Do we understand the $\eta N$ interaction from the near threshold $\eta$ photoproduction on the deuteron?

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The effects of final state interaction in incoherent $\eta$ photoproduction on deuteron are studied within a three-body approach including a realistic $NN$ potential. The results are compared with available data, and differences with other theoretical predictions are analyzed. The role of the $\eta N$ interaction and the possibility of extracting the $\eta N$ scattering parameters from this reaction are discussed.

I. INTRODUCTION

The role of final state interaction (FSI) in incoherent photoproduction of $\eta$-mesons on the deuteron was investigated in our previous work [1,2]. First we have considered in [1] the first-order approximation (FOA) where complete hadronic rescattering is included in the two-body subsystems ($NN$ and $\eta N$) only, and in the subsequent work [2] the three-body aspects of the reaction were studied. Summarizing the results we would like to conclude that the importance of final state interaction arises from two sources which are interrelated. Firstly, the impulse approximation (IA), where the FSI is ignored, predicts a very strong suppression of the cross section close to the threshold. The reason for this is a strong mismatch between the large nucleon momentum needed for $\eta$ meson production (more than 200 MeV/c) and the average available momentum in the deuteron of about 45 MeV/c. The FSI provides an efficient mechanism to compensate this momentum mismatch by the collision between the final particles. As a consequence, already the leading terms in the multiple scattering series associated with pairwise $NN$ and $\eta N$ scattering produce a very large enhancement of the $\eta$ production rate close to the threshold. This conclusion has later been confirmed by Sibirtsev et al. [3–5].

The second reason is a relatively strong $\eta N$ interaction at low-energy. Because of its amplification via the $NN$ interaction an appreciable attraction in the $\eta NN$ system is generated, which in turn leads to a virtual pole in the $S$-wave part of the three-body scattering amplitude [6–8]. In the case of the spin-singlet $\eta NN$ state ($J^\pi, T) = (0^−, 1$), which dominates the reaction close to the threshold, the pole lies on the three-body unphysical sheet not far from zero kinetic energy [7]. We would like to emphasize that the existence of a virtual state near zero energy is the reason, why the FOA cannot provide an accurate description of the reaction dynamics at low energy. This statement is corroborated by the fact, that the multiple scattering series converges slowly near the pole so that the FOA does not constitute a reliable approximation to the whole series. In particular, the FOA is unable to account for anomalies in the energy dependence of the total cross section, since singularities of the three-body scattering amplitude cannot be generated in a perturbative approach (see, e.g., the general arguments given in [9]). Indeed, as was shown in [2], only the full three-body treatment can explain, at least qualitatively, the anomalously strong rise of the experimental $\gamma d \rightarrow \eta X$ cross section just above threshold.

However, our work in [2] was only considered as a first step for taking into account the most important dynamical properties of the $\eta NN$ system, because of several simplifications for the two-body forces. This concerns primarily the $NN$-sector, especially the deuteron wave function, which in [2] was taken as a pure $^3S_1$ state, generated by the Yamaguchi potential [10]. Clearly, such a treatment is too simple for $\eta$ photoproduction where large momentum transfers come into play.

Thus the first motivation for the present paper is to overcome this shortcoming by using a realistic $NN$-potential. In Sect. II, we give the details of the two-body interactions and present the results for the total and differential cross sections. For the comparison with the inclusive data of [11] also the coherent reaction $\gamma d \rightarrow \eta d$ is calculated. As will be shown, our predictions slightly underestimate the data in almost the whole energy region from threshold up to $E_\gamma = 750$ MeV. We will discuss possible reasons for this disagreement. The second motivation of the paper is to
II. TWO-BODY INPUT AND DISCUSSION OF THE RESULTS

The transition matrix element of the reaction $\gamma d \to \eta np$ has been evaluated with inclusion of the hadronic interactions between the final particles whereas the initial electromagnetic interaction is treated perturbatively in lowest order. The general formalism of our approach is described in detail in [2]. Here we present mainly the most important two-body ingredients of the calculation.

As a basic input we need the $\eta N$ and $NN$ scattering amplitudes which were restricted to $S$-states only in view of the near threshold region. For the $\eta N$ interaction we use a conventional isobar ansatz as described in detail in [12], where the $\eta N$ channel is coupled with $\pi N$ and $\pi \pi N$ channels through the excitation of the $S_{11}(1535)$ resonance. The separable $\eta N$ scattering matrix has the usual isobar form

$$t_{\eta N}(q, q'; W) = \frac{g_\alpha(q)g_\alpha(q')}{{W - M_0 - \Sigma_\eta(W) - \Sigma_\pi(W) + \frac{i}{2}\Gamma_{\pi\pi}(W)}},$$

as a function of the invariant energy $W$. The $S_{11}(1535)$ self energies $\Sigma_\eta$ and $\Sigma_\pi$ are determined by the vertex functions $g_\alpha(q)$ as

$$\Sigma_\alpha(W) = \frac{1}{(2\pi)^2} \int_0^\infty \frac{q^2 dq}{2\omega_\alpha} \frac{|g_\alpha(q)|^2}{W - E_N(q) - \omega_\alpha(q) + i\epsilon}, \quad \alpha \in \{\pi, \eta\},$$

with $E_N$ and $\omega_\alpha$, $(\alpha \in \{\pi, \eta\})$ denoting the on-shell energies of nucleon and meson, respectively. The two-pion channel is included in a simplified manner by adding the $S_{11}(1535)$ decay width, parametrized by

$$\Gamma_{\pi\pi}(W) = \gamma_{\pi\pi} \frac{W - M_N - 2m_\pi}{m_\pi}, \quad \text{with} \quad \gamma_{\pi\pi} = 4.3 \text{ MeV}.$$  

The vertex functions are taken in a Hulthén form

$$g_\alpha(q) = g_\alpha \left[1 + \frac{q^2}{\beta_\alpha^2}\right]^{-1},$$

containing the strength of the coupling $a_\alpha$ and the range of the Hulthén form factor $\beta_\alpha$. The parameters (see Table I) were adjusted to fit the $\eta N$ scattering length

$$a_{\eta N} = (0.5 + i0.32) \text{ fm},$$

and at the same time to provide a reasonably good description of the reactions $\pi N \to \pi N$ and $\pi N \to \eta N$ (for more details see [13]). We consider this value as an approximate average of the scattering lengths provided by modern $\eta N$ scattering analyses (see, e.g., the compilation in Table I of [4]). In Sect. IV, where we study the dependence of the results on the $\eta N$ interaction strength, two other sets of $\eta N$ parameters are used. The first one, giving $a_{\eta N} = (0.25 + i0.16) \text{ fm}$ is taken from [12]. The second set is adjusted such, that the scattering length $a_{\eta N} = (0.75 + i0.27) \text{ fm}$ of the analysis [14] is reproduced. The corresponding parameters are listed in Table I.

We would like to emphasize that the latter value is obtained simply by varying the parameters of our separable ansatz (1) without using the original model of [14]. One of the consequences of this strategy is that the Born term introduced in [14] is absorbed in our approach by the resonance amplitude, which leads to an overestimation of the $S_{11}$ contribution. In this case we achieve a satisfactory description of the $\pi N \to \eta N$ experimental cross section, and the resulting $\eta N$ amplitude agrees rather well with the ones of [14]. However, unlike the first two sets, we cannot reproduce the $S_{11}$ wave of the $\pi N \to \pi N$ analysis [15], also below the $\eta N$ threshold, unless a relatively large contribution of the $S_{11}(1650)$ resonance is included.

The electromagnetic vertex of the amplitude $t_{\gamma p}$ for the elementary process $\gamma p \to \eta p$ was fixed by fitting the corresponding data of [16]. Below the $\eta N$ threshold, we have required that the pion production amplitude $\gamma p \to S_{11}(1535) \to \pi N$, taken from [17], is reasonably well described as presented in [13]. For the neutron amplitude we have used the relation

$$\text{compare our results with the work of Sibirtsev et al. [3–5]. In Sect. III we point out principal disagreements between the results of [3–5] and ours, which cannot be explained by the differences of the model ingredients. Finally, in Sect. IV we analyze the possibility of extracting the strength of the $\eta N$ interaction in incoherent $\eta$ photoproduction on the deuteron.
\[ t_{\gamma n} = -0.82 t_{\gamma p}, \]  

which is consistent with the experimental value \( \sigma_{\gamma n} = 0.67 \sigma_{\gamma p} \) extracted from \( \eta \) photoproduction on very light nuclei [18,19].

With respect to the \( NN \)-interaction, only the most important \( ^1S_0 \) state is taken into account, since, as was shown in [1], the contribution of the triplet state is insignificant, primarily due to the isovector nature of the electromagnetic excitation \( \gamma N \rightarrow S_1 \). For the \( ^3S_0 \) state we have used the separable representation BEST3 [20] for the Bonn potential. The deuteron bound state wave function was calculated using the corresponding BEST4 version for the \( ^3S_1 - ^3D_1 \) states.

The three-body integral equations were solved only for the lowest \( S \)-wave three-body configuration where the orbital angular momentum \( l = 0 \) in the two-body subsystems is coupled with angular momentum \( l = 0 \) of the third particle with respect to the pair. The remaining partial waves were treated perturbatively up to the first order in the \( S \)-wave t-matrices of \( NN \) and \( \eta N \) scattering. This approximation is well justified by the strong \( S \)-wave dominance in the \( NN \) and \( \eta N \) low energy interactions. As was shown in [2], it is the lowest \( S \)-wave three-body state, that is very sensitive to the higher-order scattering contributions, whereas the higher partial waves are well approximated by the first order terms. As already indicated above, perturbative calculation, where only pairwise \( NN \) and \( \eta N \) FSI’s are taken into account in all partial waves is referred to as first-order approximation (FOA).

The significance of FSI is demonstrated in Fig. 1 for the total cross section and in Fig. 2 for the \( \eta \) angular distribution. In order to appreciate the importance of a realistic treatment of the \( NN \) sector, the results should be compared with those presented in Figs. 8 and 9 of [2], obtained by means of the Yamaguchi potential and using a pure \( S \)-wave deuteron. As expected, for a realistic \( NN \)-interaction the FSI effect becomes smaller. The obvious reason is the \( NN \)-repulsive core which weakens the attraction in the final \( \eta NN \) system at small relative distances. At the same time the “three-body” effect remains important. For instance, at \( E_\gamma = 635 \text{ MeV} \) the three-body treatment enhances the first-order result by about a factor two. Also the steep rise right above threshold is a characteristic feature of the three-body approach which is not born out in FOA.

In order to compare our results with the inclusive measurement [11] we calculated in addition to the break-up channel the coherent cross section \( \gamma d \rightarrow \eta d \). In analogy with the incoherent process, the calculation was performed within the quasiparticle formalism of the three-body problem, as described in [2]. In this case we neglect the small contribution of the deuteron \( D \)-state, which contributes only 1% in the IA cross section. Since the coherent cross section is proportional to the modulus of the isoscalar part \( g_s \) of the transition amplitude, we have fixed it according to the relation \( \alpha = |g_s|/|g_{\gamma p \rightarrow S_{11}}| = 0.25 \) as found in [21]. This value is also close to \( \alpha = 0.22 \) of [12].

The corresponding total cross section and one angular distribution are plotted in Fig. 3. As one notices, the effect of FSI is also very pronounced. It is not, however, as large as obtained in [2], where we have used a more attractive \( \eta N \) interaction with \( a_{\eta N} = (0.75 + i 0.27) \) fm.

Adding the coherent contribution we obtain the inclusive cross section shown in Figs. 1 and 2 by solid curves. We notice that although the three-body calculation leads to a sizable improvement of the theoretical prediction just above threshold, a quantitative agreement with the experimental results has not been reached yet, the theory being too low. Moreover, above \( E_\gamma = 650 \text{ MeV} \) the inclusion of higher-order terms in the lowest partial wave acts in the opposite direction by decreasing the cross section of the FOA. In this region the three-body model exhibits an even larger deviation from the data than the FOA.

The slight disagreement with the experimental results points apparently to the fact, that the mechanism of the \( \eta \) photoproduction is more complicated and some of its important details are not properly accounted for by our calculation. In this connection, we would like to make a few comments concerning possible ways of improving the theoretical treatment. The first relates to the off-shell behavior of the \( \eta N \) scattering matrix, which may be much more complex as given by the vertices \( g_\eta(q) \) in (4). In particular, it was already noted in [22] that the simple Hulthén form factors may strongly overestimate the short-range \( \eta NN \) interaction, since the resulting separable \( \eta N \) potential is probably too attractive near the origin. This is hardly important for low-energy \( \eta d \) scattering but it is relevant for \( \eta \) photoproduction which in general is quite sensitive to the \( \eta N \) wave function at small distances. Clearly, a more realistic description of this short-range behavior would require a much more thorough treatment of the \( \eta N \) interaction, where the resonance excitation is considered microscopically with respect to the internal dynamics of hadrons.

Another comment is related to the explicit coupling to \( \pi NN \) states. It was neglected in the present calculation (except for virtual \( \pi N \) decays in the \( S_{11} \) propagator), since a correct inclusion of a pion would require substantial refinements of our three-body treatment, in particular the insertion of a variety of resonances which are excited in \( \pi N \)-collisions. If the \( \pi NN \) states are properly taken into account, then among other factors the \( \eta \) photoproduction can proceed according to the two-step scheme \( \gamma N \rightarrow \pi N \rightarrow \eta N \), where a pion, being produced by the photon, is subsequently rescattered into an \( \eta \) by the other nucleon. According to the results of [21,23], the contribution of the intermediate pion depends strongly on the role of large momentum transfers in the reaction mechanism. For example, whereas \( \pi NN \) states provide only a small fraction of the \( \eta d \) scattering cross section [2,23], they contribute
rather sizeably to coherent $\eta$ photoproduction on the deuteron [21]. It was shown in [2], where only the $S_{11}(1535)$ resonance was included into the $\gamma N \rightarrow \pi N$ amplitude, that the $\gamma d \rightarrow \eta N$ cross section is insignificantly affected if the pion rescattering mechanism is considered. However, it may well be that inclusion of other resonances into the pion photoproduction amplitude can improve the agreement between our calculation and the data [11].

Finally, it must be noted that the relation (6), fixing the neutron amplitude, is only the simplest variant matching the required relation between the elementary cross sections. Namely, as was pointed out in [21], it does not account for a possible relative phase between $t_{\gamma p}$ and $t_{\gamma n}$. This phase is predicted, e.g., if the corresponding electromagnetic vertices are extracted by fitting different isotopic channels of pion photoproduction [12,21]. Most likely, this fact is insignificant in the region of large photon energies, where the amplitudes $t_{\gamma p}$ and $t_{\gamma n}$ are added incoherently because of large relative momenta between the particles. At the same time, near threshold the interference between $t_{\gamma p}$ and $t_{\gamma n}$ increases due to a reduction of the available phase space, which requires a more sophisticated treatment of the isotopic structure of the elementary photoproduction operator.

## III. COMPARISON WITH OTHER WORK

Now we turn to the comparison of our results to those obtained by Sibirtsev et al. in a series of papers [3–5] where the FSI effects were studied within the first-order approximation. First of all, they confirmed our previous conclusion [1] with respect to the fundamental role of the final state interaction in the near threshold region. But moreover, they have claimed that already the incoherent reaction in FOA for the final state without inclusion of the coherent channel provides a good description of the experimental cross section for $\gamma d \rightarrow \eta X$ [11] with a reasonable value for the $\eta N$ scattering length. This result is in variance with our own conclusion about the necessity of a three-body treatment of the final $\eta NN$ state, and the nonnegligible contribution of the coherent reaction. These contradicting conclusions seem to be especially surprising in view of the fact that the models used in [3–5] and in [1,2] differ only in nonessential details, such as, e.g., relativistic vs. nonrelativistic parametrizations of the $\eta$ production amplitude. Therefore, we would like to point out the principal discrepancies between our results and those of [3–5] which we were unable to trace back to model differences of the two calculations.

To this end, we show in Fig. 4 our first-order calculation of the total cross section, which reproduces our previous results obtained in [1]. Small deviations are due to different parametrization of the elementary production operator and the $\eta N$ scattering amplitude. In the left panel of Fig. 5 we compare our impulse approximation (IA) with the corresponding results of Fig. 1 in [5] and those presented in [24]. The right panel shows the ratio of the IA of [5] to our IA. One readily notices that our calculation is in reasonable agreement with that of [24] and the difference of our result to the one of [24] indicates the model dependence to be expected from different parametrization of the elementary photoproduction amplitude and the deuteron wave function. But both IAs are far below the prediction of [5]. For example, at the energy $E_\gamma = 635$ MeV the latter cross section is by about a factor 3.6 larger than ours (see right panel in Fig. 5). This disagreement is especially surprising since the IA is quite insensitive to the model ingredients. Namely, if one uses a realistic deuteron wave function in conjunction with an elementary production operator fitted to the single nucleon data, than the $\gamma d \rightarrow \eta N$ cross section is fixed almost unambiguously. A little freedom associated with the choice of the invariant energy $W_{\gamma N}$ of the active $\gamma N$ subsystem does not play a role at all, since in the present calculation and in [5] this energy is taken according to the same prescription (compare the formulas (33) of [1] and (2) of [5]). In other words, a model dependence cannot explain such a strong difference between the IA results.

Turning now to the role of FSI we also find quite different effects. Namely in [5] the FSI effect decreases rapidly above the threshold and vanishes completely at $E_\gamma = 680$ MeV, so that above this point the total cross section is determined exclusively by the IA, whereas according to our findings the FSI contribution does not exactly vanish at all energies examined here. In more detail, the enhancement of the cross section due to the $NN$ interaction is reduced asymptotically to about 1.5% at $E_\gamma = 780$ MeV (see Fig. 4). For the $\eta N$ interaction we obtain an enhancement of about 20% at $E_\gamma = 635$ MeV. This effect is visibly smaller than the 45% given in Fig. 6 of [4] for the same energy [25]. This disagreement cannot be traced back to the strength of the $\eta N$ interaction because the $\eta N$ scattering length $a_{\eta N} = (0.42 \pm 0.34)$ fm used in [4] is even slightly smaller than our value given in (5). With increasing energy the influence of the $\eta N$ FSI changes sign and results, according to our calculation, in a slight reduction of about 4% of the cross section in the resonance peak (see right panel of Fig. 4). The latter effect must originate from a relatively strong absorption of $\eta$-mesons through the rescattering into pions, which is expected to be most pronounced in the resonance region. The same smearing of the resonance peak due to the strong inelasticity of the $\eta N$ interaction is observed also in heavier nuclei [26] where it manifests itself naturally much stronger. This energy dependence of the $\eta N$ FSI is not supported by the results presented in [4,5] where rescattering effects do not play any role above $E_\gamma = 680$ MeV. We would like to note that this discrepancy must be explained before any conclusion about the understanding of $\eta$...
photoproduction on the deuteron is drawn.

IV. THE ROLE OF THE $\eta N$ INTERACTION

In this last section we address the question to what extent the $\eta N$ interaction can be “extracted” from incoherent $\eta$ photoproduction on the deuteron. The idea to obtain the information on the $\eta N$ low-energy scattering parameters from this reaction was explored in [4,5]. In particular, it was found that the cross section is quite sensitive to the $\eta N$ interaction strength, making it possible to study $\eta N$ scattering by analyzing the observed single particle spectra or angular distributions.

In our opinion, such a method very likely will meet serious difficulties, and we support our scepticism by several numerical results presented below. Firstly, and this is the crucial point for the following conclusions, we do not observe any strong sensitivity of the cross section to the $\eta N$ interaction strength as is claimed in [4]. The effect on the $\eta$ angular distribution of varying the $\eta N$ parameters is demonstrated in the left panel of Fig. 6. The calculations are performed within the three-body approach as outlined above. In each case the e.m. vertex was adjusted to reproduce the elementary experimental cross section. Comparing the present result with Fig. 4 of [5] one sees that in our case the dependence on the $\eta N$ scattering length is much less pronounced. Thus, we have to conclude that the experimental discrimination between different $\eta N$ models would require extremely precise measurements.

Furthermore, such a weak sensitivity makes the extraction of the $\eta N$ scattering parameters very model dependent. The reason has to do with the off-shell behavior of the $\gamma N \rightarrow \eta N$ amplitude. This is especially critical for the region below the free $\eta N$ threshold which comes into play, when contributions beyond the IA are considered. For example, the results presented above are obtained within the so-called “spectator on-shell” choice for the invariant energy $W_{\gamma N}$ of the $\gamma N$ system, which is natural for the nonrelativistic three-body theory. In this case, as one integrates over the spectator nucleon momentum in the deuteron, $W_{\gamma N}$ covers the range $[-\infty, W - M_N]$ where $W$ is the invariant energy of the $\eta NN$ system. Although the uncertainty of the $\gamma N \rightarrow \eta N$ subthreshold behavior is not very crucial for the incoherent reaction, it makes the method of precise determination of $\eta N$ parameters much more ambiguous, than was presented in [4]. This difficulty is quite general when one tries to determine the contribution of an individual diagram to the whole amplitude.

In this connection we would like to recall the Migdal-Watson theory [27,28] which makes it possible to study the two-body interaction without regard to the particular way of embedding this interaction into the reaction amplitude. According to this theory the low-energy parameters of two-body scattering can in principle be identified by fitting the energy spectrum of the third particle close to the maximum energy value. In order to illustrate the practical applicability of this method to the $\eta N$ interaction, we show in the right panel of Fig. 6 the proton spectrum at a fixed angle $\theta_p = 18^\circ$. The curves were obtained within the three-body approach using the same three different parametrizations of the $\eta N$ sector as in the left panel. Only the upper end of the spectrum, corresponding to low relative $\eta n$ energies, is shown, since only this part is relevant. According to the Migdal-Watson theory the behavior of the proton energy distribution in this region reflects the energy dependence of the elementary $\eta n$ cross section, disregarding the terms which depend weakly on the $\eta n$ relative energy. To eliminate the effect of the kinematical boundary, the spectrum is divided by $(T_p^0 - T_p)^{1/2}$ where $T_p^0$ is the maximum proton energy. As one notes, also in this case the experimental discrimination between the $\eta N$ models seems to meet the same difficulties. Namely, the form of the spectrum is not strongly affected by the $\eta N$ interaction. The reason is that the form of the $\eta N$ elastic scattering cross section, shown in the insert in the right panel, does not sizeably vary with the value of $a_{\eta N}$.

V. CONCLUSION

The role of final state interaction in incoherent $\eta$ photoproduction on a deuteron has been investigated using a three-body model for treating the interaction in the final $\eta NN$ system. In contrast to our previous work [2] the present results were obtained by using a realistic $NN$ interaction based on the separable representation of the Bonn potential. As a summary, we would like to draw the following conclusions:

(i) As may be expected, a realistic treatment of the $NN$ sector reduces the FSI effect compared to the use of a simple Yamaguchi potential. At the same time, the influence of higher-order rescattering in the final state remains essential and must be taken into account close to the threshold.

(ii) Our three-body calculation underestimate the data [11] slightly in the energy region up to $E_\gamma = 780$ MeV. One open point, which remains to be investigated in this connection, is the short-range part of the $\eta NN$ wave function. The latter can be quite sensitive, e.g., to the contribution of the $\pi NN$ configuration which was omitted here.
(iii) There exists a principal discrepancy between our results and the ones of Sibirtsev et al. [3–5]. Most disturbing is the fact that already the simple impulse approximation of [3–5] where the model ambiguities should be small, exhibits a strong disagreement with other authors [1,24]. In this connection we would like to emphasize, that the statement, that the reaction $\gamma d \to \eta np$ is well understood within existing theoretical approaches is premature, even if a seemingly good description of the data in [5] is achieved. Further theoretical work is certainly needed.

(iv) We do not find a strong sensitivity of the $\gamma d \to \eta np$ cross section to the $\eta N$ interaction strength as was claimed in [4]. Our calculation shows that even if the off-shell uncertainty of the $\gamma N \to \eta N$ amplitude is disregarded, quite a weak dependence of the results on the $\eta N$ interaction parameters renders a precise determination of the $\eta N$ interaction practically impossible.

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| $a_{\eta N}$ [fm] | $g_{\eta}$ | $\beta_{\eta}$ [MeV] | $g_{\eta}/\sqrt{3}$ | $\beta_{\eta}$ [MeV] | $M_0$ [MeV] |
|-----------------|----------|---------------------|-----------------|-----------------|-------------|
| 0.25+0.16       | 1.43     | 654.3               | 1.57            | 379.2           | 1563        |
| 0.50+0.32       | 2.00     | 694.6               | 1.45            | 404.5           | 1598        |
| 0.75+0.27       | 2.04     | 1282.6              | 0.55            | 888.0           | 1673        |
FIG. 1. Total cross section for the reaction $\gamma d \rightarrow \eta np$. The near-threshold region is shown separately in the right panel. The dotted, dashed and dash-dotted lines correspond to the impulse approximation (IA), first-order calculation (FOA) and three-body model for the final $\eta NN$ state. The inclusive cross section $\gamma d \rightarrow \eta X$, obtained by adding the coherent cross section (see Fig.3) is shown by the solid curve. The inclusive data are taken from [11]. The dash-dotted and solid curves are indistinguishable in the left panel.

FIG. 2. Differential cross section for $\gamma d \rightarrow \eta np$. Notation as in Fig. 1.
FIG. 3. Total and differential cross sections for coherent eta photoproduction on the deuteron. The IA and three-body results are shown by the dotted and the solid curves, respectively. The data are from [18] (open triangles) and [29] (filled triangles).

FIG. 4. Total cross section of $\gamma d \rightarrow \eta np$ for two different regions of the photon energy. The dotted curve represents the IA. Successive addition of $NN$ and $\eta N$ rescattering in FOA is shown by the dashed and the solid curves, respectively. The $\gamma d \rightarrow \eta X$ data are from [11].
FIG. 5. Left panel: comparison of the IA for the total cross section from [24] (dotted curve), [5] (dashed curve), and for the present model (solid curve). The $\gamma d \rightarrow \eta X$ data are from [11]. Right panel: ratio of the IA cross section of [5] to our calculation.

FIG. 6. Left panel: The $\eta$ angular distribution calculated with different parametrization of the $\eta N$ scattering amplitude. Notation of the curves: dotted: $a_{\eta N} = (0.25+i0.16)$ fm; full: $a_{\eta N} = (0.50+i0.32)$ fm; dashed: $a_{\eta N} = (0.75+i0.27)$ fm. The $\gamma d \rightarrow \eta X$ data are taken from [11]. Right panel: Reduced proton energy spectrum, calculated at a fixed angle. Shown is the region in the vicinity of the maximum energy. The internal $\eta n$ kinetic energy $E_{\eta N}$ is indicated at the top x-axis. Insert: $\eta N$ elastic scattering cross section as a function of the energy $E_{\eta N}$. Notation of the curves as in the left panel.