Abstract

In this work, the $\eta'$ meson photoproduction near threshold is studied in the quark model framework. A pseudovector effective Lagrangian is introduced for the $\eta' NN$ coupling and the newly published data from the SAPHIR Collaboration provide good constraints to this parameter. Corrections of order $O(1/m_q^2)$ for the electromagnetic interaction vertex are taken into account, which produce corrections of order $O(1/m_q^3)$ to the transition amplitude for $\gamma p \to \eta' p$. Some low-lying resonances, $S_{11}(1535)$, $P_{13}(1720)$, and $P_{13}(1900)$ are found to have significant contributions. A bump structure around $E_\gamma \approx 2$ GeV is found arising from the $n = 3$ terms in the harmonic oscillator basis. The beam polarization asymmetries are predicted and can be tested against the forthcoming data from GRAAL.

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I. INTRODUCTION

The study of $\eta'$ meson photoproduction becomes more and more attractive in both experiment and theory. However, within a long period, only few data were available from ABBHHM [1] and AHHM [2]. According to those results, it was difficult to even establish the feature of the resonance excitations. This situation challenged theoretical attempts to study $\eta'$ photoproduction, $\gamma p \rightarrow \eta'p$. In Ref. [3], a hadronic model was proposed with a single resonance $P_{11}(2050)$ excited in this reaction, which thus accounted for the possible strong peak near threshold. However, the roles of those low-lying resonances in this reaction have not been explained. A more self-consistent attempt was made by Li [4] through a quark model approach based on the assumption that the $\eta'$ meson couples to the constituent quark via the same pseudovector coupling as that of the $\eta$. Although the available data at that time did not permit detailed investigation of this channel, the author showed that a strong peak near threshold could be produced by the low-lying resonances, especially dominated by the $S_{11}(1535)$. Meanwhile, a small bump around $E_\gamma = 2$ GeV, which comes from the $n = 3$ resonance excitations, was predicted. However, in that work, two resonances, $P_{13}(1900)$ and $F_{15}(2000)$, which are close to the $\eta'$ production threshold, had not been included as little was known about them from Ref. [7]. we are left with questions about the role played by these two resonances and more interestingly, whether the $\gamma p \rightarrow \eta'p$ can provide signals for the existences of these two resonances. This is related to the motivation of searching for “missing resonances” in the meson photoproductions. To answer such questions, more precise measurements for this reaction are needed.

The recently published data from the SAPHIR Collaboration [5] brought the possibility of further systematic study of the resonance excitations in $\gamma p \rightarrow \eta'p$. With the photon energy covering a range from threshold (1.45 GeV) to 2.6 GeV for the first time, the resonance excitations were established experimentally in $\eta'$ meson photoproduction. The steep rise and fall of the total cross section near threshold indicated the dominance of resonance excitations. In Ref. [5], the resonance structure was explained as due to two resonance excitations ($S_{11}(2090)$ and $P_{11}(2100)$). Such a prescription led to the arbitrary assumption that low-lying resonance effects could be completely neglected, and risked overlooking influences of small cross sections to the angular distributions. As an example, we recall that in the study of $\eta$ meson photo- and electro-production, the small contribution from the $D_{13}(1520)$ produces significant effects which result in the deviation of the angular distributions from the dominant $S$-wave in the low energies. Moreover, resonances other than the $S_{11}(2090)$ and $P_{11}(2100)$ have masses very close to the $\eta'$ production threshold, such as the $P_{13}(1900)$. Therefore, one should also include their threshold effects in this reaction. Based on the above considerations, we believe that a self-consistent treatment to systematically include all the resonances in the $\gamma p \rightarrow \eta'p$ reaction is required, and the quark model provides an ideal starting point [8].

In this work, we follow the same scheme as in Ref. [4] to study the $\eta'$ photoproduction at tree level in the quark model. There are a number of new features in this work: 1) The electromagnetic interaction is expanded to $O(1/m_q^2)$, which will partly account for the relativistic corrections beyond the NRCQM. 2) Two resonances, $P_{13}(1900)$ and $F_{15}(2000)$, which are attributed to representation $[70, 2, 2, 2, J]$, have been included. These two resonances are not well determined experimentally. In $\gamma p \rightarrow \eta'p$, we expect that some signals for these
two resonances can be clarified. 3) In this study, the equal velocity frame (EVF) is adopted for the Lorentz boost [14]. As shown in Ref. [9], the EVF can boost the particle momentum close to a realistic value. An ad hoc form factor is avoided in the EVF for the spatial integrals. 4) The new data have better constraints to the nucleon pole terms, that sheds light on the $\eta'NN$ coupling constant from the quark model. 5) In addition to the study of the differential and total cross sections, the beam polarization asymmetry is predicted, which can be tested by the forthcoming data from GRAAL and JLab in the near future.

In Section II, a brief introduction to this model is given. Results and analysis are presented in Section III. Conclusions are drawn in Section IV.

II. THE MODEL

The $\eta'NN$ interaction is introduced at quark level through the effective Lagrangian for the quark-meson vertex [3]:

$$L_{eff} = \sum_{j} \frac{1}{f_{\eta'}} \bar{\psi}_{j} \gamma^{\mu} \gamma_{5} \psi_{j} \partial_{\mu} \phi_{\eta'},$$

(1)

where $\bar{\psi}_{j}$ ($\psi_{j}$) represents the $j$th quark (anti-quark) field in the nucleon, and $\phi_{\eta'}$ is the meson field. It is still not elementarily clear if the $\eta'$ couples to the nucleon through a pseudoscalar or pseudovector coupling, or even both. However, as pointed out in Ref. [10], the operators for the pseudoscalar and pseudovector coupling have the same leading order expression at quark level.

We take the assumption that the $\eta'NN$ coupling satisfies the Goldberger-Treiman relation, which means that we can relate the quark-meson coupling to the $\eta'NN$ coupling through

$$g_{\eta'NN} = \frac{g_{A} M_{N}}{f_{\eta'}}.$$  

(2)

The coupling $g_{\eta'NN}$ is treated as a parameter that will be determined by the experimental data.

The electromagnetic interaction for the three-quark baryon system can be expanded to order $O(1/m_{q}^{2})$:

$$H_{em} = (\frac{\omega_{3}}{2})^{2} \sum_{j} (h_{j}^{c} + h_{j}^{i}) e^{ik \cdot r_{j}},$$

(3)

where

$$h_{j}^{c} \equiv [\epsilon_{j} \hat{R} \cdot \epsilon - \frac{e_{j}}{2m_{j}} \sigma \cdot (\epsilon \times \hat{k})$$

$$+ \frac{1}{4M_{T}} (\frac{e_{j}}{m_{j}} - \frac{e_{T}}{M_{T}}) \sigma \cdot (\epsilon \times \hat{P})],$$

(4)

and
\[ h_j^i \equiv e_j (r_j - R) \cdot \epsilon + \frac{1}{4} \left( \frac{e_j}{m_j} - \frac{e_T}{M_T} \right) \sigma \cdot \left[ \epsilon \times \left( \frac{p_j}{m_j} - \frac{P}{M_T} \right) \right], \]  

(5)

where \( \hat{k} \) is the unit vector of the photon momentum, \( \hat{k} \equiv k / \omega \); \( h_j^c \) and \( h_j^i \) are operators for the c.m. and internal motions of the three-quark baryon system, respectively. The charge and mass of the \( j \)th quark are denoted as \( e_j \) and \( m_j \), respectively, while \( e_T \) and \( M_T \) denote the total charge and mass of the baryon, respectively. The position and momentum for the \( j \)th quark are \( r_j \) and \( p_j \), respectively, while \( R \) and \( P \) are for the c.m. motion of the baryon system.

We explicitly write the \( S_{11}(1535) \) transition amplitude given by Eq. 3 as follows:

\[
\mathcal{M}_{S_{11}} = \frac{2M_{S_{11}}e^{- \frac{k^2+q^2}{4m_q^2}}}{(s - M_{S_{11}}^2 + iM_{S_{11}}\Gamma_{S_{11}})} \left[ \frac{\omega_{\eta}'}{\mu_q} - \left( \frac{\omega_{\eta}'}{M_N + E_f} + 1 \right) \frac{2q^2}{3\alpha^2} \right] 
\times \left\{ \frac{1}{6} (\omega_\gamma + \frac{k^2}{2m_q}) \sigma \cdot \epsilon \gamma 
+ \frac{\omega_\gamma}{9m_q} (\mu_p - \mu_0) |k||q| \cos(\theta) \sigma \cdot \epsilon \gamma 
+ \frac{\omega_\gamma}{36m_q} (\mu_p - \mu_0) \sigma \cdot (q \times (k \times \epsilon \gamma)) 
- \frac{\omega_\gamma}{36m_q} (\mu_p - \mu_0) \sigma \cdot (kq \cdot \epsilon \gamma) \right\},
\]

(6)

where \( \omega_{\eta}' \) denotes the meson energy in the \( \eta'N \) c.m. system, and \( \mu_q \) is the reduced mass of two quarks which equals \( m_q/2 \) here. The first term in the curly bracket is the formula with the electromagnetic interaction expanded to order \( O(1/m_q) \), i.e. from the first lines in Eq. 4 and 5. This term is the leading contribution which has been extensively used in literature \[4,6\]. Taking into account the meson interaction vertex, which is also expanded to order \( O(1/m_q) \), the transition amplitudes of the first term in the curly bracket of Eq. 6 has been rigorously expanded to order \( O(1/m_q^2) \). In other words, the electromagnetic interaction of order \( O(1/m_q^2) \) will introduce corrections at order \( O(1/m_q^3) \) to the transition amplitudes. These corrections are given by the terms proportional to \( (\mu_p - \mu_0) \) which is the anomalous magnetic moment of the proton at leading order. Here, \( \mu_p \) and \( \mu_0 \) are the proton magnetic moment and the nuclear magneton, respectively, and they are defined as

\[
\langle N_f | \sum_j \frac{e_j}{2m_j} \sigma_j | N_i \rangle \equiv \langle N_f | \mu_p \sigma | N_i \rangle,
\]

(7)

and

\[
\mu_0 \equiv \frac{e_T}{2M_T}.
\]

(8)

Note that in Eq. 8, the proton charge, \( e_T \), has been moved out of the transition amplitude.

The formalism for the \( \eta' \) production is the same as for the \( \eta \) apart from the inclusion of order \( O(1/m_q^2) \) corrections to the electromagnetic interaction. We even adopt the quark
model parameters in the $\eta$ production here, i.e. the harmonic oscillator strength $\alpha = 384.5$ MeV and $m_q = 330$ MeV. The main decay channels of the $S_{11}(1535)$ are $\pi N$ and $\eta N$, for which the branching ratios 0.45 and 0.55 are adopted, respectively. These values are the same as derived in the $\eta$ meson production [9]. It should be pointed out that at high energies the phase space will permit small branching ratio for a off-shell resonance decaying into higher mass channels. This feature will change the on-shell branching ratio values accordingly. However, at leading order, the influences from such deviations are negligibly small. Within a wide energy regions, the energy-dependence of the total decay width is still dominated by decay of the lower mass channels. Therefore, although the total width of the $S_{11}(1535)$ is occupied by the $\pi N$ and $\eta N$, the off-shell $S_{11}$ will permit small branching ratio to $\eta'N$ channel.

### III. RESULTS AND ANALYSIS

The newly published photoproduction data from SAPHIR provide for the first time the near-threshold angular distributions for $\gamma p \to \eta'p$. Although the uncertainties are still large, the resonance excitations can be recognized clearly.

We present calculations of the angular distributions in Fig. 1 for seven energies and compare with the data of Ref. [5]. Near threshold (Fig. 1a), the cross section is dominated by the $S$-wave. The quite flattened angular distribution cannot be explained by the $t$-channel meson exchanges, which are generally forward peaked. In Fig. 1a, the solid curve denotes the calculation with the middle energy $E_\gamma = 1.490$ GeV. It shows that the threshold cross section is dominated by the $S_{11}(1535)$ excitation. We illustrate the $S_{11}(1535)$ influence by switching off its contribution. Comparing the dot-dashed curve to the solid one, we see that the $S_{11}(1535)$ strongly enhances the cross section near threshold. To show the possible uncertainties arising from the photon energy, we also present the calculations with $E_\gamma$ deviating 20 MeV from the middle values, i.e. $E_\gamma = 1.470$ (dashed curve) and 1.510 GeV (dotted curve).

The $S_{11}$ excitation also accounts for the large cross section in the energy region up to $E_\gamma \approx 1.8$ GeV. In Fig. 1b and c, the forward peaking is found to be produced by the interferences between the $S$-wave and $P$-wave amplitudes. Without the $S_{11}$, the angular distribution exhibits only a weak forward peak (dot-dashed curve in Fig. 1b). It should be noted that the excitation threshold of the $P_{13}(1900)$ is very close to the $\eta'$ production threshold. However, the threshold energy is not ideal for identifying the $P_{13}(1900)$ signal due to the strong $S_{11}(1535)$ contribution. At $E_\gamma = 1.490$ GeV, the effects of the $P_{13}(1900)$ absence cannot be seen clearly. But in Fig. 1c, we show that the $P_{13}(1900)$ is important to account for the forward peaking at $E_\gamma = 1.690$ GeV. The dashed curve in Fig. 1c denotes the calculations without the $P_{13}(1900)$, which completely changes the angular distributions at forward angles. The difference between the solid and dashed curves, are found from the interferences between the $S$ and $P$-wave amplitudes (dominant ones), which result in the forward peaking. This feature might suggest that the forward peaking can serve as a signal of the $P_{13}(1900)$ excitation in $\gamma p \to \eta'p$.

In the energy region, $1.74 < E_\gamma < 2.04$ GeV, the angular distributions show some structures which might come from resonance excitations. We find that apart from the $S_{11}(1535)$, $P_{13}(1720)$ and $P_{13}(1900)$, contributions from $n = 3$ terms become important. Although the
calculation cannot account for the obvious increase at 90° in Fig. 1d, one can see that the cross section turns to level off at this energy. Qualitatively, such a structure favors a smaller $P$-wave strength. In Fig. 1d, the dashed curve denotes the angular distribution without the $P_{13}(1900)$.

The dashed curve in Fig. 1e denotes the calculations without $n = 3$ terms. The backward peak comes from the $P_{13}(1900)$ contribution. At this energy region, the $P_{13}(1900)$ cross section has dropped down significantly. However, its influence still makes sense as the dotted curve in Fig. 1e illustrates that without the $P_{13}(1900)$, the dashed curve is significantly changed. Comparing the three curves, we learn that the $P_{13}(1900)$ and the $n = 3$ terms play important roles in producing the forward peak, and simultaneously attenuate the backward peaking tendency. The same feature continues up to 2.14 GeV. As shown in Fig. 1f, without the $P_{13}(1900)$ and $n = 3$ terms (dotted curve) the forward peak cannot be reproduced.

Up to 2.44 GeV photon energies, the resonance excitations of the low-lying resonances with $n \leq 2$ will be competing with the excitations from higher harmonic oscillator shells. It becomes very complicated to distinguish the individual resonance excitations because of their small cross sections and wide mass overlaps. Here, the energy evolution turns out to be important for such a collective description. Impressively, in Fig. 1g, we find that the forward peaking character keeps in the degenerate $n = 3$ terms, while the energy evolution of the cross sections are also perfectly satisfied. It should be noted that the backward peak in Fig. 1g is produced by the $u$-channel processes.

In Fig. 2, the total cross sections are presented in comparison with the old data from Ref. [1,2], and the new data from SAPHIR [3]. Good agreement with the data is obtained over large energy regions. The solid curve shows the full calculation of our model. Explicitly, the role of the strong $S_{11}(1535)$ can be shown by comparing the results of the exclusive $S_{11}$ excitation (dashed curve), and the calculations without the $S_{11}$ (dot-dashed curve). It shows that the sharp peak near threshold is produced by the $S_{11}(1535)$ excitation. Interestingly, a bump is automatically produced around $E_\gamma \approx 2$ GeV. This structure is a unique signal of the $n = 3$ resonance contributions in the harmonic oscillator basis. As shown by the dotted curve, the second bump does not show up without $n = 3$ excitation. The discrepancy between the solid and dotted curve suggests that the interferences between the low-lying resonance excitations and the $n = 3$ terms play a key role in producing the second bump around $E_\gamma = 2$ GeV. In the data from SAPHIR, a very similar structure shows up at the same energy. But taking into account the large errors, we need more precise data to justify it.

In Fig. 2, large partial cancellations exist between the resonance excitations and the nucleon pole terms near threshold. With the $S_{11}$ excitation and the nucleon pole terms, the cross sections are over-estimated significantly in the lower energy region $E_\gamma < 2.0$ GeV. We do not show this result in order to keep the figure clear to read. Such an enhancement, however, is cancelled by the resonance terms. Removing the $S_{11}$ contribution, and comparing the results with (dot-dashed curve) and without (heavy dotted curve) the nucleon pole terms, we see that the nucleon pole terms cancel the amplitudes significantly over a large energy range. Interestingly, such a cancellation is sensitive to the $\eta^\prime NN$ coupling and we find that very little freedom is left for the nucleon pole terms. When a coupling $\alpha_{\eta^\prime} = 0.22$ is adopted, the $\eta^\prime NN$ coupling $g_{\eta^\prime NN} = 1.66$ can be derived.

Among those low-lying resonances of $n \leq 2$, only the $P_{13}(1900)$ and $F_{15}(2000)$ of rep-
representation \([70, 28, 2, 2, J]\) are above the threshold of the \(\eta'\) production. In general, for a resonance \(N^*\) above the \(\eta'\) threshold, one can relate its helicity amplitudes to its exclusive total cross section at its mass position \(M_R\):

\[
\sigma_{tot}(\gamma p \rightarrow N^* \rightarrow \eta'p) = \frac{M_N}{M_R} \frac{b_{\eta'}}{\Gamma_R} 2\{|A_1|^2 + |A_2|^2\},
\]

(9)

where \(b_{\eta'}\) is the branching ratio of the \(N^*\) decay into \(\eta'N\) channel. \(A_1\) and \(A_2\) are the two independent photon excitation helicity amplitudes of the resonance. This relation is model-independent, and thus given sufficient information for the photon interaction vertex, one would in principle be able to derive information about the meson interaction vertex for the resonance. However, at present the status of these two resonances, \(P_{13}(1900)\) and \(F_{15}(2000)\), has not been well-established. Information about their decay modes, branching ratios, as well as the photon excitation helicity amplitudes, is not available. To go as far as possible based on the present situation, we shall use the quark model calculations as input to study the excitations of these two resonances in the \(\eta'\) production. On the one hand, we calculate the exclusive total cross section \(\sigma_{tot}(\gamma p \rightarrow N^* \rightarrow \eta'p)\) in this model. From Eq. (8) we can derive the quantity \(\xi \equiv \{|A_1|^2 + |A_2|^2\}^{\frac{1}{2}}\) for three values of the branching ratio \(b_{\eta'} = 0.1, 0.2\) and 0.3. The sensitivity of \(\xi\) to the branching ratio \(b_{\eta'}\) can be shown. On the other hand, we separately calculate the helicity amplitudes \(A_1, A_2\) for \(\xi\) in the NRCQM. Assuming the NRCQM provides the leading order calculations for the resonance photo-excitations, we can determine the magnitudes of the \(\eta'N\) branching ratios for these two resonances by comparing the results of the two methods.

In Table II, the quantities derived through the two methods are shown. Several lessons can be learned here: i) For both resonances, their branching ratios to \(\eta'N\) might be as large as 20%. ii) The \(P_{13}(1900)\) becomes very interesting due to the feature that it favors a larger branching ratio to \(\eta'N\) channel, and its mass is very close to the threshold of \(\eta'\) production. iii) The NRCQM calculations of the helicity amplitudes give the same order of magnitude for \(\xi\), which might suggest that the meson coupling has been reasonably estimated.

We do not expect the results in Table II to be precise due to the present lack of data, as well as shortcomings of this approach. However, the prediction of the order of magnitude of the \(\eta'NN\) coupling can be regarded as reasonable, although so far it has not been well-determined. The \(\eta'\) meson production might be a possible channel to determine the \(P_{13}(1900)\) and \(F_{15}(2000)\) experimentally.

We extend the calculation to the beam polarization asymmetry \(\Sigma\). In terms of the helicity amplitudes, this has the form

\[
\Sigma = -\text{Re}\{H_1(\theta)H_3^*(\theta) - H_2(\theta)H_4^*(\theta)\},
\]

(10)

where \(H_{1,2,3,4}\) are the four independent helicity transition amplitudes that have been extensively discussed in the literature (See, for example Ref. [13]). The angle \(\theta\) is given by \(\mathbf{k} \cdot \mathbf{q} = |\mathbf{k}| |\mathbf{q}| \cos(\theta)\) in the c.m. system of the final state hadrons. Normalized by the differential cross section, the beam polarization asymmetry is presented in Fig. 3a for three energies, \(E_\gamma = 1.65\) (solid), 1.6 (dashed) and 1.5 GeV (dotted). Near threshold, the asymmetries are found to be small. It shows that the \(\Sigma\) is sensitive to the energy. Large negative asymmetries are produced at \(\theta \approx 140^\circ\). We find the nodal structure is governed by the \(S\) and \(P\)-wave interferences. To be more clear, we study the \(S\) and \(P\)-wave interference effects.
at $E_\gamma = 1.65$ GeV in Fig. 3b by eliminating the $P_{13}(1900)$ to produce the dashed curve. Comparing the dashed curve to the solid one, it shows that the $P_{13}(1900)$ strongly shifts the positive node to the backward angles, and enhances it significantly. Then by removing the $P_{13}(1720)$, we obtain the dotted curve in Fig. 3b, which exhibits small asymmetries and is governed by the dominant $S$-wave amplitudes. It should be noted that the influence of the $D_{13}(1520)$ becomes very weak in the $\eta'$ production. Qualitatively, a stronger $D$-wave will attenuate the negative trend and enhance the positive one. We expect that data from the GRAAL Collaboration and JLab will provide tests of this prediction.

IV. CONCLUSIONS

In summary, we studied $\eta'$ meson photoproduction near threshold by introducing an effective Lagrangian for the quark meson vertex in the quark model. The $s$-channel resonances are systematically taken into account. With the data from SAPHIR, the resonance excitations near threshold are established in this reaction. Contrary to the explanation of Ref. [5], we find that the $S_{11}(1535)$ accounts for the steep peaking near threshold. The $S$ and $P$-wave interference is found to be important in reproducing the angular distributions near threshold. This feature might be useful for identifying the $P_{13}(1900)$ signal in this channel. In this model, a bump structure is automatically produced due to the contributions of $n = 3$ terms in the harmonic oscillator basis. A similar structure appears in the data but needs to be justified with more precise measurement. The nucleon pole terms can be reasonably constrained in this approach and provide an estimate of the $\eta'NN$ coupling of $g_{\eta'NN} = 1.66$.

Expanding the photon interaction to $O(1/m_q^2)$, we obtain corrections to order $O(1/m_q^3)$ for $\gamma p \rightarrow \eta'p$. We find that near threshold such corrections do not change the calculations significantly in comparison with calculations at $O(1/m_q^2)$. This suggests that the leading order contributions based on the NRCQM contain the main characters of such a non-perturbative process. Such an empirical starting point is still helpful for us to understand the nucleon resonance internal structures.

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TABLE I. The dependence of the quantity $\xi$ on the $\eta'N$ branching ratio in this model for the $P_{13}(1900)$ and $F_{15}(2000)$. Independent calculations in the NRCQM are presented as a comparison.

|                           | $P_{13}(1900)$ | $F_{15}(2000)$ |
|---------------------------|----------------|----------------|
| $M_R$ (MeV)               | 1900           | 2000           |
| $\Gamma_R$ (MeV)         | 400            | 400            |
| $\xi_1(b_{\eta'} = 0.1)$ (GeV$^1$) | $5.10 \times 10^{-3}$ | $2.64 \times 10^{-3}$ |
| $\xi_2(b_{\eta'} = 0.2)$ (GeV$^{3/2}$) | $3.61 \times 10^{-3}$ | $1.86 \times 10^{-3}$ |
| $\xi_3(b_{\eta'} = 0.3)$ (GeV$^{1/2}$) | $2.62 \times 10^{-3}$ | $1.52 \times 10^{-3}$ |
| $\bar{\xi}_{\gamma p \to N^*}$ (GeV$^{3/2}$) | $2.84 \times 10^{-3}$ | $1.74 \times 10^{-3}$ |
| $A_{\frac{1}{2}}$ (GeV$^{3/2}$) | $-2.74 \times 10^{-3}$ | $1.04 \times 10^{-3}$ |
| $A_{\frac{3}{2}}$ (GeV$^{1/2}$) | $0.732 \times 10^{-3}$ | $-1.39 \times 10^{-3}$ |
FIGURES

FIG. 1. Angular distributions for 7 energy bins are compared to the data from SAPHIR [3]. The calculations at the middle energies of each bins are presented by the solid curves. See text for the notations for the dashed, dotted and dot-dashed curves.

FIG. 2. Total cross sections compared with the data [3,4,5]. The solid curve denotes the full calculation of this model. The dashed curve denotes the exclusive contribution from the $S_{11}(1535)$ excitation, while the dot-dashed denotes full calculation with the $S_{11}(1535)$ switched off. The dotted curve denotes calculation without $n = 3$ terms, while the heavy-dotted curve, without the $S_{11}(1535)$ and the nucleon pole terms.

FIG. 3. Beam polarization asymmetries predicted by this model. In (a), asymmetries for three energies $E_\gamma = 1.65$ (solid), 1.6 (dashed) and 1.5 GeV (dotted), are shown. In (b), the solid curve is the same as that in (a). The dashed curve denotes calculation without the $P_{13}(1900)$, while the dotted curve without both $P_{13}(1900)$ and $P_{13}(1720)$. 
