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Nucleon-nucleon wave function with short-range nodes and high-energy deuteron photodisintegration

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We review a concept of the Moscow potential of the NN interaction. On the basis of this concept, we derive by quantum inversion optical partial potentials from the modern partial-wave analysis data and deuteron properties. Point-form relativistic quantum mechanics is applied to the two-body deuteron photodisintegration. Calculations of the cross-section angular distributions cover photon energies between 1.1 and 2.5 GeV. Good agreement between our theory and recent experimental data confirms the concept of deep attractive Moscow potential with forbidden S and P states.

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I. INTRODUCTION

Opportunities to observe manifestations of quark degrees of freedom in nuclear reactions at intermediate energies have attracted the attention of the scientific community for a long time. It was noted in Ref. [1] that the most suitable subject of research here is the deuteron, because it is the simplest nucleus for which secondary rescattering has little effect on the primary process.

The deuteron photodisintegration at photon energies of $\simeq$2 GeV generates great interest among experimentalists [1–4] and theoreticians [5–9] with the main emphasis on the properties of the NN system which are beyond the scope of realistic mesonic NN potentials [6] and can be interpreted within quark concepts [5,7]. First, it was shown in papers by the Khar'kov group [6] that starting from mesonic potentials, it is possible to explain the $d\gamma \rightarrow np$ data at energies $E_\gamma > 1$ GeV only if the electromagnetic part of the theory is revised and instead of the ordinary nucleon electromagnetic form factors the essentially different ones are used with poles of third order. Second, the phenomenological theory of Reggeon poles was taken as the basis in Ref. [5] with the selection of dominant poles according to the quark string model [10]. Free parameters of these theories make it possible to describe the experimental data reasonably well. Third, also giving reasonable results, the hard rescattering model was developed [7] within a semiempirical approach, when the photon is absorbed by a quark of one of the nucleons and then the hard rescattering of this quark by another nucleon takes place. The wave function amplitude of the final np state with large relative momentum is evaluated empirically by extrapolation of the corresponding np-scattering experimental data.

In this paper, we use point-form (PF) relativistic quantum mechanics (RQM) to treat the deuteron photodisintegration in a Poincaré-invariant way. Modern development of the RQM and an exhaustive bibliography are presented in the review by Keister and Polyzou [11]. The PF is one of the three forms proposed by Dirac [12]. The other two are the front form and instant form. These forms are associated with the different possibilities for putting interactions in generators of the Poincaré group. All the forms are unitary equivalent [13], but each has certain advantages. Most of the calculations in nuclear physics have been performed in the instant and front forms. Only in recent years have important simplifying features of the PF been realized. These features are connected with the fact that in the PF all the generators of the homogeneous Lorentz group are free of interactions. Thus only in the PF does the spectator (impulse) approximation (SA) preserve its spectator character in any reference frame [14,15]. For an electromagnetic NN process, the SA implies that the NN interaction does not affect the photon-nucleon interaction, and therefore the sum of the one-particle electromagnetic current operators may be taken as an electromagnetic current operator for the system of interacting nucleons. It is supposed that the SA may be valid when the process is quick due to the large momentum transfer. General covariant PF expressions for the electromagnetic current operator for composite systems are given in Refs. [15,16]. The PF SA was applied to calculate form factors of various composite particles [17–19] with reasonable results. Our calculation of the proton-proton bremsstrahlung [20] showed that the PF SA violates the continuity equation for the NN current operator, but the violation is relatively small for the considered kinematics.

In this paper, we show that recent deuteron photodisintegration data at $E_\gamma = 1.5$–2.5 GeV [4] confirm the Moscow NN potential model [21] characterized by deep attractive partial potentials with forbidden S and P states. In this study, the Moscow partial potentials are reconstructed from the NN partial-wave analysis (PWA) data within the energy range $0 \leqslant E_{\text{lab}} \leqslant 3$ GeV [22]. This reconstruction is based on our approach to the inverse-scattering problem for optical potentials [23].

The plan of the paper is as follows. In Sec. II, we review a concept of the Moscow potential (MP) of the NN interaction. In Sec. III, we present the optical Moscow-type NN potential derived by quantum inversion [23] within the relativistic quasipotential approach [11,24]. We show that the modern PWA data of NN scattering [22] are compatible with the concept of the MP. In Sec. IV, the formalism of PF

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RQM [11,15,20] is applied to the high-energy energy deuteron photodisintegration. Results and future prospects are discussed in Sec. V. In Appendix A, we give necessary details of the calculation techniques. In Appendix B, in the PF SA we derive an expression for the momentum Q_N transferred to the nucleon and show that Q_N is not the same as the momentum transfer seen by the deuteron. The expression is a generalization of a similar expression for the elastic electron-deuteron scattering [19].

II. POTENTIALS WITH FORBIDDEN STATES IN NUCLEAR PHYSICS

In describing systems of composite X particles consisting of some y particles, it is a common approach to exclude explicit degrees of freedom of y particles. In the simplest case of the XX system, the microscopic Hamiltonian that includes all possible pair yy interactions is substituted by an effective Hamiltonian (by sum of X particle kinetic energy terms and of an effective XX potential). The common requirement is that the effective Hamiltonian would give for the XX system the same spectrum and the same corresponding relative motion wave functions as the initial microscopic Hamiltonian. In some cases, the effective Hamiltonian has redundant eigenvalues and eigenstates, which must be disregarded. These eigenstates are called forbidden states, and the effective XX potential is called then "the potential with forbidden states."

For instance, in the oscillator shell model of the potential theory of α-α scattering [25], the antisymmetric wave function of the ^{8}Be nucleus ground state (eight-nucleon configuration s^4p^4 and orbital permutation symmetry $[f]_x = [44]$) being projected onto an α-α channel results in a $4S$-wave relative motion wave function (see our review [26]). This wave function accumulates all four oscillation quanta of the system and has two nodes. Momentum distributions corresponding to such wave functions were investigated in quasielastic knock-out of α particles from p-shell nuclei by intermediate energy photons [27]. The $0S$ and $2S$ states of α-α relative motion are forbidden as far as they correspond to the lower s^8 and s^6p^2 eight-nucleon configurations, respectively, which are forbidden by the Pauli principle. Based on these considerations, a concept of the deep attractive α-α potential with $0S$, $2S$, and $2D$ forbidden bound eigenstates was elaborated [25]. According to the concept, there is no repulsive core in the α-α interaction, and α particles can penetrate into each other. Forbidden bound eigenstates take the lowest energy levels. Unforbidden eigenstates (including scattering ones) being orthogonal to the forbidden eigenstates have a nodal structure at short range. For instance, the S-wave relative motion wave function has two nodes in the region of α-α overlap. This model is substantiated by the phase shift analysis based on the generalized Levinson theorem (GLT) [28]. For example, the S-wave phase shift of α-α scattering equals 360° at zero energy, rises up to 540° at the energy slightly above the low-lying $4S$ resonance and then runs down with increasing energy within the broad energy range up to $E_{\mathrm{lab}} \simeq 200$ MeV, where the phase shift approaches the asymptotic region of small values and becomes negative due to absorption [29]. Such a picture of the S-wave phase shift behavior was confirmed by experiments performed in a broad energy range [30], while D-wave phase shift behavior shows one forbidden state. Phase shifts of higher waves do not show forbidden states (see Ref. [29] for further details).

For the NN system, the concept of the deep attractive NN potential with forbidden states appeared in 1975 [21] when we analyzed the pp-scattering phase shift data extended at the time up to $E_{\mathrm{lab}} \cong 6$ GeV. It was shown that the singlet S-wave phase shift data with an extended gap between low- and high-energy groups of data could be interpolated by a smooth curve if the empirical low-energy group were raised 180°. This interpolation demonstrates a decrease of the S-wave phase shift in the broad energy range from zero up to $E_{\mathrm{lab}} \gtrsim 5$ GeV as a manifestation of the GLT. The high-energy part ($E_{\mathrm{lab}} \simeq 3$–6 GeV) of the interpolation for the S wave remains in the asymptotic region of small values, corresponding to the Born approximation. The energy dependence of the singlet D-wave phase shift is smooth, and there is no need to raise the initial values. Calculation showed [21] that results of this analysis are described by a deep attractive NN potential with one forbidden bound S-wave state. The forbidden state has a wave function without a node. As a result, the $^1S_0$-wave scattering wave function has a short-range nodal structure instead of short-range suppression specific to a repulsive core potential (RCP). After that, a preliminary attempt was made within the concept of MP [31] to reconstruct NN potentials for the lowest partial waves (S and P) from data of the pp and pn PWA extended at the time to intermediate energies.

At the same time, the quark microscopic foundation of the MP remains the principal problem. Unlike the nuclear shell-model picture of the α-α interaction, the lowest quark configuration s^6 is not forbidden by the Pauli principle, and the corresponding $0S$-wave state of relative NN motion is not forbidden either. Microscopic quark investigations of the last two decades with various kinds of qq interactions have resulted in the following short-range properties of the NN system [32]. There is a strong mixing of different six-quark configurations in the overlap region of two nucleons. For the S-wave states, the leading configurations are s^6 and s^4p^2 with comparable weights and destructive interference. This destructive interference leads to strong short-range suppression of the NN wave function. The suppression is described effectively by an RCP [33]. The $s^4p^2[42]_x$ configuration introduced in our papers [21] and corresponding to the $2S$ state of relative NN motion (i.e., to the MP) would dominate, for instance, in the case of strong instanton-induced quark-quark interaction, but this interaction is not strong enough [34]. Further investigations [35,36] showed the existence of a source for strengthening the s^4p^2 configuration. Namely, if coupling of the NN, $\Delta\Delta$, and hidden color CC channels is taken into account within the resonating group method, then the symmetry structure of the highly dominant six-quark configuration s^4p^2 implies the existence of a node in the S-wave relative motion wave function at short distances. Such nodes are specific to the MP. In the same manner, microscopic qq interaction may give a short-range node in a P wave of a relative NN motion wave function (in the case of the dominant six-quark configuration $s^3p^3[33]_x$).

In summary, the question of which type of the potential (MP or RCP) would be equivalent to the short-range quark microscopic picture of the NN interaction is highly controversial. For any RCP, a phase equivalent supersymmetric partner with forbidden states (i.e., an MP) may be constructed [37]. Therefore, these potentials are indistinguishable for the NN PWA. Specific to the MP, the appearance of short-range nodes in S- and P-wave relative motion wave functions is a result of complicated six-quark dynamics which is yet to be clarified. The nodal behavior of the MP wave function means that the wave function is not suppressed at short range as in case of an RCP. Thus the MP produces a high-momentum component richer than an RCP. This high-momentum component may be seen in electromagnetic reactions with two nucleons. Reference [38] showed that the available MP produces a too rich high-momentum component in contradiction with the deuteron electromagnetic form factors. Thus we use the latest high-energy PWA data to refine the short-range part of the MP. In our previous papers [20,39], we showed that the hard $pp \to pp\gamma$ bremsstrahlung at moderate energies ($E_{\rm lab} \simeq 500$ MeV) is critical to the kind of potential (MP versus RCP). The available experimental data at smaller energy of $E_{\rm lab} = 280$ MeV [40] give only a preliminary indication of MP validity [20]. Our present paper strengthens this line of phenomenological research using modern deuteron photodisintegration data.

III. RELATIVISTIC OPTICAL NN POTENTIAL

We apply the method of inversion [23] to the analysis of NN data up to energies at which relativistic effects are essential. We take into account these effects in the frames of the RQM [11,15]. A system of two particles is described by the wave function, which is an eigenfunction of the mass operator $\hat{M}$. In this case, we may represent this wave function as a product of the external and internal wave functions. The internal wave function $|\chi\rangle$ is also an eigenfunction of the mass operator, and for a system of two nucleons with masses $m_1 = m_2 = m$ it satisfies the equation

$$\hat{M}|\chi\rangle \equiv [2\sqrt{\mathbf{q}^2 + m^2} + V_{\rm int}]|\chi\rangle = M|\chi\rangle, \quad (1)$$

where $V_{\rm int}$ is an operator commuting with the full angular momentum operator and acting only through internal variables (spins and relative momentum), $\mathbf{q}$ is a momentum operator of one of the particles in the center-of-mass frame (relative momentum). Rearrangement of Eq. (1) gives

$$[\mathbf{q}^2 + mV]\chi = q^2\chi, \quad (2)$$

where V acts like $V_{\rm int}$ only through internal variables, and

$$q^2 = \frac{M^2}{4} - 2m^2. \quad (3)$$

Equation (2) is identical in form to the Schrödinger equation. The formally same equation may be deduced as a truncation of the quantum field dynamics [24]. The quasicoordinate representation corresponds to the realization $\mathbf{q} = -i\frac{\partial}{\partial \mathbf{r}}$, $V = V(\mathbf{r})$.

We applied the method of inversion [23] to reconstruction of the nucleon-nucleon partial potentials

$$V(r) = (1 + i\alpha)V^{(0)}(r), \quad (4)$$

for single waves, and

$$V(r) = \begin{pmatrix} (1+i\alpha_1)V_1^{(0)}(r) & (1+i\alpha_3)V_T^{(0)}(r) \\ (1+i\alpha_3)V_T^{(0)}(r) & (1+i\alpha_2)V_2^{(0)}(r) \end{pmatrix}, \quad (5)$$

for coupled waves, where $V^{(0)}(r)$ are energy-independent real, and inelasticity parameters α depend on energy. As input data for the reconstruction, we used modern PWA data (single-energy solutions) up to 1200 MeV for isoscalar states and up to 3 GeV for isovector states of the NN system [22]. The deuteron properties were taken from Ref. [41]. These data allow us to construct Moscow-type NN partial potentials sustaining forbidden bound states. These potentials describe part of the deuteron properties and the PWA data by the construction. According to the MP concept and the GLT, some phase shift data of Ref. [22] are raised 180°. Namely, the $^1S_0$-wave phase shift and all four $^{2S+1}P_J$-wave phase shifts are equal to 180° at zero energy; the $^3S_1$-wave phase shift is equal to 360° at zero energy. The mixing parameters ϵ_1 and ϵ_2 of the MP differ in sign from those of a traditional RCP. All phase shifts for higher waves (for $L \geqslant 2$) are "small," they have zero values at zero energy. According to our model, we have fitted free parameters of the inversion solutions to get nodes at $r \simeq 0.5$ fm in S and P waves and to make central parts of the potentials close to each other and to the Gaussian shape. The energies of forbidden states are in the range 300–750 MeV.

Our calculations show that the final state interaction (FSI) in the S and P waves gives by far the largest contribution to the deuteron photodisintegration cross section compared with the FSI in other waves, so we present results of inversion only for these waves and for waves coupled to them. Some of the results presented in Figs. 1–4 (for $^1S_0$ and $^3SD_1$ waves) we presented earlier in Ref. [23].

The reconstructed potentials $V^{(0)}(r)$ are displayed in Fig. 1. The inelasticity multipliers α are displayed in Fig. 2. Figure 3 displays the reproduction of the corresponding

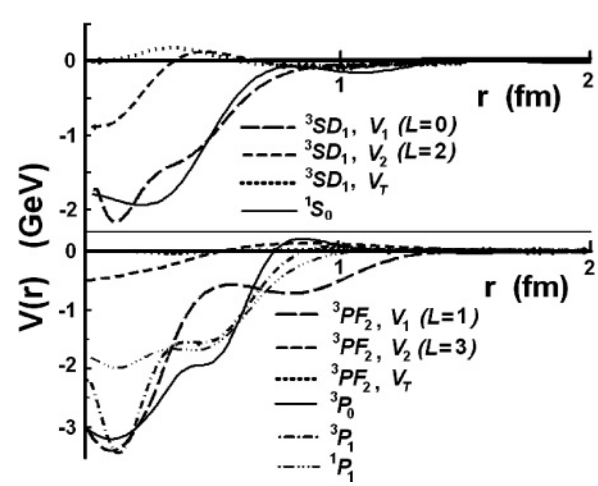

FIG. 1. Reconstructed partial potentials for lower orbital momentum (single and coupled channels).

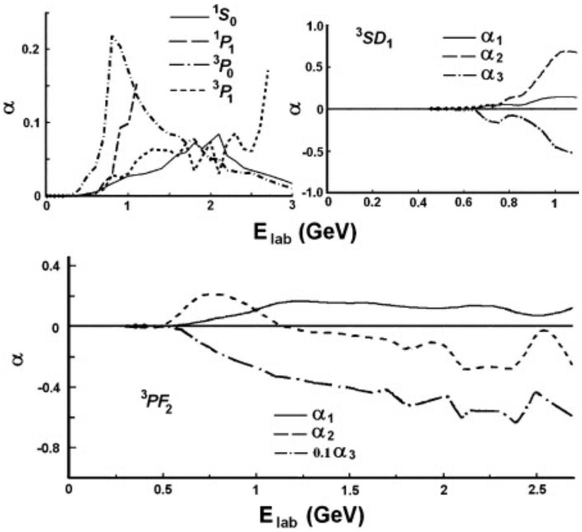

FIG. 2. Reconstructed inelasticity multipliers α for the potentials presented in Fig. 1.

phase shifts and mixing parameters. In Fig. 4, the description of inelasticity parameters is shown. All the P-wave phase shifts are positive according to the GLT. The large difference between $^3P_0$-wave and $^3P_2$-wave phase shift curves reflects a large spin-orbital interaction which is attractive for the $^3P_2$ wave as we see. These features correspond to the general properties of the MP (its large positive gradient in the region $r < 1$ fm). It is interesting to learn from Fig. 3 that among the four lowest pp phase shifts, three of them ($^1S_0, {}^3P_0$, and $^3P_2$) correspond to the MP, but the experimental data within the energy range $E_{\rm lab} = 2$–3 GeV are contradictory for the $^3P_1$-wave phase shift. It would be important to refine the PWA data in this range using modern polarization data on pp scattering. The S- and D-state wave functions of deuteron are displayed in Fig. 5. There is a node in the S-wave function at $r \simeq 0.5$ fm, and both wave functions are not suppressed at short range in contrast with wave functions produced by an RCP. For continuum S- and P-wave functions, the node radii equal 0.5–0.9 fm at the considered energies. All potentials and inelasticity multipliers (α's) can be accessed via a link to the web site [42].

It should be pointed out that in nuclear matter calculations, the NN potentials should be used in the form

$$V^{\rm nucl}(r) = V^{(0)}(r) + \lambda\langle\chi_{S,L,J}|, \tag{6}$$

where operator $\langle\chi_{S,L,J}|$ projects onto the forbidden state $|\chi_{S,L,J}\rangle$, positive constant λ tends to infinity. The forbidden state $|\chi_{S,L,J}\rangle$ may be found from Eq. (2) by some numerical method as a bound state of the partial potential $V^{(0)}(r)$ (all bound states are forbidden except the deuteron one). Constant λ is a large number, such that its further increase does not change the calculation results. This procedure orthogonalizes the nuclear wave function to forbidden two-nucleon states. Thus, we exclude the unphysical collapse of nuclear matter.

IV. DEUTERON PHOTODISINTEGRATION IN POINT-FORM RELATIVISTIC QUANTUM MECHANICS

Formalism of the PF is considered in detail in Refs. [11,15], while general covariant PF expressions for the electromagnetic current operator for composite systems are given in Refs. [15,16]. Therefore, we give only the results necessary for our calculation, in the notation of Ref. [15]. We use the algorithm of Ref. [15] to calculate the matrix elements of the electromagnetic current operator. We applied this formalism to the $pp\gamma$ process [20]. A similar approach was applied to the elastic electron-deuteron scattering [19].

We consider the pn system and neglect the difference of neutron and proton masses ($m_1 = m_2 = m$). Let p_i be the four-momentum of nucleon i, $P \equiv (P^0, \mathbf{P}) = p_1 + p_2$ the system four-momentum, M the system mass, and $G = P/M$ the system four-velocity. The wave function of two particles with four-momentum P is expressed through a tensor product of external and internal parts

$$|P, \chi\rangle = U_{12}|P\rangle \otimes |\chi\rangle, \tag{7}$$

where the internal wave function $|\chi\rangle$ satisfies Eqs. (1) and (2). The operator

$$U_{12} = U_{12}(G, \mathbf{q}) = \prod_{i=1}^{2} D[\mathbf{s}_i; \alpha(p_i/m)^{-1}\alpha(G)\alpha(q_i/m)] \tag{8}$$

is the unitary operator from the "internal" Hilbert space to the Hilbert representation space of two-particle states [15]. $D[\mathbf{s}; u]$ is the representation operator of the group SU(2) corresponding to the element $u \in$ SU(2) for the representation with the generators $\mathbf{s}$. Action of $D[\mathbf{s}; u]$ and matrices α are defined in Appendix A; $s_i = 1/2$ is the spin of a nucleon. The momenta of the particles in their c.m. frame are

$$q_i = L[\alpha(G)]^{-1} p_i, \tag{9}$$

where $L[\alpha(G)]$ is the Lorentz transformation to the frame moving with four-velocity G ($L[\alpha(G)]^{-1}$ is the inverse transformation). It is easy to verify that $\mathbf{q}_1 = -\mathbf{q}_2$.

The external part of the wave function is defined as

$$\langle G|P'\rangle \equiv \frac{2}{M'} G'^0 \delta^3(\mathbf{G} - \mathbf{G}'), \tag{10}$$

with scalar product

$$\langle P''|P'\rangle = \int \frac{d^3\mathbf{G}}{2G^0} \langle P''|G\rangle\langle G|P'\rangle = 2\sqrt{M'^2 + \mathbf{P}'^2}\, \delta^3(\mathbf{P}'' - \mathbf{P}'), \tag{11}$$

where $G^0(\mathbf{G}) \equiv \sqrt{1 + \mathbf{G}^2}$. The internal part of the wave function $|\chi\rangle$ is characterized by momentum $\mathbf{q} = \mathbf{q}_1 = -\mathbf{q}_2$ of one of the particles in the c.m. frame. Interaction appears according to the Bakamjian-Thomas procedure $\hat{P} = \hat{G}\hat{M}$, where $\hat{M}$ is the sum of the free mass operator M and the interaction V; i.e., $\hat{M} = M + V_{\rm int}$ [compare with Eq. (1)]. The interaction operator acts only through internal variables. Operators $\hat{M}$, M, $V_{\rm int}$. and V commute with spin operator S (full angular momentum) and with four-velocity operator $\hat{G}$.

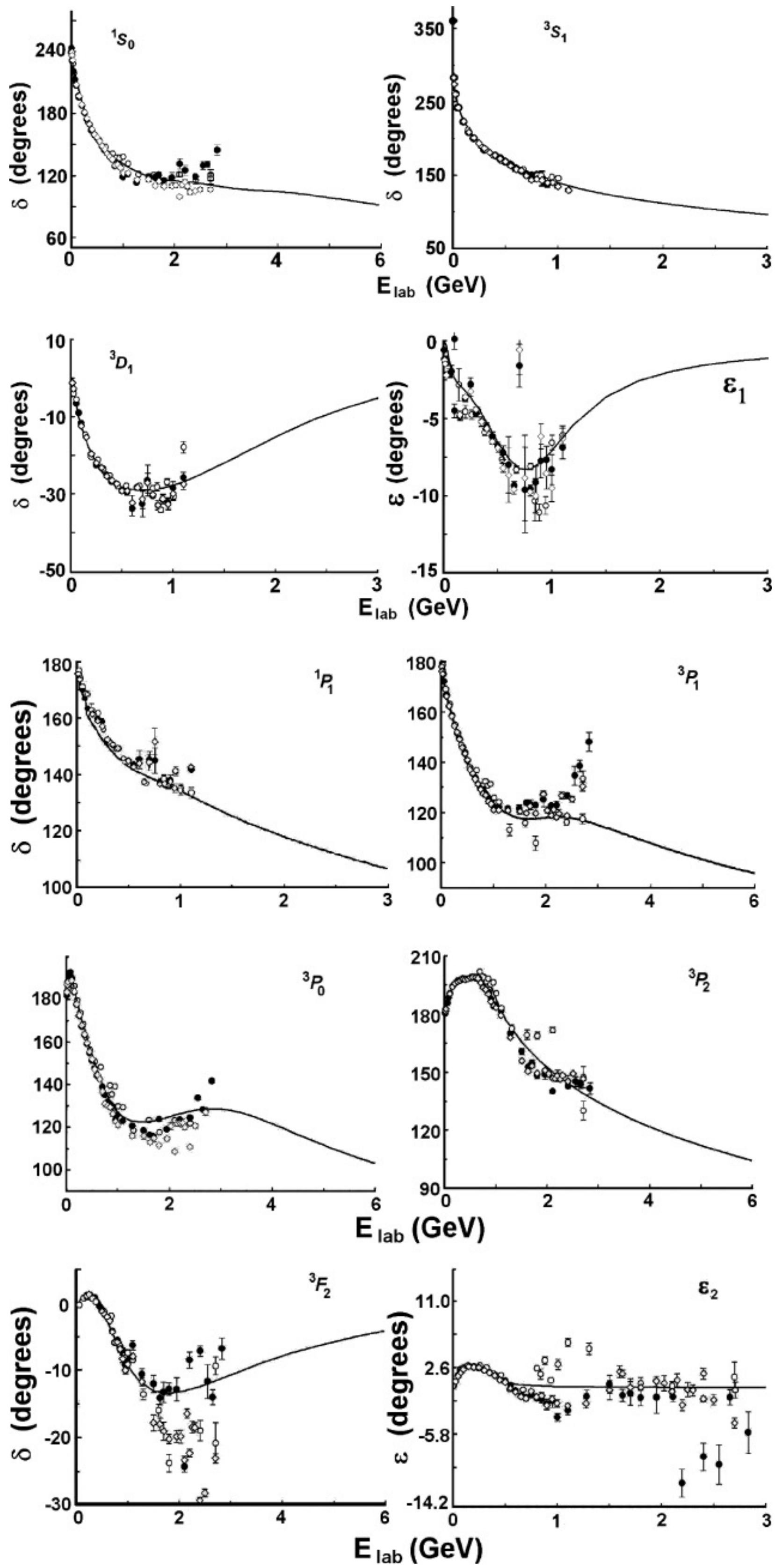

FIG. 3. Phase shifts and mixing parameters in the present optical model. PWA data are from Ref. [22]. For $^1S_0$ and $^{2S+1}P_J$ waves, the original data set from Ref. [22] is raised 180°. To leave the S matrix unchanged, we changed the sign of mixing parameters ε_1 and ε_2.

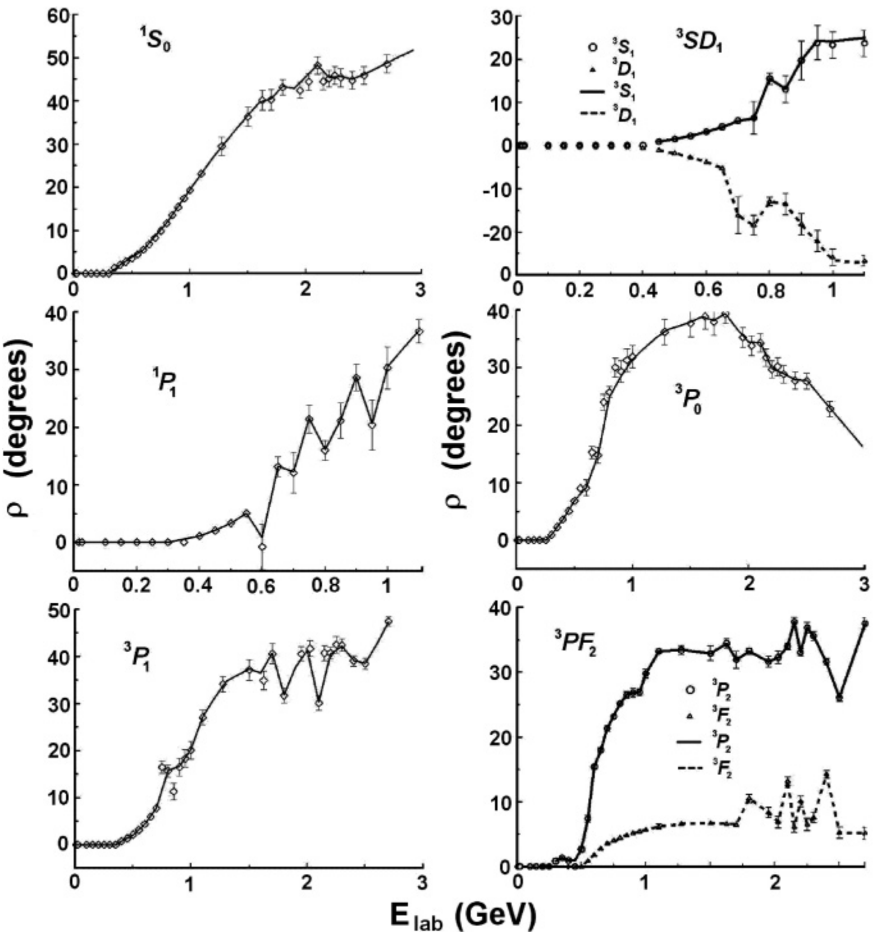

FIG. 4. Inelasticity parameters ρ in the present optical model. PWA data are from Ref. [22].

The interaction term is present in all components of total four-momentum. Generators of Lorentz boosts and generators of rotations are free of interaction. In the c.m. frame, the relative orbital angular momentum and spins are coupled together as in the nonrelativistic case. Moreover, most nonrelativistic scattering theory formal results are valid for our case of two particles [11].

The deuteron wave function $|P_i, \chi_i\rangle$ is normalized as

$$\langle P_i', \chi_i | P_i'', \chi_i\rangle = 2P_i^{0'}\delta^3(\mathbf{P}_i' - \mathbf{P}_i''). \quad (12)$$

For one-particle wave functions normalized in the same manner, the free two-particle states are normalized as

$$\begin{aligned}\langle P', \chi' | P'', \chi''\rangle &\equiv \langle p_1' | p_1''\rangle\langle p_2' | p_2''\rangle \delta_{\mu_1'\mu_1''}\delta_{\mu_2'\mu_2''} \\ &= 4w(\mathbf{p}_1')w(\mathbf{p}_2')\delta^3(\mathbf{p}_1'' - \mathbf{p}_1') \\ &\quad \times \delta^3(\mathbf{p}_2'' - \mathbf{p}_2')\delta_{\mu_1'\mu_1''}\delta_{\mu_2'\mu_2''} \\ &= 2W(\mathbf{P}')\delta^3(\mathbf{P}'' - \mathbf{P}')\frac{2w^2(\mathbf{q})}{M(\mathbf{q})} \\ &\quad \times \delta^3(\mathbf{q}'' - \mathbf{q}')\delta_{\mu_1'\mu_1''}\delta_{\mu_2'\mu_2''} \\ &= 2W(\mathbf{P}')\delta^3(\mathbf{P}'' - \mathbf{P}')\frac{M(\mathbf{q})}{2} \\ &\quad \times \delta^3(\mathbf{q}'' - \mathbf{q}')\delta_{\mu_1'\mu_1''}\delta_{\mu_2'\mu_2''}, \end{aligned} \quad (13)$$

where $w(\mathbf{p}) \equiv \sqrt{m^2 + \mathbf{p}^2}$, $M(\mathbf{q}) \equiv 2\sqrt{m^2 + \mathbf{q}^2}$, $W(\mathbf{P}) \equiv \sqrt{M^2 + \mathbf{P}^2}$, μ_i are spin projections in the c.m. frame. Multiplier $\frac{M(\mathbf{q})}{2}$ is a relativistic invariant, therefore we may normalize the internal part of the scattering state wave function in the nonrelativistic manner

$$\langle P', \chi' | P'', \chi''\rangle_{\text{n.r.}} = 2W(\mathbf{P}')\delta^3(\mathbf{P}'' - \mathbf{P}')\delta^3(\mathbf{q}'' - \mathbf{q}')\delta_{S'S''}\delta_{\mu'\mu''}, \quad (14)$$

where S and μ are full angular momentum and its projection in the c.m. frame.

The differential cross section for the $\gamma d \to np$ process is given by

$$\frac{d\sigma}{d\Omega} = \frac{q_f}{64\pi^2 M_f^2 k_c}|A_{if}|^2, \quad (15)$$

where q_f is the final asymptotic np relative momentum, k_c is the photon energy in the c.m. frame. The $d\gamma \to np$ amplitude A_{if} is defined in the same manner as the $pp\gamma$ amplitude used in Ref. [20], that is,

$$(2\pi)^4\delta^4(P_i + k - P_f)A_{if} = \sqrt{4\pi}\int d^4x \langle P_f, \chi_f | \times \varepsilon_\mu \hat{J}^\mu(x)|P_i, \chi_i\rangle e^{ikx}, \quad (16)$$

where P_i and P_f are initial and final four-momenta of the NN system correspondingly, and ε_μ is the photon polarization vector.

Following Ref. [15], we choose for calculation of the invariant amplitude A_{if} a special frame defined by the

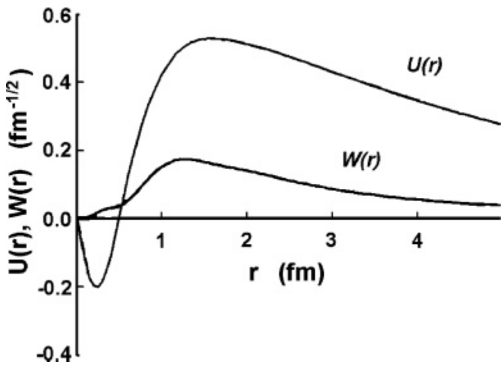

FIG. 5. Deuteron S- and D-wave functions for present version of the NN Moscow potential.

condition

$$\mathbf{G}_i + \mathbf{G}_f = 0, \tag{17}$$

where $G_i = P_i/M_i, G_f = P_f/M_f$ are four-velocities of initial and final NN c.m. frames, respectively ($\mathbf{G}_i$ and $\mathbf{G}_f$ are their three-vector parts). The initial mass M_i is the deuteron mass. The final mass M_f is the invariant mass of the final NN system. These masses are different because of absorption of a photon; therefore, the coordinate frame corresponding to Eq. (17) is not equivalent to the Breit frame where $\mathbf{P}_i + \mathbf{P}_f = 0$. Masses M_i and M_f define also the corresponding wave functions through Eqs. (3) and (2).

The matrix elements of the current operator $\hat{J}^\mu(x)$ appears especially simple in the frame defined by Eq. (17):

$$\langle P_f, \chi_f|\hat{J}^\mu(x)|P_i, \chi_i\rangle = 4\pi^{3/2}\sqrt{M_i}M_f e^{i(P_f-P_i)x} \times \langle\chi_f|\hat{j}^\mu(\mathbf{h})|\chi_i\rangle_{\text{n.r.}}, \tag{18}$$

where $\hat{j}^\mu(\mathbf{h})$ is the current operator $\hat{J}(0)$ [see Eq. (A16)] in the frame (17) expressed through $\mathbf{h}$ and $\mathbf{q}$ (see details in Ref. [15]). We use the dimensionless vector $\mathbf{h} = \mathbf{G}_f/G_f^0$, where G_f is a four-velocity of the final NN system in the frame defined by Eq. (17). This parameter may be expressed through the photon momentum $\mathbf{k}$, so that $\mathbf{h} = 2(M_iM_f)^{1/2}(M_i + M_f)^{-2}\,\mathbf{k}$, $|\mathbf{h}| \equiv h = (M_i - M_f)/(M_i + M_f) < 1$. The convenience of this parameter is illustrated in Appendix B.

The internal wave functions of the deuteron and final scattering state are normalized in the nonrelativistic manner. The deuteron wave function is

$$|\chi_i\rangle = |\chi_i\rangle_{\text{n.r.}} = \frac{1}{r}\sum_{l=0,2} u_l(r)|l, 1; 1M_J\rangle, \tag{19}$$

with normalization $\langle\chi_i|\chi_i\rangle_{\text{n.r.}} = 1$, where

$$|l, S; JM_J\rangle = \sum_m\sum_\mu |S,\mu\rangle\mathcal{Y}_{lm}(\hat{n})\mathcal{C}^{JM_J}_{lmS\mu}. \tag{20}$$

The internal wave function of the final continuum np state is

$$|\chi_f\rangle \equiv |q_f, S_f, \mu_f\rangle_{\text{n.r.}} = \sqrt{\frac{2}{\pi}}\frac{1}{q_f r}\sum_{J=0}^{\infty} \times \sum_{M_J=-J}^{J}\sum_{l=J-S}^{J+S}\sum_{l'=J-S}^{J+S}\sum_{m=-l}^{l} i^{l'}u^J_{l',l}(q_f,r) \times \mathcal{C}^{JM_J}_{lm\,S_f\mu_f}\mathcal{Y}^*_{lm}(\hat{q}_f)|l', S_f; JM_J\rangle, \tag{21}$$

with normalization $\langle\chi_{f'}|\chi_f\rangle_{\text{n.r.}} = \delta(\mathbf{q}'_f - \mathbf{q}_f)\delta_{S_f\,S'_f}\delta_{\mu_f\,\mu'_f}$. The corresponding plane wave $|\phi_f\rangle_{\text{n.r.}}$ is characterized by the spherical Bessel functions $j_l(q_f, r)\delta_{ll'}$ instead of $u^J_{l,l'}(q_f, r)$. The deuteron partial-wave functions $u_l(r)$ presented in Fig. 5 and partial waves of the final np states $u^J_{l,l'}(q_f, r)$ are calculated from Eq. (2).

We define a reduced amplitude

$$T_{fi} = \langle\chi_f|\varepsilon^*_\mu\hat{j}^\mu(\mathbf{h})|\chi_i\rangle_{\text{n.r.}}. \tag{22}$$

As a result, the differential cross section (15) can be rewritten as

$$\frac{d\sigma}{d\Omega} = \frac{\pi^2 q_f M_i}{6k_c}\sum_i\sum_f|T_{fi}|^2, \tag{23}$$

where we average over photon polarizations, spin orientations of initial deuteron and sum over spin orientations of final nucleons.

In our calculations, we approximate the above matrix element

$$\langle\chi_f|\hat{j}^\mu(\mathbf{h})|\chi_i\rangle_{\text{n.r.}} \approx \langle\phi_f|\hat{j}^\mu(\mathbf{h})|\chi_i\rangle_{\text{n.r.}} + \langle\chi_f - \phi_f|\tilde{\hat{j}}^\mu(\mathbf{h})|\chi_i\rangle_{\text{n.r.}}. \tag{24}$$

The first term is a plane-wave approximation (PlWA) and is calculated using the exact current operator (A6). In this case, the operator $\hat{\mathbf{q}}$ can be substituted by $\mathbf{q}_f$, and operator structure of $\hat{j}^\mu(\mathbf{h})$ can be presented as

$$\hat{j}^\mu(\mathbf{h}) = j^\mu(\mathbf{h}) + \delta j^\mu = \sum_{i=1,2}\big(\mathbf{B}^\mu_{1i} + \big(\mathbf{B}^\mu_{2i}\cdot\mathbf{s}_i\big) + \big(\mathbf{B}^\mu_{3i}\cdot\mathbf{s}_k\big) + \big(\mathbf{B}^\mu_{4i}\cdot\mathbf{s}_i\big)\big(\mathbf{B}^\mu_{5i}\cdot\mathbf{s}_k\big)\big)I_i(\mathbf{h}) + \delta j, \tag{25}$$

where $j^\mu(\mathbf{h})$ is the sum of the one-nucleon electromagnetic current operators (spectator approximation), and addend δj^μ restores the current conservation equation; $k = 2$, if $i = 1$, and, conversely, $k = 1$, if $i = 2$. B^μ_{1i} and $\mathbf{B}^\mu_{mi}$, $m \geqslant 2$ are vector and tensor functions of arguments $\mathbf{h}$ and $\mathbf{q}_f$. These functions are given in Appendix A, where we calculate addend δj^μ from the current conservation equation following Ref. [15] as we did for the $pp\gamma$ process in Ref. [20]. Obviously, our phenomenological quasipotential model offers no microscopic picture of the interaction that would allow us to unambiguously determine the current operator. We use the defined-below δj^μ only to estimate the violation of the current conservation equation. Assuming gauge invariance (which follows from the Poincaré invariance and the current conservation equation), we use the transverse gauge

$$\varepsilon_\mu = (0, \varepsilon), \qquad (\varepsilon\mathbf{k}) = 0. \tag{26}$$

Thus, we exclude the $j^0(\mathbf{h})$ and $j_{||}(\mathbf{h})$ [see Eq. (A21)] components of the current from Eq. (22). The Poincaré invariance is ensured by definition of the current operator $\hat{J}^\mu(x)$ through the operator $\hat{j}^\mu(\mathbf{h})$ (see details in Ref. [15]).

The use of the $(\chi(\mathbf{r}) - \phi(\mathbf{r}))$ combination in Eq. (24) accelerates the convergence of the partial-wave expansion. This term is nonzero because of the FSI of the neutron and proton. It is calculated from the first order in h approximation of the current operator $\hat{j}^\mu(\mathbf{h})$. This approximation calculated

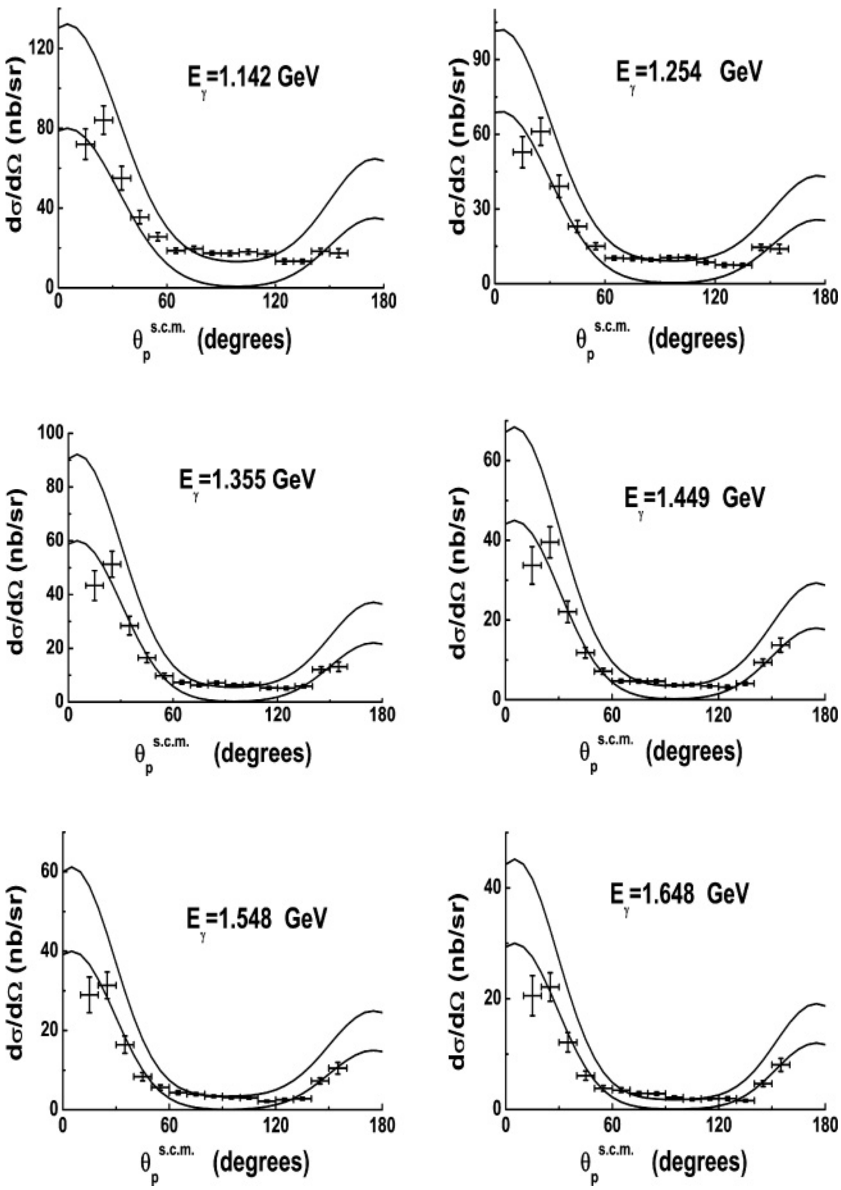

FIG. 6. Angular dependence of $d\gamma \to pn$ reaction differential cross sections for different photon energies E_γ. Our theory is compared with experimental data from Ref. [4].

in the same manner as the one for the pp system in Ref. [20] is given by

$$\begin{aligned}\hat{\mathbf{j}}(\mathbf{h}) \approx \hat{\tilde{\mathbf{j}}}(\mathbf{h}) &= \delta\mathbf{j} + \frac{\mathbf{q}}{w}\hat{g}_e^{pn}(0) - \mathbf{h}\hat{G}_e^{pn}(0)\\ &\quad + \imath\left(\frac{m}{w}[\mathbf{S}\times\mathbf{h}] + \frac{1}{w(w+m)}[\mathbf{q}\times\mathbf{h}](\mathbf{q}\cdot\mathbf{S})\right)\hat{G}_m^{pn}(0)\\ &\quad + \imath\left(\frac{m}{w}[\mathbf{T}\times\mathbf{h}] + \frac{1}{w(w+m)}[\mathbf{q}\times\mathbf{h}](\mathbf{q}\cdot\mathbf{T})\right)\hat{g}_m^{pn}(0)\\ &\quad + \imath(\mathbf{h}\cdot[\mathbf{q}\times\mathbf{S}])\mathbf{q}\left(\frac{\hat{G}_m^{pn}(0)}{mw} + \frac{\hat{G}_e^{pn}(0)}{w(w+m)}\right)\\ &\quad + \imath(\mathbf{h}\cdot[\mathbf{q}\times\mathbf{T}])\mathbf{q}\left(\frac{\hat{g}_m^{pn}(0)}{mw} + \frac{\hat{g}_e^{pn}(0)}{w(w+m)}\right)\\ &\quad - (\mathbf{h}\cdot\mathbf{q})\mathbf{q}\frac{\hat{G}_e^{pn}(0)}{mw},\end{aligned} \tag{27}$$

$$\begin{aligned}\delta\mathbf{j} &= \left(\frac{4w}{M_f+M_i} - 1 - h\right)\frac{\mathbf{q}}{w}\hat{g}_e^{pn}(0)\\ &\quad + \imath h\left\{[\mathbf{q}\times\mathbf{T}]\left(\frac{\hat{G}_m^{pn}(0)}{m} - \frac{\hat{G}_e^{pn}(0)}{w+m}\right) - 2\hat{g}_e^{pn}(0)w\mathbf{r}\right.\\ &\quad \left. + [\mathbf{q}\times\mathbf{S}]\left(\frac{\hat{g}_m^{pn}(0)}{m} - \frac{\hat{g}_e^{pn}(0)}{w+m}\right)\right\},\end{aligned} \tag{28}$$

where $\mathbf{S} = \mathbf{s}_1 + \mathbf{s}_2$, $\mathbf{T} = \mathbf{s}_1 - \mathbf{s}_2$,

$$\begin{aligned}\hat{g}_e^{pn}(0) &= G_e^p(0)I_1(\mathbf{h}) - G_e^n(0)I_2(\mathbf{h}),\\ \hat{g}_m^{pn}(0) &= G_m^p(0)I_1(\mathbf{h}) - G_m^n(0)I_2(\mathbf{h}),\\ \hat{G}_m^{pn}(0) &= G_m^p(0)I_1(\mathbf{h}) + G_m^n(0)I_2(\mathbf{h}),\\ \hat{G}_e^{pn}(0) &= G_e^p(0)I_1(\mathbf{h}) + G_e^n(0)I_2(\mathbf{h}),\end{aligned} \tag{29}$$

$$w \equiv w(\mathbf{q}) = \sqrt{m^2+\mathbf{q}^2}, \quad I_i(\mathbf{h})\chi(\mathbf{q}) = \chi(\mathbf{d}_i(\mathbf{q})), \tag{30}$$

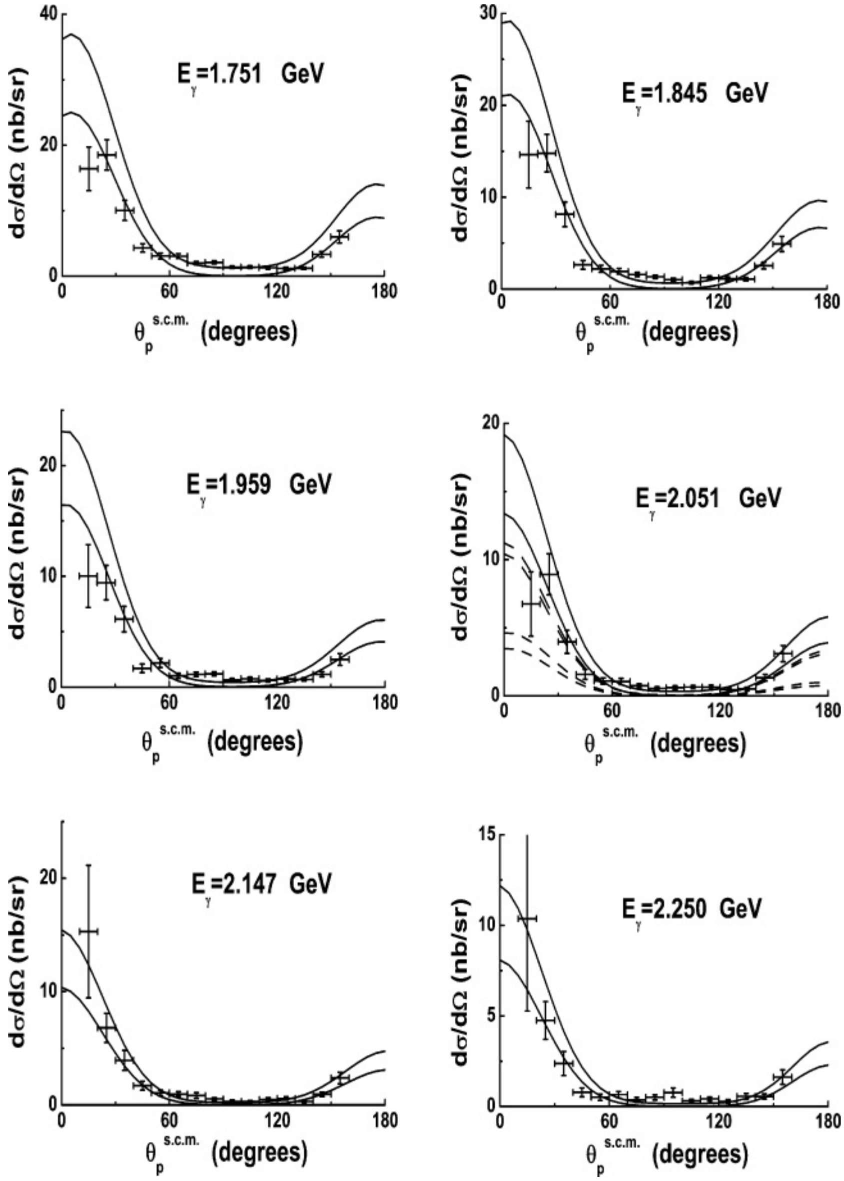

FIG. 7. Same as Fig. 6, but (for $E_\gamma = 2.051$ GeV) our theoretical results for the RCP (Paris potential [45]) are also shown (two lower dashed curves). Upper dashed curves show results of our calculations for PIWA with deuteron wave function in the initial state calculated with our MP.

$$\mathbf{d}_i(\mathbf{q}) = \mathbf{q} + (-1)^i \frac{2\mathbf{h}}{1-h^2}[w + (-1)^i(\mathbf{h}\cdot\mathbf{q})] \approx \mathbf{q} + (-1)^i 2\mathbf{h}w, \qquad (31)$$

where $G_m^n(Q_N^2), G_m^p(Q_N^2), G_e^n(Q_N^2), G_e^p(Q_N^2)$ are nucleon electromagnetic form factors parametrized according to Ref. [43]. In Ref. [19], the elastic electron-deuteron scattering was described in frames of the PF RQM. It was shown that in the PF SA, the momentum of the unstruck particle (the spectator) is unchanged, while the impulse given to the struck particle is not the impulse given to the deuteron.

Following a general approach to construction of the electromagnetic current operator for the relativistic composite system [15], we define the momentum transfer Q_i^2 to the particle i as an increment of the particle four-momentum q_i [19], that is,

$$Q_i^2 = |(q_i' - q_i)^2|. \qquad (32)$$

For interacting particles, the individual four-momenta are not defined before photon absorption as well as after it. Therefore we introduce an operator Q_i^2 corresponding to the physical quantity of the momentum transfer Q_i^2. In Appendix B, we generalize the deduction presented in Ref. [19] and show that

$$Q_1^2 = -(q_1' - q_1)^2 = 16\left(m^2 + \mathbf{q}^2 - \frac{(\mathbf{q}\cdot\mathbf{h})^2}{h^2}\right)\frac{h^2}{(1-h^2)^2}. \qquad (33)$$

This is the general expression of the $Q_1^2 = Q_2^2 = Q_N^2$ in the case of free two-particle states (for particles of equal masses); therefore, we use this expression in the PF SA for evaluation of the current operator in Eq. (25). The parameter $\mathbf{h}$ does not depend on the interaction and is specified by the relative "position" of the initial and final NN c.m. frames. In the case of two interacting particles, $\mathbf{q}$ and Q_i^2 are operators in the internal space. In impulse representation, $\mathbf{q}$ is a variable of integration [19]. It is obvious that in action on a plane wave (for PlWA), this

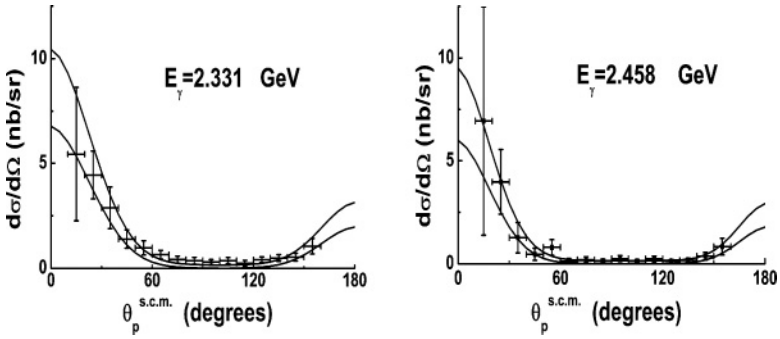

FIG. 8. Same as Fig. 6.

operator is equivalent to the multiplication by a number $Q_N^2 > 0$ if $h \neq 0$. The first order in the h approximation gives $Q_N^2 \approx 0$ for Eq. (27). Consideration similar to that of Ref. [19] gives for the PlWA in our case of the deuteron photodisintegration

$$Q_N^2 = E_\gamma^2 - (w'_n - w'_p)^2 = (2w'_n - m_D)(2w'_p - m_D), \quad (34)$$

where w'_n and w'_n are the final energies of neutron and proton in the initial c.m. frame (laboratory frame). Detailed deduction of Eq. (34) is published in Ref. [44].

V. RESULTS AND FUTURE PROSPECTS

Our theoretical description of the differential cross section of the $d\gamma \to pn$ reaction is compared with recent experiment [4] in Figs. 6–8 at a few energies around $E_\gamma = 2$ GeV. We do not use free parameters. However, there are uncertainties in our calculation. The first uncertainty is caused by uncertainty in the form factor parametrization of Ref. [43] due to errors of the experimental data on the form factors. We estimated this uncertainty at about ±15% of the results varying parameters inside the limits defined in Ref. [43]. The second uncertainty is connected with the approximation (27) used to calculate the FSI term in Eq. (24). In Eq. (27), the nucleon form factors are equal to their values at $Q_N^2 = 0$, and therefore the FSI term is overestimated. Figure 7 ($E_\gamma = 2.051$ GeV) shows the contribution of the FSI term. The PlWA term of the amplitude (24) is dominant, but the FSI is not negligible. Therefore it is desirable to estimate the second order in h correction to the approximation (27). We plan to do this estimation in the future. The third uncertainty is caused by uncertainty of the addend $\delta\mathbf{j}$ that restores the current conservation equation. To estimate this uncertainty, we calculated two curves for every energy of the photon. The lower curves correspond to calculations without addend $\delta\mathbf{j}$ in the current operator and with form factors of Ref. [43] varied to their lower limits. The upper curves correspond to full calculations and with form factors varied to their upper limits. The FSI is included for both curves.

We see good general correspondence of the theory and experiment both in absolute values and in the shape of the angular dependence of the differential cross section at various energies. Large absolute values of cross sections in our theory in comparison with results for the RCP [Fig. 7 ($E_\gamma = 2.051$ GeV)] originate mainly in the nodal character of the deuteron S-wave functions (greater weight of the high-momentum wave function components). The ability to describe both the absolute value and the angular dependence of differential cross sections confirms the detailed algebraic structure of our theory. A persistent forward-backward asymmetry is determined mainly by the angular dependence of the nucleon electromagnetic form factors according to Fig. 9 (proton knockout dominates at forward angles, and neutron knockout dominates at forward angles).

To complete this line of our investigation, we plan to make an analysis of polarization $d\gamma \to pn$ experiments and to consider the pionic radiative capture $pp \to d\pi^+$ at proper energies. Other actual problems are outlined in Ref. [20]. The first concerns the microscopic theory of the MP. As we suppose, it is connected to the short-range quark exchange between nucleons accompanied by excitations of color dipole states of two virtual baryons with very strong attraction between them. This scenario is based on the quark configuration $s^4p^2[42]_x$ in deuteron.

As a concluding remark, it should be stressed that usage of the MP instead of an RCP in the theory of complex nuclei demands accurate evaluation of $3N$ forces. Effect of these forces is much enhanced [23], as far as three nucleons without a NN core can overlap and form short-range $9q$ subsystems with large probability. Recent experiments [46] on the knock-out of nucleon from the ^{3}He nucleus may clarify the situation. In these experiments, the missing momentum is great, and recoil to the $2N$ subsystem with large relative momentum of two spectator nucleons is observed.

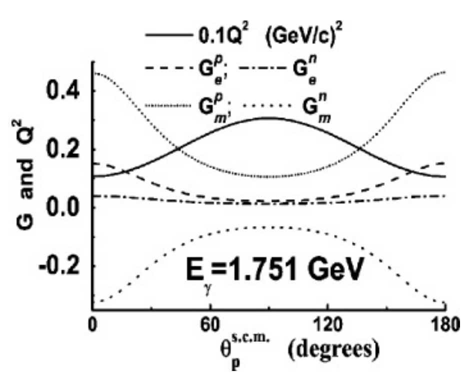

FIG. 9. Angular dependence of four-momentum transfer Q^2, and of the nucleon electromagnetic form factors for $d\gamma \to pn$ reaction calculated from Eq. (34) for the PlWA. In our calculations, we use dependance of the form factors on Q^2 according to parametrization of Kelly [43].

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APPENDIX A

In this Appendix we explain the calculation of the electromagnetic current matrix elements. The derivation is based on results of Ref. [15], where Eq. (18) and Eq. (A6) were deduced.

Let us define a matrix [47]

$$\alpha(g) = \frac{g^0 + 1 + \sigma \cdot \mathbf{g}}{\sqrt{2(g^0+1)}}, \quad \text{(A1)}$$

corresponding to a four-velocity g, where $\sigma = (\sigma_x, \sigma_y, \sigma_z)$ are the Pauli matrices. Let us define the matrix $\check{p} = \widetilde{M}(p) \equiv \sigma^\mu p_\mu$ corresponding to a four-vector p (σ^0 is 2×2 unit matrix). Operator $\widetilde{M}(p)$ transforms the four-vector p to (2×2) matrix. The inverse transformatiodn is defined as

$$\begin{aligned} p_0 &= \frac{1}{2}(\check{p}_{11} + \check{p}_{22}), \quad p_1 = \frac{1}{2}(\check{p}_{12} + \check{p}_{21}),\\ p_2 &= \frac{1}{2i}(\check{p}_{21} - \check{p}_{12}), \quad p_3 = \frac{1}{2}(\check{p}_{11} - \check{p}_{22}), \end{aligned} \quad \text{(A2)}$$

and we denote this transformation as $p = \widetilde{V}(\check{p})$. The boost $p \to L[\alpha(g)]p$ is equivalent to the matrix transformation

$$\check{p} \to \alpha(g)\check{p}\alpha(g)^+. \quad \text{(A3)}$$

It is easy to see that $L[\alpha(g)](1,0,0,0) = g$. The Poincaré group transformation $U(a,l)$ is characterized [15] by the four-shift a and four-rotation l,

$$U(\alpha, l)\varphi(g) = e^{\imath m g' a} D[s; \alpha(g)^{-1} l \alpha(g')]\varphi(g'), \quad \text{(A4)}$$

where $\varphi(g)$ is a normalized spinor function of a particle with mass p; s is the spin of the particle; and $g' = L(l)^{-1}g$. In our case of spin $s = 1/2$ particles, we deal with the fundamental representation [47], i.e., $\mathbf{s}_i \equiv \frac{1}{2}\sigma_i$ and

$$D(s; \alpha(g)^{-1} l \alpha(g')) \equiv \alpha(g)^{-1} l \alpha(g'). \quad \text{(A5)}$$

The "internal" electromagnetic current operator for a system of two particles in the SA is [15]

$$j^\mu(\mathbf{h}) = \sum_{i=1,2} (L^i)^\mu_\nu D^i_1 D^i_2 j^\nu_i(\mathbf{h}) D^i_3 K^i I_i(\mathbf{h}), \quad \text{(A6)}$$

where

$$(L^i)^\mu_\nu = L\left(L[\alpha(f)]\frac{q_i}{m_i}, L[\alpha(f')]\frac{d_i}{m_i}\right)^\mu_\nu, \quad \text{(A7)}$$

$$\begin{aligned} D^i_1 &= D[s_k; \alpha(q_k/m_k)^{-1}\alpha(f)^{-1}\alpha(f')\alpha(d_{ki}/m_k)]\\ &= \alpha_k(q_k/m_k)^{-1}\alpha_k(f)^{-1}\alpha_k(f')\alpha_k(d_{ki}/m_k), \end{aligned} \quad \text{(A8)}$$

$$\begin{aligned} D^i_2 &= D[s_i; \alpha(q_i/m_i)^{-1}\alpha(f)^{-1}\alpha(z_i)]\\ &= \alpha_i(q_i/m_i)^{-1}\alpha_i(f)^{-1}\alpha_i(z_i), \end{aligned} \quad \text{(A9)}$$

$$\begin{aligned} D^i_3 &= D[s_i; \alpha(f'_i)^{-1}\alpha(z_i)^{-1}\alpha(f')\alpha(d_i/m_i)]\\ &= \alpha_i(f'_i)^{-1}\alpha_i(z_i)^{-1}\alpha_i(f')\alpha_i(d_i/m_i), \end{aligned} \quad \text{(A10)}$$

kinematic multipliers

$$K^i = \frac{m_i w_i(\mathbf{q}_i)}{w_i(\mathbf{d}_i)}\left(\frac{M(\mathbf{d}_i)}{M(\mathbf{q})}\right)^{3/2}. \quad \text{(A11)}$$

Here, $k = 2$, if $i = 1$; and, conversely, $k = 1$, if $i = 2$. $L(G, G')$ denotes the Lorentz transformation $L[\alpha(G, G')]$, and $\alpha(G, G') \equiv \alpha[(G+G')/|G+G'|)]$; $z_i = L[\alpha(f)]q_i/m_i$, $L[\alpha(f')]d_i/m_i$. Next,

$$f = L(G, G')^{-1}G, \quad f' = L(G, G')^{-1}G' \quad \text{(A12)}$$

represent the four-velocities of the two-nucleon c.m. in the initial and final state, respectively, meaning the coordinate frame (17). The following formal aspects should be mentioned here:

$$\begin{aligned} f^2 &= f'^2 = 1, \quad \mathbf{f} + \mathbf{f}' = 0,\\ f^0 &= f^{0'} = (1 + \mathbf{f}^2)^{1/2},\\ \mathbf{h} &\equiv \mathbf{f}/f^0;\\ L(G, G') &= L(\alpha(G, G')),\\ \alpha(G, G') &= \alpha((G+G')/|G+G'|);\\ d_1 &= (w_1(\mathbf{d}_1), \mathbf{d}_1),\\ d_2 &= (w_2(\mathbf{d}_2), \mathbf{d}_2),\\ d_{12} &= L[\alpha(f')^{-1}\alpha(f)]q_2 = (w_2(\mathbf{d}_1), -\mathbf{d}_1),\\ d_{21} &= L[\alpha(f')^{-1}\alpha(f)]q_1 = (w_1(\mathbf{d}_2), \mathbf{d}_2). \end{aligned}$$

The last equations give d_i also. Index i or k of matrices α and σ means that it acts in ith or kth particle spin space and appears as in Eq. (A1) but with σ_i or σ_k correspondingly instead of σ. Let $d_1 = (\omega_1(\mathbf{d}_1), \mathbf{d}_1)$, $d_2 = (\omega_2(\mathbf{d}_2), -\mathbf{d}_2)$ and $I_i(\mathbf{h})$ $(i = 1,2)$ be operators defined by the conditions $I_i(\mathbf{h})\chi(\mathbf{q}) = \chi(\mathbf{d}_i)$.

$$\begin{aligned} g'_i &= L[\alpha(f)]\frac{q_i}{m_i},\\ g''_i &= L[\alpha(f')]\frac{d_i}{m_k},\\ f_i &= L[z_i]^{-1}g'_i,\ f'_i = L[z_i]^{-1}g''_i,\\ \mathbf{h}_i &= \frac{\mathbf{f}_i}{f^0_i},\ w_i(q) \equiv \sqrt{m_i^2 + q^2}. \end{aligned}$$

Finally, $j^\nu_i(\mathbf{h})$ is a four-current of the particle i,

$$\begin{aligned} j^0_i(\mathbf{h}) &= eF^i_e(Q^2_i),\\ \mathbf{j}_i(\mathbf{h}) &= -\frac{ie}{\sqrt{1-\mathbf{h}^2_i}} F^i_m(Q^2_i)(\mathbf{h}_i \times \mathbf{s}_i), \end{aligned} \quad \text{(A13)}$$

where vectors $\mathbf{h}_i$ are defined below, $\mathbf{s}_i \equiv \sigma_i/2$, and $Q^2_i = 4m^2_i\mathbf{h}^2_i/\sqrt{1-\mathbf{h}^2_i}$ [see also Eq. (33)].

From Eq. (A6)–(A13), it is obvious that for a plane-wave final state, when operator $\mathbf{q} = -i\nabla$ can be substituted by vector $\mathbf{q}_f$, operator $j(\mathbf{h})$ becomes an exterior product $j^\nu(\mathbf{h}) \equiv \sum_{i=1,2} A^i_{i\nu} \otimes A^i_{k\nu} I_i(\mathbf{h})$; $k = 2$, if $i = 1$, and, conversely, $k =$

1, if $i=2$. The $\mathbf{q}_f$-dependent matrix $A^k_{i\nu}$ acts in ith particle spin space, and presentation $A^k_{i\nu}=\sigma^\mu_i a^k_{i\nu\mu}\equiv a^k_{i\nu0}+2(\mathbf{s}_i\cdot\mathbf{a}^k_{i\nu})$ is valid. "Components" a_ν are extracted by the Eq. (A2) transformation

$$a^i_{\nu i}=\widetilde{V}\big((L^i)^\nu_\mu D^i_2 j^\mu_i(\mathbf{h})D^i_3 K^i\big),\qquad a^k_{\nu i}=\widetilde{V}\big(D^i_1\big),\quad i\neq k. \tag{A14}$$

Functions B of Eq. (25) are expressed as

$$B^\nu_{1i}=a^i_{i\nu0}a^k_{i\nu0},\quad \mathbf{B}^\nu_{2i}=2a^k_{i\nu0}\mathbf{a}^i_{i\nu},\quad \mathbf{B}^\nu_{3i}=2a^i_{i\nu0}\mathbf{a}^k_{i\nu},\quad \mathbf{B}^\nu_{4i}=2\mathbf{a}^i_{i\nu},\quad \mathbf{B}^\nu_{5i}=2\mathbf{a}^k_{i\nu},$$

$k=2$, if $i=1$, and, conversely, $k=1$, if $i=2$.

Now, we should take into account the current conservation equation

$$\frac{\partial\hat{J}^\mu(x)}{\partial x^\mu}=0. \tag{A15}$$

Using also the four-shift

$$\hat{J}^\mu(x)=\exp(i\hat{P}x)\hat{J}^\mu(0)\exp(-i\hat{P}x), \tag{A16}$$

we obtain the relation

$$\hat{P}_\mu\hat{J}^\mu(0)-\hat{J}^\mu(0)\hat{P}_\mu=0. \tag{A17}$$

In terms of the internal variables of the NN system, Eq. (A17) can be reduced to the matrix element

$$\langle\chi_f|M_f G_{f0}j^0(\mathbf{h})-M_f\mathbf{G}_f\mathbf{j}(\mathbf{h})-M_i G_{i0}j^0(\mathbf{h})+M_i\mathbf{G}_i\mathbf{j}(\mathbf{h})|\chi_i\rangle=0, \tag{A18}$$

which can be rewritten in the form

$$\langle\chi_f|(\mathbf{h}\cdot\hat{\mathbf{j}}(\mathbf{h}))|\chi_i\rangle=\frac{M_f-M_i}{M_i+M_f}\langle\chi_f|\hat{\mathbf{j}}^0(\mathbf{h})|\chi_i\rangle, \tag{A19}$$

as far as $\mathbf{G}_i=-\mathbf{h}G_{i0}$, $\mathbf{G}_f=\mathbf{h}G_{i0}$, $\mathbf{P}_i=M_i\mathbf{G}_i$, $\mathbf{P}_f=M_f\mathbf{G}_f$, $\hat{M}|\chi_i\rangle=M_i|\chi_i\rangle$, $\hat{M}|\chi_f\rangle=M_f|\chi_f\rangle$. The current Eq. (A6) does not satisfy Eq. (A19) and needs a modification. Following Ref. [15], we use the unique decomposition into longitudinal and transverse parts

$$\hat{\mathbf{j}}(\mathbf{h})=\hat{\mathbf{j}}(0)+\frac{\mathbf{h}}{h}\hat{j}_{||}(\mathbf{h})+\hat{\mathbf{j}}_\perp(\mathbf{h}), \tag{A20}$$

where $\mathbf{h}\mathbf{j}_\perp(\mathbf{h})=0$ and

$$\hat{j}_{||}(\mathbf{h})=\frac{1}{|\mathbf{h}|}(\mathbf{h}\cdot(\hat{\mathbf{j}}(\mathbf{h})-\hat{\mathbf{j}}(0))),\qquad \hat{\mathbf{j}}_\perp(\mathbf{h})=\hat{\mathbf{j}}(\mathbf{h})-\hat{\mathbf{j}}(0)-\frac{\mathbf{h}}{|\mathbf{h}|^2}(\mathbf{h}\cdot(\hat{\mathbf{j}}(\mathbf{h})-\hat{\mathbf{j}}(0))). \tag{A21}$$

To estimate violation of the current conservation equation, we assume that NN interaction does not change transverse and time components of operator $\hat{j}(\mathbf{h})$. Then we can reconstruct $\hat{\mathbf{j}}(0)$ and $\hat{j}_{||}(\mathbf{h})$ from Eq. (A22). In the transverse gauge (26), the longitudinal component has no effect on our calculation results; therefore, we determine only the matrix element of $\hat{\mathbf{j}}(0)$ as

$$\langle\chi_f|\hat{\mathbf{j}}(0)|\chi_i\rangle=\frac{M_f-M_i}{M_i+M_f}\langle\chi_f|\left.\frac{\partial\hat{\mathbf{j}}^0(\mathbf{h})}{\partial\mathbf{h}}\right|_{h=0}|\chi_i\rangle, \tag{A22}$$

The corresponding addend $\delta\mathbf{j}$ that restores Eq. (A19) is given in Eq. (28).

The first term in Eq. (24) (PIWA) appears as

$$\begin{aligned}\langle\phi_f|\hat{j}^\mu|\chi_i\rangle_{\text{n.r.}}&=\sqrt{\frac{2}{\pi}}\frac{1}{q_f}\sum_{J=0}^{3}\sum_{l=0,2}\sum_{m=-l}^{l}i^l C^{JM_J}_{lm\,1\mu}\\&\times\Bigg[\sum_{k=1}^{4}\sum_{i=1}^{2}\mathcal{Y}^*_{lm}(\hat{q}_i)\langle l,S;JM_J|L^\mu_{ki}\\&\times|l,1;1M_J\rangle U^i_l+\mathcal{Y}^*_{lm}(\hat{q}_f)(1-\delta_{0,\mu})\\&\times\Bigg(\sum_{i=1}^{3}\langle l,S;JM_J|K^\mu_i|l,1;1M_J\rangle U_l(q_f)\\&-2g^{pn}_e\,w(q_f)\sum_{l'=1,3}\langle l',1;JM_J|\hat{r}^\mu|l,1;1M_J\rangle\\&\times\int_0^\infty r\hat{j}_{l'}(q_f r)u_l(r)dr\Bigg)\Bigg],\\U_l(q_f)&=\int_0^\infty\hat{j}_l(q_f r)u_l(r)dr,\\U^i_l&=\frac{w(q_f)(1-h^2)}{w(q_f)(1+h^2)+(-1)^i 2(\mathbf{h}\cdot\mathbf{q}_f)}\\&\times\int_0^\infty\hat{j}_l(d_i(\mathbf{q}_f)r)u_l(r)dr,\end{aligned} \tag{A23}$$

where L^μ_{ki} $(k=1,2,3,4)=B^\mu_{1i}$, $(\mathbf{B}^\mu_{2i}\mathbf{s}_2)$, $(\mathbf{B}^\mu_{3i}\mathbf{s}_1)$, and $(\mathbf{B}^\mu_{4i}\mathbf{s}_2)$ $(\mathbf{B}^\mu_{5i}\mathbf{s}_1)$, respectively; K^μ_i represent the μ components ($\mu=1,2,3$) of the the first three ($i=1,2,3$) terms in Eq. (28).

The second term in Eq. (24) appears as

$$\begin{aligned}\langle\chi_f-\phi_f|\hat{\tilde{j}}^\mu(\mathbf{h})|\chi_i\rangle_{\text{n.r.}}&=\sqrt{\frac{2}{\pi}}\frac{1}{q_f}\sum_{L=0,2}\sum_{J=0}^{\infty}\sum_{l=J-S}^{J+S}\\&\times\sum_{l'=J-S}^{J+S}\sum_{m=-l}^{l}i^{l'}C^{JM}_{lm\,S\mu}\mathcal{Y}^*_{lm}(\hat{q}_f)\\&\times\int_0^\infty dr\langle l',S;JM|\big(u^J_{l',l}(q_f,r)\\&-\delta_{l,l'}\hat{j}_l(q_f r)\big)\hat{\tilde{j}}^\mu(\mathbf{h})u_L(r)\\&\times|L,1;1M_i\rangle.\end{aligned} \tag{A24}$$

We use further the algebraic results (A24)–(A27) of Ref. [20] and obtain the final expression for the differential cross section which is reduced to radial integrals and spherical harmonics but, unfortunately, is too unwieldy to be given here.

Now, by a few examples, we illustrate the calculation technique for the matrix elements of various components of the relativistic current operator

$$\langle l_f, S_f; J_f M_f | (\mathbf{B}^{\mu}_{3i} \cdot \mathbf{s}_1) | l_i, S_i; J_i M_i\rangle$$
$$= \left(\mathbf{B}^{\mu}_{3i} \cdot \langle l_f, S_f; J M_f | \mathbf{s}_1 | l_i, S_i; J_i M_i\rangle\right),$$
$$\langle l_f, S_f; J M_f | (s_1)_\nu | l_i, S_i; J_i M_i\rangle$$
$$= (-1)^{L_i + J_f + S_f + 1} \delta_{L_i L_f} C^{J_i M_i}_{1\nu\, J_f M_f} \sqrt{2J_i + 1} \times \begin{Bmatrix} L_f & J_f & S_f \\ 1 & S_i & J_i \end{Bmatrix} \langle S_f || s_1 || S_i\rangle,$$
$$\langle S_f || s_1 || S_i\rangle = (-1)^{S_f} \sqrt{(2S_i + 1)(6S_f + 3)/2} \times \begin{Bmatrix} 1/2 & 1/2 & S_f \\ 1 & S_i & 1/2 \end{Bmatrix}. \quad \text{(A25)}$$

$$(\mathbf{a}_1 \cdot \mathbf{s}_1)(\mathbf{a}_2 \cdot \mathbf{s}_2) = \sum_{k=0}^{2} C_k [[a_1 \times a_2]^{(k)} \times [s_1 \times s_2]^{(k)}]^{(0)},$$
$$C_k = (1, \sqrt{3}, \sqrt{5}); \quad \text{(A26)}$$

$$\langle S_f || [s_1 \times s_2]^{(k)} || S_i\rangle = \frac{3}{2}\sqrt{(2S_i + 1)(2S_f + 1)(2k + 1)} \times \begin{Bmatrix} 1/2 & 1/2 & S_f \\ 1/2 & 1/2 & S_i \\ 1 & 1 & k \end{Bmatrix}.$$
$$(\nabla \cdot \mathbf{S})\nabla_\mu = -\frac{1}{\sqrt{3}} [[\nabla \times \nabla]^{(0)} \times S]^{(1)}_\mu - \sqrt{\frac{5}{3}} [[\nabla \times \nabla]^{(2)} \times S]^{(1)}_\mu. \quad \text{(A27)}$$

$$(\mathbf{h} \cdot [\nabla \times \mathbf{S}])\nabla_\mu = -i\frac{\sqrt{6}}{3}\left(\frac{\sqrt{15}}{2}[[[\nabla \times \nabla]^{(2)} \times S]^{(2)} \times h]^{(1)}_\mu + [[[\nabla \times \nabla]^{(0)} \times S]^{(1)} \times h]^{(1)}_\mu - \frac{\sqrt{5}}{2}\left[[[\nabla \times \nabla]^{(2)} \times S]^{(1)} \times h\right]^{(1)}_\mu\right). \quad \text{(A28)}$$

$$\langle L_f, S_f = 1; J_f M_f \mid [[\nabla \times \nabla]^{(k)} \times S]^{(1)}]^{(n)}_\mu f(r) | L_i, S_i = 1; J_i M_i\rangle = C^{J_f M_f}_{J_i M_i n\mu} \begin{Bmatrix} L_f & 1 & J_f \\ L_i & 1 & J_i \\ k & 1 & n \end{Bmatrix} \times \sqrt{6(2J_i + 1)(2n + 1)} \langle L_f \| [\nabla \times \nabla]^{(k)} f(r) \| L_i\rangle, \quad \text{(A29)}$$

$$\langle L_f \| [\nabla \times \nabla]^{(2)} \frac{f(r)}{r} \| L_i\rangle = \frac{\sqrt{2L_f + 1}}{\sqrt{6} C^{L_f 0}_{L_i 020}} \frac{1}{r} \left\{ \delta_{L_i L_f} \left(-1 + \frac{3(2L_i^2 + 2L_i - 1)\sqrt{2(2L_i + 1)}}{(2L_i - 1)(2L_i + 1)(2L_i + 3)}\right)\left(\frac{d^2}{dr^2} - \frac{L_i(L_i + 1)}{r^2}\right) f(r) \right.$$
$$+ \delta_{L_i L_f - 2} \frac{3(L_i + 1)(L_i + 2)\sqrt{2(2L_i + 1)}}{(2L_i + 1)(2L_i + 3)(2L_i + 5)} \left(\frac{d^2}{dr^2} - \frac{(2L_i + 3)}{r}\frac{d}{dr} + \frac{(L_i + 3)(L_i + 1)}{r^2}\right) f(r)$$
$$\left. + \delta_{L_i L_f + 2} \frac{3L_i(L_i - 1)\sqrt{2(2L_i + 1)}}{(2L_i + 1)(2L_i - 3)(2L_i - 1)} \left(\frac{d^2}{dr^2} - \frac{(2L_i - 1)}{r}\frac{d}{dr} + \frac{L_i(L_i - 2)}{r^2}\right) f(r) \right\}. \quad \text{(A30)}$$

In these expressions, an upper index in round brackets means a tensor rank of an operator. The first rank is omitted where it is obvious ($\nabla \equiv \nabla^{(1)}$, etc.)

APPENDIX B: POINT-FORM MOMENTUM TRANSFER

In the general case, there are an initial NN state with associated initial c.m. frame (i.c.m.f.) and a final NN state with associated final c.m. frame (f.c.m.f.) Suppose that the photon momentum (momentum transfer) is along the z axis. Values of photon momentum and energy in i.c.m.f. are $|\mathbf{q}_\gamma|$ and q^0_γ, correspondingly. Momentum transfer is $Q^2 = |\mathbf{q}_\gamma|^2 - (q^0_\gamma)^2$. Let P be the total four-momentum of the NN system, M the mass of the NN system, and $G = P/M$ the system four-velocity. Index $i(f)$ means initial (final) state of the NN system. Transformation from i.c.m.f. to the special frame suggested by Lev [15] (L.s.), where

$$\mathbf{G}_f + \mathbf{G}_i = 0|_{\text{L.s.}}, \quad \text{(B1)}$$

is defined by angle $\Delta/2$ such that

$$\tanh \Delta/2 = h, \quad \text{(B2)}$$

where $\mathbf{h} = \mathbf{G}_f / G^0_f|_{\text{L.s.}}$. The Lev frame (B1) is not equivalent to the Breit frame defined by the condition $\mathbf{P}_f + \mathbf{P}_i = 0$ if $M_f \neq M_f$. For elastic electron-deuteron scattering, these frames coincide.

From this point, we may use a special derivation of Ref. [19] Eqs. (B3)–(B7) of the present paper].

The initial energies and z components of momenta in L.s. are

$$\begin{aligned} w_1 &= w\cosh\Delta/2 + q_z\sinh\Delta/2,\\ q_{1z} &= q_z\cosh\Delta/2 + w\sinh\Delta/2,\\ w_2 &= w\cosh\Delta/2 - q_z\sinh\Delta/2,\\ q_{2z} &= -q\cosh\Delta/2 + w\sinh\Delta/2, \end{aligned} \tag{B3}$$

where $\mathbf{q}$ and $w = \sqrt{\mathbf{q}^2 + m^2}$ are center-of-momentum variables, $\mathbf{q}$ is the momentum of particle one (internal variable). After photon absorption, the z component of the internal variable and corresponding energy change

$$q'_z = q_z\cosh\Delta \mp w\sinh\Delta, \tag{B4}$$

$$w' = w\cosh\Delta \mp q_z\sinh\Delta, \tag{B5}$$

where the minus (plus) sign is used when particle one (two) is struck. The final energies and momenta in L.s. will then be

$$\begin{aligned} w'_1 &= w\cosh 3\Delta/2 - q_z\sinh 3\Delta/2,\\ q'_{1z} &= q_z 3\cosh\Delta/2 - w\sinh 3\Delta/2,\\ w'_2 &= w_2, \quad q'_{2z} = q_{2z}, \end{aligned} \tag{B6}$$

other components do not change. Some hyperbolic trigonometry reveals that

$$(q'_1 - q_1)^2 = 4\left(q_z^2 - w^2\right)\sinh^2\Delta, \tag{B7}$$

and it follows from Eq. (B2) that

$$\sinh\Delta = \frac{2h}{1-h^2}. \tag{B8}$$

Since

$$q_z^2 - w^2 = -(m^2 + \mathbf{q}_\perp^2) = -\left(m^2 + \mathbf{q}^2 - \frac{(\mathbf{q}\cdot\mathbf{h})^2}{h^2}\right), \tag{B9}$$

the resulting Eq. (33) is established.

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