Effects of negative-energy components in two-body deuteron photodisintegration

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Several observables in two-body deuteron photodisintegration are investigated in the framework of the Bethe-Salpeter formalism. Apart from keeping the Lorentz covariance throughout, special attention is paid to inclusion of both the positive-energy and negative-energy partial-wave components of the deuteron state. Using the Bethe-Salpeter equation for the deuteron in the ladder approximation with one-boson exchanges as a driving term, the contribution of the negative-energy states is studied for the unpolarized differential cross section as well as for linear photon and tensor target asymmetries. These states are found to have an impact on the observables and, thus, should be taken into account in a complete theoretical treatment of the reaction in the intermediate energy regime.

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I. INTRODUCTION

An application of the Bethe-Salpeter equation (BSE) for spinor particles and Mandelstam's theory to the analysis of two-body deuteron photodisintegration was reported in previous work [1,2], where the contributions of relativistic effects to the differential cross section were estimated using the Bethe-Salpeter (BS) formalism as well as the equal-time approximation to the BSE. The analysis shows some deficiency due to not taking into account the full Dirac structure of the two-nucleon bound state. It is well known that while dealing with the BSE the problem of many coupled states arises since there are four types of solutions of the Dirac equation for a given momentum, two positive-energy states and two negative-energy states. In our case this deficiency is ascribed to making use of the BSE in the context of a relativistic separable interaction kernel while leaving out various negative-energy states [3]. The present paper is intended to shed a light on the effect including of negative-energy components of the deuteron state on the observables of deuteron photodisintegration. Their driving mechanism is the one-boson exchange (OBE) kernel, in which the pion-nucleon coupling is described with the help of axial-vector (A) theory.

The relativistic covariant BSE with a superposition of π, ω, ρ, η, and δ exchanges originally was applied to a description of low-energy nucleon-nucleon (NN) scattering (see Ref. [4] and references therein). Agreement with the experimental data is achieved for partial waves with $J>0$ considering the coupling for the pion-nucleon vertex to be an A type, i.e., a weak $N\bar{N}\pi$ coupling. The same model has been implemented for research on the deuteron electromagnetic (EM) form factors in the relativistic framework [5]. These investigations have inspired others to carry out studies of the EM properties of the two-nucleon system [6] and to develop an effective theory of strongly interacting particles at momentum transfers of a few GeV/c [7]. Moreover, a deuteron in this model has been obtained from solution of the BSE in the ladder approximation with these basic mesons plus one more meson, the σ [8]. Based on this numerical solution, extended calculations of the relativistic corrections to the deuteron static properties, such as the magnetic and quadrupole moments of deuteron, were made [9].

The work introduced in this paper is a study of two-body deuteron photodisintegration using this relativistic framework. The objective is to establish a role played by the negative-energy spinor states of the relativistic wave function of the deuteron. As energy and momentum transfers involved in the process are held below 1 GeV, the appropriate effective degrees of freedom in this energy range are the mesons and nucleons. Assuming a theoretical description in terms of these effective degrees of freedom, we study the differential cross section as well as the photon and tensor asymmetries. A consistent relativistic treatment requires knowledge of the deuteron vertex function, final-state interactions and the EM current operator. For the present work, we use the plane-wave one-body approximation (PWOA) and, thus, extend the work of Ref. [1]. The treatment of the deuteron structure is performed within the relativistic OBE model of NN interaction. The model allows us for the incorporation of the full Dirac structure of two-body bound state amplitudes.

The paper is organized as follows. In the following section we summarize the relevant expression for the polarization observables in terms of the deuteron current matrix elements between the initial and final states of the two-nucleon system. We also briefly describe the dynamical model used in this work and specify the deuteron properties calculated in the framework of the BS theory. Since the chief source of the relativistic effects comes from the relativistic wave function of deuteron, we discuss the detail of the partial-wave decomposition of the full BS amplitude for the deuteron. In Sec. IV, the general structure of the EM current matrix elements in the PWOA is shown. Here the terms generated from the negative-energy spinor states of the deuteron vertex function and the boosting of initial deuteron state from the rest frame to the c.m. frame of the reaction are taken into account. In Sec. V, we describe the numerical results for the polarization observables with an emphasis on the leading contribution coming from the dominant negative-energy triplet P state in the relativistic OBE model with the pseudovector πN coupling. We also present the comparison of our results with

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corresponding predictions calculated in the framework of the BSE with multirank separable interactions. The differences are traced back to the choice of the interaction kernel in actual calculations of the deuteron vertex function. Section VI, finally, contains a discussion of our results and conclusions.

In conclusion, it should be mentioned that the most important relativistic contribution comes from the boost on the single-particle propagator, as it was first shown for the deuteron EM form factors [5] and later confirmed for the photodisintegration of the deuteron [1]. In doing the numerical part of the work we presume that the relative-energy dependence of the vertex function of the deuteron is not so important in comparison with that in the propagator.

II. DEFINITION OF THE POLARIZATION OBSERVABLES

In this section we summarize useful formulas describing the photodisintegration process of the deuteron and the deuteron EM current matrix elements between the initial and final states of the two-nucleon ($2N$) system within the BS formalism. We confine ourselves to the description of the breakup either of tensor polarized deuterons with unpolarized photons or unpolarized deuterons with linearly polarized photon beam. We closely follow the work of Ref. [10], from where the formal expressions of all possible polarization observables in $d(\gamma,N)N$ are derived.

The treatment uses a standard coordinate system, where the incoming photon three-momentum c.m. $\mathbf{q}$ ($|\mathbf{q}|\equiv\omega$) is along the z axis, while the x axis is directed along the maximal linear polarization of the photon. The spin indices of the initial state are specified by the photon polarization $\lambda=\pm1$ and the deuteron spin projection $m_d=0,\pm1$ with respect to the z axis. The polarization of the final state is characterized by the total spin $S=0,1$ of the neutron-proton (np) pair and its z projection m_s. The spherical angles Θ_p and Φ_p define the direction of the relative three-momentum $\mathbf{p}$ of the np pair. At this point the concept of the reduced reaction t matrix elements can be introduced. It can be shown that the azimuthal Φ_p dependence of the reaction matrix $T(\Theta_p,\Phi_p)$ for two-body photodisintegration of the deuteron in the c.m. frame can be isolated. Therefore, the reaction T matrix has the following form:

$$T_{Sm_s\lambda m_d}(\Phi_p,\Theta_p)=e^{\imath(\lambda+m_d)\Phi_p}t_{Sm_s\lambda m_d}(\Theta_p). \tag{1}$$

Here the reduced t matrix depends on the polar angle Θ_p only. The general formulas for the polarization observables are expressed in terms of products of the reaction t amplitudes. The differential cross section is given by

$$\frac{d\sigma_0}{d\Omega_p}=\frac{\alpha}{16\pi s}\frac{|\mathbf{p}|}{\omega}F(\Theta_p),$$

$$F(\Theta_p)=\frac{1}{3}\sum_{Sm_Sm_d}|t_{Sm_s\lambda=1m_d}(\Theta_p)|^2, \tag{2}$$

where $\alpha=e^2/4\pi$ is fine structure constant and s is the square of the total energy of the np pair. The photon asymmetry for linearly polarized photons (Σ^l) and the target asymmetries for tensor oriented deuterons (T_{2M})($M=0,1,2$) are expressed as

$$-\Sigma^lF(\Theta_p)=\frac{1}{3}\sum_{Sm_Sm_d}t^*_{Sm_s1m_d}(\Theta_p)t_{Sm_s-1m_d}(\Theta_p), \tag{3}$$

$$T_{2M}F(\Theta_p)=\frac{\sqrt{5}}{3}\sum_{Sm_Sm_d}\mathrm{C}^{1M+m_d}_{1m_d2M}\mathrm{Re}\{t^*_{Sm_s1m_d}(\Theta_p)\times t_{Sm_s1M+m_d}(\Theta_p)\}(2-\delta_{M0}). \tag{4}$$

It is understood that the observables in Eqs. (2)–(4) depend on the photon energy ω. In most of the experiments on the deuteron photodisintegration angular distributions are measured for fixed laboratory photon energies E_γ. It should be distinguished from the c.m. energy ω. The relation between the photon energies in the two frames has the form

$$\omega=\frac{M_d}{\sqrt{s}}E_\gamma \quad \text{with} \quad s=M_d(M_d+2E_\gamma), \tag{5}$$

where M_d is the rest mass of the deuteron.

The reaction t matrix is expressed in terms of the deuteron EM current matrix element between the final and initial $2N$ states. The procedure for the calculation of the matrix elements in the framework of the BS formalism is based on Mandelstam's theory and the reduction formalism of quantum field theory. It preserves a consistency between amplitudes and EM current operator from the very outset. In the deuteron breakup one deals with the matrix elements of the current between the deuteron and asymptotically free np scattering state:

$$t_{Sm_s\lambda m_d}=\frac{1}{4\pi^3}\int d^4kd^4u\overline{\chi}_{Sm_s}(u;\hat{p}\mathbb{P})\times\epsilon_\lambda\cdot J(u,k;q,\mathbb{K})\chi_{m_d}(k;\mathbb{K}), \tag{6}$$

where J is the irreducible EM vertex, $\chi_{m_d}(k;\mathbb{K})$ is the BS amplitude for the deuteron (isospin $I=0$) with the total momentum $\mathbb{K}=(\sqrt{M_d^2+|\boldsymbol{\omega}|^2},-\boldsymbol{\omega})$, and $\overline{\chi}_{Sm_s}(k;\hat{p}\mathbb{P})$ denotes the conjugate BS amplitude of the final np pair (isospin $I=0,1$). The emission of the nucleon is confined to the x-z plane, and the momentum components in the final state given by $\mathbb{P}=(\sqrt{s},\mathbf{0})$ and $\hat{p}=(0,\mathbf{p})$, with $\mathbf{p}=(|\mathbf{p}|\sin\Theta_p,0,|\mathbf{p}|\cos\Theta_p)$. We set the relative energy $\hat{p}_0$ equal to zero, as the two outgoing nucleons are on mass shell.

The EM vertex operator J is conventionally split up into two pieces,

$$J=J^{[1]}+J^{[2]}, \tag{7}$$

where $J^{[1]}$ is a free part (a nucleon couples individually to the radiation field without interacting with another) and $J^{[2]}$ is a two-body EM vertex operator. Owing to the fact that the EM interaction does not conserve the total isospin I, the transition amplitudes related to $\Delta I=1$ interfere with those having $\Delta I=0$. For consistency, gauge independence of the reaction amplitude has to be respected. It is necessarily guaranteed by the one- and two-body Ward identities, which are derived in the BS framework [11,12]. In particular, the two-body Ward identities are used to elucidate properties of the current operator of a two-body bound state and of a asymptotically free two-body scattering state. The sufficient condition is defined by that both BS amplitudes have to satisfy the corresponding BS equations with the same interaction kernel:

$$\chi_{m_d}(k;\mathbb{K})=\frac{\imath}{4\pi^3}\int d^4u S^{(1)}\left(\frac{\mathbb{K}}{2}+u\right)\times S^{(2)}\left(\frac{\mathbb{K}}{2}-u\right)\mathcal{V}(k,u)\chi_{m_d}(u;\mathbb{K}), \quad (8)$$

$$\chi_{Sm_s}(k;\hat{p}\mathbb{P})=\chi^{(0)}_{Sm_s}(k;\hat{p}\mathbb{P})+\frac{\imath}{4\pi^3}\int d^4u S^{(1)}\left(\frac{\mathbb{P}}{2}+u\right)\times S^{(2)}\left(\frac{\mathbb{P}}{2}-u\right)\mathcal{V}(k,u)\chi_{Sm_s}(u;\hat{p}\mathbb{P}), \quad (9)$$

where $S^{(l)}$ is the free propagator of the lth nucleon and $\chi^{(0)}_{Sm_s}(k;\hat{p}\mathbb{P})$ is the amplitude for the free motion of two nucleons. It is given by the antisymmetric combination of the free Dirac spinors and isospin functions.

III. THE BOUND-STATE VERTEX FUNCTION

In this section, the deuteron vertex function is presented and reduction of the BSE and the BS amplitude for the deuteron to a practical form suitable for numerical solution is then described. The first step in the reduction of the BS amplitude is to show explicitly the spinor and angular-momentum dependence of its factors. This can be done by using either the direct product representation [13] or the matrix representation [14] of two-particle states for a given value of the total angular momentum. The latter allows us to nicely absorb the angular-momentum factors into the specification of the partial-wave states. The second step is the calculation of the matrix elements for different interaction kernels for the partial-wave equations. In the present work we cannot give formulas for all the terms in detail, since we are interested only in the most general structure of the matrix elements for the kernels.

The relativistic wave function χ_{m_d} in Eq. (6) is defined in the c.m. frame of the reaction, in which the deuteron moves with a velocity ω/E_d. In a general moving frame, the vertex function of the deuteron is given by

$$\Gamma_{m_d}(k;\mathbb{K})=\left[S^{(1)}\left(\frac{\mathbb{K}}{2}+k\right)S^{(2)}\left(\frac{\mathbb{K}}{2}-k\right)\right]^{-1}\times\chi_{m_d}(k;\mathbb{K}). \quad (10)$$

As the BSE (8) is solved for the deuteron at rest, we need the relation between vertex functions in the laboratory and c.m. frames. It has the form

$$\Gamma_{m_d}(k;\mathbb{K})=\Lambda^{(1)}(\mathcal{L})\Lambda^{(2)}(\mathcal{L})\Gamma_{m_d}(\mathcal{L}^{-1}k;\mathbb{K}_{(0)}), \quad (11)$$

where Λ is the operator for spin $\frac{1}{2}$ particles corresponding to the boost $\mathcal{L}$ between the two frames: $\mathbb{K}^\mu=\mathcal{L}^\mu_{\ \nu}\mathbb{K}^\nu_{(0)}$ with $\mathbb{K}_{(0)}=(M_d,\mathbf{0})$.

Now we can proceed with the partial-wave expansion of the vertex function $\Gamma_{m_d}(k;\mathbb{K}_{(0)})$. In analyzing eigenstates of the BSE for a given total angular momentum, the total parity $\mathcal{P}$ and exchange quantum (the exchange parity) numbers $\mathcal{P}_{12}$ are used. The first one represents the spacial and intrinsic parities, and the second describes the symmetry of the states under interchange of all coordinates [13]. For on-shell coupled channels $^3S_1$-$^3D_1$ the relativistic wave function consists of eight partial-wave states. We label our spin-angular-momentum basis states involving the Dirac u and v spinors as $\Gamma_{m_d}(\mathbf{k},\alpha)$. Index α runs over all eigenstates of $\mathcal{P}$ and $\mathcal{P}_{12}$, for which the total parity is conserved and symmetry under interchange of particles does not change. The expression for the deuteron vertex function is composed of terms that have the form

$$\Gamma_{m_d}(k_0,\mathbf{k};\mathbb{K}_{(0)})=\sum_{\alpha=1}^{8} g(k_0,|\mathbf{k}|,\alpha)\Gamma_{m_d}(-\mathbf{k},\alpha)\zeta^0_0, \quad (12)$$

where $g(k_0,|\mathbf{k}|,\alpha)$ is the radial function for the corresponding α channel and $\zeta^{I_3}_I$ denotes the normalized eigenstate of the total isospin I and its projection I_3. In the spectroscopic notation, α is labeled as $\alpha={}^{2S+1}L^{\rho}_{J=1}$, where L is the orbital angular momentum, S is the spin, and ρ is the projection of the total energy spin of the $2N$ system. The ρ spin is analogous to the usual σ spin except that instead of operating on spin-up and spin-down states, the ρ spin operates on the positive- and negative-energy states in the same manner. In the deuteron we have the eight symmetrical states:

$$1:{}^3S_1^+,\quad 2:{}^3D_1^+,\quad 3:{}^3S_1^-,\quad 4:{}^3D_1^-,$$
$$5:{}^1P_1^e,\quad 6:{}^3P_1^o,\quad 7:{}^1P_1^o,\quad 8:{}^3P_1^e. \quad (13)$$

The superscript ρ indicates triplet energy spin states as $\rho=(+)$, $(-)$, and (e), and singlet energy spin state is labeled as (o). The radial functions g for the first six states in Eq. (13) are even in the relative energy, while the last two are odd. The separation into even and odd channels clearly exhibits the symmetry of the $2N$ states due to the generalized Pauli principle. Two positive-energy states in Eq. (13) have the corresponding counterparts in conventional nonrelativistic physics, while the rest are of purely relativistic origin.

The BSE for the radial functions are found by the partial-wave decomposition of the kernel in Eq. (8) and by carrying out the angular integration. The decomposition yields the set of the eight coupled two-dimensional integral equations in the relative energy k_0 and in absolute value of the relative three-momentum $|\mathbf{k}|$:

$$g(k_0,|\mathbf{k}|,\alpha)=\frac{\imath}{4\pi^3}\int\int dq_0 d|\mathbf{q}|\sum_{\beta,\gamma}\mathcal{V}(k_0,|\mathbf{k}|,\alpha;q_0,|\mathbf{q}|,\beta)\times G_0(q_0,|\mathbf{q}|,\beta,\gamma)g(q_0,|\mathbf{q}|,\gamma), \quad (14)$$

where $\mathcal{V}(k_0,|\mathbf{k}|,\alpha;q_0,|\mathbf{q}|,\beta)$ and $G_0(q_0,|\mathbf{q}|,\beta,\gamma)$ are the partial-wave decomposition of the interaction kernel $\mathcal{V}$ and the two-particle spinor propagator $G_0=S^{(1)}S^{(2)}$, respectively. The coupling between positive- and negative-energy spinor states occurs directly through the interaction kernel matrix $\mathcal{V}(k_0,|\mathbf{k}|,\alpha;q_0,|\mathbf{q}|,\beta)$ and indirectly through the propagator matrix $G_0(q_0,|\mathbf{q}|,\beta,\gamma)$. The term G_0 has a simple form independent of angle and spin variables, as it depends only on the ρ-spin indices. The structure of the matrices $\mathcal{V}(k_0,|\mathbf{k}|,\beta;q_0,|\mathbf{q}|,\gamma)$ is complicated. Relevant expressions for the pseudoscalar and scalar exchanges can be found in Ref. [13] and for the axial-vector, vector, and tensor exchanges in Ref. [16]. As mentioned above, the radial functions are determined by the set of the eight coupled equations (14). For the first time the deuteron vertex function was calculated for the BSE with a superposition of π, η, ϵ, δ, ρ, and ω exchanges that describe the effective NN forces [5]. The parameters of the relativistic OBE kernel (the meson coupling constants, the masses of the exchange mesons and the cutoff parameters) are constrained by all the low-energy NN data and the binding energy of the deuteron. This investigation was followed by other authors, and the BSE (14) was solved numerically by performing the Wick rotation, $q_0\to\imath q_4$ [15]. Additional variables to the set of the OBE parameters in Ref. [5] are the mass ($m_\sigma=571$ MeV) and the coupling constant of the σ meson. In momentum space, the form of the partial-wave components of the deuteron vertex function, which results from this relativistic OBE model, is discussed fully in Ref. [9].

There is no simple way to compare the BS amplitude and the Schrödinder wave function, which describe the same system, namely, the deuteron. In order to make a comparison possible, a few approximations can be made with respect to the exact equations. One makes use of separable potentials in the context of the BS equation, disregarding the negative-energy spinor states. This is done using a relativistic covariant generalization of the nonrelativistic separable Graz-II potential for the NN system in the coupled $^3S_1^+$ and $^3D_1^+$ states [3]. This turns out to be a solvable model for the deuteron, which was employed in the previous work [1]. The method allows a rigorous relativistic calculation of the differential cross section of the deuteron photodisintegration, including recoil effects.

In this paper we are also interested in the question of how strongly the observables of the deuteron breakup are sensitive to different inputs of the kernel in the BSE. We will compare the deuteron vertex function obtained from the BSE in the ladder approximation with a kernel modeled by the sum of one-boson exchanges and a kernel in the separable form. The positive-energy components, which result from the two models, are shown in Fig. 1. The momentum space functions are plotted against the absolute value of the relative three-momentum $|\mathbf{k}|$, as the relative energy variable k_0 is set equal to zero.

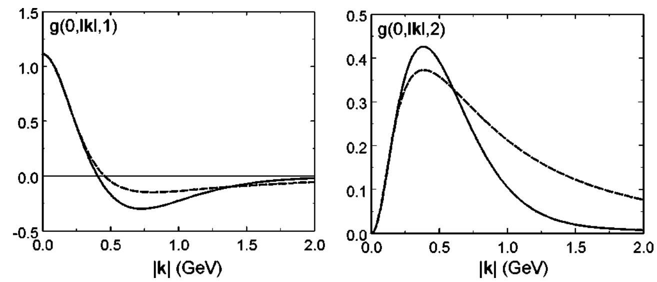

FIG. 1. The momentum dependence of the deuteron vertex function for the coupled $^3S_1^+$-$^3D_1^+$ channels is shown for the relative-energy variable $k_4=0$. The left panel shows the S-wave vertex function $g(0,|\mathbf{k}|,1)$, the right panel shows the D-wave vertex function $g(0,|\mathbf{k}|,2)$. Solid and dashed curves depict the behavior of the vertex functions for the relativistic OBE and separable models, respectively.

It is seen that the behavior essentially depends on the type of interaction model used for the description of the $2N$ system. The radial functions in both channels are similar at low momenta $|\mathbf{k}|\leqslant 250$ MeV, while "tails" of the tensor components at momenta above 500 MeV are much different [28]. Moreover, in the region $0.5\leqslant|\mathbf{k}|\leqslant 1$ GeV the $^3S_1^+$ component for the OBE model has a more pronounced dip than the other. These differences of the momentum components will have an effect on the observables.

In the case of the relativistic OBE interaction, the contribution of the negative-energy spinor components should be discussed. Their importance can be estimated using the normalization condition for the partial-wave components of the deuteron vertex function. The definition and normalization of these is discussed fully in Ref. [9]. They are normalized according to

$$\frac{1}{N_d}\sum_{\alpha=1}^{8}\int_0^{+\infty}dk_4\int_0^{+\infty}d|\mathbf{k}||\mathbf{k}|^2\omega_\rho[\phi(k_4,|\mathbf{k}|,\alpha)]^2=1, \quad (15)$$

where N_d is an appropriate normalization factor, $\omega_\rho=\frac{1}{2}(\rho_1+\rho_2)E_k-(M_d/2)$ with $\rho_{1/2}=\pm\frac{1}{2}$ and $E_k=\sqrt{m^2+|\mathbf{k}|^2}$. The relation between the radial components of the vertex function $g(\alpha)$ and those of the BS amplitude, denoted as $\phi(\alpha)$, has the form

$$g(1,2)=G_0(+,+)^{-1}\phi(1,2) \quad (16)$$

for the S and D positive-energy states,

$$g(3,4)=G_0(-,-)^{-1}\phi(3,4) \quad (17)$$

for the S and D negative-energy states,

$$g(5,7)=\frac{G_0(e,e)\phi(5,7)-G_0(e,o)\phi(6,8)}{G_0(e,e)^2+G_0(e,o)^2} \quad (18)$$

for the two P states even in the relative energy, and

$$g(6,8)=\frac{G_0(e,o)\phi(5,7)+G_0(e,e)\phi(6,8)}{G_0(e,e)^2+G_0(e,o)^2} \quad (19)$$

TABLE I. Deuteron properties in the framework of the BS formalism, compared with experiment.

	OBE	Graz-II	Empirical	Reference(s)
Binding energy ε_d (MeV)	2.2250	2.2254	2.224 575(9)	[3,9,17]
Asymptotic D/S ratio $\rho_{D/S}$	0.02497	0.02691	0.0256(4)	[3,15,17]
Quadrupole moment Q_d(fm^2)	0.2678 [a]	0.2774 [a]	0.2859(3)	[3,9,17]
Magnetic moment μ_d ($e/2m$)	0.8561 [a]	0.8512 [a]	0.857 506(1)	[3,9,17]
D-state probability P_D (%)	5.10	5.0		[3,15]
Pseudoprobability P_- (%)	-0.0050			[9]
Pseudoprobability $P_{\rm even}$ (%)	-0.0920			[9]
Pseudoprobability $P_{\rm odd}$ (%)	-0.0230			[9]

[a]Without meson current contributions and with relativistic corrections.

for the two P states odd in the relative energy. In Eqs. (16)–(19) the partial-wave projections of the two-nucleon propagator G_0 in ρ-subspace are given by

$$G_0(+,+)=\frac{1}{\left(\frac{M_d}{2}-E_k\right)^2+k_4^2},$$

$$G_0(-,-)=\frac{1}{\left(\frac{M_d}{2}+E_k\right)^2+k_4^2}, \tag{20}$$

$$G_0(e,e)=\frac{1}{2}[G_0(+,+)+G_0(-,-)],$$

$$G_0(e,o)=\frac{1}{2}[G_0(+,+)-G_0(-,-)]. \tag{21}$$

The normalization condition (15) defines the probability of the corresponding channel. According to Ref. [5], each can be understood as a measure of the effective charge of the partial-wave state. The corresponding numerical values are shown in Table I. In actual calculation, the upper limits of twofold integrations in Eq. (15) were limited to 3 GeV. It is seen that the probabilities for the relativistic components of the vertex function are negative, and they will be referred to as "pseudoprobabilities." As a consequence, inclusion of the negative-energy spinor states reinforce the contribution of the two positive-energy states to the baryon charge. The smallness of the pseudoprobabilities of the negative-energy states is explained by the use of the πN vertex in A theory. The dominant relativistic component is the triplet $^3P_1^o$ channel with $P_6=-0.08\%$. The relativistic wave functions in the $^3S_1^-$ and $^3D_1^-$ channels are negligible for most purposes. In the case of the πN interaction due to pseudoscalar coupling, values of the pseudoprobabilities of the negative-energy channels are expected to be quite different. Qualitative analysis shows that the relativistic wave function for the $^3S_1^-$ channel may have the largest pseudoprobability, while probabilities of all other channels involving (v) spinor degrees of freedom are of very small percentages [6].

Static properties of the deuteron in the two models discussed above are also shown in Table I. The experimental values are cited according to Table XVII of Ref. [17]. In both models the deuteron binding energy, taken to be 2.225 MeV, was treated as a constant in solving the BSE. As the Graz-II interaction favors the D-state percentage of 5% [3], we have focused our attention on the vertex function calculated for this case. The difference between the D-state percentages is about 0.1%. The deuteron quadrupole Q_d and magnetic μ_d moments were calculated using relativistic formulas without accounting for two-body current contributions. The magnetic moments agree with the empirical value, but the quadrupole moment and the asymptotic D/S ratio $\rho_{D/S}$ are both slightly low. In the OBE model the total relativistic correction to the quadrupole moment, determined by the large components of the deuteron vertex function $Q_d=0.2690$ fm^2, is negative, i.e., $\delta Q_d=-0.0012$ fm^2, reinforcing the discrepancy with the experimental value [9]. The asymptotic D/S ratio was calculated explicitly using formula

$$\rho_{D/S}=\left.\frac{g(k_4,|\mathbf{k}|,2)}{g(k_4,|\mathbf{k}|,1)}\right|_{k_4=0,E_k=M_d/2}. \tag{22}$$

Extrapolation of the radial functions to the unphysical value of $|\mathbf{k}|^2=(M_d^2/4)-m^2$ was carried out by expansion of the numerator and denominator in Eq. (22) in the Taylor series up to terms of the third order.

It is worthwhile to compare the percentages of the purely relativistic components of the deuteron vertex function of the two OBE models [5,15] discussed above. Deuteron relativistic wave functions resulting from these models were obtained using the BSE in the ladder approximation with the same superposition of the basic meson exchanges. But the OBE model of Ref. [15] incorporates one extra meson, the σ. The difference in the percentages of the spin-singlet and spin-triplet P channels reaches one order of magnitude. Compare the total strength of these states $P_-=-2.5\times10^{-2}$ given in Table II of Ref. [5] with $P_-=-1.1\times10^{-1}$ of Ref. [9]. The former value arises from the spin-singlet even and odd P states, while the major contribution to the latter value is due to the spin-triplet even P states. A source of such incompatibility may reside in the form of matrix elements for the one-boson exchanges in the BSE (14). It should be noted that there is a direct coupling between the partial-wave states that are even and odd in relative energy for the axial-vector, vector, and tensor exchanges.

The case of the scalar exchange is different: mixing between even and odd states does not occur. Moreover, such effects as turning off the $q^\mu q^\nu/\mu_V^2$ term in propagators for vector mesons and a tiny change of the πN coupling in A theory may influence quite sensitive changes in strengths of the negative-energy spinor channels [7].

IV. THE REACTION t AMPLITUDE

Having determined the properties of the vertex function of the deuteron at rest, we can calculate the reaction t amplitudes (6), and as a result the differential cross section (2) and photon and the target asymmetries (3) and (4). In PWOA we have (all details of the derivation are given in Ref. [1])

$$\begin{aligned} t_{Sm_s\lambda m_d} = & \sum_{l=1,2} \overline{\chi}_{Sm_s}(\mathbf{p})\zeta_s^+ \Lambda(\mathcal{L})\Gamma_\lambda^{(l)}(q^2=0) \\ & \times S^{(l)}[\tfrac{1}{2}\mathbb{K}_{(0)} - (-1)^l k_l]\Gamma_{m_d}(k_l;\mathbb{K}_{(0)}) \\ & - \sum_{l=1,2}(-1)^S \overline{\chi}_{Sm_s}(-\mathbf{p})\zeta_v^+ \Lambda(\mathcal{L})\Gamma_\lambda^{(l)}(q^2=0) \\ & \times S^{(l)}[\tfrac{1}{2}\mathbb{K}_{(0)} + (-1)^l k_l]\Gamma_{m_d}(-k_l;\mathbb{K}_{(0)}). \end{aligned} \tag{23}$$

where the four-momenta $k_l = \mathcal{L}^{-1} p_l$ and $p_l = p + (-1)^l (q/2)$ are the relative four-momenta of nucleons in the deuteron in the rest and c.m. frames; $\overline{\chi}^{(0)}_{Sm_S} = \chi^{(0)+}_{Sm_S}\gamma_0^{(1)}\gamma_0^{(2)}$ is the conjugate of the $2N$ continuum amplitude, which is composed of terms that have the form $\chi_{Sm_s}(\mathbf{p})\zeta_s$ and $\chi_{Sm_s}(-\mathbf{p})\zeta_v$, where $\chi_{Sm_s}(\mathbf{p})$ is given by combinations of the free Dirac positive-energy spinors

$$\chi_{Sm_s}(\mathbf{p}) = \sum_{\lambda_1\lambda_2} C^{Sm_s}_{\frac{1}{2}\lambda_1\frac{1}{2}\lambda_2} u_{\lambda_1}(\mathbf{p}) u_{\lambda_2}(-\mathbf{p}), \tag{24}$$

and $\zeta_s = \zeta_0^0 + \zeta_1^0$, $\zeta_v = \zeta_0^0 - \zeta_1^0$ are isospin singlet and triplet functions, respectively. It should be noted that the energy component of the four-momentum p_l is equal to one-half the photon energy in the c.m. frame, $p_{0l} = (-1)^l(\omega/2)$. This follows from conservation of four-momenta and assuming two free nucleons in the final state.

Since the deuteron vertex function is calculated in its rest frame, it has to be boosted to the c.m. frame of the final np pair. In Eq. (23) the $\Lambda(\mathcal{L}) = \Lambda^{(1)}(\mathcal{L})\Lambda^{(2)}(\mathcal{L})$ is the operator for spin $\frac{1}{2}$ particles corresponding to this Lorentz transformation $\mathcal{L}$:

$$\Lambda^{(l)}(\mathcal{L}) = \left(\frac{E_d + M_d}{2M_d}\right)^{1/2} \left[1 - \frac{\gamma_0 \boldsymbol{\gamma}\cdot\boldsymbol{\omega}}{E_d + M_d}\right]^{(l)}. \tag{25}$$

The matrix $\mathcal{L}$ defined as $\mathbb{K} = \mathcal{L}\mathbb{K}_{(0)}$ is the boost transformation of the initial BS amplitude from the rest frame to the c.m. frame, in which the deuteron moves with a velocity ω/E_d. Here we only need the boost opposite to the z axis with momentum ω. From this we find

$$\begin{aligned} k_{l0} &= \frac{\omega}{M_d} p_z + (-1)^l \frac{\sqrt{s}}{M_d}\frac{\omega}{2}, \\ k_{lx} &= p_x, \quad k_{ly} = p_y, \\ k_{lz} &= \frac{\sqrt{s}-\omega}{M_d} p_z + (-1)^l \frac{\sqrt{s}}{M_d}\frac{\omega}{2}. \end{aligned} \tag{26}$$

In the breakup of the deuteron, we deal with the half off-mass-shell photon-nucleon vertex Γ_λ, since the knocked-out nucleon is on the mass shell. As a consequence of gauge invariance, $\Gamma^{(l)}$ takes on the on-shell form at the real photon point [18]

$$\Gamma_\lambda(q^2=0) = \epsilon^\mu_\lambda \gamma_\mu \frac{1+\tau_3}{2} + \frac{\imath}{2m}\sigma_{\mu\nu}\epsilon^\mu_\lambda q^\nu \frac{\kappa_s + \kappa_v \tau_3}{2}, \tag{27}$$

where $\kappa_s = \kappa_p + \kappa_n$ and $\kappa_v = \kappa_p - \kappa_n$ with the anomalous part of the proton (neutron) magnetic moments in units of the nuclear magneton $1/2m$ denoted as $\kappa_{p(n)}$.

We further proceed by introducing the matrix representation for each partial-wave channel of the deuteron vertex function in Eq. (13) and for the BS amplitude of np pair in Eq. (24). In evaluating the isospin part of the matrix elements (23), the reaction T amplitudes $t^a_{Sm_s\lambda m_d}$ for the isoscalar $\Delta I = 0$ $(a=1)$ and isovector $\Delta I = 1$ $(a=2)$ transitions are cast into terms composed of γ-matrix traces. Due to the symmetry consideration, the first two terms in Eq. (23) are identical to the latter two. Further evaluation of the reaction t amplitudes for a given set of the polarization indices λ, m_d, and m_s was carried out with the help of the formulas manipulating language REDUCE [19]. The resulting have the general form

$$\begin{aligned} t^a_{Sm_s\lambda m_d} = \sum_{l=1,2}\sum_{\alpha=1}^{8} & [\tilde{\Gamma}^a_{Sm_s\lambda m_d}(\omega,\mathbf{k}_l;l,\alpha) S^{(l)}_{\rho_1}(k_{l0},|\mathbf{k}_l|) \\ & \times g(k_{l0},|\mathbf{k}_l|;\alpha) + \tilde{\Gamma}^a_{Sm_s\lambda m_d}(\omega,-\mathbf{k}_l;l,\alpha) \\ & \times S^{(l)}_{\rho_2}(k_{l0},|\mathbf{k}_l|) g(-k_{l0},|\mathbf{k}_l|;\alpha)], \end{aligned} \tag{28}$$

where $S^{(l)}_{\rho_i=\pm 1}$ are the positive- and negative-energy parts of the single-particle propagator

$$S^{(l)}_{\rho_i}(k_{l0},|\mathbf{k}_l|) = \frac{1}{E_{k_l} - \rho_i\left[\dfrac{M_d}{2} - (-1)^l k_{l0}\right]}. \tag{29}$$

The factors $\tilde{\Gamma}^a_{Sm_s\lambda m_d}(\omega,\mathbf{k}_l;l,\alpha)$ absorb the spin-angular part of the reaction amplitude. A REDUCE code is set up for the explicit analytical evaluation of the spin-angular factors with free polarization indices, given in square brackets of Eq. (28), for each isospin a, particle l, and state α numbers.

V. RESULTS

A. Contributions from relativistic components of the deuteron vertex function

It is time to compute the value of the matrix element in Eq. (28). Next we can proceed with studying various relativistic effects in the differential cross section and the four polarization observables. In view of the discussion following Eq. (14), we meet with a difficulty of finding the deuteron vertex function for real values of the relative energy variable k_0. The problem lies in the method of obtaining numerical results from the BS equation in the bound-state region. It is the Wick transformation of the equation resulting in the analytical continuation of the c.m. vertex function to the imaginary axis with $k_4 = -\imath k_0$. Rotation back to the real k_0 axis is rather hard to do, even if feasible, since a numerical procedure seems to be very unstable and ill defined. It is possible to calculate the vertex function along the real values of k_0 (below the first branch points) by expanding it in a Taylor series around $k_4 = 0$. Unfortunately, our examination shows that the derivatives of the radial components $g(k_4, |\mathbf{k}|; \alpha)$ of the vertex function with respect to k_4 cannot be calculated with a controlled accuracy. The explanation for this is that the solution χ of the BSE is known only at dicretized mesh points. As it was noted in Ref. [5], derivatives of χ should be calculated by taking the derivatives of the kernel with respect to initial momenta $\mathcal{V}'$ and then iterating the integral equation over, rather than by interpolation of the vertex function. In Ref. [1] it was shown that a good approximation of the exact result is the zeroth order approximation (BS-ZO) for the BS amplitude. The approximation amounts to keeping k_0 dependence of the single-particle propagator in Eq. (29) and setting k_{0l} in $g(k_{0l}, |\mathbf{k}|_l, \alpha)$ equal to zero. The Lorentz boost on the spin and momentum arguments of the deuteron vertex function and the single-particle propagator is retained. This approximation also implies that the retardation effects are far less important than the boost on the single-particle propagator due to recoil. A minor drawback of the BS-ZO approximation is loss of the channels that are odd in the relative energy, namely, ${}^1P_1^o$ and ${}^1P_1^e$. We may hope it will produce a small effect on the observables, since the percentage of the odd P states, $P_{\rm odd} \approx -0.02\%$, is roughly one order of magnitude less than that of the even P states, $P_{\rm even} \approx -0.1\%$. As for the model with the separable interaction, we calculate contributions arising from the coupled ${}^3S_1^+$-${}^3D_1^+$ channels of the deuteron vertex function, because all the negative-energy spinor states are disregarded in this case.

The size of the angular distributions for the differential cross section $d\sigma_0/d\Omega_p$ and the linear photon asymmetry Σ^l at three different values of laboratory photon energies are shown in Fig. 2. The BS-ZO calculation in the OBE model is the solid curve. The dotted line shows the effect of shutting off all the negative-energy spinor states of the deuteron vertex function. This leaves us with only large components. The dot-dashed curve calculated should be compared with the dotted one. The former depicts the calculation in the separable model, where relativistic effects apart from those generated by the negative-energy partial waves are accounted for. As seen in Fig. 2, shapes of the angular distributions are formed by the positive-energy components of the deuteron vertex function. The negative-energy components increase markedly the differential cross section, but leave its global structure untouched. Comparing the dotted and dot-dashed lines, we find discrepancy between results that the two models of the deuteron yield. Apparent explanation is supplied by

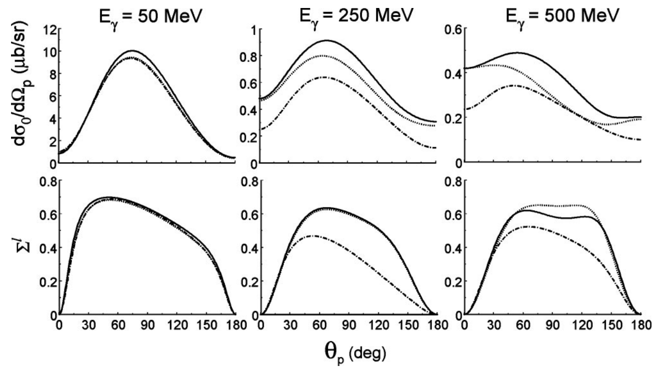

FIG. 2. The differential cross section and the photon asymmetry for linearly polarized photons in the plane-wave one-body approximation at different laboratory photon energies E_γ. Curves: solid line, the BS-ZO approximation takes into account the four negative-energy spinor states ${}^3S_1^-$, ${}^3D_1^-$, ${}^1P_1^e$, and ${}^3P_1^o$; dotted line, the BS-ZO approximation takes into account two positive-energy spinor states ${}^3S_1^+$ and ${}^3D_1^+$. For comparison, the result with the separable vertex function is shown.

Fig. 1. As follows from Eq. (2), the differential cross section is proportional to the sum of squares of the partial-wave components multiplied by the overall kinematic factor. At $E_\gamma = 50$ MeV, the absolute value of the relative momentum $|\mathbf{k}_l|$ in Eq. (26) lies in the range 200–265 MeV. Since the separable and OBE vertex functions are almost identical in this momentum interval, the results are identical. A miniscule rise of the differential cross section in the OBE model is due to the smallness of the negative-energy channels in this momentum range. At $E_\gamma = 500$ MeV, the absolute value of $|\mathbf{k}_l|$ is within the momentum interval 450–950 MeV. Now the ${}^3S_1^+$ partial-wave component in the OBE model has deeper minimum with respect to the other. We conclude that the momentum behavior of the partial-wave components of the deuteron vertex function is significant. Moreover, the ${}^1P_1^e$ and ${}^3P_1^o$ states have persistent superiority over the positive-energy states at relative three-momenta comparable with the rest mass of the nucleon. As the photon energy increases, the contribution of the latter becomes more pronounced. Effects produced by ${}^3S_1^-$ and ${}^3D_1^-$ states are completely negligible. Shutting off the negative-energy states has a minor effects, particularly, on the photon asymmetry in a wide photon energy range. We conclude that the relativistic effect due to the presence of the negative-energy spinor state components is important. The dominant contribution comes from the ${}^3P_1^o$ channel. Our results for the photon asymmetry shows that this relativistic effect almost cancels for medium laboratory photon energies.

We further discuss three observables associated with the tensor oriented deuterons. Figure 3 shows the size of the relativistic effect discussed above. The calculation with inclusion of the six channels of the deuteron vertex function in the OBE model is the solid curve. We can see that negative-energy spinor states have an effect on T_{2M}, particularly at the forward and backward proton c.m. angles. Again, the global structure of the observables is untouched. The dot-dashed line should be compared to the dotted one. For the

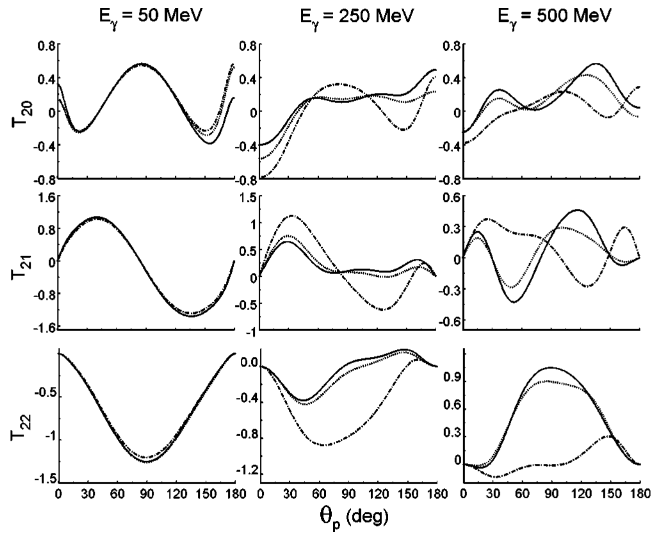

FIG. 3. The tensor target asymmetries. Notations are the same as in Fig. 2.

photon energies about 500 GeV, we find that the observables are more sensitive to the input of the momentum distributions in the deuteron than to the presence of the $^3P_1^o$ state. The dissimilarity between the OBE and separable models clearly shows up in the observable T_{22}. It is explained by the fact that there are fewer interference terms between amplitudes $t_{Sm_s\lambda m_d}$ with different m_d. For $M=2$ the Clebsch-Gordan coefficient in the Eq. (4) selects a product of the reaction t amplitudes with $m_d=-1$ and $m_d=+1$. Conversely, the asymmetry T_{20} is less changed with this respect.

In total, our results show that the relativistic effects due to the dominant negative-energy spinor state $^3P_1^o$ of the vertex function is significant for the differential cross section in two-body deuteron photodisintegration. They are largely canceled in the linear photon and the tensor asymmetries.

B. Supplementary remarks

Finally, in this section we add extra remarks on our above analysis and discuss other related works. The strongest limitations in the present study of the deuteron photodisintegration concern the neglect of the two-body EM current and final-state interaction (FSI). Future investigation of the two-body EM contribution to the reaction t amplitude is compelling, but the omission of the FSI is presumably not too relevant for intermediate photon energies. It by no means underestimates a specific role played by the NN interaction in the final state. The deuteron photodisintegration by low-energy photons, except when $E_\gamma \leqslant 10$ MeV, is dominated by the contribution to the reaction t amplitude from isovector spin-triplet final states. The most numerically significant ones are the partial states ^3P_J $(J=0,1,2)$. In view of the interference of spin-triplet transition amplitudes, it is also important to treat properly isoscalar spin-triplet final states, e.g., $^3S_1$-$^3D_1$ and $^3D_2$ [20]. When omitting the FSI, our analysis suggests that the effects due to negative-energy components of the deuteron vertex function become pronounced, as the photon energy increases. It is an expected result, since the high-energy photon probes the deuteron structure at shorter distances. This relativistic correction seems to be one of the important short-range contributions associated with the virtual production of an extra $N\bar{N}$ pair. As a matter of fact, authors of Ref. [21] investigate a three-dimensional reduction of the covariant transition amplitude for the deuteron electrodisintegration near threshold. It is found that in the relativistic impulse approximation, the P-wave component of the BS amplitude is responsible for the pair current contribution of nonrelativistic theories [22]. However, recent calculations of Refs. [23,24] consider the effects of negative-energy components of nucleon spinors on the proton-proton bremsstrahlung cross section. The analysis, based on the Blankenbecler-Sugar-Logunov-Tavkhelidze (BSLT) formalism in the impulse approximation, shows that substantial cancellations occur between the effects of negative-energy components in the single-scattering diagrams (the photon is emitted from external lines of the NN scattering amplitude) and the rescattering diagrams (the photon is emitted from internal nucleon lines). The net relativistic effect in the bremsstrahlung processes due to negative-energy components is still of the order of 20%. This conclusion raises the question of whether such cancellations may occur in the process of the deuteron photodisintegration. Although the deuteron photodisintegration amplitude is similar to the time reverse of the bremsstrahlung one, with two nucleons in the initial state being a bound state rather than in a scattering one, there are few substantial differences. As said above, a decrease of the NN interaction in the final state is expected at higher photon energies E_γ, since high E_γ in the initial state is directly related to the energy of the np pair. Moreover, there can be an additional indication of the small role played by negative-energy components in some partial channels in the low energy NN scattering. Authors of Ref. [4] show that with the pseudovector πN coupling the deviation between the BS equation and the Blankenbeckler-Sugar equation that couples only positive-energy states, is not appreciable in the isovector spin-triplet ^3P_J channels. These conclusions somewhat point out that the cancellations found in the bremsstrahlung analysis cannot strongly suggest similar cancellations in the deuteron photodisintegration amplitude, at least in the whole momentum range.

The second remark is of a technical character. It concerns the Lorentz boost formulas, given by the Eq. (26). The reaction t amplitude in the Eq. (23) is initially written in the c.m. frame of the np pair. When omitting the FSI, the relative-energy variable in the deuteron vertex function is equal to one-half of the photon energy in the c.m. frame, $p_{0l} = (-1)^l(\omega/2)(l=1,2)$. Thus, the boost formulas in Eq. (26) incorporate the value of the time component of the relative four-momentum, which follows from conservation of four-momentum at the $NN\gamma$ vertex and of the total four momentum. In numerical calculations with one-body EM current we set the time component of the vertex function $k_{0l} = \mathcal{L}^{-1}p_{0l}$, it is needed in the rest frame of a deuteron, equal to zero, leaving both bound nucleons equally off-shell. As discussed above, this approximation implies that the retardation in the vertex function is numerically not able to have a great effect on observables. It should be compared to the relativistic description of electron-deuteron scattering, based on the three-dimensional reduction of the BSE using the BSLT quasipotential formalism [25]. Here, at some stage, the Lorentz transform is stated separately from the energy-momentum conservation. That is clear, since the initial deuteron-bound state is replaced by the BSLT vertex function, and the relative-energy variable k_{0l} in the single-particle propagator is restricted by momentum conservation (the spectator particle is on-shell).

Finally, we will briefly discuss our results including only the positive-energy components of the deuteron vertex function. It concerns differences between the two deuteron vertex functions; the first one is the solution of the BSE with the OBE interaction kernel and the second is obtained for the separable interaction. Results for the differential cross section, vs Fig. 2, are distinct for these two types of interactions. Apart from explanations given in the Sec. V A, which directly stem from observing momentum distributions in Fig. 1, we should mention a major problem. It is that even a refined separable potential, relying on a heuristic ansatz with parameters determined only from low-energy experimental data and on-shell observables, may differ significantly from a potential model derived from dynamic concepts [26]. Care should be taken that the off-shell behavior of the separable kernel is not left to arbitrariness. Large deviations usually occur in off-shell entities, e.g., wave functions. In our case, both interactions describe well the static properties of the deuteron and fits the *NN* phase shifts in the coupled ${}^3S_1^+$-${}^3D_1^+$ channels. However, these constraints determine only asymptotic behavior of the deuteron vertex function and its first derivative in momentum space as $|\mathbf{k}|\to 0$, where $\mathbf{k}$ is the relative three-momentum. Discernibility between two vertex functions develops in the range of medium and high relative three-momenta. It is now clear that results from the application of these two kernels to the two-body problem are rather sensitive to such aspects.

VI. CONCLUDING REMARKS

In this paper we have studied the influence of the relativistic effects on the polarization observables in two-body deuteron photodisintegration within the framework of the BS formalism. Energy transfers involved in the process are held below 1 GeV. We have looked at the sensitivity of our results to the contribution arising from the existence of the relativistic components of the deuteron vertex function. These components comply with requirements of the relativistic covariance, since the full Dirac structure of u and v spinors should be taken into account in the partial-wave analysis of the full BS amplitude of the deuteron. Relativistic effects due to these components are usually ignored in calculations.

The role of the negative-energy spinor components is examined using the BS equation with an OBE interaction kernel. The parameters of the OBE kernel are fitted to produce static properties of the deuteron and the low-energy *NN* phase shifts. Moreover, the importance in particularities of the behavior of the positive-energy components is tested as well. We compare the results obtained in the OBE model of the *NN* interaction with those in the multirank separable model. Unfortunately, our results do not properly account for the role of the relative-energy variable k_0 of the deuteron vertex function. A fundamental obstacle here is posed via the numerical treatment of the BS equation that uses the Wick rotation to produce a transformed integral equation with the standard Euclidean metric that is more susceptible to solution than the nonrotated equation with its Lorentz metric. However, this technique prevents a direct application of the two-body vertex function in the Euclidean space to some physical process. Considering reliable possibilities of making approximate calculations, we employ the zeroth order approximation for the radial part of the deuteron vertex function. It allows us to take into account four out of six negative-energy spinor states as well as all other very important relativistic effects.

The conclusions of this paper are summarized as follows:

(1) We find that inclusion in our relativistic analysis of the negative-energy states of the bound nucleons in the deuteron leads to a sizeable increase (up to 10%) of the differential cross section in a wide range of the proton c.m. angles. This tendency becomes more pronounced at higher photon energies.

(2) Angular distributions of the beam-target asymmetries are less influenced by these states and, thus, show minor changes. However, modifications become noticeable in the tensor asymmetries in the forward and backward directions.

(3) The results show rather strong dependence of the polarization observables to different inputs of the deuteron vertex function in the considered region of the photon energies.

Our numerical results are obtained in the plane-wave approximation with one-body EM current operator. Since the one-body current is not conserved, the choice of gauge for the radiation field becomes very important. Using the transverse gauge in combination with the one-body current operator leads to too small a value for the deuteron EM current matrix elements. This is the reason why we do not make a comparison with the experimental data. The two-body nucleonic current as well as meson exchange currents should be called into play to restore the gauge independence of the reaction amplitude and comply with the Siegert limit [11,27]. Therefore, a consistent treatment of the EM interaction is of prime importance. Corresponding numerical work will be done in the future and results will be reported in a forthcoming paper.

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