triggering was used. Specifically, data acquisition during the g13a experiment required a coincidence between the start counter and the time-of-flight systems in at least two sectors of the CLAS. During the g13b run period, the DAQ required a coincidence between the start counter and the time-of-flight systems in at least one of the CLAS sectors. These requirements were chosen to minimize the background from accidental hits while maximizing the data acquisition rate for charged-particle events [99].

Chapter 4

Detector Calibration

The raw information from the different detector systems of the CLAS and Hall B consists of ADC and TDC channel IDs and digital numbers that are converted into real physics quantities, such as energy and time, via a process referred to as cooking. This CPU-intensive, time-consuming process involves, among other things, reconstructing particle tracks from the drift chamber hits that are necessary for the determination of the particle momenta. Other valuable information, like timing, hit position, and photon energy assignment, is also translated from raw data during this process. Thus, detector calibration plays a vital role in the cooking process, since only a good calibration can provide usable values for these physics quantities.

4.1 Overview

Due to the instability over time of the electronics connected to the CLAS detector, the response of each detection system varies throughout an experiment such as g13. Instabilities include, among other things, gains of photomultiplier tubes that drift, individual electronic channels, or even entire modules that become inoperational for some time, and problematic drift chamber wires that break. Therefore, calibration of the CLAS has to be done in a way such that these instabilities are accounted for. On the other hand, there are in total $\sim 6 \times 10^5$ g13 data files, which is too large a number to individually calibrate each file. Therefore, for initial calibration, the running period was divided into subperiods in which the performance of the detector was relatively stable (based on notes taken during the experiment) and overall experimental conditions were constant. For each such subperiod, one run was chosen to be calibrated (“golden” run) and then the calibration constants were applied to all the other runs from the corresponding subperiod. The calibration constants for each detector system are adjusted in an iterative procedure. A team of collaborators put in typically several hundreds of hours to complete this time-consuming task. In the g13 experiment, Edwin Munevar from the George Washington University carried out the drift chamber calibration and Paul Mattione from Rice University carried out the elec-

tromagnetic calorimeter calibrations. Two Ph.D. students from Idaho State University, Chuck Taylor and Danny Martinez, carried out the final start counter and time-of-flight calibrations, respectively. The tagger calibration was carried out by the author.

In general, the calibration of the different systems is done in a certain order, since measurements of one detector system often depend on measurements of a different system. In addition, as mentioned above, the calibration process is iterated a number of times until all systems are considered to be satisfactorily calibrated. Initially, only the “golden” runs are used to calibrate the CLAS. The next iteration includes dividing each subperiod into smaller sub-subperiods, improving the calibration constants even further. This process is then repeated until the calibration of all runs reaches a satisfactory level. In the case of the tagger calibration, a typical subrange had a size of 10 runs.

The importance of a well calibrated set of data will become more obvious in Chap. 5, where the determination of various quantities characterizing the produced particles is crucial for correctly determining physics observables. For example, particle identification is done by comparing the particle’s velocity calculated using information from the time-of-flight system with the velocity of the particle calculated using drift chamber information (see Sec. 5.3.2). Also, timing information from the tagger system, along with timing information from either the start counter or the time-of-flight system, is needed to correctly identify the photon that initiated the reaction and reconstruct the complete kinematics of an event. Therefore, poor calibrations of the systems mentioned above would cause, for example, misidentification of particles or assignment of the wrong photon to the event reconstructed within the CLAS. The sections below outline the calibration procedures for the various detector systems.

4.2 Timing Calibration

The calibrations described here refer to calibrations of various systems that are used to extract particle timing information. The timing information from various detectors in the CLAS, such as the start counter and the time-of-flight system, as well as from detector systems outside the CLAS, such as the tagger, is essential to correctly identify the reaction of interest. Common features that need the attention of the calibrators for all

these systems are described below.

4.2.1 TDC timing calibration

The channel-to-time calibrations of all TDCs is performed using a special pulser run. This is done via a dedicated pulser system of each detector, and during this run, pulsed logic signals are delivered simultaneously to all TDC channels, in sets of 50. The time interval between the pulses in each set is increased by approximately 2 ns [101]. The information collected from this run is then compared to the information fed by the pulser, allowing the determination of the TDC response over a wide range of timings. Specifically, the response for each TDC is determined by studying the raw data with respect to the pulsed data and fitting this with a predefined function. In the case of the time-of-flight TDCs, the predefined functions are quadratic polynomials:

$$t = c_0 + c_1T + c_2T^2, \quad (4.1)$$

where T is the raw time in units of TDC channels and t is the converted time in nanoseconds. The polynomial fit identifies the three constants, c_0 , c_1 , and c_2 , for each time-of-flight TDC that are used for the channel-to-time conversion.

4.2.2 Time-walk corrections

The TDCs in the CLAS are fed with the logic output signals from leading-edge discriminators (LED). The logic signal in LEDs is generated at the moment the analog pulse crosses a pre-determined threshold. In leading-edge discriminators, a problem arises when the input analog signals can have a wide range of amplitudes. Since the rise time of the analog pulse depends primarily on the detector characteristics, it is approximately constant for all signals with different amplitudes. Therefore, pulses with different amplitudes will cross the pre-determined threshold at different times, introducing a time difference, Δt , in the output signal of the LED. Figure 4.1 shows how two analog pulses with different amplitudes cross the pre-determined threshold at different times, introducing the amplitude-dependent time-difference, referred to as “time walk”. The “time-walk” correction is usually obtained by fitting the correlation between timing and pulse heights

Figure 4.1. Example of time walk. Two signals with identical rise times but different pulse heights reach the threshold value in the discriminator at different times.

for various signals with a pre-defined function. For example, the predefined function in start-counter time-walk calibration has the form:

$$\Delta t^w = W_0 + \frac{W_1}{A - W_2}, \quad (4.2)$$

where A is the ADC channel, and W_0 , W_1 , W_2 and are the time-walk calibration constants that are determined by the fit.

4.2.3 Event start time

As mentioned before, timing information is crucial for particle identification and event reconstruction. For photon experiments, timing calibration becomes more difficult since the accelerator RF time is not directly used. Instead, times from the tagger detector system are used and the tagger calibration becomes very important. The RF signal, which drives the CEBAF accelerator, indicates the time of each electron beam bucket. The

CEBAF accelerator delivers the electron beam in buckets separated by 2.004 ns. The electron that produced the photon that initiated an event in the target can be identified after the tagger time calibration, where the tagger time is correlated with the RF time. Once this is done, subsequent detector systems are calibrated with respect to the tagger system, ensuring that timing is measured by each system from a common zero time, t_0 .

4.3 Tagger Calibration

As mentioned above, the tagger calibration constitutes a crucial part of the data handling since a good calibration allows the substitution of the tagger time with the more accurate RF time. An error of one RF bunch in the tagger calibration, i.e. 2.004 ns delays on the tagger time, may introduce variations of several hundreds of MeV in the particle reconstructed mass. Figure 4.2 shows the T-counter TDC timing spectrum of the tagger against the RF signal, indicating the 2-ns interval between consecutive beam buckets.

Figure 4.2. Tagger TDC timing spectrum, showing the timing resolution of the tagger T-counters (~ 110 ps). The signals in the T-counter TDCs are timed against the RF signal that accelerated the electron beam. The 2 ns beam bucket structure is also indicated. The small peaks just before the main ones, corresponds to spills from the million times more intense beam being sent to the other experimental Halls [94].

For the proper calibration of the tagger system, there are four sets of constants

that need to be adjusted. These constants are used for the channel-to-time conversion. The times and hit positions are then used by the tagger reconstruction software to calculate photon energies and ultimately to correctly identify the photon that initiated the reaction in the target. Specifically, for an electron hit to be reconstructed as valid, it has to satisfy the following requirements: hits must be found in both the left and the right TDCs of one T-counter (unless during calibration one of the TDC is flagged as bad); these hits must match in time with a hit in the TDC of one E-counter (referred to as time matching); the time-matched detectors must be on an electron trajectory allowed by the tagger-magnet optics (referred to as geometric matching). The final stage of electron reconstruction involves re-binning of the hit positions which takes into account hits from adjacent E-counters that overlap (see Sec. 3.2.2). Due to the overlap between counters, electron hits in adjacent counters are reconstructed as a single hit (one electron going through two overlapping E-counters) if their timing is close enough. The re-binning procedure effectively increases the energy resolution of the hodoscope.

A different set of constants is needed to calibrate each of the following:

- TDC slopes
- Base peak position
- RF fine adjustment
- Global tagger-time offsets with respect to the CLAS

A description of the above calibrations follows.

All of the constants depend on each other, and therefore must be determined in a specific order. First, the TDC slopes are determined. The TDC slope is the constant c_1 from Eq. (4.1). The E-counter TDC slopes are fixed at 500 ps/channel, whereas the T-counter TDC slopes are close to 50 ps/channel. The latter vary slightly from counter to counter, and the goal of this first tagger calibration step is to find the correct slope for each T-counter TDC, both left and right. The slopes for the left/right TDCs on the same T-counter also need to be balanced so that the reconstructed electron timing does not depend on the position of the hit. The left/right TDC balance is visualized by plotting the raw TDC time difference between the left and right photomultiplier tube of each individual

Figure 4.3. Tagger calibration plots showing the raw TDC timing difference $t_{left}^6 - t_{right}^6$ as a function of their sum, $t_{left}^6 + t_{right}^6$, for T-counter 6 in a run from g13b. The left plot corresponds to a poorly calibrated run and the right plot shows the same correlation after the run was calibrated.

T-counter, $t_{left}^i - t_{right}^i$, as a function of their sum, $t_{left}^i + t_{right}^i$. Poor calibration of the TDC slope balance between the left and right TDC can affect the resolution of each individual T-counter. Figure 4.3 shows the difference $t_{left}^i - t_{right}^i$ as a function of $t_{left}^i + t_{right}^i$ for T-counter 6, of a well calibrated run (right) and a poorly calibrated run (left).

The base-peak-position calibration is to find the centroid of the peak in the TDC spectra for each individual TDC. The raw timing spectrum of each TDC consists of a sharp peak over a continuum. The peak corresponds to self-triggered hits (*i.e.* hits in the T-counter that produced the trigger), whereas the continuum corresponds to hits when other T-counters produced the trigger. By convention, the amount of time to which the sharp peak corresponds is subtracted from the TDC value in calculating the tagger time, so the absolute value of the tagger time is close to zero relative to the trigger time. The same procedure is also applied to the E-counters to ensure correct identification of the reconstructed electron, based on time coincidence.

The next step in the tagger calibration procedure involves the fine tuning of each T-counter to the RF signal. The purpose of this calibration is to correlate the timing of the photon at the target as measured by the T-counters with the more precise RF information. The RF time, t_{RF} , should be understood as the time of an undetermined RF bucket relative to the trigger. Specifically, the RF time is given by the time of the photon being produced

(relative to the trigger), t_{pho} , plus a timing offset corresponding to an integer number of beam buckets k_{event} , which varies from event to event:

$$t_{RF} = t_{pho} + k_{event} \times 2.004 \text{ ns.} \quad (4.3)$$

A good fine-tuning calibration allows the determination of k_{event} on an event-by-event basis, through the calculation of the tagger–RF offset parameter C_i (one for each of the 121 tagger channels). The determination of k_{event} then enables the value of interest t_{pho} to be determined from the RF time, t_{RF} . The tagger–RF offset parameter C_i can be decomposed into an overall phase offset, C_i^{RF} , for each tagger channel i relative to the RF signal, plus an integer number of beam buckets, k_i , which depends on the tagger channel in question:

$$C_i = C_i^{RF} + k_i \times 2.004 \text{ ns.} \quad (4.4)$$

Therefore, by constructing the time difference, $t_i - t_{RF}$, between the tagger channel in question and the RF time, the overall phase offset C_i^{RF} can be extracted:

$$\begin{aligned} t_i - t_{RF} &= t_i - t_{pho} - k_{event} \times 2.004 \text{ ns} \\ &= C_i - k_{event} \times 2.004 \text{ ns} \\ &= C_i^{RF} + (k_i - k_{event}) \times 2.004 \text{ ns.} \end{aligned} \quad (4.5)$$

Specifically, the 121 offsets C_i^{RF} for each tagger channel are extracted from Gaussian fits to the distributions $(t_i - t_{RF}) \bmod 2.004 \text{ ns}$ (see Fig. 4.4). In order to calculate the C_i offset, the integer k_i needs to be determined (see Eq. (4.4)). To determine k_i for each tagger channel, a timing comparison between the channel of interest and a reference detector is made. The reference detector could be any detector whose timing is correlated to the production of the tagged photon. In g13 the start counter was used as the reference detector. The timing offset, C_i^{ST} , between the tagger channel time, t_i , and the time of a hit in the start counter, t_{ST} , is first determined:

$$C_i^{ST} = t_i - t_{ST}. \quad (4.6)$$

This is done by Gaussian fits to the time difference distribution, $t_i - t_{ST}$ (see Fig. 4.7). Next, the difference between the offset constants C_i^{ST} and C_i^{RF} is used to determine k_i for

each tagger channel:

$$C_i^{ST} - C_i^{RF} = K_{det} - k_i \times 2.004 \text{ ns}, \quad (4.7)$$

where K_{det} is a detector-specific offset between the two constants. Specifically, K_{det} is determined from a fit to the distribution of $(C_i^{ST} - C_i^{RF}) \bmod 2.004 \text{ ns}$ summed over all channels, and k_i is determined from Eq.(4.7). Using the values of C_i^{RF} and k_i , the integer k_{event} is calculated from Eq. (4.5).

Figure 4.4 shows the difference between the tagger channel time, t_i , and the RF time, t_{RF} for a particular T-counter, before (left) and after (right) calibration. As mentioned above these plots are used for the extraction of the offset constant C_i^{RF} . Figure 4.5

Figure 4.4. Individual T-counter fine tuning. The plots correspond to the difference between the tagger channel time, t_i , and the RF time, t_{RF} for tagger channel 71 before (left) and after (right) calibration for a given run. The fit corresponds to a Gaussian that is used for the extraction of the mean and thus the adjustment of the RF fine tuning constants.

shows the time difference, $t_i - t_{RF}$ as a function of tagger channel for a given run before and after calibration. As mentioned above, a key part of the tagger calibration is to make sure that the calibration has the same quality for all runs adjacent to the golden run. If the quality of the calibration is not the same, then the run range around the golden run is split in two and the calibration is redone until its quality reaches a satisfactory level. The quality of the calibration is generally visualized by producing a timeline for each T-counter. The timelines show the mean value of the time difference between the tagger channel time

Figure 4.5. Difference between the tagger channel time and the RF time, $t_i - t_{RF}$, as a function of tagger channel for a given run. A proper calibration of the fine-tuned constants aligns all T-counters at 0 ns. The line at 0 ns is shown only to guide the eye.

Figure 4.6. Timelines for T-counters between 73 and 84. The calibration was done for every 10 runs and the run index on the X axes corresponds to one randomly chosen run from the calibrated run interval of ten runs. One can see that the mean values of $t_i - t_{RF}$ are well within the tolerance limit of 30 ps.

and the RF time, $t_i - t_{RF}$, as extracted from the Gaussian fit shown in Fig. 4.4, as a function of run number. Figure 4.6 shows these timelines after the final calibration iteration for g13b, for T-counters 73–84 (T-counters that typically detect electrons that produce highly polarized photons). A tolerance of ± 30 ps is set to the time difference $t_i - t_{RF}$ for the determination of C_i^{RF} . Figure 4.7 shows the time difference $t_i - t_{ST}$ as a function of tagger

Figure 4.7. Time difference between tagger and start counter as a function of T-counter, with T-counter 95 clearly mis-calibrated.

channel that is used to determine the offset constant C_i^{ST} mentioned above.

In addition to determining the calibration constants, the calibrator also has the responsibility to flag bad left or right T-counter photomultiplier tubes. As mentioned before, the tagger reconstruction software requires hits in both left and right photomultiplier tubes attached to each T-counter. Thus, when one of the two photomultiplier tubes is unstable or dead, the electron fails to be reconstructed unless the bad photomultiplier tube is flagged as bad. Once a photomultiplier tube is flagged as bad, the reconstruction software ignores the flagged photomultiplier tube and reconstructs the electron from a hit only in the other photomultiplier tube. On one hand, this reduces the resolution of a specific T-counter, but on the other hand, it allows for electrons to be reconstructed that otherwise would have been lost. In g13, photomultiplier tubes were flagged as bad by studying the number of hits recorded by each T-counter. Specifically, if the events recorded in photomultiplier tubes

Figure 4.8. T-counter hits as recorded in left (left plot) and right (right plot) photomultiplier tubes as a function of T-counter. The difference between the events recorded in the first 19 T-counters and the T-counters 20-61 is due to the smaller size of the paddles in the former (see Sec. 3.2.2). Flagged photomultiplier tubes for this run are left-48, left-61, and right 49.

of one end in adjacent T-counters differed by more than 50% (see events in T-counter 65 and 66 left photomultiplier tubes), then the photomultiplier tube with the smaller event count was flagged as bad. Figure 4.8 shows the events recorded in the left photomultiplier tube (left plot) and in the right photomultiplier tube (right) as a function of T-counter for a given run.

4.4 Start Counter Calibration

The start-counter calibration affects the reconstructed event time in the CLAS, and, thus, affects the correlation between the event in the CLAS and the tagged photon that produced that event. As mentioned in Sec. 3.2.3, the start counter consists of 24 scintillator paddles, and one of the most important steps of the start-counter calibration is to align the times of the 24 paddles to each other. First, the channel-to-time TDC conversion is done using data from the pulser. The raw TDC spectra are fitted with a function

$$t_{ST} = C_T T, \quad (4.8)$$

where t_{ST} is the start-counter time in nanoseconds, T is the TDC channel, and C_T is the channel-to-time calibration constant determined from the fit. Following the TDC calibration, the hit time in the start counter is aligned to the RF-adjusted tagger time. In this step, several effects are corrected for, such as time delays, time-walk corrections, and the time it takes the light to travel through the scintillator paddle to the light guide. The time-delay constant corrects for the time difference between the RF time and the start counter time for each paddle. The time-walk corrections are described in Sec. 4.2.2, and finally, the correction for the position of the hit takes into account the non-linear geometry of each paddle.

Figure 4.9. Time difference between the event vertex time from the start counter and the event vertex time from the RF signal as a function of distance along paddle 2. The black points show the Δt peak positions in each d -bin, and the red line is the empirical-second-degree polynomial fitted to the distribution. The fit parameters determined from the fit are the calibration constants used to correct for the finite time-propagation of light through each scintillator.

Figure 3.11 shows the start counter, indicating the non-linear geometry of each paddle, which consists of a straight rectangular “leg” followed by a tapered forward “nose”. This design results in a non-linear relation between the hit position in the scintillator and the light-propagation time to the light guide. The empirical function used for determining the calibration constants, k_0 , and k_1 , and, thus, correcting for the time it takes the light to travel through the scintillator paddle to the light guide is,

$$\Delta t = \frac{z_0}{V_{eff}} + k_0 z_1 + k_1 z_1^2 + K_{RF}. \quad (4.9)$$

This function takes into account the effective velocity for light propagation, V_{eff} , through the “leg”. The variable z_0 corresponds to the distance of the hit position from the light guide along the “leg” part of the paddle and the variable z_1 corresponds to the distance from the “leg”–“nose” junction along the “nose”. At this point, the constant K_{RF} is treated as a free parameter. Figure 4.9, shows the time difference between the event vertex time as calculated using start-time information and RF information as a function of the distance along a start counter paddle. The black line corresponds to the fit of the function of Eq. (4.9) that is used to extract the calibration constants. The non-linear behavior of the “nose” of the start counter is clearly visible.

Figure 4.10. Event vertex time difference Δt^c for the four paddles in sector 6 of the CLAS, after the timing of each paddle is corrected for light propagation time along the scintillators. The average resolution extracted from the Gaussian fit is ~ 300 ps.

The final alignment of the start counter time to the RF time is done by using the propagation-time-corrected time difference, Δt^c , between the start-counter event vertex time and the RF event vertex time. Specifically, for each start counter paddle, this time difference is fitted with a Gaussian extracting the mean value which corresponds to the calibration constant K_{RF} . Figure 4.10 shows the event distributions over Δt^c for the four paddles in sector 6 of the CLAS. The width of the distributions indicates the ~ 300 -ps start-counter time resolution.

The time-walk corrections, performed at initial stages of the calibration, are established by studying the time walk as a function of the ADC channel. The latter is a

Figure 4.11. Time-walk, Δt^w , as a function of ADC channel. The upper left plot shows the time-walk for protons detected in paddle 5, where the upper right shows the time-walk for pions detected in paddle 6. The lower two plots show the time-walk for detected protons and pions in paddles 7 and 8 respectively. The black points correspond to the mean values of the time-walk distributions in each ADC channel extracted by Gaussian fits. The red line is the fit to the dependence of the mean value on the ADC channel used to establish the time-walk calibration constants.

measure of the pulse height. Specifically, the mean values of the time-walk distributions are obtained for each ADC channel by Gaussian fits (see Eq. (4.2) for function used). Then, the distribution of the mean values as a function of the ADC channels are fitted with an empirical function to determine the calibration constants. Figure 4.11 shows these fits. To have data covering a wide range of pulse heights, both pion and proton signals are used. Since pions are minimally-ionizing particles, they produce small pulse heights, while protons, due to their larger mass and lower momenta, deposit more energy and produce larger pulse heights. This difference between proton and pion pulse heights is indicated in Fig. 4.11. A more detailed description of the start-counter calibration can be found in Ref. [102]

4.5 Drift Chamber Calibration

Drift chamber calibration is focused on correctly reconstructing the particle tracks through the six superlayers in each sector. An initial track reconstruction is being done during the data collection of the experiment. In this first stage, individual tracks are fit only to wire hit positions in a procedure known as “hit-based” tracking. Specifically, wire hits are combined into track segments within individual superlayers and these segments are then grouped together using a “link table” to form tracks across the three regions within one sector. Due to the comparatively small size of the drift cells and the large number of wire layers, the particle momentum can already be reconstructed with a resolution $\sigma_p/p \sim 5\%$ using the hit-based tracking.

During the second stage of track reconstruction, timing information from other detectors of the CLAS is used to determine the electron drift times from the location of the hit to the sense wire in order to reconstruct precisely the hit position within each layer. Once the drift time is calculated, the drift position is established using pre-determined tables. These tables were generated from a GEANT 3 simulation and they take into account the position of each layer in the CLAS magnetic field and the angle of approach of the particle producing the hit. The hit positions in each drift cell are then fit in a procedure known as “time-based” tracking that yields the final track parameters. The final reconstruction results in a tracking resolution $\sigma_p/p \sim 0.5\%$. Several corrections need to be determined for the calculation of the drift time in the time-based tracking. Specifically,

the time-based tracking involves time-walk corrections, corrections for fixed time-delays for individual wires, channel-to-time calibration of the drift-chamber TDCs, and calculations of the flight time of the particle from the event vertex to the hit position in each layer.

Figure 4.12. Drift chamber residuals as a function of drift time. The residual is the difference between the distance of closest approach from the fitted track to the sense wire and the distance calculated from the drift time. This distribution is used to adjust the calibration constants for correctly identifying the hit position within a layer.

During the drift-chamber calibration, extra attention is paid to the dependence of the fit residuals of each superlayer on the drift time. A fit residual is the difference between the absolute value of the distance of closest approach (DOCA), D_{abs} , and the calculated distance, D_{calc} , from the sense wire to the track. The distance of closest approach is the distance from the fitted track to the sense wire, whereas the calculated distance is calculated from the drift time as well as the electric-field map within the layer. Figure 4.12 shows the residuals as a function of drift time for superlayer 3 of sector 1. This distribution is then fitted with a predefined function extracting calibration constants that are used by the “time-based” reconstruction. A more detailed description of the drift-chamber calibration and track reconstruction can be found in Ref. [103].

4.6 Time-of-Flight Calibration

The time-of-flight system is calibrated both for the time of particle arrival at a particular paddle as well as for the energy deposited by the particle in the respective paddle. The first step in the energy calibration of the time-of-flight system involves determining the pedestals for each ADC. These correspond to the response of ADC (in channel value) in the absence of a real hit, and represent the residual current present in the ADCs. The pedestal information is usually collected during a pulser run taken after any configuration changes. The pedestal information is then loaded in the time-of-flight ADCs, so that the ADCs can operate in pedestal-subtracted mode. The conversion for the time-of-flight ADC values to energy in MeV is:

$$E = \frac{10(A - P)}{A_{NMIP}}, \quad (4.10)$$

where A is the ADC value of the hit in channels, P is the pedestal, and A_{NMIP} is the pedestal-subtracted ADC value corresponding to a normally incident minimum ionizing particle (NMIP), which deposits 10 MeV in the scintillator material.

After TDC channel-to-time conversion (see Sec. 4.2.1), the time-of-flight calibration involves timing corrections. The timing corrections include three main adjustments: the time-walk corrections, the left-right balance correction between the TDC times from each end of the scintillator, and the alignment of the paddle hit times to the common zero time. The nature of the TDC calibration and the time-walk correction is explained in Sections 4.2.1 and 4.2.2. For the case of the time-of-flight system, the time-walk and the left-right balance corrections are established by using data collected with the pulser, where the signals pulse simultaneously at the central part of each scintillator paddle and at a diode. The diode provides the reference time to which the TDC time of the scintillator is compared. The corresponding pulse height, on the other hand, which is varied using a neutral density filter, is measured by the ADC. Then a fit of the measured TDC time as a function of the ADC value determines the time-walk corrections. The fit used to determine the three w_i for the time-of-flight time-walk correction is:

$$\Delta t^w = f^w \left(\frac{600}{V_T} \right) - f^w \left(\frac{A - P}{V_T} \right), \text{ with} \quad (4.11)$$

$$f^w(x) = \begin{cases} \frac{w_2}{x^{w_3}} & \text{if } x < w_0 \\ \frac{w_2}{w_0^{w_3}}(1 + w_3) - \frac{w_2 w_3}{w_0^{w_3+1}} & \text{if } x > w_0 \end{cases} \quad (4.12)$$

In Eq. (4.12), A corresponds to the ADC channel, V_T corresponds to the discriminator threshold in ADC channels, and P to the pedestal value. The pulser data is also used for the left-right balance correction. Specifically, calibration constants are adjusted such that the PMTs at either end of the scintillator paddle give the same timing, since the light that produces these events originates from the center of the paddle.

The next step in the time-of-flight calibration involves determining the attenuation length of each counter, which allows a correct determination of the energy deposited in the paddle, independent of the hit position. This is done in two steps. First, a measurement of the energy deposited in each counter, independent of the hit position, is established. Then, a comparison between the energy measured by each end of the paddle is done as a function of the hit position. The hit position is established by the difference in time measured by both ends of a particular paddle, using the effective velocity of the light propagating in the scintillator. The effective velocity is calculated using the dependence of the hit distance relative to the edge of each counter, established from the drift chambers by tracking, on the time measured in the TDC at that end of the counter. Typical values of the effective velocity in the scintillator paddles are ~ 16 cm/ns.

Finally, the times of the paddles are corrected such that all paddles in the time-of-flight system are aligned with each other. Figure 4.13 shows the time-of-flight paddle alignment for all paddles in Sector 5 before and after calibration. The alignment corrections are obtained from Gaussian fits to event distributions over the difference between the measured time-of-flight and the calculated time-of-flight for each paddle. The mean values of the fits are used for the correction. With a good calibration, the typical time resolution of the scintillation counters ranges from 150 ns to 250 ns between the smaller, forward-region counters and the longer, backward-region counters. A more detailed description of the time-of-flight calibration can be found in Ref. [101].

4.7 Electromagnetic Calorimeter Calibration

The final stage of detector calibration involves the calibration of the electromagnetic calorimeters. Since these were not required for the reconstruction of a proton and, thus, not used in this analysis, they are mentioned only for completeness. The main purpose

Figure 4.13. Time-of-flight paddle to paddle alignment. The plots show the difference in measured and calculated time as a function of paddle number for Sector 5, before (left) and after (right) calibration. The black points correspond to the mean values of the time difference as extracted from a Gaussian fit.

of the electromagnetic calorimeter is to detect and identify neutral particles, and specifically, to identify photons and neutrons. The neutral-particle identification is based on the time-of-flight of the particle as measured by the calorimeter. For this reason, the main calibration of the electromagnetic calorimeter is the time calibration. The latter involves timing corrections that are established using electron and charged pion data for which the time-of-flight could be easily obtained from other detector systems. The electromagnetic calorimeter timing calibration involves TDC calibration, time-walk corrections, as well as an overall detector offset. ADCs connected to the electromagnetic calorimeter are also calibrated so that the energy deposited by a particle can be reconstructed.

Chapter 5

Analysis I: Event Selection

5.1 Overview

The Bank Object System (BOS) format is used for the data acquired in Hall B during an experiment. The “raw” BOS files that contain all the detector-specific information such as ADC and TDC channels and detector status flags are kept on the JLab data storage silo. As mentioned in Chap. 4, during the cooking process, the “raw” information is converted to physical quantities using calibration parameters via the simple event builder (SEB) reconstruction package [104]. The SEB routines find groups of geometrically matched tracks and hits in the CLAS coordinate system and reconstruct the trigger particle and trigger time, and also determine the particles’ ID. New banks are then created that contain physical quantities. These are added to the data stream during cooking. These new banks containing information on the reconstructed events (such as momenta, particle ID, and timing information) are used in the subsequent physics analysis. Specifically, information from the BOS files that was of interest for the analysis of deuteron photodisintegration was saved in a format that allowed the use of an object-oriented framework for large-scale data analysis, named ROOT [105]. ROOT, which was developed at CERN, allowed the utilization of standard libraries for data-analysis procedures, including fitting routines and histogramming of results.

This chapter describes in detail the procedure to select the events of interest as well as to identify and reconstruct the reaction $\gamma + d \rightarrow p + n$. The general idea of the procedures followed is common for all experiments in Hall B, however, the implementation and adjustment of this procedures to deuteron photodisintegration events were carried out by the author.

5.2 Event Reconstruction

The first step in the event reconstruction took place during cooking. This first step determined a basic set of physical quantities for each event, such as the charge of

the particle, momentum, and velocity. A more sophisticated event reconstruction followed, with its steps described in detail in Secs. 5.3 – 5.4. An outline and a brief description of the basic event reconstruction follows.

5.2.1 Charge and momentum

The charge of the particle that has produced a track in the CLAS was determined from the direction of curvature of the reconstructed track through the toroidal magnetic field. The track was reconstructed using information from the drift chambers as outlined in Sec. 3.2.3. According to the Lorentz force law, a force, $\mathbf{F}$, will act on a charged particle with charge q , traveling through a magnetic field, $\mathbf{B}$, with velocity $\mathbf{v}$:

$$\mathbf{F} = q\mathbf{v} \times \mathbf{B}. \quad (5.1)$$

Using the curvature radius of the track, $\mathbf{r}$, and the magnetic field map of the CLAS, the momentum of the particle is obtained,

$$\mathbf{p} = q\mathbf{r} \times \mathbf{B}. \quad (5.2)$$

As mentioned in Sec. 3.2.3, the toroidal magnetic field is always approximately perpendicular to the velocity of the particles, maximizing the momentum sensitivity with smaller changes in the radius of curvature.

5.2.2 Velocity

The velocity of the particle is reconstructed from its time of flight, t_f . The particle's time of flight corresponds to the time between the event start time at the event vertex and the hit in the time-of-flight detector. The event start time is determined by the reconstruction software by comparing the hit times in the start counter to the RF electron beam bunch times. The distance traveled by the particle, d_f , between the event vertex and the hit position in the time-of-flight detector is determined from the reconstructed track length in the drift chambers. The velocity of the particle, as a fraction of the speed of light, c , is thus,

$$\beta = \frac{d_f}{c \cdot t_f}. \quad (5.3)$$

5.2.3 Mass

The rest mass of the detected particle, m_0 , is reconstructed from the two independent measurements of its momentum and its velocity,

$$m_0^2 = \frac{1 - \beta^2}{\beta^2} p^2. \quad (5.4)$$

The particle identification performed by the SEB package depends on the value of the reconstructed mass. Specifically, the particle identification performed by the SEB routines is determined from the best match between the velocity reconstructed using Eq. (5.3), β_{meas} , and the velocity calculated from the reconstructed momentum and an assumption about the particle's mass, β_{calc} . Specifically, β_{calc} is calculated as

$$\beta_{calc} = \sqrt{\frac{p^2 c^2}{m_i^2 c^4 + p^2 c^2}}, \quad (5.5)$$

where m_i is the nominal mass of the assumed particle. For the purposes of this analysis a more elaborated method for particle identification was used (see Sec. 5.3.2). This was done to better select protons and to reduce any ambiguities in the particle identification procedure.

5.2.4 Event vertex

The event vertex was reconstructed by a least-squares fit to the reconstructed tracks, minimizing the distance of the beam line to each track. The algorithm used is described in Ref. [106]. The multi-track vertex fitting routine assumes that the magnetic field is negligible in the vicinity of the target and that the particles' tracks in the area between the vertex and the start counter can be approximated by straight lines. In the case where only one track is reconstructed in an event, the event vertex was determined to be the point of closest approach of that track to the beam line. The beam line was determined during calibration by Gaussian fits to the reconstructed event vertex of events with multiple tracks. Figure 5.1 shows the Gaussian fits to the x and y event vertices used to establish the beam line position, and Fig. 5.2 shows the extracted mean and standard deviation of x (blue) and y (red) event vertex position as a function of run number.

Figure 5.1. Distribution of horizontal (left) and vertical (right) components of the vertex position of multi-track events, for all events from one representative run in g13b. The Gaussian fit is used to extract the mean and the standard deviation of the beam line position.

Figure 5.2. Mean (left), and standard deviation (right) of the x (blue) and y (red) components of the beam-line position as a function of run number. The mean values were used for the vertex reconstruction of events where one particles is detected.

5.3 Selecting $\gamma + d \rightarrow p + n$ Events

An initial filter was applied to the data via a process referred to as skimming, in order to reduce the file size, making the data analysis easier to handle. First, events with only one positively-charged particle, and no neutrals, were kept. Including events with one neutral and one positive track, while increasing our statistics, increased the background significantly. Since the neutron detection efficiency of the CLAS is relatively small, compared to that of the proton ($\sim 30\%$) [107], it was decided to exclude those events from the analysis.

An initial cut on the difference $\Delta\beta$ was applied as part of an initial particle identification,

$$\Delta\beta = \beta_{meas} - \beta_{calc} = \beta_{meas} - \sqrt{\frac{p^2 c^2}{m_p^2 c^4 + p^2 c^2}}, \quad (5.6)$$

where m_p corresponds to the nominal mass of the proton. The width of this cut was chosen to be $\sim 7\sigma$ from the mean value. This corresponds to events with $-0.1 < \Delta\beta < 0.1$, which selected the events that were kept for further analysis. Figure 5.3 shows the $\Delta\beta$ distribution of events with only one positively-charged particle detected.

The next cut applied was to select events where the square of the missing mass fell within $\sim 5\sigma$ from the square of the nominal value of the neutron mass. The 5σ width of the cut, which corresponds to about $0.6 \text{ (GeV}/c^2)^2$, was sufficiently broad to reduce the file size significantly (by a factor of ~ 1000) while retaining enough data from the background events. The latter was important for the background subtraction studies (see Sec. 5.10). The square of the missing mass of the reaction,

$$\gamma + d \rightarrow p + X, \quad (5.7)$$

was calculated using four-vector momentum conservation,

$$\begin{aligned} \mathcal{P}_\gamma + \mathcal{P}_d &= \mathcal{P}_p + \mathcal{P}_X \\ \mathcal{P}_X^2 = M_X^2 &= (\mathcal{P}_\gamma + \mathcal{P}_d - \mathcal{P}_p)^2, \end{aligned} \quad (5.8)$$

where $\mathcal{P}_i$ corresponds to the four-momentum of the particle in question. Figure 5.4 shows the missing-mass-squared distribution of the single-track events, with the red vertical line

indicating the nominal mass squared of the neutron, $0.88278 \text{ (GeV}/c^2)^2$. This shows that a number of the selected events correspond to deuteron photodisintegration events, with the neutron selected by the application of missing-mass-squared cuts.

Figure 5.3. $\Delta\beta$ distribution for events with one positively-charged particle.

Figure 5.4. The square of the missing mass of the reaction $\gamma + d \rightarrow p + X$ calculated using four momentum conservation. The red vertical line indicates the nominal mass of the neutron.

The missing-mass-squared cut reduced the size of the data even further without losing any signal events. Specifically, that application of the broad, particle ID and missing-mass-squared cut, the total resulting file size was reduced from ~ 140 Tb to ~ 10 Gb, allowing the data to be kept on a local machine for prompt recovery.

5.3.1 Event vertex cuts

Following the initial skimming, cuts on the event vertex were applied to ensure that the reconstructed event vertex location is consistent with particles being produced within the target. As mentioned in Sec. 3.2.5, the 40-cm-long liquid-deuterium target used for the g13 experiment had its center placed 20 cm upstream of the CLAS center. Figure 5.5 shows the z -component of the reconstructed event vertex. The cuts applied to the z -component of the vertex are indicated by the red vertical lines positioned at $z = -39.5$ cm and at $z = -1.3$ cm. The peak at $z = 0$ cm corresponds to the thin metal window of the target.

Figure 5.6 shows the x - (left) and the y - (right) components of the reconstructed event vertex. This indicates that the events are produced within the target with a very small

Figure 5.5. The z component of the reconstructed event vertex. The red lines correspond to the cuts applied to select particles that originated within the target.

Figure 5.6. The x - (left) and y -component (right) of the reconstructed event vertex.

contribution of events produced outside the target that can be removed with a selection cut.

5.3.2 Particle identification

Particles were identified by applying momentum-dependent $\Delta\beta$ cuts. The $\Delta\beta$ was calculated using Eq. (5.6), assuming that the particle was a proton with $m_p = 0.938 \text{ GeV}/c^2$. First, deuteron photodisintegration events were chosen by selecting events with $0.68 (\text{GeV}/c^2)^2 < m_X^2 < 1.1 (\text{GeV}/c^2)^2$ as indicated by the magenta lines in Fig. 5.7. Then, the 2-D distribution of $\Delta\beta$ as a function of the particle's momentum, p , was split into a series of 1-D projections along the momentum axis in bins of $50 \text{ MeV}/c$. For true protons, the distribution of $\Delta\beta$ is centered around zero.

Figure 5.7. $\Delta\beta$ as a function of M_X^2 . The magenta lines correspond to the cuts applied to select deuteron photodisintegration events.

A Gaussian was fitted to the $\Delta\beta$ distribution for each momentum bin extracting the mean, $\mu_{\Delta\beta}$, and standard deviation, $\sigma_{\Delta\beta}$. These parameters were used to apply a cut, for each momentum bin, discarding all events outside 3σ from the mean $\mu_{\Delta\beta}$. Specifically, for event-by-event data analysis, the momentum-dependent $\Delta\beta$ cuts were parametrized by fitting a fourth-order polynomial to the extracted positions $\mu_{\Delta\beta} \pm 3\sigma_{\Delta\beta}$ as a function of momentum.

Figure 5.8 shows the $\Delta\beta$ distribution for four momentum bins. The red line corre-

Figure 5.8. $\Delta\beta$ for proton candidates. The blue lines correspond to the $\pm 3\sigma$ cuts determined by fitting each distribution with a Gaussian. The plots correspond to momentum bins $p = 0.725 \text{ GeV}/c$, $p = 1.325 \text{ GeV}/c$, $p = 1.775 \text{ GeV}/c$, and $p = 1.975 \text{ GeV}/c$.

Figure 5.9. $\Delta\beta$ as a function of momentum. The red lines define the $\pm 3\sigma$ cuts and are determined by fitting each momentum bin with a Gaussian, as described in the text.

Figure 5.10. $\Delta\beta$ as a function of momentum for proton candidates. The magenta lines show the $\pm 3\sigma$ cuts determined by fitting each momentum bin with a Gaussian.

sponds to the Gaussian fit used to extract the $\mu_{\Delta\beta} \pm 3\sigma_{\Delta\beta}$ positions (shown by the blue lines). Figure 5.9 shows the fourth-order polynomial (red line) that was used to fit the $\mu_{\Delta\beta} + 3\sigma_{\Delta\beta}$ (left) and $\mu_{\Delta\beta} - 3\sigma_{\Delta\beta}$ (right) values as a function of momentum. This procedure yielded the momentum-dependent $\Delta\beta$ cuts. Figure 5.10 shows the $\Delta\beta$ distribution as a function of momentum with the missing-mass-squared cut applied (see Fig 5.7). The magenta lines correspond to the final momentum-dependent $\Delta\beta$ cuts used for particle identification. The extra bands shown in Fig. 5.10 are a result of assigning the particle the wrong mass (*i.e.* particles other than protons), accidental tracks, or selection of the incorrect electron-beam bunch for the event start-time calculation. Accidental tracks are particles that originated from different physics events than the event that caused the trigger. The $\Delta\beta$ cuts for $p < 0.7$ GeV/ c and for $p > 2.0$ GeV/ c were extrapolated due to poor statistics in those momentum bins. The change in the width of the $\Delta\beta$ distribution with momentum is primarily due to the uncertainty of β_{calc} . The uncertainty of the momentum determination, $\sigma_p \sim 0.005 \times p$, is propagated to the uncertainty of β_{calc} as shown in Eq. (5.9).

$$\sigma_{\beta_{calc}} = \frac{m_p^2 c^5}{(m_p^2 c^4 + p^2 c^2)^{3/2}} \sigma_p, \quad (5.9)$$

Figure 5.11 shows the calculated mass using Eq. (5.4), before (red) and after (yellow) the

Figure 5.11. Mass squared as measured for positive particles, before (red) and after (yellow) the particle identification cuts. The magenta line indicates the nominal mass of the proton. The small offset of the invariant proton mass peak from its true position is due to energy losses in the target and the CLAS.

particle identification cuts. The magenta line indicates the nominal mass of the proton.

5.4 Identifying the Incident Photon

The next step in the event reconstruction involved identifying the photon that initiated the reaction. Typically, an average of 14 “good” electron hits were recorded in the tagger for each trigger event (see Fig. 5.12). A “good” photon corresponds to a photon whose associated electron T-counter hit passed a set of consistency checks in the tagger (see Sec. 4.3). These consistency checks were designed to strongly reduce background and random electron hits not associated with a bremsstrahlung event. As mentioned before, the tagger reconstruction software assigned a photon for each electron hit. The time of the photon is based on information from the T-counters, and its energy is based on the electron E-counter hit position. The time assigned to the photon was the time of the photon at the center of the g13 target with respect to the global start time.

Figure 5.12. Number of electron hits in the tagger for each trigger event.

The photon identification was done by studying the time correlation between the calculated arrival time of the photon at the event vertex, t_γ (from tagger timing), and the event vertex time of the reaction, t_v , calculated using information from the CLAS. The event vertex time was calculated using the timing information from the start counter propagated back to the event vertex,

$$t_v = t_{ST} - \frac{d}{c\beta_{calc}}, \quad (5.10)$$

where t_{ST} is the time measured by the scintillation counters of the start counter with respect to the global start time, d is the length of the track from the particle vertex to the start counter paddle obtained from tracking, and β is the particle's velocity as a fraction of the speed of light calculated using the drift chamber momentum and the nominal mass of the proton. The time of the photon at the event vertex was calculated using the photon time from the tagger propagated to the event vertex,

$$t_\gamma = t_T + \frac{z + 20}{c}, \quad (5.11)$$

where t_T is the photon arrival time at the center of the target, and z is the distance of the event vertex along the beam axis measured relative to the CLAS center. The x and y

offsets of the event vertex, which are due to the finite beam spot size (~ 2 cm), are neglected as they are comparable to the vertex resolution.

The coincidence time, Δt , between the event vertex time calculated using information from the CLAS and the tagger information for all photons in an event,

$$\Delta t = t_v - t_\gamma, \quad (5.12)$$

is shown in Fig. 5.13. The peak centered at $\Delta t = 0$ ns corresponds to photons that are in coincidence with the reaction detected in the CLAS. The small neighboring peaks at 2 ns intervals correspond to the bunched nature of the electron beam. Specifically, the photons in these neighboring peaks originate from other beam bunches (not the one that initiated the reaction) being recorded in the 20 ns time window opened by the event's trigger.

An initial selection is done by choosing the photon that has the smallest coincidence time, $|\Delta t|$. Figure 5.14, shows the result of this pre-selection. Photons in the peak centered at zero (with coincidence time $|\Delta t| < 1.0$ ns) were considered as primary candidates for the photons that initiated the reaction. This value was chosen to reflect photons that come from the correct electron-beam bunch.

Figure 5.13. Coincidence time between the proton and all “good” photons, after particle identification.

Figure 5.14. Coincidence time between the proton and the “good” photon with the smallest time coincidence $|\Delta t|$, after particle identification.

Several events were identified in which two photons with a time coincidence $|\Delta t| < 1.0$ ns were detected (see Fig. 5.15). These photons occurred from background hits in the T-counters that happened in the same narrow time window defined by the central beam bucket. These events were discarded from further analysis since it was difficult to

unambiguously determine, using timing cuts, which photon caused the interaction. Selection of the wrong photon, using energy considerations, could artificially alter the yield distributions, which would then result to a false asymmetry. Events with zero photon with coincidence time $|\Delta t| < 1.0$ ns correspond to accidental events and they were also discarded from further analysis. Figure 5.16 shows the missing-mass-squared distributions of events

Figure 5.15. Number of photons in each event that have a coincidence time $|\Delta t| < 1.0$ ns. For approximately 78% of all events the photon was unambiguously determined, and only these events were selected for further analysis.

having only one photon with coincidence time $|\Delta t| < 1.0$ ns (red histogram), compared to the missing-mass-squared distribution of events having two photons with coincidence time $|\Delta t| < 1.0$ ns (yellow histogram). The left yellow histogram corresponds to the missing-mass-squared distribution calculated using the energy of the photon with the smallest coincidence time whereas the right yellow histogram corresponds to the missing-mass-squared distribution calculated using the energy of the photon with the second smallest coincidence time. From these plots one can see that events having two photons with coincidence times $|\Delta t| < 1.0$ ns are primarily background events. While the contribution at this stage is $\sim 16\%$, after background subtraction (see Sec. 5.10) the statistical contribution of events having two photons with coincidence times $|\Delta t| < 1.0$ ns reduces to $\sim 1\%$ of the total signal events.

Figure 5.16. Missing-mass-squared distribution of events having only one photon with coincidence time $|\Delta t| < 1.0$ ns (red histograms) compared to the missing-mass-squared distribution of events with two photons with time coincidence $|\Delta t| < 1.0$ ns (yellow histograms). The left yellow histogram shows the missing-mass-squared distribution calculated using the energy of the photon with the smallest coincidence time, whereas the right yellow histogram shows the missing-mass-squared distribution calculated using the energy of the photon with the second smallest coincidence time. One can see that events with two “good” photons are dominated by background. These events are excluded in further analysis.

5.5 Beam Trips

Another condition was applied to select events that were produced when the electron beam was stable. During the g13 experiment, the electron beam supplied by the CEBAF accelerator often tripped for a time ranging from few seconds to ~ 20 minutes, typically due to beam steering migration. The longer trips were often a result of hardware problems such as magnet trips and vacuum leaks. A proper photon flux determination over time intervals of beam trips is very hard. Even though this affects cross-section measurements more than asymmetry determinations, it was decided to exclude data collected during periods of time where the electron beam was unstable.

Events collected during a beam trip were flagged by using scaler counts from different subsystems in the CLAS. The trip status was obtained by studying the mean number of physics events per scaler interval (collected every 10 s). Intervals with the largest deviation from the average were rejected (assigned a non-zero status flag) if they contained 90% or less of the physics events of the mean. This process was iterated until scaler regions stopped being rejected. More details on the routines can be found in Refs. [108] and [109].

5.6 Fiducial Cuts

Fiducial cuts were applied to exclude events in which the detected particle fell in the region of the CLAS where the acceptance changes rapidly. This typically occurs near the edges of the CLAS drift chambers, where charged particles often escape detection or fail track reconstruction. Typically, particles hitting the support frames or the cryostats of the torus magnet fail track reconstruction. In addition, the magnetic field close to the torus magnet varies rapidly with position and is not mapped very accurately. Particle tracks reconstructed in these regions are characterized by large systematic uncertainties, which in turn, propagates to a large systematic uncertainty in their momentum determination.

The fiducial cuts were determined by studying the azimuthal and polar distributions of the reconstructed tracks. To account for the shape of the drift chambers, the azimuthal angle of the reconstructed track was studied as a function of the polar angle, as shown in Fig. 5.17. The six active sectors of the CLAS are clearly seen as regions with a

Figure 5.17. Proton azimuthal distribution as a function of the proton polar angle. The six independent tracking sectors of the CLAS are clearly evident. The low statistics areas correspond to the regions where the coils of the torus magnet are positioned.

high yield of particles. On the other hand, regions containing the torus magnet result in strips of low statistics. These are centered at $\phi = 30^\circ, 90^\circ, 150^\circ, 210^\circ, 270^\circ$, and 330° .

The functional form of the fiducial cuts used to remove regions of rapidly changing

acceptance in these distributions is,

$$\theta = -\frac{1}{c_0(|\phi - c_3| - c_1)} + c_2, \quad (5.13)$$

where the c 's are the function parameters. These parameters were determined manually so that the cuts removed almost all of the inefficient reconstruction regions at the edges of the drift chambers. The acceptance of the CLAS is momentum-dependent, since the radius of curvature of a charged-particle trajectory depends on the particle's momentum (the magnetic field produced by the torus magnet was set to bend positively-charged particles inward, towards the beam line). However, the difference in the CLAS acceptance between low-momentum (~ 0.6 GeV/ c) and high-momentum protons (~ 2.5 GeV/ c) in deuteron photodisintegration events was small (a shift of the order of 2° in the polar angle was observed between high- and low-momentum protons) and was not reflected in the fiducial cuts.

The CLAS acceptance also depends on the z -component of the event vertex, since particles with smaller polar angles could be detected only if they were produced at smaller z -vertex positions. The effect of the z -vertex position on the functional form of the fiducial cuts was also quite small and it too was ignored in the determination of the function pa-

Figure 5.18. Proton azimuthal distribution as a function of the proton polar angle after fiducial cuts. The left plot shows the angular distribution of protons in the lab frame, whereas the right plot shows the proton angular distribution in the center-of-mass frame.

rameters. In the functional form of the cut, Eq. (5.13), the six sectors of the drift chambers were assumed to be symmetric about their centers. Deviations from radial symmetry of different sectors of the drift chambers were very small and thus neglected in the application of these cuts. Figure 5.18 shows the proton polar angle as a function of the azimuthal angle after the application of the fiducial cuts. The left plot shows the distribution in the lab frame, whereas the right plot shows the distribution in the center-of-mass frame.

5.7 Energy-Loss Corrections

As charged particles pass through matter, they lose energy by ionizing or exciting the electrons in the matter. The mean rate of energy loss is described by the Bethe-Bloch formula [110],

$$-\frac{dE}{dx} = Kz^2 \frac{Z}{A} \frac{1}{\beta^2} \left[\frac{1}{1} \ln \frac{2m_e c^2 \beta^2 \gamma^2 T_{max}}{I^2} - \beta^2 - \frac{\delta(\beta\gamma)}{2} \right], \quad (5.14)$$

where A is the atomic mass of the absorber, Z is the atomic number of the absorber, I is the mean excitation energy, T_{max} is the maximum kinetic energy which can be imparted to a free electron in a single collision, and K is a constant that depends on the classical electron radius, Avogadro's number and the mass of the electron.

A software package, called *eloss*, was developed to correct for the energy lost due to the passage of particles through matter [111]. As mentioned before, the momentum of a charged particle in the CLAS is determined by tracking it in the magnetic field with the drift chambers. The drift-chamber reconstruction software returns a value for an “effective” momentum (without energy loss) between the Region-1 and the Region-3 drift chambers. The *eloss* software package determines the path length of the particles in each of the materials along the track, starting with the air gap between Region-1 drift chambers, start counter, scattering chamber, and ending with the target cell wall, and target material between the target wall and the event vertex. Then, it corrects for the energy lost through each material depending on the particle's measured momentum.

Figure 5.19 shows the energy loss of protons as a function of their measured momentum. Figure 5.20 shows the one-dimensional projection of the energy loss. On average, the protons lose about 4 MeV along their path towards the drift chambers. The peak

Figure 5.19. Energy loss as a function of measured momentum for protons. See text for explanation of the observed structure.

Figure 5.20. Proton energy loss through different materials before the particle enters the drift chambers. See text for explanation of the structure at 1.6 MeV.

observed near 1.6 MeV in Fig. 5.20 corresponds to particles produced outside the target. The v -shaped structure (gap) centered at $p_{meas} \sim 0.85$ GeV/ c in Fig. 5.19 is due to the CLAS acceptance. Specifically, this gap occurs due to the fact that particles produced at forward angles that also have low momenta are bent inwards, towards the beam line, and are not detected.

5.8 Photon Energy Corrections

A correction to the photon energy was applied due to a small gravitational sag in the tagger scintillator paddles [112]. This sag resulted in reconstructing electrons with systematically larger energies than their actual energy. This error propagates to the determination of the photon energy, and, specifically, results in tagging the photon with an energy smaller than its actual energy. Since the electron energy is determined by the hit position in the tagger E-channel, the applied correction was a tagger E-channel-dependent multiplicative constant. Specifically, the multiplicative constant was calculated (see Ref. [112]) to correct the photon energy with the values shown in Fig. 5.21. The maximum correction occurs at E-channel 200 and corresponds to a change of about 12 MeV for a 2.655-GeV electron beam.

5.9 Momentum Corrections

Additional corrections to particle momenta were applied after the *eloss* corrections. These corrections account cumulatively for energy loss in the drift chambers, drift chamber misalignments, and differences between the actual toroidal field and the field map used for track reconstruction. The momentum corrections were determined by performing kinematic fitting of the reaction $\gamma + d \rightarrow p + p + \pi^-$ [113]. The fit was constrained by requiring four-momentum conservation, with the magnitudes of the track momenta treated as unknowns to remove their bias from the kinematic fit.

The momentum corrections were determined by comparing the kinematically-corrected momentum with the *eloss*-corrected momentum. The comparison was done by taking the difference, Δp_i , between the kinematically- and *eloss*-corrected momenta. The average difference, Δp , between the fitted and *eloss*-corrected momentum was studied for

Figure 5.21. The photon-energy multiplier that corrects for the gravitational sag in the photon tagger as a function of the tagger E-channel.

each type of particle in each sector of the CLAS. In order to take into account misalignments of the drift chambers and magnetic-field map deviations, the difference Δp was studied as a function of θ and ϕ . The energy loss in the drift chambers was determined by studying the difference Δp as a function of *eloss*-corrected momentum. Predefined functions were used to determine the momentum corrections with respect to θ , ϕ , and p . These fits were iterated until the momentum corrections converged to a mean Δp of zero. Figure 5.22 shows the average difference between the fit and the detected momentum, Δp , as a function of (A) p , (B) ϕ , and (C) θ for sector 1 of the CLAS for protons and π^- . Figure 5.23 shows the same distributions after the momentum corrections. More information regarding the kinematic-fitting procedure can be found in Ref. [112].

5.10 Background Subtraction

From the missing-mass-squared distribution shown in Fig. 5.4 it is evident that there is a significant background that dilutes the deuteron photodisintegration events. Since this can result in false asymmetries, it was crucial to separate the signal events from the background events. To separate the signal from the background, the probabilistic event-

Figure 5.22. Average difference between kinematically- and *eloss*-corrected momentum, Δp , as a function of (A) p , (B) ϕ , and (C) θ for sector 1 of the CLAS for protons (left column) and π^- (right column), used to determine momentum corrections.

Figure 5.23. Average difference between kinematically- and $eloss$ -corrected momentum, Δp , as a function of (A) p , (B) ϕ , and (C) θ for sector 1 of the CLAS for protons (left column) and π^- (right column), after applying momentum corrections.

weighting method was used [114]. The aim of the procedure is to explicitly identify the chance that a given event, e_i , is a signal event, Q_i , or a background event, $(1 - Q_i)$. The first step in this procedure was done by defining a *distance measure* that was used to identify how close together two events were in non-reference coordinates. In the case of deuteron photodisintegration, the reference coordinate corresponds to the square of the missing mass, m_X^2 , since the functional form of the signal and the background as a function of the missing mass squared is known *a priori*. The non-reference coordinates can be any coordinates that are expected to affect the reference coordinate distribution, *i.e.* photon energy, angles, etc. Since the photon energy is binned in 200-MeV wide bins, the variation of the missing-mass-squared distribution across a photon-energy bin is quite small. For this reason, the non-reference coordinates used for the *distance measure* in this study were the proton polar and azimuthal angles, θ and ϕ .

For a given non-reference coordinate, ξ_k , the maximum difference between any pair of events i and j is r_k . Then the distance between these two events is,

$$d_{ij} = \sum_k \left[\frac{(\xi_k)_i - (\xi_k)_j}{r_k} \right]^2, \quad (5.15)$$

where the sum is over all the non-reference coordinates. This translates to,

$$d_{ij} = \left[\frac{\cos \theta_i - \cos \theta_j}{2} \right]^2 + \left[\frac{\phi_i - \phi_j}{2\pi} \right]^2. \quad (5.16)$$

While the definition of distance requires taking the square root of d_{ij} , doing so does not change anything in the procedure. On the other hand, working with the *distance measure*, as defined in Eq. (5.16), saves significant computational time that would otherwise be required for the computation of a large number of square roots. For each event in the data set, e_i , the distances to all other events in the data set, d_{ij} ($j = 1, 2 \dots n$), are computed. Then, a pre-defined number of events, N_d , is retained, which are closest to e_i . These are referred to as the closest neighbors to the event of interest e_i , and they define a dynamic bin width. For this analysis, 200 events were retained as closest neighbors. Using the N_d retained events, the distribution of their reference coordinate, ξ_r , was fitted to the known signal and background shapes. The signal was described by a Gaussian,

$$g(m_X^2, A, \mu, \sigma) = A e^{-\frac{1}{2} \left(\frac{m_X^2 - \mu}{\sigma} \right)^2}, \quad (5.17)$$

where A is the amplitude of the signal, and μ and σ are the mean and standard deviation, respectively. The background is due to several physics channels, such as $\gamma + d \rightarrow p + n + \pi^0$, $\gamma + d \rightarrow p + p + \pi^-$, and accidental events, making its exact shape hard to characterize. For this reason, the background shape was chosen as the shape that resulted in the best fit of the missing-mass-squared distribution, which in this case consisted of two exponentials,

$$b(m_X^2, A_1, A_2, B_1, B_2) = A_1 e^{A_2 m_X^2} + B_1 e^{B_2 m_X^2}. \quad (5.18)$$

Other background shapes have been studied (1^{st} through 4^{th} -order polynomials) without much success in adequately fitting the missing-mass-squared distributions. In several kinematic bins, the 3^{rd} -order polynomial resulted in good fits, however, the results for the background contribution were unphysical (giving negative background contributions under the nominal mass of the neutron).

Figure 5.24 shows a fit of the signal plus background distribution of the 200 closest neighbors for a specific event. The closest neighbors were chosen among the events that passed the selection criteria mentioned before (trip flags, particle ID, photon selection, fiducial cuts), in addition to a selection of the coherent photon peak (see Chap. 6).

Figure 5.24. Missing-mass-squared distribution for the 200 closest neighbors of an event. The red line shows the fit of the Gaussian (for the signal) and the two exponentials (for the background). The fit parameters were used to determine the probability, Q_i , that the specific event corresponds to a signal event.

The fit parameters, A , μ , σ , A_1 , A_2 , B_1 , and B_2 , determined from the fit, were

used to identify the signal and background functions, $G(m_X^2)$, and $B(m_X^2)$, respectively. The missing-mass-squared value of event e_i , $m_{X_i}^2$, was used to calculate the values of the signal and background of the event, $G(m_{X_i}^2)$, and $B(m_{X_i}^2)$, respectively. The Q -value of event e_i was then calculated,

$$Q_i = \frac{G(m_{X_i}^2)}{G(m_{X_i}^2) + B(m_{X_i}^2)}. \quad (5.19)$$

As mentioned before, the Q -value of the event represents the probability that the event of interest is a signal event.

Figure 5.25. Weighted missing-mass-squared distribution (yellow) for signal events (left) and background events (right), compared to the unweighted missing-mass-squared distribution (red). The distributions show events in the photon energy bin 1.7–1.9 GeV that pass all the selection criteria.

This procedure was repeated for all events that pass the selection criteria discussed above, separately for *Para* and *Perp* events¹. The Q -value was then used as a weighting function for any distribution (missing mass squared, momentum, polar and azimuthal angle, etc.), to separate the signal from the background. Figure 5.25 shows the missing-mass-squared distribution of the *Para* events in the photon energy bin 1.7 – 1.9 GeV that passed the selection criteria (red distribution). The yellow distributions show events weighted with the Q -value (left), and with the $(1 - Q)$ -value (right), which represent the signal and the background, respectively.

¹*Para/Perp* events correspond to events collected when the photon polarization vector was parallel/perpendicular to the Hall-B floor

Different dynamic bin widths were studied with consistent results. The dynamic bin of 200 closest neighbors was chosen since it corresponds to a bin width similar to that used for the extraction of the azimuthal asymmetry. Specifically, for events in the lowest photon energy bin ($E_\gamma = 1.1 - 1.3$ GeV), the dynamic bin size of 200 closest neighbors corresponds to about 5° in polar and azimuthal angles for events produced at forward angles, and to about 10° for events produced at backward angles. On the other hand, for events in the highest photon-energy bin ($E_\gamma = 2.1 - 2.3$ GeV), the dynamic bin width of 200 closest neighbors corresponds to about 10° in polar and azimuthal angles for events produced at forward angles, and to about 20° for events produced at backward angles.

Chapter 6

Analysis II: Photon Polarization

6.1 Overview

Linearly polarized photons in Hall B are produced using the coherent bremsstrahlung technique [91]. As discussed in Sec. 3.2.1, the electron beam scatters coherently from a crystal radiator, resulting in enhancement at specific energies over the $\sim 1/E_\gamma$ photon-energy spectrum observed with an amorphous radiator (carbon). A common practice for studying the photon-energy spectrum of coherent bremsstrahlung is to plot the *enhancement* of the spectrum, which is defined as the ratio of the spectrum obtained with the crystal radiator to that obtained with an amorphous radiator (see Fig. 6.1 for a typical enhancement plot). While this technique enhances the coherent bremsstrahlung spectrum over the incoherent spectrum, it also eliminates any channel-to-channel fluctuations caused by variations in the widths and efficiencies of the tagger spectrometer.

Figure 6.1. Enhancement plot calculated from the ratio of the spectrum obtained with the crystal radiator to that obtained with an amorphous radiator. The different peaks correspond to coherent contributions due to coherent bremsstrahlung off different reciprocal lattice vectors.

In g13b the photon polarization vector was rotated between two orthogonal ori-

entations: 1) parallel to the Hall-B floor, referred to as *Para*, and 2) perpendicular to the Hall-B floor, referred to as *Perp*. The orientation of the crystal lattice plane, which defined the photon polarization vector, was adjusted by setting the azimuthal angle of the crystal. The relative position of the main coherent peak on the photon-energy axis was set by adjusting the angle between the crystal lattice and the electron-beam direction. As mentioned in Sec. 3.2.1, the photon polarization depends on the coherent-edge position and on the electron-beam energy. Table 6.1 shows the different electron energies used for each coherent-edge position in g13b. Data with both *Para* and *Perp* orientations of the photon polarization were collected for each these settings. A description of the method used to

Table 6.1. Different Photon energy settings used during g13(b).

E_γ (GeV)	E_e (GeV)
1.1-1.3	3.302, 3.914, 4.192
1.3-1.5	4.065, 4.475
1.5-1.7	4.065, 4.748
1.7-1.9	5.057
1.9-2.1	5.057, 5.157
2.1-2.3	5.157

determine the linear photon polarization follows. This method was primarily developed by a collaborator [115]. The author has contributed significantly in studying and optimizing this method. The determination of the degree of photon polarization for the g13 experiment was also carried out by the author.

6.2 Analytic Bremsstrahlung Calculation

Coherent bremsstrahlung has been investigated theoretically in great detail since interference effects of high-energy bremsstrahlung produced in crystals were first suspected by Williams in 1934 [116]. Typical features, like the appearance of maxima and minima as a function of the crystal orientation and an intensity enhancement with increasing electron energy, were predicted by an intensity formula developed by Ferretti in 1950 [117]. Überall showed that the photon beam resulting from the bremsstrahlung process in a crystal has a strong linear polarization for certain orientations of the crystal lattice with respect to the

incident electron beam [118, 119]. This has been verified through numerous experiments and the theory was significantly improved to accommodate experimental features, like beam collimation, beam divergence, etc. Experimental evidence for the validity of coherent bremsstrahlung theory can be found in Refs. [120, 121, 122]. Reference [123] covers in detail the arguments and calculations for the interference effect between the electrons and the lattice structure of the crystal.

Appendix A describes in detail the analytic bremsstrahlung calculation and the features added to describe beam and collimation characteristics. Specifically, Appendix A shows the function that characterizes the enhancement distributions and determines, in this way, the degree of photon polarization. In general, there are five parameters that determine the form of the enhancement and, thus, the polarization. These are:

- θ – Angle between the beam and the crystal planes defined by the 022 direction
- σ – Gaussian smearing of θ to account for beam divergence, multiple scattering, beam-spot size etc.
- θ_r – Relative angle of collimation in units of characteristic angle
- σ_r – Smearing factor for collimation around θ_r via cumulative distribution function C_g
- $I_2^0, I_4^0, \dots, I_G^0$ – Amplitudes of the discrete coherent peaks

Details and definitions of these parameters are shown in Appendix A. These parameters that are used to determine the degree of polarization are obtained by fitting the enhancement distribution, as discussed in detail in Sec. 6.4.

6.3 Enhancements

As mentioned before, the enhancement is calculated by dividing the photon-energy spectrum obtained using the crystal radiator by that obtained from an amorphous radiator. This is done to primarily remove any channel-to-channel variations in the counter efficiencies. Since counter efficiencies can be time-dependent, the amorphous data run used

to divide each coherent data run was chosen such that the time difference between the collection time of the two runs was the smallest.

To account for the effect of collimation on the photon polarization, event-by-event data were used. These were taken at a rate of ~ 4 kHz; therefore, it takes a long time to build up an enhancement spectrum with reasonable statistics (eg. > 5 min). Figure 6.2 shows the photon-energy spectrum obtained from the coherent-edge setting $E_\gamma = 1.5$ GeV (lower panel), and the photon-energy spectrum obtained from an amorphous radiator (upper panel).

The coherent-edge position was set by adjusting the angle between the crystal radiator and the beam line using the goniometer. However, due to small oscillations of the crystal radiator, it was necessary to determine the coherent-edge position from the data. This was done by fitting the edge position with a 3^{rd} -order polynomial. The position with the biggest negative slope was chosen as the coherent-edge position (see Fig. 6.3). In several cases, the difference between the set value and the measured value of the coherent-edge position was as large as 100 MeV. Figure 6.4 shows the typical drift in coherent peak position for the settings where the coherent edge was nominally set at 1.3, 1.5, 1.7, and 2.1 GeV. These fluctuations significantly affect the photon polarization and needed to be taken into account. However, the time needed to build up an enhancement spectrum with reasonable statistics using the event-by-event data was much bigger compared to the time scale where the coherent peak was drifting in position. For this reason, the coherent-edge position was determined from E-counter scalers, which on one hand, don't reflect the effect of collimation, but on the other hand, provide an order-of-magnitude better statistics and can be used to determine the coherent-edge position every 2 seconds. The polarization was then determined using the event-by-event data for the different coherent-edge positions. Figure 6.5 shows how the coherent-edge position drifted around the nominal set value as a function of event number for a *Para* run (left) and a *Perp* run (right). Figure 6.6 shows the 2-MeV-wide binning used for each coherent-edge position to determine the photon polarization.

Details on the fitting procedure for determining the photon polarization are presented in the section below [115].

Figure 6.2. Photon-energy spectrum obtained from an amorphous radiator (upper panel) and a diamond radiator (lower panel). The peaks from the coherent spectrum that correspond to highly polarized photons are clearly seen.

Figure 6.3. Determination of the coherent-edge position using a 3^{rd} -order polynomial.

Figure 6.4. Drift in the coherent peak position for data taken with the diamond radiator with the coherent edge nominally set at 1.3, 1.5, 1.7, and 2.1 GeV. The different lines correspond to the enhancement plots produced every 2 s.

Figure 6.5. Drift in the coherent peak position as a function of event number for a *Para* run with the coherent edge nominally set at 2.3 GeV (left) and for a *Perp* run with the coherent edge nominally set at 1.9 GeV (right).

Figure 6.6. Coherent-edge bins for *Para* event-by-event data collected with the coherent edge nominally set at 1.5 GeV. For each 2-MeV-wide bin in coherent-edge position the photon polarization was determined.

6.4 Photon Polarization Determination

The photon polarization is determined by fitting the enhancement function shown in Eq. (A.17). The fit parameters are then used to calculate the photon polarization distribution using Eq. (A.19). These functions take into account several beam and detector characteristics, like beam collimation, beam-spot size and divergence, fluctuations in the crystal position, etc., as discussed in Sec. 6.2. At this point, the procedure followed for determining the photon polarization had to address additional characteristics of the g13 data.

One of the problems was that the enhancement from the event-by-event data was far from a smooth distribution. This was primarily due to the low statistics in the amorphous data used for the calculation of the enhancement. On several occasions, drifts in the efficiencies of the tagger counters over time resulted in a non-smooth enhancement distribution. These problematic enhancements were more common if there was a long time gap between the data-taking period of the polarized and amorphous data.

Another problem with the enhancement calculation was that the scaling factor was quite difficult to determine. A point or a region with no enhancement, and thus no coherent contribution, had to be chosen to determine the scale factor for setting the baseline to 1. The chosen point or region is sensitive to statistics, and for this reason, the scaling factors were poorly determined for enhancements resulting from low-statistics coherent-edge bins. In addition to this, leakage of the secondary coherent peaks into the primary coherent peak needed to be taken into account in the scaling factor determination. Fortunately, in the g13 experiment, a very tight collimation of the beam was chosen to increase the photon polarization, which resulted in a relatively good separation of the primary and secondary coherent peaks.

Furthermore, the enhancement fit functions depend on at least 5 parameters (in the case where only the main coherent peak is fitted). The tight beam collimation used for the g13 experiment resulted in a relatively good separation between the different coherent peaks; therefore, fitting only the coherent peak produced by the 022 lattice vector was justified. Despite this, the fitting routine that attempted to minimize the 5 free parameters often failed to produce a good fit, and in several cases, the values of the extracted fit

parameters were unphysical. This was primarily observed in fitting bins with low statistics.

A procedure was developed in an attempt to address all the above problems. In this procedure, the enhancements binned in 2-MeV-wide coherent-edge position bins (see Fig. 6.6) were used, as described below.

6.4.1 Initial fitting

As mentioned before, the coherent-edge position is determined using the E-counter scaler information and passed to the event-by-event data. Initially, a coherent-edge position bin with high statistics is chosen to fit the enhancement distribution with the fit function (see Eq. (A.17)), treating all fit parameters as unknown. Figure 6.7 shows the enhancement distribution (upper panel) of *Para* data with coherent edge 1.501 GeV fitted with the enhancement function of Eq. (A.17), and the degree of photon polarization (lower panel) as calculated from Eq. (A.19) using the parameters determined from the minimizing routine of the fit.

It is worth noting here that the photon polarization is quite large at points where the enhancement is relatively small. It will become evident later that the photon polarization is more sensitive to the points where the enhancement is relatively small. This can also be observed from studying Fig. 6.7. From these, the small difference in the enhancement between two adjacent E-counters closer to the left magenta line corresponds to a large difference in the photon polarization, whereas a larger difference in the enhancement between two adjacent E-counters closer to the peak corresponds to a smaller difference in photon polarization. The highest photon polarization occurs at points where the enhancement is maximum.

The parameters that describe the beam collimation are then fixed to the values extracted from fitting the high-statistics coherent-edge bin. Specifically, the values of θ_r and σ_r are used for fitting the rest coherent-edge bins. This addressed, to some degree, the problem that occurred from fitting the enhancement function on low-statistics enhancement distributions. In general, the fitting routine produces meaningful results after fixing two of the five free parameters, becoming much more reliable. Figure 6.8 shows fits on enhancement distributions (upper two plots) for *Para* event-by-event data after fixing θ_r

Figure 6.7. Enhancement (upper panel) for *Para* event-by-event data with coherent-edge position 1.500–1.502 GeV (red points) and the degree of photon polarization (lower panel) determined from the enhancement fit. The blue line in the upper panel corresponds to the fitted enhancement function described by Eq. (A.17). The magenta lines can be used to guide the eye in comparing the enhancement and photon polarization plots.

and σ_r , along with the degree of polarization (lower two plots) determined from fitting these enhancements.

Figure 6.8. Enhancement and determined polarization for two different coherent-edge bin, $E_\gamma^{\text{edge}} = 1.521$ GeV (left plot) and $E_\gamma^{\text{edge}} = 1.483$ GeV (right plot). The enhancement function shown here has the parameters θ_r and σ_r fixed to the values extracted from the high-statistics coherent-edge bin.

6.4.2 Rescaling

The initial scaling of the event-by-event enhancement distribution was done by looking at photon-energy bins located above the coherent-edge position. Specifically, the integral of 11 consecutive bins between the primary coherent edge and secondary coherent peak is calculated. The integral calculation is repeated for the next 11 bins shifted by 1 bin

with respect to the previous integral. The smallest bin integral is then used to calculate the scaling factor for the initial scaling. In general, the region between the main coherent edge and the secondary coherent peak is where the statistics are poorest and where any effect of tagger channels having drifted between the coherent and amorphous data are strongest. Thus, rescaling becomes very important in the photon polarization determination.

The results from the initial enhancement fitting were used to address the issue of setting the baseline to 1. Specifically, the enhancements of the event-by-event data of each coherent-edge bin were rescaled by the ratio of integrals of the fitted enhancement, ξ_{data} , and the event-by-event enhancement, ξ_{data} ,

$$k_{rescale} = \frac{\int_{low}^{high} \xi_{fit}}{\int_{low}^{high} \xi_{data}}. \quad (6.1)$$

The limits of integration were predefined to cover most of the coherent peak. These are shown by the magenta lines in the figures above. The statistical fluctuations of the individual points of the enhancement distribution are expected to average out in the calculation of the integral. Another problem arises, however, in using this method to rescale the enhancement distributions. In several cases, the enhancement fit was not satisfactory, which resulted in the calculation of a bad rescaling constant. This occurred more often in low-statistics coherent-edge bins. To prevent a bad rescaling, fit parameters extracted from good fits (*i.e.* fits with small χ^2) were used for rescaling all enhancement distributions. Specifically, these parameters were plotted as a function of coherent-edge position and fitted with a second-degree polynomial, as shown in Fig. 6.9. This was done in order to study their dependence on coherent-edge position and rescale the enhancement distributions that had inadequate enhancement fits. Figure 6.10 shows the enhancement distribution before (left) and after (right) rescaling as compared to the fitted enhancement. Specifically, the enhancement fit shown in the left plot (initial enhancement distribution) of Fig. 6.10 corresponds to the function of Eq. (A.17), with its parameters determined from fitting the event-by-event enhancement data. On the other hand, the enhancement fit shown in the right plot (rescaled enhancement distribution) corresponds to the enhancement whose parameters are determined from fitting the extracted parameters, as shown in Fig. 6.9.

Figure 6.9. Parameters θ and I_2 extracted from fitting the enhancement distributions of different coherent-edge bins. The red line corresponds to a 2^{nd} -degree polynomial fitted to these values which is then used for rescaling the enhancement distributions.

Figure 6.10. Enhancement distribution before (left) and after (right) rescaling. The blue line on the left plot corresponds to the fitted enhancement on the event-by-event data. The green line on the right plot corresponds to the enhancement determined using parameters determined from fitting the parameters of the different coherent-edge enhancement fits. The event-by-event enhancement data on the right plot are rescaled according to Eq. (6.1)

6.4.3 Corrections on the photon polarization

After the rescaling of the event-by-event data is performed, differences between the enhancement fit and the event-by-event enhancement have to be accounted for. The photon polarization determination is more sensitive above the coherent-edge position and well below it, in regions where the enhancement is small relative to its maximum. Clearly, if the enhancement is greater than the fitted function, the polarization should be increased on the basis of that difference.

To determine the correction to the polarization based on the difference between

the measured and calculated enhancement, Eqs. (A.1) and Eq. (A.14) are used. Under the assumption that the incoherent contribution, $I^i(x)$, is reasonably described by the distribution of the data collected using an amorphous radiator (*i.e.* $I^i(x) \simeq kI^{amo}(x)$), Eq. (A.14) simplifies to,

$$\xi = 1 + \frac{I^c}{I^i}. \quad (6.2)$$

Substituting for I^i/I^c in terms of ξ in Eq. (A.1) gives

$$P_g = -\Phi(x) \left(1 - \frac{1}{\xi}\right). \quad (6.3)$$

Differentiating Eq. (6.3) with respect to ξ gives

$$\frac{\partial P_g}{\partial \xi} = -\Phi(x) \frac{1}{\xi^2}. \quad (6.4)$$

Therefore, a small change, $\Delta\xi$, in the enhancement, ξ , will result in a polarization correction

$$\Delta P_g = -\frac{\Phi(x)}{\xi^2} \Delta\xi. \quad (6.5)$$

The above can also be written in terms of a fractional change, using Eq. (6.3),

$$\frac{\Delta P_g}{P_g} = \frac{1}{\xi - 1} \frac{\Delta\xi}{\xi}. \quad (6.6)$$

Figure 6.11 shows the enhancement fit compared with rescaled enhancement of the event-by-event data (left plot) and the polarization determined from the fit (green line on right plot). The blue points on the right plot correspond to the corrected polarization using Eq. (6.6). It is evident from Eq. (6.6) that corrections are generally bigger in bins where the enhancement is small. The corrections in these bins usually fluctuate, reflecting the statistical fluctuations in the data and the variations in the tagger efficiencies of the tagger channels between the accumulation of the coherent and incoherent data sets used to make the enhancement.

6.4.4 Accounting for statistical fluctuations

In order to remove the statistical and tagger-efficiency fluctuations, the corrected photon polarization for each E-counter channel is plotted as a function of coherent-edge position. Then the corrected polarization is fitted with a 3rd-degree polynomial in an

Figure 6.11. The left plot corresponds to the rescaled enhancement of event-by-event data compared to the fitted enhancement function. The right plot shows the polarization determined using the fit parameters of the enhancement function of the left plot. The blue points correspond to the corrected polarization using Eq. (6.6).

attempt to remove the statistical fluctuations and variations in the tagger efficiencies, as shown in Fig. 6.12.

Figure 6.12. Corrected photon polarization as a function of coherent-edge position for E-counters 339, 326, 319, and 312 respectively. The red line corresponds to a 3rd-degree polynomial, fitted on the corrected polarization, used to remove the corrected-polarization dependence on the statistical fluctuations and variations in the tagger efficiencies.

Figure 6.13. Polarization of a low-statistics coherent-edge bin as a function of photon energy. The green line corresponds to the polarization determined from fitting the event-by-event enhancement distribution, the blue points correspond to the corrected polarization which takes into account the differences between the data and the fit, and the magenta line corresponds to the polarization determined for each individual E-counter to account for any statistical fluctuation or variations in the tagger efficiencies (see Fig. 6.12).

The polarization is then determined for each E-counter using the polynomial fit. Figure 6.13 shows the determined polarization for a specific coherent-edge bin. The green line corresponds to the polarization determined using the fit parameters of the enhancement fit function, the blue points correspond to the corrected polarization taking into account the differences between the data and the fit, and finally, the magenta line corresponds to the polarization determined from fitting each E-counter with a polynomial to smooth out any statistical fluctuations or variations in the counter efficiencies in the data. This method of determining the polarization does not work very well for E-counters located far from the nominal coherent edge. This is because these channels are usually associated with small statistics, resulting in a large photon polarization correction, and in addition, lacking sufficient information to study the E-counter dependence on the coherent-edge position.

6.4.5 Photon polarization tables and usage

Using the information provided by the individual E-counter polynomial fits, polarization tables are produced for each coherent-edge bin in steps of 1 MeV. The method described above was used to determine the photon polarization for each E-counter bin (about 9 MeV wide) and for each coherent-edge position (in steps of 1 MeV). On average, 100 tables are produced for each photon energy containing the information needed to calculate the average photon polarization. Table 6.2 shows a portion of the table produced for a specific photon-energy setting (*Para* events, coherent-edge nominal position: 1.5 GeV, coherent-edge bin: 1.493 GeV), with the polarization of each E-counter shown in the last column.

Table 6.2. Part of a photon polarization table showing the polarization for each E-counter for the coherent-edge bin 1493 GeV.

$\mathbf{E}_{id}$	$\mathbf{E}_{\gamma}$	ξ	σ_{ξ}	ξ_{fit}	$\mathbf{P}_{fit}$	$\mathbf{P}_{corr}$	$\mathbf{P}_{smooth}$
384	993.1	0.852	0.004	1.003	0.002	0	0
...	...	...	...	...	...	...	...
332	1393.2	11.359	0.033	10.869	0.832	0.833	0.831
331	1401.9	11.851	0.034	11.508	0.838	0.838	0.836
330	1410.7	12.699	0.036	12.128	0.843	0.844	0.844
329	1419.5	13.333	0.039	12.706	0.847	0.848	0.847
...	...	...	...	...	...	...	...
321	1489.8	8.666	0.028	9.344	0.820	0.812	0.803
320	1498.6	7.098	0.036	7.583	0.802	0.787	0.784
319	1507.4	5.581	0.025	6.059	0.772	0.750	0.753
...	...	...	...	...	...	...	...

The goal is to be able to calculate the average value of the degree of polarization for each photon-energy setting. The average degree of photon polarization can be usually calculated in two different ways. In the first method, the different coherent-edge tables are combined by calculating the average photon polarization per tagger E-counter. This is done by weighting the polarization in each E-counter by the number of events in the

corresponding coherent-edge bins (see Fig. 6.6),

$$\bar{P}_{E_{id}} = \frac{\sum_j^K N_{events}^j P_{E_{id}}^j}{\sum_j^K N_{events}^j}, \quad (6.7)$$

where j runs over the K 1-MeV-wide coherent-edge polarization tables, N_{events}^j is the number of events in the specific coherent-edge bin, and $P_{E_{id}}^j$ is the photon polarization of the specific E-counter for the coherent-edge bin j . This will result in one photon-polarization table for each photon setting, independent of the coherent-edge position.

The second method involves calculating the average polarization using information from the reaction of interest (in this case, $\gamma + d \rightarrow p + n$). For each event in the reaction of interest (*i.e.* events that pass all the selection criteria discussed in Chap. 5, such as PID, photon selection, fiducial cuts etc.), the polarization of the selected photon is retrieved from the lookup tables depending on the event's coherent-edge position and photon energy. This way, one can calculate the average polarization for a specific reaction by taking into account only events of interest.

As mentioned before, the tables produced from this method are more reliable for bins that are closer to the nominal coherent-edge position. The rescaling of the low-statistics bins with coherent-edge positions far from the nominal position was determined by extrapolating the fit parameters to the lower (or higher) coherent-edge positions, due to inadequate fits. Because of this, only cuts on events with coherent-edge position less than $|\pm 3\sigma|$ of the nominal position were used for further analysis (see Fig. 6.14).

In addition to the above cuts, events with photon polarization larger than 0.50 were kept for further analysis. This is because photon polarization less than 0.50 corresponds to small enhancements that become more sensitive to statistical fluctuations and variations in the tagger-counter efficiencies. As mentioned before, the photon polarization is more stable close to the coherent peak, where the variations of photon polarization between adjacent E-counters are small. By selecting events that fall close to the coherent peak, one reduces the systematic uncertainty of the determined polarization, since the photon polarization of events with photon energies far from the coherent peak is more sensitive to the predicted beam collimation and beam divergence. Therefore, for the purposes of this analysis, events with photon energies that fall between their corresponding coherent-

Figure 6.14. Coherent-edge position fitted with a Gaussian to determine the 3σ cuts for selecting events with reliable photon determination.

edge position and 200 MeV below their corresponding coherent-edge position were kept for further analysis. Table 6.3 shows the average photon polarization determined for events that pass the analysis cuts mentioned in Chap. 5 and the cuts mentioned above. The uncertainty in the degree of photon polarization is determined from an independent study of systematics (see Chap. 8) and is of the order of 7%.

Table 6.3. Average photon polarization determined using events from deuteron photodisintegration as described in the text. The number of events corresponds to the number of events that pass all analysis cuts. These events were used for the calculation of the average polarization.

E_{γ}^{edge} (GeV)	E_e (GeV)	P_{γ} (<i>Para</i>)	Number of events (<i>Para</i>)	P_{γ} (<i>Perp</i>)	Number of events (<i>Perp</i>)
1.3	3.302	0.73	16928	0.73	19627
1.3	4.192	0.80	27034	0.77	13887
1.5	4.065	0.81	13529	0.80	14392
1.5	4.475	0.78	14238	0.77	12505
1.7	4.065	0.75	11260	0.75	6095
1.7	4.748	0.78	6402	0.80	9658
1.9	5.057	0.68	8643	0.71	9144
2.1	5.057	0.76	1894	0.76	1663
2.1	5.157	0.75	950	0.75	888
2.3	5.157	0.71	1523	0.71	1535

Chapter 7

Determination of the Azimuthal Asymmetry

7.1 Overview

The beam asymmetry, Σ , was extracted from the differential cross section of deuteron photodisintegration. The differential cross section of Eq. (2.14) reduces to

$$\frac{d\sigma}{d\Omega} = \left(\frac{d\sigma}{d\Omega} \right)_0 (1 + P_\gamma \Sigma \cos[2\phi]) \quad (7.1)$$

in the case of a linearly polarized photon beam. In Eq. (7.1) P_γ corresponds to the degree of linear polarization (see Chap. 6), whereas ϕ corresponds to the azimuthal angle between the photon polarization vector and the reaction plane (see Fig. 2.1). This chapter presents the method used by the author to extract Σ from the data.

7.2 Choice of Bin Size

The beam asymmetry is a function of both the incident photon energy, E_γ , and the polar angle, θ . It is common to describe the observables of a two-body reaction in the center-of-mass frame since the $\theta_{c.m.}$ dependence of the observables is symmetric between the two final-state particles. The boost velocity to the center-of-mass frame, β , is calculated to be

$$\vec{\beta} = \frac{\vec{p}_\gamma}{|\vec{p}_\gamma| + m_d}, \quad (7.2)$$

using the fact that in the center-of-mass frame the momentum of the deuteron is equal and opposite to that of the incident photon. The choice of photon-energy bin width was determined to be 200 MeV at the initial stages of the experiment. Specifically, data were collected for six different positions of the coherent peak, between 1.1 and 2.3 GeV (see Chap. 3). The same bins were used for the determination of Σ . In principle, smaller bins could have been used but that would increase the statistical uncertainties in Σ , especially in the higher photon-energy bins where statistics are quite low. In addition, a large photon-

energy bin allowed us to use smaller bin widths in $\theta_{c.m.}$. Figures 2.3 and 2.5 show that the available theoretical predictions for the azimuthal asymmetry are much more sensitive to the dependence of Σ on $\theta_{c.m.}$ than on the photon energy (see Fig. 2.13).

On the other hand, the bin width of the proton angle, $\theta_{c.m.}$, was chosen such that the statistical uncertainties in Σ were comparable between different bins. This resulted in binning $\theta_{c.m.}$ in relatively small bins ($\sim 5^\circ$) for data collected at smaller incident photon energies and forward angles, and much bigger bins (up to $\sim 40^\circ$) for data at backward angles and large incident photon energies.

7.3 Σ Determination Method

The azimuthal asymmetry can be determined in various ways. For example, one can determine the asymmetry from fitting the differential cross section with a $\cos 2\phi$ function. To do so, a very careful study is needed that will account for the detector acceptance and efficiencies. This also involves knowing *a priori* the exact value of the unpolarized differential cross section. Figure 7.1 shows the ϕ distribution of the yield for a single photon-energy bin. From this, it is evident that the distribution is strongly dominated by the acceptance of the CLAS, although small differences can be observed between the *Para* and *Perp* data. The six fiducial regions of CLAS are also clearly visible in the data. A careful study of the CLAS acceptance is not easy and results in additional systematic uncertainties in the determined asymmetry.

Since acceptance effects are not polarization-dependent, dividing the polarized *Para* and *Perp* distributions by the distributions collected with an amorphous radiator (*Amo*) would remove most of the acceptance effects. From the cross section (see Eq. (7.1)), the distribution of deuteron photodisintegration events as a function of ϕ is uniform for *Amo* data. On the other hand, the distribution for *Para* and *Perp* events has the form

$$N(\phi)^{\parallel,\perp} \sim F^{\parallel,\perp} (1 \pm P_{\gamma}^{\parallel,\perp} \Sigma \cos[2\phi]) A(\phi), \quad (7.3)$$

where $F^{\parallel,\perp}$ corresponds to the incident photon flux and $A(\phi)$ to the CLAS acceptance. By taking the ratio between polarized and amorphous data, the acceptance $A(\phi)$ cancels out, thereby eliminating any associated systematic errors. The resulting distribution has the

Figure 7.1. ϕ distribution of protons in events collected with a parallel (left) and perpendicular (right) to the Hall B floor beam polarization.

form

$$\frac{N(\phi)^{\parallel,\perp}}{N(\phi)^a} = \frac{F^{\parallel,\perp}}{F^a} (1 \pm P_\gamma^{\parallel,\perp} \Sigma \cos[2\phi]). \quad (7.4)$$

By fitting the ratio of ϕ distributions between the polarized and the amorphous data with a function of the form $A(1 + B \cos[2\phi])$, one can determine Σ from the fit parameter B using the determined photon polarization P_γ .

Another way of determining Σ from the data is by taking the ratio between a linear combination of the polarized data. Specifically, by calculating the quantity $\frac{N(\phi)^{\parallel} - N(\phi)^{\perp}}{N(\phi)^{\parallel} + N(\phi)^{\perp}}$, the acceptance effects cancel out and the asymmetry can be easily extracted from fitting. Using Eq. (7.3), this ratio results in

$$\frac{N(\phi)^{\parallel} - N(\phi)^{\perp}}{N(\phi)^{\parallel} + N(\phi)^{\perp}} = P \Sigma \cos[2\phi] \quad (7.5)$$

in the case where the incident photon flux is well known and the data can be scaled reliably to make $F^{\parallel} = F^{\perp}$, and under the general case of having $P_\gamma^{\parallel} = P_\gamma^{\perp} = P$. The asymmetry, Σ , can then be determined by fitting this distribution with the function $B \cos[2\phi]$.

A complication arises, however, due to the accuracy with which the orientation of the photon polarization plane is known. Ideally this is known accurately enough (within less than a degree). However, there is likely to be some systematic offset ϕ_0 , which is required to make the $\cos[2\phi]$ distribution consistent with the data. This systematic offset is expected to be the same between the *Para* and *Perp* data, since the photon polarization

orientation was obtained from accurate changes in the orientation of the crystal radiator.

In the more general case, the average photon polarization of *Para* data is not equal to that of *Perp* data. Furthermore, the incident photon fluxes for *Para* and *Perp* are not known accurately enough to scale the data, and thus $F^{\parallel} \neq F^{\perp}$. Taking this into account, the ratio between the polarized data becomes

$$\frac{N(\phi)^{\parallel} - N(\phi)^{\perp}}{N(\phi)^{\parallel} + N(\phi)^{\perp}} = \frac{F_R - 1 + \frac{F_R P_R + 1}{P_R + 1} 2\bar{P}\Sigma \cos[2(\phi - \phi_0)]}{F_R + 1 + \frac{F_R P_R - 1}{P_R + 1} 2\bar{P}\Sigma \cos[2(\phi - \phi_0)]}, \quad (7.6)$$

where $F_R = F^{\parallel}/F^{\perp}$, $P_R = P_{\gamma}^{\parallel}/P_{\gamma}^{\perp}$, and $\bar{P} = (P_{\gamma}^{\parallel} + P_{\gamma}^{\perp})/2$. This can be fitted with a function with four free parameters, A , B , C , and D ,

$$Y(\phi) = \frac{A - 1 + \frac{AB+1}{B+1} 2C \cos[2(\phi - D)]}{A + 1 + \frac{AB-1}{B+1} 2C \cos[2(\phi - D)]}. \quad (7.7)$$

The statistical uncertainty of the quantity $Y(\phi)$ is calculated by propagating the uncertainties of $N(\phi)^{\parallel}$ and $N(\phi)^{\perp}$, which follow a Poisson distribution (*i.e.* $\Delta N(\phi) = \sqrt{N(\phi)}$). It can be easily shown that the statistical uncertainty of $Y(\phi)$ is given by

$$\Delta Y = \frac{2}{(N^{\parallel} + N^{\perp})^2} \sqrt{N^{\parallel} N^{\perp} (N^{\parallel} + N^{\perp})}. \quad (7.8)$$

From fitting the function shown in Eq. (7.7) to the data the following quantities could be extracted:

- A : Ratio of *Para* and *Perp* fluxes, F_R
- B : Ratio of *Para* and *Perp* polarizations, P_R
- C : Product of average polarization and asymmetry, $\bar{P}\Sigma$
- D : Offset of the photon polarization vector, ϕ_0

For the purposes of this analysis, the azimuthal asymmetry was determined using the ratio of the difference between *Para* and *Perp* events to their sum (as indicated by left hand side of Eq. (7.6)). This choice was made since the amorphous data collected in g13 was about an order of magnitude less than the polarized data. The asymmetry determined using the

amorphous data would thus result in large statistical uncertainties, as well as systematic uncertainties from unreliable fitting, and therefore that method was abandoned.

We can use Eq. (7.8) to obtain the statistical uncertainty on the azimuthal asymmetry. An estimate on this uncertainty can be obtained by assuming that $F_R = 1$, and $P_R = 1$ (i.e. $P^{\parallel} = P^{\perp} = P$). With these assumptions in place, the statistical uncertainty of the azimuthal asymmetry can be given by (if uncertainties in the degree of polarization are not taken into account)

$$\Delta\Sigma = \frac{\Delta Y}{P} = \frac{2\sqrt{N^{\parallel}N^{\perp}(N^{\parallel} + N^{\perp})}}{P(N^{\parallel} + N^{\perp})^2}. \quad (7.9)$$

We can simplify the above further if the experiment is designed to collect approximately the same amount of *para* and *perp* data (i.e. $N^{\parallel} \approx N^{\perp} \approx N_{tot}/2$). Doing so, the statistical uncertainty of the azimuthal asymmetry can be estimated by

$$\Delta\Sigma = \frac{1}{P\sqrt{N_{tot}}} . \quad (7.10)$$

Equation (7.10) shows that the statistical uncertainty of the azimuthal asymmetry can be reduced by either collecting more data (increasing N_{tot}) or improving the experimental setup to obtain higher degree of photon polarization (increasing P). Specifically, for reducing the statistical uncertainty by a factor of two one can either increase the amount of data collected by a factor of four or double the degree of photon polarization.

A common quantity in polarization experiments is the *figure of merit* (*fom*), which is defined as

$$fom = P^2 L, \quad (7.11)$$

where L corresponds to the luminosity (proportional to the incident photon flux). This quantity is inversely proportional to the amount of time necessary to make a measurement of a desired precision and it is used to estimate the degree of photon polarization and luminosity needed. Because of the tight beam collimation needed in linearly-polarized photon-beam experiments, which makes the determination of the incident photon flux almost impossible to calculate, the *fom* was not used in its traditional form in the development stages of the g13 experiment.

Figure 7.2. ϕ distribution of events from the $\bar{\gamma} + d \rightarrow p_s + p + \pi^-$ reaction, with invariant-mass value $W = 1.6 - 1.9$ GeV [124]. Data were integrated over all θ to increase statistics. The offset phase ϕ_0 was extracted from the fit with a function given in Eq. (7.7). Points with big uncertainties correspond to the areas of the CLAS between the different sectors that were removed by applying the fiducial cuts.

7.3.1 Optimization of Fit Parameters

Several choices were made to optimize the fit parameters. The polarization-plane systematic offset, ϕ_0 , depends upon the accuracy of the alignment of the diamond radiator in the goniometer. Since no changes have been made during the g13 experiment to this alignment, the offset ϕ_0 is expected to be constant throughout the experiment. Therefore, ϕ_0 can be determined in a separate analysis (using a different reaction) and this value can be used in fitting the deuteron photodisintegration data. The high-statistics reaction $\bar{\gamma} + d \rightarrow p_s + p + \pi^-$ was used to determine this offset [124]. Specifically, the ϕ_0 value was obtained from fits to high-statistics ϕ distributions using Eq. (7.7). In order to reduce the statistical uncertainty, the ratio between *Para* and *Perp* data was integrated over all θ . A photon-energy bin with a positive asymmetry was chosen to improve the fit [124]. Figure 7.2 shows the distribution used to determine ϕ_0 overlaid with the function $Y(\phi)$. The value of ϕ_0 was determined to be $0.125^\circ \pm 0.172^\circ$. Even though this is consistent with zero, the value of 0.125° was used for fixing ϕ_0 in the fit function.

The polarization ratio $P_R = P_\gamma^\parallel / P_\gamma^\perp$ can be determined using the polarization tables (see Chap. 6). This ratio can also be determined from fitting the distribution of the data as mentioned before. Ideally, one should determine this ratio using the two independent methods, studying in this way the systematic uncertainties in the polarization value. However, extracting the polarization ratio from the data proved to be more challenging, especially when fitting low-statistics bins. Fitting the distributions by treating P_R as a free parameter produced unrealistic results, and it was thus decided to constrain the polarization ratio to the value extracted from the polarization tables. A high-statistics channel with a well measured beam asymmetry can be used to extract this ratio and study the systematics of the determined polarization indicated in the tables (this is discussed in more detail in the following chapter).

The incident photon flux should be constant for a given photon-energy bin, and independent of the reaction of interest. For this reason, a high-statistics reaction can be used to determine the flux ratio between *Para* and *Perp* data. However, to be able to use the determined values from the high-statistics channel, the data chosen, as well as the analysis cuts used, needed to be the same. To remove this complication, an integrated θ bin of deuteron photodisintegration events produced the statistics needed to accurately determine the flux ratio F_R . The value of F_R determined from the integrated θ bin was then used for the individual θ -bin fitting with the function of Eq. 7.7.

The above constraints on ϕ_0 , P_R , and F_R reduced the free parameters of the fit of Eq. (7.7) to one: $\bar{P}\Sigma$. A study that focused on the quality of the fit using simulated data was done by a collaborator and can be found in Ref. [125]. Studies on the systematics of the fit and the analysis cuts, along with the results, are presented in the following chapter.

Chapter 8

Systematic Studies

Systematic studies aim to identify any biases in the measurements. There are several sources of systematic errors in measurements and a lot of effort has been put into identifying the ones that contribute the most. This chapter describes the studies performed by the author to estimate the systematic uncertainties that result from various sources.

8.1 Quality of Fit

The quality of the fit was determined from the χ^2 values. A reliable fit results in a confidence that the observable is correctly determined, especially since the azimuthal asymmetry is determined from fitting the ratio of the differences of azimuthal distributions between the *Para* and *Perp* data to their sum (see Fig. 8.1). From the fits, the χ^2 per degree of freedom was calculated. These values are shown in Fig. 8.2. The distribution peaks

Figure 8.1. Ratio $(N_{||} - N_{\perp})/(N_{||} + N_{\perp})$ for data in photon-energy bin $E_{\gamma} = 1.2$ GeV and $90^{\circ} < \theta_{c.m.} < 95^{\circ}$. The red line corresponds to the fit function of Eq. (7.7) used to determine the asymmetry.

Figure 8.2. The χ^2 per degree of freedom distribution as calculated from fitting the ϕ distribution in each photon energy and $\theta_{c.m.}$ bin.

close to χ^2 per degree of freedom of 1, which indicates a good overall fit to the individual bins. To obtain adequate fits on the ratio of the ϕ distributions, the $\theta_{c.m.}$ -bin widths were increased for higher photon-energy bins to account for the drop in the differential cross section ($\frac{d\sigma}{dt} \sim s^{-11}$).

8.2 Estimating Systematic Uncertainties

Many different studies were performed to determine the systematic uncertainties associated with the azimuthal asymmetries. Table 8.1 summarizes the different sources of systematics and how these were estimated. Detailed descriptions for the individual studies are presented in the following subsections. These were used to calculate the overall uncertainty of each individual point.

The systematic studies performed here were based on statistical means. As shown in Table 8.1, the systematic uncertainties were determined by varying one selection criterion at a time (e.x. missing-mass-squared cut) and studying how this affects the final result on the azimuthal asymmetry. In several cases, the selection criterion changed our statistical sample, introducing statistical variations in the comparison between the azimuthal asymmetries. These statistical variations were accounted for by quoting an upper limit on

Table 8.1. Summary of the systematic uncertainties in the determination of the azimuthal asymmetry of $\vec{\gamma} + d \rightarrow p + n$.

Source	Study
Photon polarization – P_γ	Reference [126]
Offset – ϕ_0	$\phi_0 = -0.047^\circ$ vs $\phi_0 = 0.297^\circ$
Flux ratio – F_R	F_R^{min} vs F_R^{max}
ϕ -bin Width	$\Delta\phi = 12.5^\circ$ vs $\Delta\phi = 25.0^\circ$
Particle ID cuts	3σ vs 2σ $\Delta\beta$ cuts
Missing-mass-squared cuts	3σ cut vs 2σ cut
Fiducial Cuts	nominal vs $\sim 3^\circ$ tighter cuts
Background subtraction	Probabilistic event weighting vs Bin scaling

the estimated systematic uncertainties using the mean and RMS values of the difference between the azimuthal asymmetries. More details on this procedure can be found in the following subsections.

8.2.1 Photon Polarization

One of the main systematic uncertainties in the determination of the azimuthal asymmetry was due to the determination of the photon polarization. This was estimated to be less than 7% [126] by studying the high-statistics channel $\vec{\gamma} + p \rightarrow p + \pi^0$ (using data from a different experiment). This channel allowed us to extract the azimuthal asymmetry for each E-counter, using the *Para* and *Perp* data independently. Specifically, due to the typical drift in the coherent-edge position, the coherent peak between two settings had a significant overlap region. A comparison between the overlap-region asymmetries determined using data from different nominal coherent-edge positions resulted in an estimate of the uncertainty in the photon polarization. Figure 8.3 shows the ratio of the azimuthal asymmetries extracted using the 1.1–1.3 and 1.3–1.5 GeV photon-energy bins as a function of E-counter. The top panel corresponds to the ratio of the asymmetries using the *Para* data and the bottom panel to the ratio of asymmetries using the *Perp* data.

The effect of the photon polarization uncertainty on the extracted value of Σ was twofold. First, the uncertainty on the photon polarization propagated in the polarization ratio P_R , which was used as a fixed parameter in the fit. This affected the fit quality and

Figure 8.3. Ratio of azimuthal asymmetries using the 1.1–1.3 and 1.3–1.5 GeV photon-energy bin for each E-counter. The blue line corresponds to a 0^{th} -order fit to the ratios with its extracted parameter indicated on the plot. The upper panel corresponds to the ratio of *Para* data and the bottom to the ratio of *Perp* [126].

resulted in an error in the calculation of the free-fit parameter that was used to determine Σ . This effect was investigated by comparing the asymmetries, Σ , determined from the fit when the photon-polarization ratio was fixed at its maximum and minimum values, according to

$$\Delta P_R = P_R \sqrt{\left(\frac{\Delta P_{\parallel}}{P_{\parallel}}\right)^2 + \left(\frac{\Delta P_{\perp}}{P_{\perp}}\right)^2} \text{ and thus,} \quad (8.1)$$

$$P_R^{min} = P_R - \Delta P_R = P_R(1 - 0.07 \times \sqrt{2}) \quad (8.1)$$

$$P_R^{max} = P_R + \Delta P_R = P_R(1 + 0.07 \times \sqrt{2}). \quad (8.2)$$

Figure 8.4 shows the azimuthal asymmetry Σ determined using the minimum (blue points) and maximum (red points) values of P_R . The data used for the two cases of P_R in determining Σ are exactly the same, and thus, any difference between the asymmetries is entirely due to the uncertainty of the determination of P_R .

To estimate the effect that the uncertainty of P_R has on the asymmetry Σ , the

Figure 8.4. Azimuthal asymmetry for data in the photon-energy bin 1.1–1.3 GeV, as a function of $\theta_{c.m.}$. The blue points correspond to Σ extracted using P_R^{min} , whereas red points correspond to Σ extracted using P_R^{max} .

differences between the Σ determined using the minimum and maximum values of P_R were plotted as a function of $\theta_{c.m.}$ for each photon-energy bin (see Fig. 8.5). The mean difference was determined by fitting these data with a 0th-order polynomial, indicated by the black lines in Fig. 8.5. The blue lines correspond to the uncertainty of the mean difference calculated using the deviations of the points from the average value. The exact values of $\overline{\Delta\Sigma}$, along with their uncertainties, are shown in each panel of Fig. 8.5.

The uncertainty of the degree of photon polarization also results in a $\sim 5\%$ uncertainty in the calculation of the average polarization $\bar{P}$,

$$\Delta\bar{P} = \frac{1}{2} \sqrt{(\Delta P_{\parallel})^2 + (\Delta P_{\perp})^2}, \quad (8.3)$$

by using $P^{\perp} \approx P^{\parallel} \approx \bar{P}$. Since Σ is determined from the fit parameter $C = \bar{P}\Sigma$ (see Sec. 7.3), which itself is associated with an uncertainty, the total uncertainty on Σ had to be calculated by propagation of errors,

$$\Delta\Sigma = \Sigma \sqrt{\left(\frac{\Delta C}{C}\right)^2 + \left(\frac{\Delta\bar{P}}{\bar{P}}\right)^2}. \quad (8.4)$$

Figure 8.5. Difference between the asymmetries determined from the two values of P_R as a function of $\theta_{c.m.}$. The different panels correspond to data from different photon-energy bins. The values of $\overline{\Delta\Sigma}$ indicated in each panel correspond to the mean differences along with their uncertainty.

In general, the uncertainties due to P_R and $\bar{P}$ are highly correlated and should be treated together. However, it is evident that the uncertainty in the average polarization ($\sim 5\%$) has a much bigger effect on Σ than the uncertainty in the polarization ratio P_R ($\sim 0.1\%$), and thus, the latter can be neglected.

It is possible to reduce the uncertainty of the degree of photon polarization to about 3%. In order to do so, a high-statistics channel (e.g. single-pion production) is needed to be analyzed using the g13 data. The idea behind this is to try to determine Σ using data within each photon-energy bin by applying different photon-energy cuts. Because of the drift in coherent-edge position, the asymmetry determined for each E-counter will overlap within adjacent photon-energy cuts. Comparing these would allow us to correct the degree of photon polarization such that it is self-consistent [115].

8.2.2 ϕ_0 offset

In addition to the ratio P_R , the fit was constrained by the parameter ϕ_0 . As mentioned in Chap. 7, this constant accounts for any offset between the photon polarization vector and the x -axis for the *Para* orientation, or the y -axis for the *Perp* orientation. The value of this offset was determined using a high-statistics channel to be $0.125^\circ \pm 0.172^\circ$ [124]. Since the accuracy of this parameter affects the quality of the fit, the asymmetry Σ was obtained for the upper and lower limits of ϕ_0 , in an attempt to estimate the magnitude of its effect. This was done by comparing the Σ values obtained using the two values of ϕ_0 of 0.297° and -0.047° . A comparison between the Σ values for the 1.1–1.3 GeV photon-energy bin is shown in Fig. 8.6.

Figure 8.7 shows the differences $\Delta\Sigma$, as a function of $\theta_{c.m.}$, for all six photon-energy bins. Since these differences did not show any variation with $\theta_{c.m.}$, the mean difference $\overline{\Delta\Sigma}$ was determined by fitting the data with a 0^{th} -order polynomial. The mean values along with their uncertainties are indicated in each photon-energy bin in Fig. 8.7. On average, these differences correspond an uncertainty of less than 0.2% on Σ . Since the effect of the ϕ_0 offset should be independent of the photon-energy and angular bin, the average over all bins was calculated and used as the value of the systematic uncertainty.

Figure 8.6. Azimuthal asymmetry as a function of $\theta_{c.m.}$ for data in the lowest photon-energy bin. The blue points correspond to Σ obtained using $\phi_0 = -0.047^\circ$ whereas the red points correspond to Σ obtained using $\phi_0 = 0.297^\circ$.

8.2.3 Flux ratio

The third fit parameter that was treated as a constant was the photon flux ratio, F_R . Several studies have been performed to examine the effect of this quantity on the determined Σ . Ideally, the photon flux would be known from some independent monitor. However, determining the photon flux using an independent monitor is very challenging, and of course just like any other measurement, it would also have its own uncertainties. In this work, the flux ratio was determined by the same procedure used for determining Σ (see Sec. 7.3). In order to get a more reliable value for F_R , the function $Y(\phi)$ was fitted to the data integrated over all $\theta_{c.m.}$ bins. In this way, the fit was less sensitive to any statistical fluctuations. Table 8.2 shows the flux ratios obtained from fitting the ratios of azimuthal distributions in each photon-energy bin.

To study the sensitivity of Σ to the flux ratio F_R , Σ was determined using the upper $(F_R + \Delta F_R)$ and lower $(F_R - \Delta F_R)$ values of F_R in each photon-energy bin. Figure 8.8 shows a comparison between these two results for the lowest photon-energy bin. The

Figure 8.7. Differences between the Σ obtained using the maximum and minimum values of the ϕ_0 offset. The different panels correspond to data from different photon-energy bins. The values of $\Delta\Sigma$ indicated in each panel correspond to the mean differences along with their uncertainty.

Table 8.2. Photon-flux ratios determined from fitting the integrated- $\theta_{c.m.}$ bin for the six different photon-energy bins.

E_γ Bin (GeV)	F_R	ΔF_R
1.1 – 1.3	1.293	± 0.010
1.3 – 1.5	1.032	± 0.009
1.5 – 1.7	1.119	± 0.013
1.7 – 1.9	0.930	± 0.015
1.9 – 2.1	1.233	± 0.038
2.1 – 2.3	0.944	± 0.039

differences $\Delta\Sigma$ in the six different photon-energy bins are shown in Fig. 8.9. Since there was no apparent $\theta_{c.m.}$ dependence on the differences, these points were fitted with a 0^{th} -order polynomial to determine the mean value. The mean value $\overline{\Delta\Sigma}$, along with their uncertainties, are indicated in the individual panels in Fig. 8.9. On average, these differences correspond to less than 1% of Σ .

Additional studies have been performed that examined the sensitivity of Σ on the photon-flux ratio. Specifically, in the low photon-energy bins where there are enough statistics, the flux ratio was determined for each bin in $\theta_{c.m.}$. The differences between the flux ratio determined from an integrated $\theta_{c.m.}$ bin and the flux ratios for each $\theta_{c.m.}$ bin were, in general, within the statistical uncertainty of the integrated- $\theta_{c.m.}$ flux ratio. A comparison between the determined asymmetries obtained by treating F_R as a fixed parameter and the ones obtained treating F_R as a free parameter in the fit was done. The differences between the two Σ were comparable to the differences quoted in Fig. 8.9.

8.2.4 Fiducial cuts

The effect the fiducial cuts (see Sec. 5.6) had on Σ was examined by varying the ϕ -dependent θ cut. Specifically, about a 3° tighter cut was applied, as shown in Fig. 8.10. The asymmetries were determined for all photon-energy bins and a comparison between the two cuts was done. Figure 8.11 shows a comparison between the asymmetries in the photon-energy bin 1.1–1.3 GeV (blue squares: nominal; red circles: $\sim 3^\circ$ tighter).

The differences between the two Σ values did not show any angular dependence.

Figure 8.8. A comparison between the Σ obtained using the maximum and minimum values of the photon-flux ratio, as a function of $\theta_{c.m.}$ (data from the 1.1–1.3 GeV photon-energy bin). The blue points correspond to Σ obtained using $F_R^{\min}$, whereas the red points correspond to Σ obtained using $F_R^{\max}$.

Just as in the previous studies, the mean $\Delta\Sigma$ was determined by fitting these differences with a 0^{th} -order polynomial. Figure 8.12 shows the differences in Σ from the individual photon-energy bins.

Unfortunately, the 3° -tighter fiducial cut on the highest photon-energy bin reduced the statistics to a point where the $Y(\phi)$ fit became unreliable. This is the reason the 3° -tighter cut has a larger effect on the asymmetry of the most forward-going particles (see Fig. 8.11). The differences between the points includes statistical fluctuations, as well as systematic effects. However, it is extremely difficult to separate the two sources. To reduce the effect on any statistical fluctuations, it was decided to compare the asymmetries for the different fiducial cuts before any background subtraction. Figure 8.13 shows the differences $\Delta\Sigma$ between the asymmetries obtained with a nominal and a 3° -tighter fiducial cut, before any background subtraction. These plots clearly indicate that the differences in the asymmetries for the forwardmost $\theta_{c.m.}$ bin were the largest. To better study the effect that the fiducial cuts might have on these $\theta_{c.m.}$ bins, the differences in the asymmetries

Figure 8.10. Lab-frame phase space of the protons in sector 2. The black line in the left panel corresponds to the nominal fiducial cut, whereas the black line in the right panel corresponds to the 3°-tighter fiducial cut applied to study the systematic uncertainties of this cut.

results of this study are shown in Fig. 8.14.

The mean differences $\overline{\Delta\Sigma}$ of the latter two studies were significantly lower than the ones obtained by comparing the asymmetries obtained using the nominal and the 3°-tighter fiducial cut after background subtraction. However, the latter was used as an upper limit of the systematic uncertainty of the fiducial cuts. This upper limit corresponds to an uncertainty of about 3% of Σ .

8.2.5 ϕ -bin width

As discussed in Sec. 7.3, Σ is obtained by fitting the proton azimuthal distributions with a $\cos 2\phi$ function. These distributions are binned in ϕ with a pre-determined number of bins. A comparison between the Σ obtained by varying the bin width was carried out. Figure 8.15 shows the distribution $\frac{N(\phi)^{\parallel}-N(\phi)^{\perp}}{N(\phi)^{\parallel}+N(\phi)^{\perp}}$ in one $\theta_{c.m.}$ and photon-energy bin. The left panel corresponds to the ϕ distribution binned in 29 ϕ bins ($\Delta\phi = 12.5^\circ$), whereas the right panel corresponds to the ϕ distribution binned in 17 ϕ bins ($\Delta\phi = 25.0^\circ$).

Figure 8.11. A comparison between the Σ obtained using the nominal and the $\sim 3^\circ$ -tighter fiducial cuts as function of $\theta_{c.m.}$ (data from the 1.1–1.3 GeV photon-energy bin). The blue points correspond to Σ obtained using F_R^{min} , whereas the red points correspond to Σ obtained using F_R^{max} .

Figure 8.16 shows a comparison between the asymmetries obtained with the two ϕ -bin widths for the lowest photon-energy bin. The differences between these points was used to determine the systematic uncertainty that results from different ϕ binning. Figure 8.17 shows the differences $\Delta\Sigma$ between the asymmetries obtained using the different ϕ -bin widths, for each photon energy bin. Since there was no evident $\theta_{c.m.}$ dependence in the differences $\Delta\Sigma$, the average difference $\overline{\Delta\Sigma}$ (indicated in each panel in Fig. 8.17) was assigned to the systematic uncertainty of the ϕ -bin width. On average, the ϕ -binning has less than 7% systematic uncertainty on Σ .

Different groups [127, 128] have shown that the ϕ -bin method for determining the azimuthal asymmetry is quite robust. In fact they have shown, using simulated data, that the size of the ϕ -bin width has a very small systematic effect on the value of the azimuthal asymmetry. Our study is strongly influenced by statistical variations and for this reason the value quoted here really represents an upper limit of the systematic uncertainty. A smaller variation in the ϕ -bin width (lets say by $\sim 20\%$ of the nominal size of the bin)

Figure 8.12. Differences between the Σ obtained using the nominal and the 3° -tighter fiducial cuts. The different panels correspond to data from different photon-energy bins. The mean values of $\overline{\Delta\Sigma}$ along with their uncertainties are indicated in each panel.

Figure 8.13. Difference between the asymmetries obtained using the nominal and the 3°-tighter fiducial cut, before any background subtraction. The different points correspond to the differences between the asymmetries in different photon-energy bins.

Figure 8.14. Difference between the asymmetries obtained using the nominal and the 3°-looser fiducial cut, before any background subtraction. The different points correspond to the differences between the asymmetries in different photon-energy bins.

Figure 8.15. ϕ distributions $\frac{N(\phi)^{\parallel} - N(\phi)^{\perp}}{N(\phi)^{\parallel} + N(\phi)^{\perp}}$ using two different ϕ -bin widths: 12.5° (left panel), and 25.0° (right panel). The red line corresponds to the function $Y(\phi)$ that is fitted to these distributions in order to determine Σ .

would have given us a better estimate on the upper limit, and might have reduced the systematic uncertainty to a few percent.

8.2.6 Particle ID cuts

The systematic effect that the particle-identification cut had on the asymmetry was studied by varying the momentum-dependent $\Delta\beta$ cut (see Sec. 5.3.2 for details on the particle identification method). Specifically, a comparison between the asymmetries using

Figure 8.16. A comparison between the Σ obtained using the different ϕ -bin widths. The blue points corresponds to Σ obtained using $\Delta\phi = 12.5^\circ$, whereas the red points correspond to Σ obtained using $\Delta\phi = 25.0^\circ$.

a 3σ and a 2σ $\Delta\beta$ cut was done. Figure 8.18 shows the $\Delta\beta$ distribution as a function of proton momentum, with the 3σ (red line) and 2σ (black-dashed line) cuts overlaid on the distribution.

Ideally, the asymmetries Σ after background subtraction should have been compared for the two different cuts. However, the background-subtraction procedure was very time-consuming and for this reason the comparison was made between Σ before any background subtraction. Figure 8.19 shows the comparison between the asymmetries for the lowest photon-energy bin, obtained from the two different particle-identification cuts before any background subtraction.

Figure 8.20 shows the difference $\overline{\Delta\Sigma}$ (indicated in each panel in Fig. 8.20). Since there wasn't any evident $\theta_{c.m.}$ dependence, the mean difference was calculated for each photon-energy bin and used as the systematic uncertainty of the particle identification cut. On average the particle-identification cut has a 7% systematic uncertainty on Σ .

Figure 8.17. Difference between the asymmetries determined from using the different ϕ bin widths. The different panels correspond to data from different photon-energy bins. The values of $\Delta\Sigma^{\Delta\phi}_{syst}$ indicated in each panel correspond to the values of the parameters determined from fitting the data with a 0th-order fit.

Figure 8.18. $\Delta\beta$ distribution as a function of proton momentum. The red line corresponds to the 3σ cut used for the nominal particle identification procedure, and the black line corresponds to the 2σ cut used for studying the systematic effects that cut has on the determined asymmetry.

Figure 8.19. A comparison between the Σ obtained using the different particle-identification cuts, before any background subtraction. The blue points corresponds to Σ obtained using a 3σ cut, whereas the red points correspond to Σ obtained using a 2σ cut on the $\Delta\beta$ distribution.

Figure 8.20. Difference between the asymmetries determined from using the different particle-identification cuts. The different panels correspond to data from different photon-energy bins. The values of $\Delta\Sigma$ indicated in each panel correspond to the values of the parameters determined from fitting the data with a 0^{th} -order fit.

8.2.7 Missing-mass-squared cuts

In this analysis, the reaction was reconstructed by applying on the square of the missing-mass cuts. The nominal missing-mass-squared cut was determined by fitting the missing-mass-squared distribution with a Gaussian plus two exponentials. The 3σ cuts were determined from the fit parameters. Figure 8.21 shows the missing-mass-squared distribution of the reaction $\vec{\gamma} + d \rightarrow p + X$. The 2σ and 3σ cuts are indicated by the vertical lines.

Figure 8.21. Missing-mass-squared distribution of the reaction $\vec{\gamma} + d \rightarrow p + X$. The vertical lines correspond to the 2σ and 3σ cuts applied to study their effect on the obtained asymmetry.

A study was performed to examine the effect of a missing-mass-squared cut on the determined asymmetry Σ . Figure 8.22 shows a comparison between the asymmetries determined with a 3σ missing-mass-squared cut (red points) and a 2σ missing-mass-squared cut (blue points) for data in the lowest photon energy bin.

The differences between the Σ obtained in the two cases were used to estimate the value of the systematic uncertainty due to this cut. Figure 8.23 shows these differences $\Delta\Sigma$ for each photon-energy bin. On average, the missing-mass-squared cut has a 5% systematic uncertainty on Σ .

Figure 8.22. A comparison between the Σ obtained using a 3σ missing-mass-squared cut (red points) and a 2σ missing-mass-squared cut (blue points). The asymmetries shown here correspond to data in the lowest photon-energy bin.

8.2.8 Background subtraction

The systematic effect that the background subtraction had on the asymmetries was studied by examining the differences in the asymmetries determined using two different background-subtraction methods. As mentioned in Sec. 5.10, the background-subtraction method used for the final results was the probabilistic event weighting. In this, the probability of an event being a signal or background event was calculated using the missing-mass-squared distribution of its 200 closest neighbors. The number of closest neighbors was chosen such that adequate fits on the missing-mass-squared distributions would be obtained.

The second method used for background subtraction corresponds to the same principle as the probabilistic event weighting, with a slightly different approach. Specifically, in this method, the missing-mass-squared distribution for each $\theta_{c.m.}$ and ϕ bin was fitted with the same functional form as the one used in the probabilistic weighting. From these fits, a scaling factor was calculated that corresponds to the ratio between the integral of the signal (Gaussian) and the total (Gaussian plus two exponentials) between the 3σ missing-mass-

Figure 8.23. Difference between the asymmetries determined from a 3σ missing-mass-squared cut and a 2σ missing-mass-squared cut. The different panels correspond to data from different photon-energy bins. The values of $\Delta\Sigma$ indicated in each panel correspond to the mean value of $\Delta\Sigma$.

Figure 8.24. Missing-mass-squared distributions for different ϕ bins, within a given photon-energy and $\theta_{c.m.}$ bin. The fitted functions were used to calculate the scaling factors, $s2t$, indicated in the plots.

squared cuts. This factor was used to scale the proton yields in each bin in $\theta_{c.m.}$ and ϕ , thus removing “any” background contribution. Figure 8.24 shows the missing-mass-squared distributions for six different ϕ bins in a given $\theta_{c.m.}$ and photon-energy bin, fitted with the functional form of a Gaussian plus two exponentials. The yellow region corresponds to the region where the Gaussian and the total function were integrated to calculate the scaling factor.

Unfortunately, the fits in the nominal $\theta_{c.m.}$ and ϕ bins failed to produce reasonable results in the higher photon-energy bins and at backward angles. This is because the number of events in those bins was much smaller than the number of events required to have an adequate missing-mass-squared distribution. To account for this, the comparison between the asymmetries was done for larger $\theta_{c.m.}$ -bin widths, in an attempt to increase the statistics and produce adequate missing-mass-squared fits. In principle, this makes the comparison between the two methods more reliable since, in the probabilistic event weighting, the bin widths varied to always include 200 events. Figure 8.25 shows a comparison between the asymmetries determined using the probabilistic event-weighting and the bin-scaling background-subtraction methods.

The systematic uncertainty of the background subtraction was determined by studying the differences between the asymmetries obtained using the two different methods. Figure 8.26 shows these differences as a function of $\theta_{c.m.}$ for all photon-energy bins. Since there was no evident $\theta_{c.m.}$ dependence on these differences, the systematic uncertainty was obtained by determining the mean value of these differences in each photon-energy bin. On average the systematic uncertainty of the background subtraction corresponds to about 7% of Σ .

It is evident that our choice of two exponentials for describing the background might introduce systematic uncertainties in the background subtraction procedure. Because of this, our study could benefit from simulated data in the effort to better understand the actual shape of the background. This requires precise knowledge on the cross sections of the reactions that contribute to our background, along with precise knowledge of the detector efficiency. Unfortunately, there are no available cross-section data on the background reactions. Because of this, our study of the systematics associated with the choice of the background shape cannot benefit from simulated data.

Figure 8.25. Comparison between the asymmetries determined using the probabilistic event weighting (blue) and bin scaling (red). The asymmetries shown here correspond to data in the lowest photon-energy bin.

8.2.9 Photon-energy bins

The 200 MeV width of the photon-energy bins used for this analysis was chosen such that there were adequate statistics in each bin to study the $\theta_{c.m.}$ dependence of the azimuthal asymmetry. Since the g13 experiment took data for six different nominal coherent-edge positions, it made it easier to bin the photon energies with these nominal coherent-edge positions. In general, the photon flux is relatively constant close to the coherent edge, giving equal weight to all events in the bin.

In the method used for obtaining the asymmetry from the data, an assumption was made for the average photon energy. Specifically, when the ratio of *Para* and *Perp* is constructed to obtain the asymmetry, we assumed that the average photon energy was the same for both *Para* and *Perp* data. This was checked by calculating the average photon energy in each photon-energy bin for the two photon polarizations. Only events that passed all analysis selection criteria were used for this calculation. Table 8.3 shows the average photon energy for the two orientations of the photon polarizations for the six

Figure 8.26. Difference between the asymmetries determined from the two background subtraction methods (see text). The different panels correspond to data from different photon-energy bins. The values of $\overline{\Delta\Sigma}$ indicated in each panel correspond to the mean value of $\Delta\Sigma$.

Table 8.3. Average photon energy for *Para* and *Perp* data in each photon-energy bin.

E_γ bin (GeV)	<i>Para</i> $\bar{E}_\gamma$ (GeV)	<i>Perp</i> $\bar{E}_\gamma$ (GeV)
1.1 – 1.3	1.206	1.208
1.3 – 1.5	1.403	1.403
1.5 – 1.7	1.600	1.601
1.7 – 1.9	1.793	1.798
1.9 – 2.1	1.991	1.991
2.1 – 2.3	2.201	2.200

different photon energy bins. The differences between the average photon energy of *Para* and *Perp* data are of the order of a few MeV, which is well within the resolution of the tagger spectrometer.

8.3 Summary

Table 8.1 summarized the methods used to study the systematic uncertainties associated with the azimuthal asymmetry of deuteron photodisintegration. In general, these systematic uncertainties can be split into four categories. Systematic uncertainties which are correlated for all photon energies and angles are referred to as overall systematic. Angle correlated systematics correspond to the systematic uncertainties that are the same for all angles but different between photon energy bins. Photon energy correlated systematics correspond to the systematic uncertainties that are the same for all photon energies but different between different angles. Finally, point-to-point systematic uncertainties correspond to the uncertainties that differ between all angles and all photon energies. Table 8.4 shows the different categories and the associate sources of systematic uncertainties that correspond to each category.

As mentioned before, the systematic uncertainties were determined using statistical means. In addition, the variations between the different point-to-point systematic within one photon-energy bin were quite small. For this reason, the point-to-point systematic uncertainties were assigned with an estimate of the upper limit for each photon-energy bin (one value of the systematic uncertainty for each photon energy bin). This upper limit is

Table 8.4. Categories of systematic uncertainties associated with deuteron photodisintegration.

Category	Source of Systematic
Overall	Photon polarization ϕ_0 offset
Photon-energy correlated	Flux ratio
Angle correlated	—
Point-to-point	ϕ -bin width Particle ID cuts Missing-mass-squared cuts Fiducial cuts Background subtraction

determined by taking into account the variations of each point from the mean value (i.e. $|\overline{\Delta\Sigma}| + \sigma_{\overline{\Delta\Sigma}}$). More precise studies on the choice of cuts and the determination of systematic uncertainties, using simulated data, can and will be done for the publication.

Tables 8.5 and 8.6 show the results of these systematic studies. The next chapter presents the results obtained from this study for the azimuthal asymmetry of deuteron photodisintegration, incorporating the systematic studies performed, as described in this chapter.

Table 8.5. Systematic uncertainties for photon-energy bins $E_\gamma = 1.1 - 1.7$ GeV for the reaction $\vec{\gamma} + d \rightarrow p + n$. The values $|\overline{\Delta\Sigma}|$ correspond to the mean differences of the azimuthal asymmetries and $\sigma_{\overline{\Delta\Sigma}}$ correspond to the RMS values that are calculated using the differences between the points and the mean.

Photon Energy Bin (GeV)	Source of Systematic	$ \overline{\Delta\Sigma} $	$\sigma_{\overline{\Delta\Sigma}}$
1.1 – 1.3	Photon polarization	5%	—
	ϕ_0 offset	0.0002	—
	Flux ratio	0.0002	0.0009
	ϕ -bin width	0.015	0.017
	Particle ID cuts	0.0015	0.0092
	Missing-mass-squared cuts	0.003	0.015
	Fiducial cuts	0.0	0.0095
	Background subtraction	0.0011	0.0071
1.3 – 1.5	Photon polarization	5%	—
	ϕ_0 offset	0.0002	—
	Flux ratio	0.0006	0.0013
	ϕ -bin width	0.0065	0.0072
	Particle ID cuts	0.004	0.012
	Missing-mass-squared cuts	0.003	0.015
	Fiducial cuts	0.0017	0.0089
	Background subtraction	0.0	0.012
1.5 – 1.7	Photon polarization	5%	—
	ϕ_0 offset	0.0002	—
	Flux ratio	0.0010	0.0011
	ϕ -bin width	0.025	0.046
	Particle ID cuts	0.0	0.022
	Missing-mass-squared cuts	0.011	0.032
	Fiducial cuts	0.0	0.011
	Background subtraction	0.004	0.011

Table 8.6. Systematic uncertainties for photon-energy bins $E_\gamma = 1.7 - 2.3$ GeV for the reaction $\vec{\gamma} + d \rightarrow p + n$. The values $|\overline{\Delta\Sigma}|$ correspond to the mean differences of the azimuthal asymmetries and $\sigma_{\overline{\Delta\Sigma}}$ correspond to the RMS values that are calculated using the differences between the points and the mean.

Photon Energy Bin (GeV)	Source of Systematic	$ \overline{\Delta\Sigma} $	$\sigma_{\overline{\Delta\Sigma}}$
1.7 – 1.9	Photon polarization	5%	—
	ϕ_0 offset	0.0002	—
	Flux ratio	0.0013	0.0019
	ϕ -bin width	0.015	0.019
	Particle ID cuts	0.007	0.023
	Missing-mass-squared cuts	0.005	0.019
	Fiducial cuts	0.0	0.015
	Background subtraction	0.0	0.015
1.9 – 2.1	Photon polarization	5%	—
	ϕ_0 offset	0.0002	—
	Flux ratio	0.0017	0.0054
	ϕ -bin width	0.012	0.044
	Particle ID cuts	0.027	0.025
	Missing-mass-squared cuts	0.020	0.022
	Fiducial cuts	0.018	0.027
	Background subtraction	0.027	0.038
2.1 – 2.3	Photon polarization	5%	—
	ϕ_0 offset	0.0002	—
	Flux ratio	0.004	0.013
	ϕ -bin width	0.009	0.049
	Particle ID cuts	0.017	0.040
	Missing-mass-squared cuts	0.021	0.089
	Fiducial cuts	0.062	0.047
	Background subtraction	0.041	0.023

Chapter 9

Results and Discussion

9.1 Overview

Chapters 5–7 discuss the analysis procedure we followed in order to obtain the final results for the azimuthal asymmetry of the reaction $\vec{\gamma} + d \rightarrow p + n$. Detailed descriptions of the theoretical predictions along with previous measurements of the azimuthal asymmetry are presented in Chap. 2. Chapter 8 addresses the systematic uncertainties involved in the determination of the azimuthal asymmetry using the g13 data. Here we discuss the final results, comparing them with the theoretical model predictions, as well as with previous measurements. Tabulated results can be found in Appendix B.

9.2 Comparison with Previous Measurements

Previous measurements exist for the azimuthal asymmetry, Σ , of the reaction $\vec{\gamma} + d \rightarrow p + n$ [69, 70] in the energy range $E_\gamma = 0.8 - 1.6$ GeV and only for $\theta_{c.m.} = 90^\circ$ [69, 70]. Figure 9.1 shows the result obtained from this work compared to these previous measurements. The red and cyan points in Fig. 9.1 correspond to results from Yerevan [69, 70], whereas the black points correspond to the results of the present work. The $\theta_{c.m.}$ bin width used in our data varied from 5° widths for the lowest photon-energy bins to 30° widths for the highest photon-energy bins, centered at $\theta_{c.m.} = 90^\circ$, in an attempt to increase the statistics in the higher photon-energy bins in which the cross section of deuteron photodisintegration is sufficiently smaller ($d\sigma/dt \sim s^{-11}$). For the Yerevan results, the photon-energy bin width varies from 40 MeV to 140 MeV depending on E_γ . The photon-energy bin width for the CLAS data was kept constant at 200 MeV. The results from this work significantly increase the kinematic coverage of the previous data set (see Figs. 9.2 and 9.3 and discussion in next section). Specifically, at $\theta_{c.m.} = 90^\circ$, the data set for the azimuthal asymmetry is extended to 2.2 GeV, whereas new data are added for center-of-mass angles between 25° and 160° and energies 1.2 to 2.2 GeV. This latter kinematic domain is completely unexplored in terms of polarization measurements. The results of

Figure 9.1. Azimuthal asymmetry as a function of the incident photon energy, for $\theta_{c.m.} = 90^\circ$. The red and cyan points correspond to results from Yerevan [69, 70], whereas the black points to the results from the present work. Statistical and systematic uncertainties of the CLAS data are shown with the black and red error bars respectively.

this work agree well with the published results from Yerevan. Figure 9.1 also indicates that the g13(b) data set allowed us to significantly improve the precision of the existing data at these lower photon energies.

9.3 Angular and Energy Distributions of Σ

Figures 9.2 and 9.3 show the results for the azimuthal asymmetry, Σ , determined from this work. Specifically, Fig. 9.2 shows the angular dependence of the asymmetry for the six different photon-energy bins (between 1.1 and 2.3 GeV), whereas Fig. 9.3 shows the energy dependence of the asymmetry for five angles ($\theta_{c.m.} = 35^\circ, 60^\circ, 90^\circ, 120^\circ$, and 145°) in the center-of-mass frame (see Tables in Appendix B for bin width of each point). The bin width in $\theta_{c.m.}$ varies between 5° for forward angles and small photon energies, to 30° for backward angles and large photon energies. The photon-energy bin width is 200 MeV.

In Figure 9.2, at the lowest photon-energy bin, the azimuthal asymmetry exhibits a minimum near $\theta_{c.m.} = 90^\circ$. This minimum shifts to larger angles as the photon energy

Figure 9.2. Azimuthal asymmetry as a function of the proton angle in the center-of-mass frame for six photon-energy bins. Statistical and systematic uncertainties of the CLAS data are shown with the black and red error bars respectively.

increases and vanishes for $E_\gamma \geq 1.8$ GeV, where the asymmetry exhibits a maximum at $\theta_{c.m.} = 90^\circ$ (see lower right panel of Fig. 9.2). The azimuthal asymmetry for all photon-energy bins also changes sign, being negative for small $\theta_{c.m.}$ and positive for intermediate $\theta_{c.m.}$ (see Fig. 9.2). In addition, the asymmetry at backward angles, and specifically for $\theta_{c.m.} = 145^\circ$, is positive for $E_\gamma \leq 1.7$ GeV and changes sign, becoming negative, for $E_\gamma \geq 1.7$ GeV (see Fig. 9.3). The energy dependence of the azimuthal asymmetry at $\theta_{c.m.} = 90^\circ$ indicates a transition from lower to higher values between 1.6 GeV and 2.0 GeV

Figure 9.3. Azimuthal asymmetry as a function of photon energy for five different angles in the center-of-mass frame. Statistical and systematic uncertainties of the CLAS data are shown with the black and red error bars respectively.

(see Fig. 9.3). This might be an indication of a change in the production mechanism. As mentioned before, the data for the azimuthal asymmetry of deuteron photodisintegration are the only polarization data available that cover photon energies $E_\gamma = 1.1 - 2.3$ GeV and all angles.

9.4 General Remarks

Figures 9.4–9.6 show the available data for the differential cross section, $d\sigma/dt$, and polarization observables Σ , P_y , $C_{x'}$, and $C_{z'}$ as a function of photon energy for $\theta_{c.m.} = 90^\circ$.

It should be pointed out here that the results on observables other than the azimuthal asymmetry that are presented in this section correspond to studies of other experiments in which the author was not involved. The discussion presented here corresponds to a contextual interpretation by the author. This was done in an attempt to look for common patterns among these observables that would indicate regions of interest for further exploration. No assumptions on the validity of these other observables along with their uncertainties is made. For more details on these other observables one should referred to the published papers (can be found in Chap. 2).

The azimuthal asymmetry, $\Sigma(90^\circ)$, seems to be constant between 0.8 and 1.6 GeV as is the induced polarization, P_y . On the other hand, the available data on polarization transfers, $C_{x'}$ and $C_{z'}$, indicate that there might be a maximum around $E_\gamma = 1$ GeV. The differential cross section at $\theta_{c.m.} = 90^\circ$ agrees well with the predictions of dimensional scaling. The polarization transfers, $C_{x'}$ and $C_{z'}$ are small for larger photon energies (close to, but not equal to zero). These are close to the values predicted by HHC ($C_{x'} = C_{z'} = 0$). The results for the azimuthal asymmetry obtained from this thesis work indicated a large positive asymmetry. Even though we cannot use the azimuthal asymmetry results to test the HHC rules (HHC does not constrain the values of Σ), the results on the other

Figure 9.4. Differential cross section, $d\sigma/dt$, for $\theta_{c.m.} = 90^\circ$ as a function of photon energy. Details on the data and curves can be found in Fig. 2.6.

Figure 9.5. Induced proton polarization, P_y , and azimuthal asymmetry, Σ , for $\theta_{c.m.} = 90^\circ$ as a function of photon energy.

polarization observables rule out the onset of pQCD at few-GeV photon energies, and the perturbative interpretation of dimensional scaling.

The energy dependence of the asymmetry at $\theta_{c.m.} = 90^\circ$ shows an increase for Σ

Figure 9.6. Polarization transfers, $C_{x'}$ and $C_{z'}$, for $\theta_{c.m.} = 90^\circ$ as a function of photon energy.

at $E_\gamma > 1.5$ GeV, which may indicate a change in the production mechanism for these incident photon energies. In addition, the azimuthal asymmetry at $\theta_{c.m.} = 90^\circ$ peaks (or reaches a plateau) at 2 GeV, and it is quite interesting that none of the other observables shows such a behavior.

Figure 9.7. Differential cross section, $d\sigma/dt$, for $E_\gamma \sim 2$ GeV as a function of proton angles in the center-of-mass-frame. Details on the data and curves can be found in Fig. 2.8.

Figures 9.7–9.9 show the available data for the differential cross section, $d\sigma/dt$, and polarization observables Σ , P_y , $C_{x'}$, and $C_{z'}$ as a function of proton angle in the center-of-mass frame, for $E_\gamma \sim 2$ GeV. The data indicate a change in sign for almost all polarization observables. In fact, the induced polarization and polarization transfer $C_{x'}$ seem to be changing sign at intermediate angles, whereas the azimuthal asymmetry is negative for both forward and backward angles. The polarization transfer $C_{z'}$ is compatible with zero at backward angles. The sign of the polarization observables provides evidence on the relation between different helicity amplitudes in the interference terms.

9.5 Comparison of Σ with Theoretical Predictions

Figure 9.10 shows the azimuthal asymmetry as a function of proton angle in the center-of-mass frame for six different photon-energy bins between $E_\gamma = 1.1$ and 2.3 GeV. The red dashed-dotted line shows the predictions of the QGSM [50], whereas the blue

Figure 9.8. Induced proton polarization, P_γ , and azimuthal asymmetry, Σ , for $E_\gamma \sim 2$ GeV as a function of proton angles in the center-of-mass-frame.

lines correspond to predictions of the HRM model [47]. The thick and thin solid curves correspond to the “on-shell” version of the model and represent predictions based on large and small helicity non-conservation, respectively (see Chap. 2 for details). The dashed and dotted blue lines represent two versions of the “off-shell” approximation. Figure 9.11 shows the energy dependence of the azimuthal asymmetry for $\theta_{c.m.} = 90^\circ$ compared to the two models.

The angular dependence of the azimuthal asymmetry at fixed E_γ is not well represented by the two models. The QGS predicts a maximum of the asymmetry at $\theta_{c.m.} = 90^\circ$, which is observed at the higher photon-energy bins, however, its functional form does not coincide with what is observed from the data. On the other hand, the HRM predicts the behavior of the asymmetry at the low photon-energy bins (minimum at $\theta_{c.m.} = 90^\circ$), but

Figure 9.9. Polarization transfers, $C_{x'}$ and $C_{z'}$, for $E_\gamma \sim 2$ GeV as a function of proton angles in the center-of-mass-frame.

fails to do so at the higher energies. In the energy range $E_\gamma = 1.1 - 1.7$ GeV, the angular distributions of the azimuthal asymmetry show rich structure (two peaks) which qualitatively supports the QGS, which also exhibits rich structure (although three peaks). The

Figure 9.10. Azimuthal asymmetry as a function of the proton angle in the center-of-mass frame for the six photon-energy bins between $E_\gamma = 1.1 - 2.3$ GeV. The energy bin width is 200 MeV, and the angular bin width varies between 5° , for the low photon-energy bins, to 40° , for the high photon-energy bins. The error bars show the statistical (black) and the systematic (grey) uncertainties. The red dashed-dotted line shows the predictions of the QGSM [50], whereas the blue lines correspond to predictions of the HRM model [47]. The thick and thin solid curves correspond to the “on-shell” version of the model and represent predictions based on large and small helicity non-conservation, respectively. The dashed and dotted blue lines represent two versions of the “off-shell” approximation

cross-section data show a smooth decrease with angle (see Fig. 2.8), therefore this rich structure may be due to strong interference effects.

Figure 9.11. Azimuthal asymmetry as a function of the incident photon energy, for $\theta_{c.m.} = 90^\circ$. The blue lines correspond to the predictions of the QGSM (solid line) [50], and two versions of HRM (dotted lines) [47]. The red and cyan points correspond to results from Yerevan [69, 70], whereas the black points represent the results of this work. Statistical and systematic uncertainties of the CLAS data are shown with the black and red error bars respectively.

In general, the data rule out the HRM model calculations. It is interesting, however, that the HRM model reproduces the minimum at $E_\gamma = 1.2$ GeV and $\theta_{c.m.} = 90^\circ$, given that it is typically expected to be valid at higher energies. The fact that the HRM calculation does not reproduce the peak at ~ 2 GeV (see Fig. 9.11) may mean that modeling the process as a factorization of pn elastic scattering, the deuteron wave function and hard scattering is not applicable to the azimuthal asymmetry Σ . It might be the case that any structures in the energy dependence of the azimuthal asymmetry arise from structures in the underlying $pn \rightarrow pn$ amplitude. If this is the case, the factorization to include the $pn \rightarrow pn$ amplitude is not sufficient to explain the peak at 2 GeV, or the extrapolation of this amplitude misses some important physics.

The QGSM, just like the HRM, does not predict any structure in the energy dependence of Σ at $\theta_{c.m.} = 90^\circ$. In fact, it predicts asymmetries that decrease at a constant rate as the photon energy increases. Even though this model is expected to be valid for typically higher energies, it does not predict the transition from lower to higher values of the asymmetry observed at $\theta_{c.m.} = 90^\circ$ between $E_\gamma = 1.6$ GeV and 2.0 GeV.

Chapter 10

Summary and Outlook

10.1 Overview

The determination of the azimuthal asymmetry, Σ , for deuteron photodisintegration in the energy region $E_\gamma = 1.1 - 2.3$ GeV has been presented in the previous chapters. The importance of this reaction in achieving a comprehensive description of the underlying dynamics was argued. Specifically, the determination of the azimuthal asymmetry in this energy region, where the transition from the meson-nucleon to the quark-gluon picture degree of freedom happens, is essential in the effort to theoretically describe this transition region. The work done by the author to calibrate a specific detector system, select deuteron photodisintegration events, study the degree of photon polarization, and finally determine the azimuthal asymmetry and any systematic uncertainties associated with it, was comprehensively explained.

This work was conducted at Thomas Jefferson Accelerator Facility and used data collected during the g13 experiment. In this experiment, a linearly polarized photon beam was produced using the coherent bremsstrahlung technique (see Chap. 3). The experimental setup (electron-beam energies, beam collimation, *etc.*) produced photons with degrees of polarization that reached $\sim 80\%$. Deuteron photodisintegration events were selected by timing cuts on the detected positive particles and cuts on the square of the missing mass calculated using the information from the CLAS detector (see Chap 5). The azimuthal asymmetry was determined using the ϕ -bin method as presented in Chap. 7. In addition, a lot of attention was given to studying the systematic uncertainties involved in the determination of Σ . Finally, the results obtained from the work were presented along with a brief discussion and comparison with previous measurements on the differential cross section, single and double polarization observables (tabulated results are presented in Appendix. B).

This chapter's aim is to discuss the implications of the current results in terms of extension of this work, motivation for future experiments, and the impact of this result on

our understanding of nuclear structure and hadron interactions.

10.2 Discussion of the Current Results

The present results increase the kinematic coverage of previous measurements up to 2.3 GeV and to “all” angles ($\theta_{c.m.} = 25^\circ - 160^\circ$). This latter kinematic domain was completely unexplored in terms of polarization measurements. The determination of the azimuthal asymmetry in the energy region $E_\gamma = 1.1 - 2.3$ GeV compliments the previous measurements done on the differential cross section and recoil-proton polarizations, providing additional information for the underlying dynamics. A study of the functional form of the azimuthal asymmetry, both its angular dependence and its energy dependence, can give important insight into the mechanism of deuteron photodisintegration, which will lead to a better understanding of the transition region from the hadronic picture of nuclear interactions to the quark-gluon picture.

The results presented in Chap. 9 show that the azimuthal asymmetry is large (reaches ~ 0.9 for the photon-energy bin $E_\gamma = 1.7 - 1.9$ GeV) and positive for intermediate angles. Even though our results of the azimuthal asymmetry cannot be used to test predictions of the HHC rules, since they do not constrain the values of Σ , the results on other polarization observables contradict the predictions of the HHC and rules out the onset of pQCD in this energy region. As discussed in Chap. 9, the energy dependence of the asymmetry (especially at proton angles $\theta_{c.m.} = 90^\circ$) may indicate a change in the production mechanism between 1.5–2.1 GeV.

The available theoretical models failed to adequately describe the results presented in this thesis. Even though these models have free parameters that can be adjusted to fit the results obtained, it is unlikely that they will be able to describe the structures observed in the angular dependence of the azimuthal asymmetry. Both models predict a positive asymmetry for all angles in the energy range of interest. However, the results indicate that the asymmetry is negative for small angles over the entire photon-energy range, and change sign for large angles, being positive for small energies and negative for large. This can provide some insight on the relative phase between specific helicity amplitudes.

In summary, deuteron photodisintegration has been well studied experimentally

over the last two decades. Measurements done on the differential cross section did not provide the essential information that was needed to understand the underlying dynamics, and the determination of polarization observables proved to be necessary. The available theoretical models are successful in describing the angular and energy behavior of the differential cross section. In addition, the predictions of the models for the polarization observables, P_y , $C_{x'}$, and $C_{z'}$, are inadequate for examining their validity. On the other hand, the comparison between the theoretical predictions and the results obtained from this work clearly indicate that the available models fail in describing the underlying dynamics. The differences between the theoretical predictions and our results need to be explored in more detail. Theoretical work that would aim to describe these results is essential in our understanding of nuclear structure and the dominant mechanism of nuclear interactions in the transition region.

10.3 Motivation for Future Work

It is likely that the validity of both models occurs at higher energies. It was discussed in Chap. 9 that the QGSM predicts the maximum observed in the azimuthal asymmetry that occurs at $\theta_{c.m.} = 90^\circ$ and higher photon energies. However, at higher photon energies, the results obtained from this work suffer from large statistical and systematic uncertainties that might hinder information on the underlying dynamics. Reducing the statistical and systematic uncertainties of the azimuthal asymmetry can show hidden structures in the angular and energy dependence of Σ that can prove important in testing available and new theoretical models. Therefore, the available results for the azimuthal asymmetry can benefit from more precise measurements, especially at higher photon energies ($E_\gamma > 1.9$ GeV), and also at backward angles.

Data on the differential cross section are available for energies up to ~ 6 GeV. On the other hand, information on polarization observables only exists up to ~ 2 GeV. For this reason, extending measurements on polarization observables (single and double) to complement the cross-section measurements can be very important for future theoretical advancements. The determination of the azimuthal asymmetry is arguably easier to extend to higher energies, and this can be performed at JLab using the upgraded facility that would accelerate electrons up to 12 GeV. This new information on polarization observables can

become useful for the development of new theoretical models and improvements in existing models, which are essential in our understanding of nuclear reactions. The importance of these developments is evident from the available results on polarization observables and is crucial in our understanding of the mechanisms of deuteron photodisintegration.

10.4 Final Remarks

To conclude, a measurement of the azimuthal asymmetry of deuteron photodisintegration has been performed in the energy region $E_\gamma = 1.1 - 2.3$ GeV and presented in the previous chapters. The asymmetry was determined using the ϕ -bin method described in Sec. 7.3. The use of the tagger spectrometer and the CLAS enabled the determination of the azimuthal asymmetry, increasing the kinematic coverage to all angles, by determining the photon polarization on an event-by-event basis, reducing significantly the systematic uncertainties involved. The data presented here provide the only data of Σ available in this energy region and all angles, and these can be used to improve our understanding and range of the transition region. These results also motivate significant theoretical work that is necessary for the interpretation of the results and their implications on the mechanism responsible for reactions in this energy region.

Appendix A

Analytic Bremsstrahlung Calculation

This appendix presents an overview of the analytic bremsstrahlung calculation. The theoretical calculation that describes the spectrum of linearly polarized photons is adjusted in such a way to describe the experimental condition of the g13 experiment. The experimental characteristics, such as beam collimation and divergence, or even diamond characteristics, can affect the coherent photon-energy spectrum and have to be taken into consideration in determining the photon polarization. This appendix presents the adjustments done to the theoretical calculation to better describe the data collected during the g13 experiment. More details on this can be found in Ref. [129]

The intensities and polarizations in the coherent bremsstrahlung theory are described as functions of fractional energy, $x = E_\gamma/E_{beam}$. The intensity of a single, isolated, coherent peak is represented as a sum of the relative intensities of the coherent and incoherent contributions, I^c , I^i . The polarization, $P_g(x)$, is described as a function of the fractional energy, x , and the discontinuity energy, x_g^d (fractional energy of coherent edge), as follows:

$$P_g(x) = -\frac{\Phi_g(x)}{1 + \frac{I^i}{I^c}}, \quad \text{where} \quad (\text{A.1})$$

$$\Phi_g(x) = \frac{2Q_g^2 x^2}{1-x} \left[1 + (1-x)^2 - \frac{4Q_g^2 x^2}{1-x} \left(\frac{1-x}{Q_g x} - 1 \right) \right]^{-1}, \quad \text{and} \quad (\text{A.2})$$

$$Q_g = \frac{1 - x_g^d}{x_g^d} \quad (\text{A.3})$$

where $\Phi(x)$ is a function of x and x_g^d and represents the upper limit of the polarization. As mentioned in Sec. 3.2.1, the linearly polarized photons are produced at very small aperture angles. Thus, tight beam collimation reduces the incoherent contribution, and therefore, the ratio I^i/I^c , causing the polarization to increase towards its upper limit. In coherent bremsstrahlung theory, the enhancements have a very sharp falloff at the coherent edge, and for this reason the coherent edge is referred to as the discontinuity energy. In reality, though, the intensity distribution is smeared due to beam divergence, spread in beam momentum, crystal thickness, impurities in the crystal etc., resulting in a coherent edge

with a finite falloff range. This smeared intensity distribution can be characterized by a sum of many $\Phi(x)$ functions over a range of discontinuities ($x_g^d \pm \Delta x^d$). In order to predict the measured I^i/I^c ratio for a set of runs, this smearing has to be taken into account through modeling the beam collimation and angular spread of the generated photons.

The photon polarization is determined by fitting the experimental enhancement distributions with the theoretical prediction (fitting procedure is discussed in detail in Sec. 6.4). For a realistic fit function, one needs to incorporate the smearing due to beam collimation and divergence, coherent bremsstrahlung contributions from different reciprocal lattice vectors of the crystal, and the incoherent contribution from the crystal. The coherent contributions from each reciprocal lattice vector g , and the corresponding polarizations are analytically well defined [123]. The intensity from a reciprocal lattice vector g as a function of the fractional energy x , is given by,

$$I_g(x) = I_g^0 \chi_g(x), \quad \text{where} \quad (\text{A.4})$$

$$\chi_g(x) = \frac{x Q_g^2}{1-x} \left[1 + (1-x)^2 - \frac{4 Q_g^2 x^2}{1-x} \left(\frac{1-x}{Q_g x} - 1 \right) \right]. \quad (\text{A.5})$$

In Eq. (A.4) I_g^0 is the peak intensity for vector g , and Q_g in Eq. (A.5) is defined in Eq. (A.3). The corresponding polarization that results from the vector g is given by Eqs. (A.1)–(A.2). For each vector g , the discontinuity x_g^d depends on the angles between the beam and the crystal lattice and when restricted to the reciprocal lattice vectors $G = 022, 044, \dots 0gg$ it simplifies to give a one-dimensional angular dependence on θ according to

$$x_g^d = \frac{1}{\frac{k}{g E_{beam} \theta} + 1}, \quad (\text{A.6})$$

where $k = m_e a / 4\sqrt{2}\pi = 26.56$ MeV, and a is the diamond lattice constant equal to 923.7 (unitless). Equation (A.6) allows us to substitute this expression for x_g^d into Eq. (A.3) for Q_g , representing Φ and χ as functions of the fractional energy, x , the reciprocal lattice vector, g , and the angle between the electron beam and the crystal lattice, θ .

Due to the collimation, there is a lower limit on x_g , defined as follows,

$$x_g^c(g, \theta) = \frac{x_g^d}{1 + \theta_r^2 (1 - x_g^d)}, \quad (\text{A.7})$$

where θ_r is the relative angle between the beam and the collimator opening. The polar-

ization determination is limited by kinematics below x_g^d (where the cross section drops to zero) and by the geometrical collimation above x_g^c , (*i.e.* $x_g^c \leq x \leq x_g^d$).

To account for the beam collimation, we introduce a cumulative distribution function,

$$C(x, g, \theta) = \frac{1}{2} \left[1 + \operatorname{erf} \left(\frac{x - x_g^c}{\sigma_r \sqrt{2}} \right) \right], \quad (\text{A.8})$$

where σ_r accounts for the smearing of the two-dimensional beam spot across the collimator, and $\operatorname{erf}(x)$ corresponds to the error function defined as

$$\operatorname{erf}(x) = \frac{2}{\sqrt{\pi}} \int_0^x e^{-t^2} dt. \quad (\text{A.9})$$

This smearing affects the intensity of the coherent contribution from a single lattice vector as shown by Eq. (A.10),

$$I_g(x) = C(x, g, \theta) \times I_g^0 \chi(x, g, \theta). \quad (\text{A.10})$$

The intensity as a function of the fractional energy from reciprocal vectors 022, 044, ... 0GG results from the sum of contributions from individual lattice vectors,

$$\begin{aligned} I^c(x, G, \theta) &= \sum_{g=2,4,\dots}^G I_g(x, g, \theta) \\ &= \sum_{g=2,4,\dots}^G C(x, g, \theta) \times I_g^0 \chi(x, g, \theta). \end{aligned} \quad (\text{A.11})$$

Using Eq. (A.11), the total $\phi_{tot}(x)$ function can be written as the mean of the individual $\Phi(x, g, \theta)$ contributions weighted by their respective strengths,

$$\phi_{tot}(x, G, \theta) = \frac{\sum_{g=2,4,\dots}^G \Phi(x, g, \theta) \times I_g(x, g, \theta)}{\sum_{g=2,4,\dots}^G I_g(x, g, \theta)}, \quad (\text{A.12})$$

where $\Phi(x, g, \theta)$ is defined in Eq. (A.2) with $x_g^d(\theta)$ given in Eq. (A.6).

As mentioned before, the total intensity is the sum of the coherent and incoherent contributions from the crystal,

$$I_{tot}(x, \theta) = I^i(x) + I^c(x, \theta). \quad (\text{A.13})$$

The *enhancement*, ξ , is calculated in an attempt to remove any channel-to-channel variations in the tagger counter efficiencies, as well as the $1/x$ shape from the incoher-

ent bremsstrahlung photon-energy spectrum. This is done by dividing the experimental photon-energy spectrum obtained from a crystal radiator by the photon-energy spectrum obtained from an amorphous radiator,

$$\begin{aligned}\xi_{tot}(x, G, \theta) &= \frac{I_{tot}(x, \theta)}{kI_{amo}(x)} \\ &= \frac{I^i(x) + I^c(x, \theta)}{kI_{amo}(x)},\end{aligned}\tag{A.14}$$

where k is a constant used to normalize the baseline of the enhancement to a height of ~ 1 . Anything above the baseline corresponds to the enhancement of the coherent over the incoherent contributions to the spectrum. Figure A.1 shows the enhancement spectrum of data collected in the $E_\gamma=1.3\text{--}1.5$ GeV photon-energy bin. The different peaks correspond to coherent contributions that originate from electrons scattering off different reciprocal lattice vectors of the crystal.

Figure A.1. Enhancement spectrum for data collected in the $E_\gamma=1.3\text{--}1.5$ GeV photon-energy bin. The different peaks correspond to coherent contributions due to coherent bremsstrahlung off different reciprocal lattice vectors. For the analysis, data that result from the 022 crystal plane are used since this plane results in the highest statistics and highest polarization.

Assuming that the incoherent contribution, I^i , from the crystal has the same shape as the actual incoherent spectrum from an amorphous radiator, I_{amo} , over the photon-

energy range of interest (*i.e.* over the main coherent peak), the constant k is chosen to set the baseline to 1 at some well defined point where $I^c = 0$,

$$\begin{aligned}\xi_{tot}(x, G, \theta) &= 1 + \frac{I^c(x, G, \theta)}{kI_{amo}(x)} \\ &= 1 + \frac{\sum_{g=2,4,\dots}^G I_g^0 \times C(x, g, \theta) \chi(x, g, \theta)}{kI_{amo}(x)}.\end{aligned}\quad (\text{A.15})$$

Equation (A.1) can be rewritten as,

$$P_{tot}(x, G, \theta) = -\phi(x, G, \theta) \times \frac{I^c(x, G, \theta)}{I_{tot}(x, G, \theta)}.\quad (\text{A.16})$$

In reality, the coherent edge has a finite fall-off width due to beam divergence, crystal thickness, etc., as discussed before. This smearing is taken into account by assuming a Gaussian distribution in θ and summing over the different intensity distributions. Specifically, the smeared enhancement spectrum $\xi_{tot}^s(x, G, \theta, \sigma)$ can be represented as a normalized sum over many enhancement spectra, each weighted accordingly,

$$\xi_{tot}^s(x, G, \theta, \sigma) = \frac{\int_{\theta-3\sigma}^{\theta+3\sigma} e^{\frac{(\theta'-\theta)^2}{2\sigma^2}} \left[1 + \frac{\sum_{g=2,4,\dots}^G I_g^0 \times C(x, g, \theta') \chi(x, g, \theta')}{kI_{amo}(x)} \right] d\theta'}{\int_{\theta-3\sigma}^{\theta+3\sigma} e^{\frac{(\theta'-\theta)^2}{2\sigma^2}} d\theta'}.\quad (\text{A.17})$$

Using the same principle, the polarization for any given fractional energy, x , is a normalized, weighted sum of the polarization contributions from the allowed range of θ positions. The contribution from each θ position is the product of the weight for that θ point and the total intensity at the point. From this, the total polarization, taking into account the smearing in θ is,

$$P_{tot}^s(x, G, \theta, \sigma) = -\frac{\int_{\theta-3\sigma}^{\theta+3\sigma} I_{tot}(x, G, \theta') \times e^{\frac{(\theta'-\theta)^2}{2\sigma^2}} \times \phi_{tot}(x, G, \theta') \frac{I^c(x, G, \theta')}{I_{tot}(x, G, \theta')} d\theta'}{\int_{\theta-3\sigma}^{\theta+3\sigma} I_{tot}(x, G, \theta') \times e^{\frac{(\theta'-\theta)^2}{2\sigma^2}} d\theta'},\quad (\text{A.18})$$

which reduces to,

$$P_{tot}^s(x, G, \theta, \sigma) = -\frac{\int_{\theta-3\sigma}^{\theta+3\sigma} e^{\frac{(\theta'-\theta)^2}{2\sigma^2}} \times \phi_{tot}(x, G, \theta') \times I^c(x, G, \theta') d\theta'}{\int_{\theta-3\sigma}^{\theta+3\sigma} I_{tot}(x, G, \theta') \times e^{\frac{(\theta'-\theta)^2}{2\sigma^2}} d\theta'}.\quad (\text{A.19})$$

Equation (A.17) is used to fit the enhancement spectrum. The extracted fit param-

eters from this fit can be used in Eq. (A.19) to determine the corresponding polarization as a function of the fractional energy, x . The parameters that determine the form of the enhancement and the polarization are:

- θ – Angle between the beam and the crystal planes defined by the 022 direction
- σ – Gaussian smearing of θ to account for beam divergence, multiple scattering, beam-spot size etc.
- θ_r – Relative angle of collimation in units of characteristic angle
- σ_r – Smearing factor for collimation around θ_r via cumulative distribution function C_g
- $I_2^0, I_4^0, \dots, I_G^0$ – Amplitudes of the discrete coherent peaks

The enhancement and polarization functions can be easily translated to be functions of the measured photon energy by substituting $E_\gamma = xE_{beam}$. By doing so, it becomes straightforward to fit the calculated enhancement with the enhancement function and determine the fit parameters. The procedure to calculate the enhancement from the data and determine the photon polarization is outlined in Sec. 6.4.

Appendix B

Tabulated Results, and Figures

This appendix presents the data along with the associated uncertainties of the azimuthal asymmetry for all photon-energy and $\theta_{c.m.}$ bins. The systematic uncertainties quoted here are calculated by adding the different sources of systematic uncertainties quoted in Tables 8.5-8.6 in quadrature¹.

Table B.1. Azimuthal asymmetry for $E_\gamma = 1.1 - 1.3$ GeV

$\theta_{c.m.}$ (deg)	Σ	$\Delta\Sigma^{stat}$	$\Delta\Sigma^{syst}$
25 – 30	-0.140	0.028	0.044
30 – 35	-0.077	0.026	0.043
35 – 40	0.056	0.027	0.043
40 – 45	0.155	0.027	0.044
45 – 50	0.318	0.028	0.046
50 – 55	0.378	0.026	0.047
55 – 60	0.420	0.028	0.048
60 – 65	0.440	0.028	0.048
65 – 70	0.401	0.028	0.047
70 – 75	0.437	0.032	0.048
75 – 80	0.386	0.034	0.047
80 – 85	0.313	0.031	0.046
85 – 90	0.281	0.030	0.045
90 – 95	0.235	0.031	0.045
95 – 100	0.240	0.033	0.045
100 – 105	0.236	0.030	0.045
105 – 110	0.329	0.032	0.046
110 – 115	0.363	0.038	0.047
115 – 120	0.295	0.044	0.045
120 – 125	0.380	0.066	0.047
125 – 135	0.514	0.047	0.050
135 – 145	0.532	0.051	0.050
145 – 160	0.381	0.078	0.047

¹ $\Delta\Sigma^{syst} = \sum_i^N (|\overline{\Delta\Sigma}|^i)^2 + \sum_i^N \left(\sigma_{\overline{\Delta\Sigma}}^i\right)^2$, where i loops over the different sources of systematic uncertainties.

Table B.2. Azimuthal asymmetry for $E_\gamma = 1.3 - 1.5$ GeV

$\theta_{c.m.}$ (deg)	Σ	$\Delta\Sigma^{stat}$	$\Delta\Sigma^{syst}$
25 – 30	-0.187	0.030	0.035
30 – 35	-0.158	0.027	0.035
35 – 40	-0.164	0.029	0.035
40 – 45	0.003	0.030	0.034
45 – 50	0.093	0.033	0.034
50 – 55	0.170	0.032	0.035
55 – 60	0.301	0.032	0.037
60 – 65	0.370	0.033	0.039
65 – 70	0.383	0.034	0.039
70 – 75	0.464	0.035	0.041
75 – 80	0.474	0.044	0.041
80 – 85	0.373	0.038	0.039
85 – 95	0.233	0.027	0.036
95 – 105	0.183	0.029	0.035
105 – 115	0.124	0.034	0.035
115 – 125	0.240	0.047	0.036
125 – 135	0.371	0.064	0.039
135 – 145	0.247	0.055	0.036
145 – 160	0.076	0.058	0.034

Table B.3. Azimuthal asymmetry for $E_\gamma = 1.5 - 1.7$ GeV

$\theta_{c.m.}$ (deg)	Σ	$\Delta\Sigma^{stat}$	$\Delta\Sigma^{syst}$
25 – 30	-0.200	0.037	0.091
30 – 35	-0.219	0.034	0.091
35 – 42.5	-0.234	0.030	0.091
42.5 – 50	-0.035	0.033	0.090
50 – 57.5	0.108	0.033	0.090
57.5 – 65	0.291	0.033	0.091
65 – 70	0.359	0.041	0.092
70 – 75	0.467	0.039	0.093
75 – 80	0.625	0.048	0.095
80 – 85	0.561	0.047	0.094
85 – 95	0.345	0.035	0.092
95 – 105	0.254	0.040	0.091
105 – 120	0.183	0.042	0.091
120 – 132.5	0.254	0.068	0.091
132.5 – 145	0.278	0.057	0.091
145 – 160	0.173	0.065	0.090

Table B.4. Azimuthal asymmetry for $E_\gamma = 1.7 - 1.9$ GeV

$\theta_{c.m.}$ (deg)	Σ	$\Delta\Sigma^{stat}$	$\Delta\Sigma^{syst}$
25 – 35	-0.202	0.034	0.060
35 – 45	-0.195	0.036	0.060
45 – 55	-0.031	0.039	0.059
55 – 65	0.159	0.039	0.059
65 – 75	0.455	0.039	0.063
75 – 85	0.670	0.043	0.068
85 – 95	0.713	0.042	0.069
95 – 110	0.608	0.046	0.066
110 – 125	0.292	0.066	0.061
125 – 140	-0.036	0.088	0.059
140 – 160	-0.155	0.072	0.059

Table B.5. Azimuthal asymmetry for $E_\gamma = 1.9 - 2.1$ GeV

$\theta_{c.m.}$ (deg)	Σ	$\Delta\Sigma^{stat}$	$\Delta\Sigma^{syst}$
25 – 37.5	-0.147	0.057	0.121
37.5 – 50	-0.050	0.063	0.121
50 – 62.5	0.172	0.065	0.121
62.5 – 75	0.506	0.063	0.124
75 – 100	0.880	0.048	0.127
100 – 125	0.892	0.065	0.129
125 – 160	-0.222	0.091	0.122

Table B.6. Azimuthal asymmetry for $E_\gamma = 2.1 - 2.3$ GeV

$\theta_{c.m.}$ (deg)	Σ	$\Delta\Sigma^{stat}$	$\Delta\Sigma^{syst}$
25 – 45	-0.164	0.062	0.201
45 – 65	-0.128	0.069	0.201
65 – 95	0.466	0.068	0.202
95 – 125	0.577	0.099	0.203
125 – 160	0.077	0.112	0.201

Table B.7. Azimuthal asymmetry for $\theta_{c.m.} = 35^\circ$

E_γ (GeV)	$\theta_{c.m.}$ (deg)	Σ	$\Delta\Sigma^{stat}$	$\Delta\Sigma^{syst}$
1.1 – 1.3	32.5 – 37.5	-0.013	0.026	0.043
1.3 – 1.5	30 – 40	-0.160	0.020	0.035
1.5 – 1.7	30 – 40	-0.247	0.025	0.091
1.7 – 1.9	30 – 40	-0.236	0.033	0.060
1.9 – 2.1	25 – 45	-0.066	0.047	0.121
2.1 – 2.3	25 – 45	-0.152	0.058	0.201

Table B.8. Azimuthal asymmetry for $\theta_{c.m.} = 60^\circ$

E_γ (GeV)	$\theta_{c.m.}$ (deg)	Σ	$\Delta\Sigma^{stat}$	$\Delta\Sigma^{syst}$
1.1 – 1.3	57.5 – 62.5	0.435	0.028	0.048
1.3 – 1.5	55 – 65	0.330	0.023	0.038
1.5 – 1.7	55 – 65	0.271	0.029	0.091
1.7 – 1.9	55 – 65	0.159	0.039	0.059
1.9 – 2.1	50 – 70	0.248	0.052	0.121
2.1 – 2.3	50 – 70	-0.147	0.067	0.202

Table B.9. Azimuthal asymmetry for $\theta_{c.m.} = 90^\circ$

E_γ (GeV)	$\theta_{c.m.}$ (deg)	Σ	$\Delta\Sigma^{stat}$	$\Delta\Sigma^{syst}$
1.1 – 1.3	87.5 – 92.5	0.243	0.030	0.044
1.3 – 1.5	85 – 95	0.233	0.027	0.036
1.5 – 1.7	85 – 95	0.345	0.035	0.092
1.7 – 1.9	85 – 95	0.713	0.042	0.068
1.9 – 2.1	77.5 – 102.5	0.95	0.046	0.130
2.1 – 2.3	75 – 105	0.794	0.063	0.205

Table B.10. Azimuthal asymmetry for $\theta_{c.m.} = 120^\circ$

E_γ (GeV)	$\theta_{c.m.}$ (deg)	Σ	$\Delta\Sigma^{stat}$	$\Delta\Sigma^{syst}$
1.1 – 1.3	117.5 – 122.5	0.316	0.054	0.045
1.3 – 1.5	115 – 125	0.240	0.047	0.036
1.5 – 1.7	115 – 125	0.180	0.061	0.090
1.7 – 1.9	115 – 125	0.106	0.085	0.059
1.9 – 2.1	110 – 130	0.739	0.086	0.126
2.1 – 2.3	100 – 135	0.606	0.102	0.204

Table B.11. Azimuthal asymmetry for $\theta_{c.m.} = 145^\circ$

E_γ (GeV)	$\theta_{c.m.}$ (deg)	Σ	$\Delta\Sigma^{stat}$	$\Delta\Sigma^{syst}$
1.1 – 1.3	140 – 150	0.460	0.060	0.048
1.3 – 1.5	137.5 – 152.5	0.167	0.045	0.035
1.5 – 1.7	137.5 – 152.5	0.150	0.055	0.090
1.7 – 1.9	130 – 160	-0.138	0.059	0.059
1.9 – 2.1	130 – 160	-0.240	0.094	0.121
2.1 – 2.3	130 – 160	-0.089	0.118	0.201

Figure B.1. Azimuthal asymmetry as a function of the proton angle in the center-of-mass frame, for the photon-energy bin $E_\gamma = 1.1 - 1.3$ GeV. Dashed-dotted line shows the predictions of the QGSM [50].

Figure B.2. Azimuthal asymmetry as a function of the proton angle in the center-of-mass frame, for the photon-energy bin $E_\gamma = 1.3 - 1.5$ GeV. Dashed-dotted line shows the predictions of the QGSM [50].

Figure B.3. Azimuthal asymmetry as a function of the proton angle in the center-of-mass frame, for the photon-energy bin $E_\gamma = 1.5 - 1.7$ GeV. Dashed-dotted line shows the predictions of the QGSM [50].

Figure B.4. Azimuthal asymmetry as a function of the proton angle in the center-of-mass frame, for the photon-energy bin $E_\gamma = 1.7 - 1.9$ GeV. Dashed-dotted line shows the predictions of the QGSM [50].

Figure B.5. Azimuthal asymmetry as a function of the proton angle in the center-of-mass frame, for the photon-energy bin $E_\gamma = 1.9 - 2.1$ GeV. Dashed-dotted line shows the predictions of the QGSM [50].

Figure B.6. Azimuthal asymmetry as a function of the proton angle in the center-of-mass frame, for the photon-energy bin $E_\gamma = 2.1 - 2.3$ GeV. Dashed-dotted line shows the predictions of the QGSM [50].

Figure B.7. Azimuthal asymmetry as a function of the proton angle in the center-of-mass frame, for the photon-energy bin $E_\gamma = 1.1 - 1.3$ GeV. The thick and thin solid curves correspond to the “on-shell” version of the model and represent predictions based on large and small helicity non-conservation, respectively. The dashed and dotted lines represent two versions of the “off-shell” approximation [47].

Figure B.8. Azimuthal asymmetry as a function of the proton angle in the center-of-mass frame, for the photon-energy bin $E_\gamma = 1.3 - 1.5$ GeV. The thick and thin solid curves correspond to the “on-shell” version of the model and represent predictions based on large and small helicity non-conservation, respectively. The dashed and dotted lines represent two versions of the “off-shell” approximation [47].

Figure B.9. Azimuthal asymmetry as a function of the proton angle in the center-of-mass frame, for the photon-energy bin $E_\gamma = 1.5 - 1.7$ GeV. The thick and thin solid curves correspond to the “on-shell” version of the model and represent predictions based on large and small helicity non-conservation, respectively. The dashed and dotted lines represent two versions of the “off-shell” approximation [47].

Figure B.10. Azimuthal asymmetry as a function of the proton angle in the center-of-mass frame, for the photon-energy bin $E_\gamma = 1.7 - 1.9$ GeV. The thick and thin solid curves correspond to the “on-shell” version of the model and represent predictions based on large and small helicity non-conservation, respectively. The dashed and dotted lines represent two versions of the “off-shell” approximation [47].

Figure B.11. Azimuthal asymmetry as a function of the proton angle in the center-of-mass frame, for the photon-energy bin $E_\gamma = 1.9 - 2.1$ GeV. The thick and thin solid curves correspond to the “on-shell” version of the model and represent predictions based on large and small helicity non-conservation, respectively. The dashed and dotted lines represent two versions of the “off-shell” approximation [47].

Figure B.12. Azimuthal asymmetry as a function of the proton angle in the center-of-mass frame, for the photon-energy bin $E_\gamma = 2.1 - 2.3 \text{ GeV}$. The thick and thin solid curves correspond to the “on-shell” version of the model and represent predictions based on large and small helicity non-conservation, respectively. The dashed and dotted lines represent two versions of the “off-shell” approximation [47].

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