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Q2 dependence of the S11(1535) photocoupling and evidence for a P-wave resonance in η electroproduction

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$Q^2$ dependence of the $S_{11}(1535)$ photocoupling and evidence for a $P$-wave resonance in $\eta$ electroproduction


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New cross sections for the reaction $ep \rightarrow e'\eta p$ are reported for total center-of-mass energy $W = 1.5$–2.3 GeV and invariant squared momentum transfer $Q^2 = 0.13$–3.3 GeV$^2$. This large kinematic range allows the extraction of new information about response functions, photocouplings, and $\eta N$ coupling strengths of baryon resonances. A sharp structure is seen at $W \sim 1.7$ GeV. The shape of the differential cross section is indicative of the presence of a $P$-wave resonance that persists to high $Q^2$. Improved values are derived for the photocoupling amplitude for the $S_{11}(1535)$ resonance. The new data greatly expand the $Q^2$ range covered, and an interpretation of all data with a consistent parametrization is provided.

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I. INTRODUCTION

Photoproduction and electroproduction experiments on the nucleon provide a clean probe of nucleon structure because quantum electrodynamics is well understood. As a result, the matrix elements for $\gamma N \rightarrow N^*, \Delta^*$ transitions, commonly called the photocoupling amplitudes, are sensitive to the nucleon and $N^*$ quark-level wave function. These amplitudes have traditionally been calculated using quark models [1–3], but recently progress has been made in applying the techniques of lattice QCD [4–6]. Experimental measurements are currently being made of a number of different baryon resonances in several different final states. For a review of the current status, see Refs. [7,8].

Disentangling the wide and overlapping states that populate reaction data has been a long-lasting problem. In the mass region above the total center-of-mass (c.m.) energy $W$ of 1.5 GeV, many overlapping baryon states are present, and some are not well known. The reaction $ep \rightarrow e'\eta p$ is especially clean, since processes involving $\eta N$ final states couple only to isospin $\frac{1}{2}$ resonances, simplifying the analysis. A prominent peak in the total cross section is seen for $\eta$ production at $W = 1.535$ GeV in both $\gamma N$ and $\pi N$ experiments. This is widely interpreted as the excitation of a single resonance, the spin $\frac{1}{2}$, negative parity, isospin $\frac{1}{2}$ $S_{11}(1535)$ [9]. This state has a branching ratio to $\eta N$ of 45–60% compared to at most a few percent [9,10] for other states. This is a very interesting and unusual pattern.

$\eta$ photoproduction experiments have reaffirmed the strong energy dependence and $S$-wave (isotropic) character close to threshold [11]. Using polarized photons [12], new values for $\eta N$ decay branching ratios of other resonances have been determined through interference with the dominant $S_{11}(1535)$, $N^*$, and unusual pattern.
Electroproduction cross sections can be used to extract the photocoupling amplitude for nonzero values of the squared momentum transfer $Q^2$ from the electron to the resonance. Using $\eta$ electroproduction [13–19], an unusually flat $Q^2$ dependence of the photocoupling amplitude was found for the $S_{11}(1535)$ in contrast to the nucleon form factors and photon coupling amplitudes of other established resonances, e.g., $P_{13}(1232)$ [20]. Although previous $\eta$ angular distributions were largely isotropic at all $Q^2$, no detailed response functions were extracted because of the poor angular coverage using traditional magnetic spectrometers. Although the $Q^2$ dependence was clearly different than for other resonances, the results were comprised of many different experiments whose results appeared to be inconsistent with each other. An analysis by Armstrong et al. [19] showed that much of the inconsistency was due to different assumptions about $S_{11}(1535)$ properties used by the individual experiments.

In our previous publication [21], we presented results on $\eta$ electroproduction based on the first data taken with the Continuous Electron Beam Accelerator Facility (CEBAF) Large Acceptance Spectrometer (CLAS) [22] at Jefferson Lab. We extracted the photocoupling amplitude $A_{1\perp}$ for the $\delta_{11}(1535)$ over the range $0.25 < Q^2 < 1.5$ GeV$^2$ from our data. In addition, we observed the indication of a structure at $W \approx 1.7$ GeV in the total cross section which is also seen as a change in the shape of the differential cross section at the same energy.

The energy region around $W \approx 1.7$ GeV has received significant attention lately. At the same energy, a CLAS $\pi^+\pi^-$ electroproduction experiment [23] found excess strength beyond theoretical predictions based on previous data. This excess strength was tentatively identified as a $P$-wave resonance; either the decay properties of $P_{13}(1720)$ change significantly or there is a new spin $3/2$ state. A recent $\eta$ photoproduction experiment at Bonn [24] provides a comprehensive set of cross section data from near threshold to well beyond the resonance region. In the same paper [24], a partial wave analysis of these and other $\eta$ electroproduction data finds strong excitation of a $J = 3/2^+$ state at 1775 MeV which they identify with the $P_{13}(1720)$. Although previous analyses [9] found weak evidence for any $\eta N$ decay of resonances at $W \approx 1.7$ GeV, the new data is of much higher quality than the older data.

The data and analysis reported here use a data set taken with the same apparatus (CLAS) as used in our first publication [21]. The new data have an order of magnitude more $\eta$ events than in that previous paper and a much larger kinematic range. Therefore, the new values presented here supersede the previously published data. Our reach in $Q^2$ (0.13–3.3 GeV$^2$) is more than twice as large as in our first publication. This allows a large extension of the $Q^2$ range where we can extract the photocoupling amplitude $A_{1\perp}$ of the proton to $S_{11}(1535)$ transition. We also more precisely determine the nonisotropies in the differential cross section and show evidence for a significant contribution to $\eta$ electroproduction due to a $P$-wave resonance with a mass around 1.7 GeV.

The paper first presents some formalism needed to understand the measurement and its analysis, followed by details of the experiment. We then present results for the inclusive and exclusive analyses. Discussion of these results with a Breit-Wigner model and conclusions complete the paper.

II. FORMALISM

The kinematics for the $ep \rightarrow e'\eta p$ reaction are shown in Fig. 1. It can be characterized in terms of the squared four-momentum transfer between the electron and proton ($-Q^2$) carried by the virtual photon ($\gamma^*$), the invariant mass of the $\gamma e^- p$ system ($W$), and the scattering angles of the final state $\eta$ in the rest frame of the $\gamma e^- p$ system ($\theta^*, \phi^*$). These angles are also the decay angles of the resonance in its rest frame. We use the superscript $*$ for quantities evaluated in this frame. The fivefold unpolarized differential cross section for the $ep \rightarrow e'\eta p$ process at a specific energy $E$ may be expressed as the product of the transverse virtual photon flux $\Gamma_\gamma$ in the Hand convention [25] and the c.m. cross section for virtual photoproduction of the $p\eta$ pair:

$$\frac{d^5\sigma}{dWdQ^2d\Omega_\eta} = \Gamma_\gamma(E, W, Q^2) \frac{d^2\sigma}{d\Omega_\eta}(\gamma e^- p \rightarrow \eta p).$$

The cross section for the virtual reaction $\gamma e^- p \rightarrow \eta p$ is written by convention to explicitly display the dependence on $\phi^*$, that is,

$$\frac{d^2\sigma}{d\Omega_\eta}(\gamma e^- p \rightarrow \eta p) = \sigma_T + \epsilon\sigma_L + \sqrt{2}\epsilon(1+\epsilon)\sigma_{LT}\cos\phi^*_\eta + \epsilon\sigma_{TT}\cos2\phi^*_\eta,$$

where $\epsilon$ is the longitudinal degree of polarization of the virtual photon and is given by

$$\epsilon = \left[1 + 2q^2(Q^2)\tan^2\left(\frac{\Psi}{2}\right)\right]^{-1},$$

where $q$ is the magnitude of the three-momentum of the virtual photon and $\Psi$ is the electron scattering angle. Since $\epsilon$ is invariant under collinear transformations, $q$ and $\Psi$ may be expressed either in the laboratory or c.m. frame. The component cross sections can be expressed in terms that are
related to the spin of the photon and target, very similar to the
definition of response functions [8]. In an unpolarized experi-
ment, there are four independent terms. The cross sections for
transverse and longitudinal photons are represented by \( \sigma_T \)
and \( \sigma_L \), respectively. In addition, \( \sigma_{LT} \) is a contribu-
tion due to the interference between transverse and longitudinal
amplitudes, and \( \sigma_{TT} \) describes the interference between amplitudes for
the two different transverse polarizations, either aligned or
anti-aligned with the spin of the target photon. All four of
these terms depend on \( W, Q^2 \), and \( \cos \theta^* \).

To identify individual baryon resonances, the cross section
should be decomposed into partial wave amplitudes. These
amplitudes are most often labeled by the electromagnetic
multipole notation [26]. Multipoles are commonly labeled
\( E_{l\pm}, M_{l\pm}, \) and \( S_{l\pm} \), where \( l \) is the orbital angular momentum of
the final \( \eta p \) system and \( \pm \) denotes whether the total
angular momentum is \( l \pm \frac{1}{2} \). \( E \) and \( M \) refer to electric
and magnetic transitions involving transverse virtual photons,
while the longitudinal \( S \) transitions involve longitudinal
photons.

The response functions and multipoles have contributions from
underlying resonant and nonresonant reaction mechanisms. When evaluated at the peak of the resonance, the
multipole is expressed in terms of both the photocoupling
amplitude and the hadronic decay properties of the individual
resonances in a commonly accepted way [27]. The photocou-
pling amplitudes are labeled by the \( \gamma N \) total helicity (\( \frac{1}{2} \) or \( \frac{3}{2} \)) and
the virtual photon polarization (transverse or longitudinal)
and depend on the invariant squared momentum transfer to the
resonance (\( Q^2 \)). The shape of the resonance determines the
\( W \) dependence of the resonant part of the multipole. Spin-\( \frac{1}{2} \)
resonances will be described by one transverse amplitude
\( A_\frac{1}{2} \) and one longitudinal amplitude \( S_\frac{1}{2} \). In terms of multipo-
les, an \( S_{11} \) resonance has an \( E_{0+} \) (electric dipole) and an
\( S_{0+} \) transition; a \( P_{11} \) has \( M_{1-} \) and \( S_{1-} \) transitions. Higher spin
resonances will be described by both \( A_\frac{1}{2} \) and \( A_\frac{3}{2} \) photocoupl-
ings (and both \( E \) and \( M \) multipoles). Extraction of the multi-
ple amplitudes from the cross section data, see, e.g., [28],
is not unique, because more than one bilinear combination of
multipoles have identical angular distributions. We therefore
choose simplified methods (discussed below) to analyze the
data.

The differential cross sections can be calculated with a
model of resonance production/decay and the nonresonant
processes. This cannot yet be done from a fundamental field
theory such as quantum chromodynamics (QCD). Instead,
models are used that have parameters determined from data.
The \( \eta \)-MAID [29] model uses an isobar model [30] to construct the
cross section for \( \eta \) photo- and electroproduction; parameters
are fit by comparisons with previous results [11,21,31]. We
have calculations from the MAID code for our kinematics. To
further understand our data, we also do Legendre polynomial
fits to the angular distributions. Both these results are described
in Sec. IV B.

To analyze our angle-integrated cross sections, we make a
further simplification which is possible because the \( S_{11}(1535) \)
resonance is dominant near threshold. Therefore, we ignore the
nonresonant amplitude. If one can isolate the contribution of a
single \( S_{11} \) resonance to the \( E_{0+} \) multipole, the cross section
takes the simple form

\[
\frac{d\sigma}{d\Omega_\eta} = \frac{p_\eta^* W}{m_p K} |E_{0+}(W)|^2,
\]

where \( K = (W^2 - m_p^2)/(2m_p) \) is the equivalent real photon
energy, \( p_\eta^* \) is the momentum of the outgoing \( \eta \) in the \( S_{11} \) rest
frame, and \( m_p \) is the proton mass. The longitudinal multipole
\( S_{0+} \) does, in principal, contribute, and we do not have the
data to make the separation. However, \( S_{0+} \) has been found to
be small [17], and it was therefore ignored in previous
analyses [11,13,14,16–19,21]. An isobar model analysis of
\( \eta p, \pi^0 p, \) and \( \pi^+ n \) CLAS electroproduction data [32] confirms
the assumption of a small longitudinal component. In this
analysis, the value of \( S_{\frac{1}{2}} / A_{\frac{1}{2}} \) is about 15–20\%; this translates to
a few percent contribution to the cross sections measured
with. With the assumption that a single resonance dominates
the cross section and \( S_{\frac{1}{2}} \) is small, \( A_{\frac{1}{2}} \) for \( \gamma p \rightarrow S_{11}(1535) \) can be
determined from Ref. [27] as

\[
A_{\frac{1}{2}} = \sqrt{\frac{2\pi}{K m_p^2 p_\eta^*}} \Im(E_{0+}(W_R))
\]

where \( E_{0+} \) refers only to the contribution from the resonance
which is evaluated at the peak of the resonance. If there are
other contributions to \( E_{0+} \), a model is needed to extract the
resonance contribution. This formula contains terms related to
the final state decay of the \( S_{11} \): \( \Gamma_{\eta} \), the total width of the
\( S_{11}(1535) \); \( b_{\eta} \), the branching fraction into the \( \eta p \) final state;
and \( p_\eta^* \) are all calculated at the mass of the \( S_{11} \) (\( W_R \)). Our
current lack of knowledge of these parameters leads to a model
dependence in the extracted values for \( A_{\frac{1}{2}} \). This prompted
Benmerrouche et al. [33] to propose using a quantity for each
resonance with less model dependence, that is,

\[
\xi_{\frac{1}{2}} = \frac{m_p^2 K b_{\eta} \Gamma_{\eta}}{W_R^2 p_\eta^* \Gamma_{\eta}} A_{\frac{1}{2}}
\]

\( \xi_{\frac{1}{2}} \) depends on the matrix element for the initial state
\( \gamma N \rightarrow N^* \) transition, while \( \xi_{\frac{1}{2}} \) is proportional to the prod-
uct of the matrix elements for the \( \gamma N \rightarrow N^* \) and \( N^* \rightarrow \eta p 
\) transition. For \( S_{11}(1535) \), \( \xi_{\frac{1}{2}} = \sqrt{2\pi \Im(E_{0+}(W_R))} \). Al-
though \( \xi_{\frac{1}{2}} \) is more closely related to experimental values, \( A_{\frac{1}{2}} \) is
more easily determined from calculations, e.g., using quark
models. Whichever quantity is used, the model dependence
still exists when comparing calculations to experiment.

We use the same resonance parametrization in all our
calculations. The relativistic Breit-Wigner form is taken from
previous \( \eta \) photoproduction work [34,35] and extended to
nonzero angular momentum as

\[
\frac{d\sigma_{BW}}{d\Omega_\eta}(W) = \frac{p_\eta^*}{4q^2} |T_{BW}^\ell(W)|^2,
\]

\[
T_{BW}^\ell(W) = \frac{a W_R \Gamma_{\eta}}{(W^2 - W_R^2 - i W_R \Gamma_{\eta_{tot}}},
\]

where \( a \) is a constant that contains the photocoupling am-
plitude and kinematic factors, \( q^* \) is the photon three-momentum
in the resonance rest frame, \( \Gamma_{\eta} \) is the partial width for...
TABLE I. Summary of kinematic ranges of previously published data compared with this experiment. This experiment is an extension of Ref. [21]. Values given are the maximum ranges for each experiment.

<table>
<thead>
<tr>
<th>Experiments</th>
<th>W (GeV)</th>
<th>( Q^2 ) (GeV(^2))</th>
<th>( \cos \theta^* )</th>
<th>( \phi^* ) (deg)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Daresbury [13]</td>
<td>1.51 (\rightarrow) 1.55</td>
<td>0.15 (\rightarrow) 1.5</td>
<td>not given</td>
<td>not given</td>
</tr>
<tr>
<td>Bonn [14]</td>
<td>1.51 (\rightarrow) 1.56</td>
<td>0.2 (\rightarrow) 0.4</td>
<td>(-0.766 \rightarrow 0.939)</td>
<td>(-0)</td>
</tr>
<tr>
<td>DESY [15]</td>
<td>1.5 (\rightarrow) 1.7</td>
<td>0.22 (\rightarrow) 1.0</td>
<td>(-1 \rightarrow 1)</td>
<td>0 (\rightarrow) 180</td>
</tr>
<tr>
<td>DESY [16]</td>
<td>1.49 (\rightarrow) 1.58</td>
<td>0.6, 1.0</td>
<td>(-1 \rightarrow 1)</td>
<td>15 (\rightarrow) 90</td>
</tr>
<tr>
<td>Bonn [17]</td>
<td>1.44 (\rightarrow) 1.64</td>
<td>0.4</td>
<td>(-0.643 \rightarrow 0.866)</td>
<td>(-40 \rightarrow 40)</td>
</tr>
<tr>
<td>DESY [18]</td>
<td>1.49 (\rightarrow) 1.8</td>
<td>2.0, 3.0</td>
<td>(-1 \rightarrow 1)</td>
<td>0 (\rightarrow) 120</td>
</tr>
<tr>
<td>JLab [19]</td>
<td>1.48 (\rightarrow) 1.62</td>
<td>2.4, 3.6</td>
<td>(-1 \rightarrow 1)</td>
<td>0 (\rightarrow) 360</td>
</tr>
<tr>
<td>JLab [21]</td>
<td>1.5 (\rightarrow) 1.86</td>
<td>0.375 (\rightarrow) 1.5</td>
<td>(-1 \rightarrow 1)</td>
<td>0 (\rightarrow) 360</td>
</tr>
<tr>
<td>This experiment</td>
<td>1.5 (\rightarrow) 2.3</td>
<td>0.13 (\rightarrow) 3.3</td>
<td>(-1 \rightarrow 1)</td>
<td>0 (\rightarrow) 360</td>
</tr>
</tbody>
</table>

\[
N^* \rightarrow \eta p \text{ decay, and } \Gamma_{\text{tot}} \text{ is the total width,}
\]

\[
\Gamma_{\eta} = B_{l}(p_{\eta}^*) \Gamma_{R},
\]

\[
\Gamma_{\text{tot}} = \frac{0.5 p_{\eta}^*}{p_{\eta}^* + B_{l}(p_{\eta}^*)} + \frac{0.4 p_{\eta}^*}{p_{\eta}^* + B_{l}(p_{\eta}^*)} \Gamma_{R},
\]

where \( \Gamma_{R} \) is the bare width and \( B_{l}(p^*) \) is a Blatt-Weisskopf penetration factor [36]. If \( \ell = 0 \), this factor is equal to unity. The momentum ratios [35] approximately account for proper phase space effects for the various final states (\( \pi N \), \( \eta N \), and \( \pi\pi N \), where the phase space factors for the \( \pi\pi N \) final state are ignored). They are weighted according to estimates of the branching fractions to each final state. This form has been successful in matching data but is not unique.

### III. DETECTOR AND ANALYSIS

The CLAS facility [22] was designed for the efficient detection of multiparticle final states. The data used for this measurement were taken in 1999 at electron beam energies of 1.5, 2.5, and 4.0 GeV. A cylinder of liquid hydrogen was used as the target. Two different targets were used, 5.0 and 3.8 cm long. Toroidal magnet coils separate CLAS into six largely identical sectors, each covering roughly 54° in azimuthal angle \( \phi \) (with smaller coverage at smaller polar angle). Tracking drift chambers (DCs) in CLAS measure angles and momenta of charged particles for laboratory polar angles in the range \( 8° < \theta < 142° \). Outside the DCs, scintillation counters (SCs) provide time-of-flight measurements with which we can separate the charged hadrons into pions, kaons, and protons. For laboratory angles \( \theta < 48° \), threshold Cherenkov counters (CCs) and electromagnetic calorimeters (ECs) distinguish electrons from charged hadrons.

For this analysis, events were selected with an identified electron and proton. Since the momentum four-vectors of the beam and target are known, the four-vector for the putative \( \eta \) can be determined from these two final state particles. A fiducial cut on these particles was applied to avoid the regions near the magnetic coils and the edges of the CC where the acceptance is changing rapidly. The momentum of the electron was required to be above 400 MeV in order to be well above the trigger threshold.

Cross sections were calculated as a function of \( Q^2 \) and \( W \) for the angle-integrated data analysis and as a function of \( Q^2 \), \( W \), \( \cos \theta^* \), and \( \phi^* \) for the differential data analysis. Cross sections are determined in a standard way by determining the yield in each of many bins, correcting for detector acceptance, and normalizing by the beam intensity measured with a Faraday cup and the calculated target thickness.

As discussed in Sec. II, the extraction of resonance properties comes from an analysis of the \( \cos \theta^* \), \( \phi^* \) distributions at specific values of \( W \) and \( Q^2 \). Distributions of these variables covered by the apparatus are determined by geometry. The large acceptance of CLAS guarantees almost complete coverage in \( \cos \theta^* \) and \( \phi^* \). The beam energies of the experiment coupled with CLAS provided data at a wide range of \( Q^2 \) and \( W \). We compare the kinematic range for the new experiment with that available for the previously published \( \eta \) electroproduction data in Table I.

Events were divided into separate kinematic bins as detailed in Sec. IV. For each bin, the \( \eta \) yield was determined by fitting the distribution of missing mass recoiling against the outgoing \( e^-p \) system. An example fit in one bin is shown in Fig. 2.

The fit is the sum of a signal at the \( \eta \) mass and a background function. The signal shape has a radiative tail and is corrected for experimental resolution; the background function is a polynomial. We use the data to determine both shapes. Both functions must then be modified by the geometric acceptance for this reaction, because it has a rapid variation with respect to the kinematic parameters. This method is an extension of what was used in the previous CLAS data analysis [21].

The shape of the signal was modeled in two steps to reproduce all features seen in the data. It is first described by a \( \delta \) function at the \( \eta \) mass (\( m_\eta \)) plus an exponential above \( m_\eta \) representing the radiative tail,

\[
S(m) = (1 - f)\delta(m - m_\eta) + f\Theta(m - m_\eta)e^{-m/\alpha(m - m_\eta)}.
\]

The fraction of events in the radiative tail (\( f \)) and a parameter describing the slope in the exponential (\( \alpha \)) were determined for each \( W-Q^2 \) bin from Monte Carlo generator events containing radiative effects [37,38]. This signal shape was then convoluted with a Gaussian representing the experimental resolution to obtain the final signal shape (an analytic function) used to fit
mass spectrum for each bin. In the figure, the dot-dashed curve is the function modified by the acceptance function is then fit to the missing Carlo program. The sum of the line (right scale) shows the acceptance that is calculated with a Monte Carlo program. The acceptance has significant variation across the fit region. After requiring the scattered electron and background acceptance function. We multiply our simple background function with the calculated acceptance function to obtain the final background used in the fit. Examples of all three of these curves are shown in Fig. 2.

In a small number of bins where the cross section is low, statistical fluctuations in the background can lead to a best-fit value for the number of $\eta$’s that is negative. In this case, we follow the suggestion of the Particle Data Group (PDG) [9] and report a negative value with error bars for the cross section. This provides sufficient information for constraints from these bins to be combined with nearby bins in comparing with theoretical predictions.

Acceptance for the $ep \rightarrow e\eta p$ reaction was calculated using a GEANT-based Monte Carlo simulation [39]. The event generator included radiative effects using the peaking approximation [37,38], and the cross sections have been corrected for radiation. When making a major improvement in published cross sections, development of an appropriate event generator is important. We use the data as a guide; the final cross section is dominated by $S$ and $P$ waves with the $S_{11}(1535)$ the dominant structure seen. An iterative procedure matching analytic Monte Carlo to real data was used to develop the event generator. The same fitting procedure used on data was applied to Monte Carlo events to calculate acceptance. The acceptance has significant variation across the bins with a maximum value of about 60%. When approaching the kinematic limit, the acceptance falls off rapidly. At the higher values of $W$, the proton goes forward where there is a hole in the CLAS acceptance. This causes problems for $\phi_\eta \sim 180^\circ$. We only report results in bins where the acceptance is greater than 3% and where it is not changing rapidly.

A detailed study of potential sources of systematic uncertainty was made. Since the $\eta$ peak shape and the background shape included various parameters, all were studied. The parameters were varied within the error bars determined in the fit for each cross section value. Additional tests were made for variations in particle identification and in the momentum scale. The momentum uncertainty arose from uncertainty in the details of both the magnetic field map and the alignments of the various tracking chambers. Sensitivity to momentum determination was largest close to threshold and falls off with increasing $W$. Since most of the $\eta$ events are produced near threshold, this source dominates the average systematic uncertainty. Cross sections were recalculated with slightly tighter fiducial cuts, and this variation was considered as a systematic uncertainty estimate. A variety of underlying physics models were used for evaluating the systematic on the radiative correction: using a single $S$-wave resonance or two, varying the mass and width of a single $S$ wave, or including a $P$-wave resonance. The quoted systematic uncertainty on the radiative correction includes these effects, but is dominated by Monte Carlo statistics in the calculation. The total systematic uncertainty for each bin in $W$, $Q^2$, and c.m. scattering angles was the sum of all the components added in quadrature. The average total systematic uncertainty for the angle-integrated cross sections was 3.3%, 3.9%, and 7.1% for data at 1.5, 2.5, and 4.0 GeV, respectively. The corresponding average estimated systematic uncertainties for the differential cross sections were 5.1%, 5.2%, and 7.6%. The breakdown by source

\[
D_{bkg}(M_X) = b_0(2\sqrt{\Delta m^\prime \Delta m} - \Delta m),
\]

where $\Delta m = m_{\text{max}} - M_X$ and $b_0$ is the overall strength. One example of this function is shown in Fig. 2. At the highest beam energy, a slightly more complicated function was used. Both forms contained one parameter ($\Delta m^\prime$) that was determined from our data by fitting to a polynomial in $W$. As with the peak shape function parameters, these fit parameters were included in the systematic uncertainty determination.

The main structure in the background fit function comes from the variation of the geometric acceptance of the detector as a function of missing mass. We found that proper modeling of this acceptance was very important. We determined this acceptance using a separate Monte Carlo program that generated $ep \rightarrow epX$ events with the $X$ mass thrown randomly across the fit region. After requiring the scattered electron and proton to be in the fiducial volume of CLAS, we compared the generated and accepted events in order to calculate the

FIG. 2. Sample missing mass $M_X$ spectrum for $ep \rightarrow e\eta p$. The bin shown is for $W = 1.535$ GeV and $Q^2 = 0.6$ GeV$^2$. The dashed line (right scale) shows the acceptance that is calculated with a Monte Carlo program. The sum of the $\eta$ signal shape and the raw background function modified by the acceptance function is then fit to the missing mass spectrum for each bin. In the figure, the dot-dashed curve is the raw background function $D_{bkg}$ from this fit, while the dotted curve shows those same values when multiplied by the acceptance function. The solid curve shows the full fit.
TABLE II. Summary of systematic uncertainties for the angle-integrated and differential cross section analysis. $M_b$, $\sigma_b$, $f$, $\alpha$, and "radiative corr." describe the $\eta$ missing peak mass shape; $b_{\text{bgd}}$ is the background missing mass function; other entries parametrize various detector properties. See text for details.

<table>
<thead>
<tr>
<th>Sys error source</th>
<th>Angle-integrated (int)</th>
<th>Differential (diff)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$M_b$</td>
<td>0.03%</td>
<td>0.06%</td>
</tr>
<tr>
<td>$\sigma_b$</td>
<td>0.4%</td>
<td>0.7%</td>
</tr>
<tr>
<td>$f$</td>
<td>1.3%</td>
<td>1.4%</td>
</tr>
<tr>
<td>$\alpha$</td>
<td>0.1%</td>
<td>0.1%</td>
</tr>
<tr>
<td>$b_{\text{bgd}}(M_b)$</td>
<td>0.1%</td>
<td>0.1%</td>
</tr>
<tr>
<td>Fiducial cut</td>
<td>0.6%</td>
<td>2.3%</td>
</tr>
<tr>
<td>Radiative corr.</td>
<td>1.0%</td>
<td>1.0%</td>
</tr>
<tr>
<td>Momentum scale</td>
<td>2.6%</td>
<td>4.4%</td>
</tr>
<tr>
<td>Total</td>
<td>3.9%</td>
<td>5.2%</td>
</tr>
</tbody>
</table>

for the 2.5 GeV data (the set from which the largest number of data points come) is given in Table II. The estimated systematic uncertainties for individual data points were seldom larger than the estimated statistical uncertainty.

IV. RESULTS

A. Angle-integrated cross sections

To get the angle-integrated cross sections, the events were binned in $W$ and $Q^2$, as shown in Table III. The 1.5 GeV beam energy data cover the $Q^2$ range from 0.13 to 0.4 GeV$^2$, while the upper two beam energies cover 0.6–3.3 GeV$^2$. Each bin is labeled by its centroid. Results are tabulated in the CLAS database [40]. These cross sections are presented in Fig. 3. The prominent peak at $W \sim 1.5$ GeV is primarily populated through intermediate excitation of the $S_{11}(1535)$ resonance. Fits to a Breit-Wigner relativistic form with an energy-dependent width, Eq. (8), are used to fit the low $W$ region. Various model calculations [33] in the past have found

TABLE III. Binning details for the angle-integrated cross sections. For each $Q^2$ bin, we show the minimum and maximum values of $Q^2$, the energy of the electron beam for the data set, and the maximum value of $W$ probed. The $W$ bin width in all cases was 10 MeV.

<table>
<thead>
<tr>
<th>$Q^2_{\text{min}}$ (GeV$^2$)</th>
<th>$Q^2_{\text{max}}$ (GeV$^2$)</th>
<th>$E_{\text{beam}}$ (GeV)</th>
<th>$W_{\text{max}}$ (GeV)</th>
</tr>
</thead>
<tbody>
<tr>
<td>0.13</td>
<td>0.2</td>
<td>1.5</td>
<td>1.66</td>
</tr>
<tr>
<td>0.2</td>
<td>0.3</td>
<td>1.5</td>
<td>1.64</td>
</tr>
<tr>
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<td>0.4</td>
<td>1.5</td>
<td>1.61</td>
</tr>
<tr>
<td>0.6</td>
<td>0.8</td>
<td>2.5</td>
<td>2.00</td>
</tr>
<tr>
<td>0.8</td>
<td>1.0</td>
<td>2.5</td>
<td>1.90</td>
</tr>
<tr>
<td>1.0</td>
<td>1.2</td>
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</tr>
<tr>
<td>1.2</td>
<td>1.4</td>
<td>2.5</td>
<td>1.69</td>
</tr>
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<td>1.3</td>
<td>1.7</td>
<td>4.0</td>
<td>2.30</td>
</tr>
<tr>
<td>1.7</td>
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<td>2.13</td>
</tr>
<tr>
<td>2.5</td>
<td>2.9</td>
<td>4.0</td>
<td>1.93</td>
</tr>
<tr>
<td>2.9</td>
<td>3.3</td>
<td>4.0</td>
<td>1.72</td>
</tr>
</tbody>
</table>

a small nonresonant contribution to the cross section, and none is needed here. The simple shape describes the low $W$ region well, but there are deviations for $W > 1.6$ GeV, presumably because of interference between $S_{11}(1535)$, $S_{11}(1650)$, and nonresonant processes. Although the higher mass resonance is very near to the state we seek to describe, all analyses [9] find a very small $\eta N$ branching fraction for $S_{11}(1650)$. Therefore, we restrict the fit to $W$ values less than 1.6 GeV. Two previous experiments [16,17] performed longitudinal/transverse separations in the late 1970s. Their results are consistent with no longitudinal component, albeit with large uncertainties. For the results presented here, the different beam energies have insufficient overlap in $W$ and $Q^2$ to separate these components. Under the assumption that the cross section is dominated by a single resonance and that $S_{11}$ is small, we can relate the $A_{1/2}$ to the peak cross sections extracted from the fit [see Eqs. (4) and (5)]:

$$A_{1/2}(Q^2) = \frac{W_R \Gamma_R}{2 m_b b_\eta} \sigma(W_R, Q^2)^{1/2}.$$  

Consistent with Armstrong et al. [19], a value of the full width of 150 MeV and an $S_{11} \rightarrow \eta N$ branching ratio of 0.55 were used. The results of this determination of $A_{1/2}$ are shown in Fig. 4 along with some previous results converted to be consistent with our choice of $\Gamma_R$ and $b_\eta$. The extracted values for $A_{1/2}$ for this experiment are also listed in Table IV. The precise normalization of $A_{1/2}$ depends on the choice of parameters for the contributing resonances, which are, as yet, not well determined. For instance, using the range of values listed in PDG for $\Gamma_R$ and $b_\eta$ leads to an 11%
systematic uncertainty on $A_1$. While these uncertainties affect the absolute value of $A_1$, the shape of the $Q^2$ dependence is much better determined. More detailed understanding of this state is required to better determine absolute values of $A_1$ and estimate the model dependence of those values. Using the choice of $\Gamma$ and $b_\eta$ described above, our extracted values are consistent with the previous values at low $Q^2$ [13,14,16,17], but with smaller uncertainties. At high $Q^2$, there is moderate disagreement between the previously published results of Brasse et al. [18] (at $Q^2 = 2.0$ and 3.0 GeV$^2$) and Armstrong et al. [19] (at 2.4 and 3.6 GeV$^2$). Our results match up nicely with Armstrong et al. and provide a precise determination of the shape of the $Q^2$ dependence of $A_1$ from low $Q^2$ up to their high $Q^2$ determinations.

The literature has various theoretical calculations of the photocoupling amplitude within the constituent quark model (CQM). Matching the slow falloff with $Q^2$ has been difficult. We show two recent calculations [2,41]. Aiello, Giannini, and Santopinto [41] use a hypercentral CQM and emphasize the importance of the three-body quark force. Although this prediction gives the best agreement with our data of all the calculations, it falls off more rapidly with $Q^2$ than the data. The Capstick and Keister calculation [2] starts with the more traditional CQM but uses relativistic dynamics in a light-front framework. Although the two calculations use different approaches, the CQM is not well defined, and many other results are given in the literature.

### B. Differential cross sections

For larger bins in $W$ and $Q^2$ (see Table V), we extract differential cross sections vs center-of-mass scattering angles of the $\eta (\cos \theta^* \phi^*)$. Each bin is labeled by its centroid. Results from this experiment are tabulated in the CLAS database [40]. For the $Q^2 = 0.8$ GeV$^2$ bin, Fig. 5 shows sample cross sections for four $W$ bins. The first two $W$ bins, $W = 1.52$ and 1.54 GeV, are at the peak of the $S_{11}(1535)$ resonance. They show a dominant isotropic component due to the $S_{11}(1535) \rightarrow \eta p$ process, but deviations from isotropy can be seen, especially at large $\phi^*$. By $W = 1.66$ GeV, the nons isotropy is quite evident. The cross section falls monotonically as a function of $\cos \theta^*$, with the cross section for forward $\eta$ production consistent with zero. As $W$ increases, this feature changes dramatically. At $W = 1.72$ GeV, the forward-backward asymmetry of the distributions has reversed, with forward $\eta$ production favored, while backward production is close to zero.

The $\eta$-MAID model [29], based on the MAID formalism [30], has been developed for $\eta$ electro- and photoproduction. This is an isobar model using a relativistic Breit-Wigner $W$ dependence with form factors. Eight PDG $3\pi$ and $4\pi$ resonances of mass less than 1.8 GeV and nonresonant processes are included at the amplitude level. They fit the photoproduction data [11,12,42] and the $Q^2$ dependence of the total cross section from electroproduction data [19,21] in
the \( S_{11}(1535) \) region. The results of a calculation implementing this model are included in Fig. 5. These calculations roughly match the observed cross sections. However, the angular dependence predicted by \( \eta \)-MAID does not agree with our data at \( W \) above the \( S_{11}(1535) \) region. The model was not fit to the differential cross sections of our previous work and the \( Q^2 \) dependence of the higher mass resonances, e.g., \( D_{13}(1675) \), was taken from a quark model calculation rather than from data.

For each \( W \) and \( Q^2 \) bin, the differential cross sections are fit to a form that comes from an expansion of the response functions from Eq. (2) in terms of associated Legendre polynomials \( P^m_\ell(\cos \theta^*_\eta) \), that is,

\[
\frac{d^2\sigma}{d\Omega_\eta} = \sum_{\ell=0}^{\infty} A_\ell P^0_\ell(\cos \theta^*_\eta) + \sum_{\ell=1}^{\infty} B_\ell P^1_\ell(\cos \theta^*_\eta) \cos \phi^*_\eta \\
+ \sum_{\ell=2} C_\ell P^2_\ell(\cos \theta^*_\eta) \cos 2\phi^*_\eta.
\tag{13}
\]

The parameters \( B_\ell \), \( C_\ell \), and \( A_\ell \) depend on \( Q^2 \), \( W \), and \( \epsilon \). They represent bilinear sums over contributing multipole amplitudes. Truncating to \( \ell \leq 3 \), we determined the parameters in Eq. (13) by a fit to the data. This truncation is motivated by three effects. (i) The lightest known \( N^* \) resonance with \( \ell > 3 \) is the \( G_{17}(2190) \); if the dominant effects on the differential cross section arise from interference with the dominant \( \ell = 0 \) partial wave, terms above \( \ell = 3 \) should be negligible. (ii) Fits to the \( \eta \)-MAID predicted cross sections yield negligible contributions for terms higher than \( \ell = 3 \). (iii) Good fits to the data are obtained with the truncated sum.

Results are shown vs \( W \) in Figs. 6–8. The quoted uncertainties contain both statistical and systematic uncertainties. We repeated the fit taking into account shifts in the cross section for each of the sources of systematic uncertainty studied in Sec. III. The total systematic uncertainty on the extracted parameters is the sum in quadrature of all individual sources. We normalize our fitted \( A_1 \), \( A_2 \), etc., to the isotropic term \( A_0 \) in order to more clearly show the \( W \) and \( Q^2 \) dependence of the shape of the differential cross section. For the ratios, the resulting uncertainty is dominated by the uncertainty on the numerators.

The isotropic component \( A_0 = \sigma_{tot}/4\pi \) shows the same features as the angle-integrated cross sections: a dominant peak from the \( S_{11}(1535) \) with additional structure above \( W = 1.6 \) GeV. The other prominent term in the fit is \( A_1 \), which represents the slope of the differential cross section vs \( \cos \theta^*_\eta \). A structure in the \( W \) dependence of \( A_1 \) was first seen in our previous publication [21] and is also seen by the GRAAL photoproduction experiment [42]. By examining the ratio \( A_1/A_0 \) in the new data, we can study this structure in more detail. Two features stand out in this ratio:

(i) The ratio \( A_1/A_0 \) is large and makes a rapid change from negative to positive values at \( W \approx 1.7 \) GeV.

(ii) This structure is roughly independent of \( Q^2 \) up to 2.5 GeV².

The simplest description for \( A_1 \) is in terms of interference between \( S \) and \( P \) waves. In that case, the rapid change in \( A_1 \) between \( W = 1.66 \) and 1.72 GeV could be caused by one of the waves passing through a resonance. There are two \( P \)-wave resonances in this region, \( P_{13}(1710) \) and \( P_{13}(1720) \). The former is rated 3\# by the PDG [9], but its properties are very difficult to extract from data and is therefore controversial [43]. The latter is rated 4\# by the PDG, but is also poorly understood [10]. Fits to CLAS \( \pi^+\pi^- \) electroproduction data [23] provided evidence that the existing baryon structure at \( W \sim 1.7 \) GeV should be changed. Their fits prefer either a greatly reduced \( \rho N \) decay branch for the existing \( P_{13}(1720) \) resonance or a new \( \frac{3}{2}^- \) state. In the present data, we cannot couple to a \( T = 3/2 \) state.
and are unable to distinguish between $P_{11}$ and $P_{13}$ states; we choose to use only a $P_{11}$ state. If one describes the cross section using only $S_{11}$ and $P_{11}$ partial waves, then

$$
A_1/A_0 = \frac{2\Re(E_0^* + M_{1-})}{|E_0| + |M_{1-}|^2}.
$$

(14)

In this case, the rapid shift from backward to forward peaked cross sections would be due to a rapid change in the relative phase of the $E_{0+}$ and $M_{1-}$ multipoles, because one of them is passing through resonance. The observation that this structure in $A_1/A_0$ is approximately $Q^2$ independent would then imply that $S_{11}$ and $P_{11}$ partial waves have a similar $Q^2$ dependence.

The values of $B_\ell$ shown in Fig. 7 are consistent with zero. These parameters measure the $\sigma_{LT}$ component of $\frac{d\sigma}{d\Omega}$, indicating that longitudinal amplitudes are not significant for this reaction (as was assumed in Sec. IV A). The $C_\ell$ parameters

FIG. 6. Results from fitting the angular distribution data of this experiment to Eq. (13). Coefficients of the $\phi^*$ independent terms are shown, i.e., those that contribute to $\sigma_T + \epsilon_0 \sigma_L$. Contributions from both statistical and systematic sources are displayed. The dashed line is the $\eta$-MAID prediction [29]; solid line is a four resonance fit to these terms.

FIG. 7. Same as Fig. 6, but showing the parameters corresponding to $\sigma_{LT}$. For the four resonance fit, these parameters are all zero.
in Fig. 8 measure the $\sigma_{TT}$ component of $\frac{d\sigma}{dQ^2}$. They are small, indicating that the $A_{3/2}$ components are also small for these values of $Q^2$.

To better understand the content of the $\eta$-MAID model, we also fit the parameters in Eq. (13) to the predicted cross sections from that model. The extracted parameters are also included in Figs. 6–8 as dashed lines. The prediction has a broader $S_{11}$ peak than is seen in our data. Some structure is predicted in $A_1$ arising from the $P_{11}(1710)$, but the size of this effect is not nearly enough to match our data. At high $W$, $\eta$-MAID predicts a negative $A_1$, in contrast to the significant positive value we observe. The model value of $b_{\eta N}$ for $P_{11}(1710)$ is much larger than the PDG value. Our data indicate the model value is incorrect. $\eta$-MAID contains many sources of $D$-wave contributions: $D_{13}(1520)$, $D_{13}(1700)$, and $D_{15}(1675)$ in addition to nonresonant amplitudes. This produces a value for $A_2/A_0$ that matches our data for $W < 1.6$ GeV, where the $D_{13}(1520)$ is the leading contribution. At larger $W$, agreement is poor. The model value of $b_{\eta N}$ for $D_{15}(1675)$ is much larger than the PDG value; our data indicate this is incorrect. The prediction for $A_2/A_0$ is near zero, as are our measurements.

Predictions for the $\sigma_{LT}$ and $\sigma_{TT}$ terms are consistent with our measurements. $C_2$ is the only term that is not negligible in $\eta$-MAID for our values of $Q^2$. It arises from the $A_{3/2}$ amplitudes of the $D$-wave resonances interfering with the larger $S$-wave amplitude. Our data agree with this general trend, but the effect is small compared to the uncertainties.

To gain further understanding of the resonance content of our data, we did an additional fit to the differential cross section data using relativistic Breit-Wigner resonances according to Eqs. (7)–(10). We fit the extracted parameters, up to $W = 1.8$ GeV, to a sum of four amplitudes for the following resonances: $S_{11}(1525)$, $S_{11}(1650)$, $P_{11}(1710)$, and $D_{13}(1520)$. This set of resonances was determined empirically as the minimal set required to fit the general features of our data. Although the properties of $P_{11}(1710)$ are very uncertain, it is an important contributor to this fit. We label it as $P_{11}$, but we cannot distinguish between $P_{11}$ and $P_{13}(1720)$ in our data set; specifically, a $P_{13}$ resonance would also give a rapid energy dependence in either $A_2$ or $C_2$ which we are unable to exclude with current statistical accuracy. The transverse response function $\sigma_T$ is modeled, i.e., the $A_i$ parameters. For the resonances with small contributions [$S_{11}(1650)$ and $D_{13}(1520)$], we fixed the resonance parameters to values obtained elsewhere. Masses and widths were set to average values from the Particle Data Group. For the $D_{13}(1520)$, we used the $Q^2$ dependence of $\eta$-MAID. Following the assumption of the single quark transition model [44], the ratio of the strength of the $S_{11}(1650)$ to that of the $S_{11}(1535)$ was taken to be independent of $Q^2$. Motivated by the $Q^2$ independence of $A_1/A_0$ in our data, we assumed the $P_{11}(1710)$ had the same $Q^2$ dependence as the $S_{11}$ states. This left 12 variables in the fit: the masses and widths of the $S_{11}(1535)$ and $P_{11}(1710)$, the relative strengths of the $S_{11}(1650)$ and $P_{11}(1710)$ to that of the $S_{11}(1535)$, an overall strength of the $D_{13}(1520)$, and the absolute strength of the $S_{11}(1535)$ in each of the five $Q^2$ bins. We view this as a simple fit. Our results should not be interpreted as a precise determination of resonance parameters, but rather as an indication of the dominant components needed in any future theoretical work.

The results of this fit are also shown in Fig. 6. The fit yields a reasonable, though not perfect description of our data. The isotropic term $A_0$ is described by the dominant $S_{11}(1535)$ peak, modified by the smaller $S_{11}(1650)$. The deviation from a simple Breit-Wigner is described as a combination of destructive interference between the $S_{11}(1535)$ and $S_{11}(1650)$, and a small contribution from the $P_{11}(1710)$. Including the $S_{11}(1650)$ results in an extracted value of $A_2$ for the $S_{11}(1535)$ which is 7% higher than that obtained with a single Breit-Wigner. The fitted width of the $P_{11}(1710)$ is 100 MeV, which is consistent with the central (but very uncertain) PDG value. We cannot isolate the $P_{11}(1710)$ photocoupling from that state’s branching ratio into $\eta N$; we can only quote a ratio of $\xi$ values [Eq. (6)]. The extracted value of $\xi_{1710}/\xi_{1535}$ is 0.22, which is about twice as large as in $\eta$-MAID, and nearly an order of magnitude larger than that extracted from parameters of the $P_{11}(1710)$ in the PDG. The $D_{13}(1520)$ primarily affects the quadratic term $A_2$. Including this resonance is enough to give a reasonable description of the $W$ dependence of $A_2$. Our data do not require significant contributions from higher $D$-wave states present in the $\eta$-MAID model.

We fit the structure in $A_1/A_0$ with a smooth $S$ wave and a rapidly changing $P$ wave. One could also describe this structure in terms of a new $S$-wave resonance interfering
with the $P$-wave component as in the model of Saghai and Li [45]. However, the amplitudes for the new resonance and the $P$-wave component must both fall off slowly with $Q^2$ to reproduce the data.

We also fit the $\phi$ dependence of the differential cross sections directly to Eq. (2) in order to obtain $\sigma_T + \epsilon \sigma_L$, $\sigma_{LT}$, and $\sigma_{TT}$ as a function of $W$, $Q^2$, and $\cos \theta^*$. We choose to fit the $\phi^*$ dependence in terms of the parallel/perpendicular asymmetry ($\text{Asym}_{TT}$) and the parallel/anti-parallel asymmetry ($\text{Asym}_{LT}$).

$$\text{Asym}_{TT} = \frac{\sigma_\parallel - \sigma_\perp}{\sigma_\parallel + \sigma_\perp}, \quad (15)$$

where

$$\sigma_\parallel = \frac{1}{2} (\sigma(\phi = 0) + \sigma(\phi = \pi)) \quad (16)$$

and

$$\sigma_\perp = \frac{1}{2} (\sigma(\phi = \pi/2) + \sigma(\phi = 3\pi/2)). \quad (17)$$

$$\text{Asym}_{LT} = \frac{\sigma(\phi = 0) - \sigma(\phi = \pi)}{\sigma(\phi = 0) + \sigma(\phi = \pi)} \quad (18)$$

For photoproduction, $\sigma_L$ and $\sigma_{LT}$ do not contribute. A common polarization parameter is the parallel/perpendicular asymmetry, $\Sigma$. It is defined by $\Sigma = \text{Asym}_{TT}/\Sigma_T$; note also that $\Sigma = \sigma_{TT}/\sigma_T$. In electroproduction, the possible presence of a longitudinal term makes the relationships more complicated:

$$\text{Asym}_{TT} = \frac{\epsilon \sigma_{TT}}{\sigma_T + \epsilon \sigma_L}, \quad (19)$$

$$\text{Asym}_{LT} = \frac{\sqrt{2} \epsilon (\epsilon + 1) \sigma_{LT}}{\sigma_T + \epsilon \sigma_L + \epsilon \sigma_{TT}}. \quad (20)$$

The data were analyzed in terms of these three response function combinations. A treatment of systematic uncertainties similar to that used for differential cross sections was applied. Figure 9 shows the values extracted from these fits for the same $W$ and $Q^2$ shown in Fig. 5. Error bars display the systematic and statistical uncertainties. The quantity $\sigma_T + \epsilon \sigma_L$ shows the same features we discussed earlier. The values for $\text{Asym}_{TT}$ are consistent with zero in all distributions, but the size of the estimated error bars are a strong function of $W$. For $W > 1.6$, the total cross section is smaller than where the $S_{11}$ resonance dominates. Extraction of meaningful values for the $\phi$ dependence in this manner is therefore difficult.

V. CONCLUSIONS

Our extractions of $A_1$ for the excitation of the $S_{11}(1535)$ cover a large range and match up well with Armstrong’s results [19] at higher $Q^2$. It should be noted again that there are significant model dependencies on describing the mass, width, and branching ratio into $\eta p$. These uncertainties lead
to significant systematic uncertainties on the absolute scale of $A_1$. These uncertainties are common to all points currently determined, so the shape of the distribution is well determined. It becomes a significant challenge for theory to reproduce this shape. No existing model is able to describe the full range.

Knowledge of the $N^*$ resonances in the region $W \sim 1700$ MeV is presently weak, because the quality of older $\pi N \rightarrow \pi \pi N$ and $n N \rightarrow \eta N$ data is poor. The coupling of known $P$-wave resonances to $\eta N$ is thought to be very small. In this experiment, rapid energy dependence in the strength in the $P$-wave for coupling to $\eta N$ final states is found. With a simple resonance model, we are able to describe these data with significant coupling of a $P$-wave resonance to $\eta N$. As with $S_{11}(1535)$, the falloff of this coupling must be very slow.

Although we can describe our measurements in terms of the $P_{11}(1710)$, we cannot distinguish between that or the $P_{13}(1720)$ with these data. Either resonance could produce the effect seen in $A_1/A_0$. The $P_{11}(1710)$ is more poorly understood than the $P_{13}(1720)$, so it is easier to accommodate our data by altering the partial widths of the $P_{11}$ rather than the $P_{13}$. A large $P_{13}$ could also produce effects in other terms. For instance, interference with a $D$ wave would give a small contribution to $A_1$, but it would not be significant compared to our uncertainties. A $P_{13}$ resonance would have an $A_1$ photo-excitation as well as $A_2$. Our determination of $A_1/A_0$ is sensitive to the $A_2$ amplitude; a significant $A_2$ amplitude could also lead to large effects in $\sigma_{TT}$.

Evidence for possible alterations in $N^*P$-wave resonances at masses of about 1.7 GeV is accumulating. In addition to what is found in this experiment, double-pion production experiments in this same mass range \cite{23,46} are also unable to be described with models using existing information. Since different models are used to describe the different data sets, it is important to use a common model to describe the combined measurements from these (and other) reactions. Such a program may allow us to accurately determine the properties of the $P$-wave resonances in this region.

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\begin{thebibliography}{99}


\end{thebibliography}
[38] L. Mo and Y. Tsai, Rev. Mod. Phys. 41, 205 (1969).