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Polarization observables in hard rescattering mechanism of deuteron photodisintegration

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Abstract

Polarization properties of high energy photodisintegration of the deuteron are studied within the framework of the hard rescattering mechanism (HRM). In HRM, a quark of one nucleon knocked-out by the incoming photon rescatters with a quark of the other nucleon leading to the production of two nucleons with high relative momentum. Summation of all relevant quark rescattering amplitudes allows us to express the scattering amplitude of the reaction through the convolution of a hard photon–quark interaction vertex, the large angle p–n scattering amplitude and the low momentum deuteron wave function. Within HRM, it is demonstrated that the polarization observables in hard photodisintegration of the deuteron can be expressed through the five helicity amplitudes of NN scattering at high momentum transfer. At 90° CM scattering HRM predicts the dominance of the isovector channel of hard pn rescattering, and it explains the observed smallness of induced, P_y and transfered, C_x polarizations without invoking the argument of helicity conservation. Namely, HRM predicts that P_y and C_x are proportional to the ϕ_5 helicity amplitude which vanishes at $\theta_{cm} = 90°$ due to symmetry reasons. HRM predicts also a nonzero value for C_z in the helicity-conserving regime and a positive Σ asymmetry which is related to the dominance of the isovector channel in the hard reinteraction. We extend our calculations to the region where large polarization effects are observed in pp scattering as well as give predictions for angular dependences.

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1. Introduction

Hard photodisintegration of the deuteron provides a unique tool for studying the role of quarks and gluons in nuclear interactions. During the last decade several experiments have been performed [1–6] which indicated strongly the importance of quark–gluon degrees of freedom in these reactions starting at $E_\gamma \geqslant 1$ GeV.

First QCD based predictions for high momentum transfer photodisintegration of the deuteron were done within minimal Fock component approximation [7,8] in which it is assumed that only minimal number of

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partonic constituents dominate in large angle hard two-body scattering. Within this approximation the energy dependences of the set of fixed angle hard two-body reactions can be predicted according to the counting rule: $d\sigma/dt \sim s^{-(n_1+n_2+n_3+n_4-2)}$, in which n_i is the number of fundamental constituents in the particle i which is involved in the reaction. This prediction has been confirmed experimentally practically for all two-body reactions for fixed angle hard scattering kinematics in which $-t, -u \geqslant 2\ \mathrm{GeV}^2$.

For high momentum transfer $\gamma + d \to p + n$ reaction the above counting rule predicts an energy dependence, $\sim s^{-11}$ [9], which was confirmed experimentally for photon energies starting at 1 GeV [1–3].

The minimal Fock component approximation can be proven rigorously within perturbative QCD (pQCD) in which the masses of interacting current quarks are neglected. Thus the experimental success of the minimal Fock component approximation raised the expectations that the observed energy dependences indicate the onset of pQCD regime. This was an important question since there were several arguments [10,11] against the application of pQCD in the considered energy range as well as the attempts to describe the absolute cross sections of hard two-body exclusive reactions within leading twist pQCD have been largely unsuccessful (see, e.g., [12,13]) underestimating the observed cross sections by several orders of magnitude.[1]

Since, in QCD the interaction is realized through the exchange of vector gluons, in pQCD (due to vanishing quark masses) the helicity of interacting particles should be conserved. Therefore, as an independent check of the onset of pQCD one can investigate the effects of hadronic helicity conservation (HHC).

The experiments which are aimed at the studies of polarization observables in hard reactions are best suited for HHC studies. The first experiments were performed for elastic pp scattering. While in wide range of hard scattering kinematics the pp data generally are in agreement with HHC, in some instances the striking disagreement is observed [15]. For example, in $\vec{p} + \vec{p} \to p + p$ scattering at $\theta_{\mathrm{cm}} = 90^\circ$ and $P_{\mathrm{Lab}} = 11.75$ GeV [15] the measurements demonstrated that protons polarized transverse to the scattering plane have four times larger probability to scatter with spins parallel than antiparallel to each other. This number is considerably larger than HHC predication of two [17,18]. Several theoretical approaches have been proposed to describe the observed enhancement of the polarization effects (see, e.g., [17–21]), however, the experimental evidence is very limited for meaningful progress in understanding the mechanism of HHC violation.

Since the onset of energy scaling in the cross section of deuteron photodisintegration is observed already at $E_\gamma \geqslant 1$ GeV and $\theta_{\mathrm{cm}} = 90^\circ$, the measurement of polarization observables at the same kinematics will suit ideally for HHC studies. There were several recent studies [5,6,22,23] in polarization properties of high energy deuteron photodisintegration. With JLAB building up a systematic experimental program on deuteron photodisintegration with polarization measurements one may expect a wealth of the new data within next several years [5,16].

In this Letter we study several polarization observables in hard photodisintegration reaction of the deuteron within the recently developed model of hard rescattering (HRM) [24]. HRM is based on the assumption that hard photodisintegration of the deuteron proceeds through two steps: at first, the incoming photon knocks-out a quark from one nucleon in the deuteron which then makes a hard rescattering with a quark of the second nucleon in the deuteron. This assumption allows us to express the disintegration amplitude through the convolution of the deuteron wave function, hard photon–quark interaction amplitude and the amplitude of hard pn scattering. The latter was estimated using the experimental pn scattering data. HRM provides also a convenient framework for calculation of the polarization observables of photodisintegration reaction, expressing them through the helicity amplitudes of pn scattering. In the next sections within HRM we calculate several polarization observables which are currently investigated experimentally. HRM gives rather different insight on observed regularities in polarization measurements and makes several predictions whose verification can advance our understanding the dynamics of hard photodisintegration.

[1] The smallness of the calculated cross sections does not rule out completely the relevance of pQCD regime, since one may expect a sizable effects from unaccounted hidden color component of hadronic wave functions [14].

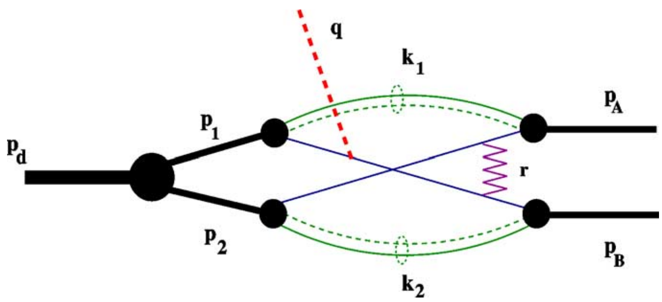

Fig. 1. Typical quark-rescattering diagram.

2. Hard rescattering mechanism

We are considering a reaction

$$\gamma + d \to p + n \tag{1}$$

in which the *polarizations* of γ and/or p are measured. The hard scattering is defined by a requirement that $-t, -u \geqslant 2\ \text{GeV}^2$, where $t = (q - p_p)^2 = (p_n - p_d)^2$, $u = (q - p_n)^2 = (p_p - p_d)^2$ and q, p_d, p_p and p_n are four-momenta of incoming photon, target deuteron, outgoing proton and neutron, respectively.

Within HRM [24] it is assumed that final two high-p_T nucleons are produced due to hard rescattering of a quark, knocked out by incoming photon from one nucleon, with a quark in other nucleon. As a result the sum of diagrams similar to the one presented in Fig. 1 gives the main contribution to the scattering amplitude of the reaction (1).

We start with analyzing the scattering amplitude corresponding to the diagram of Fig. 1:

$$\begin{aligned}
&\langle \lambda_A, \lambda_B | A | \lambda_\gamma, \lambda_D \rangle \\
&= \sum_{\substack{(\eta_1,\eta_2),(\xi_1,\xi_2),\\(\lambda_1,\lambda_2)\zeta}} \int \left\{ \frac{\psi_N^{\dagger\lambda_B,\eta_2}(p_B, x_2', k_{2\perp})}{1 - x_2'} \bar{u}_{\eta_2}(p_B - k_2)\left[-igT_c^F \gamma^\nu\right] \frac{i[u_\zeta(p_1 - k_1 + q)\bar{u}_\zeta(p_1 - k_1 + q)}{(1 - x_1)s'(\alpha_c - \alpha + i\varepsilon)} \right. \\
&\quad \left. \times \left[-ie_q \epsilon_\mu^{\lambda_\gamma} \gamma^\mu\right] u_{\xi_1}(p_1 - k_1) \frac{\psi_N^{\lambda_1,\xi_1}(p_1, x_1, k_{1\perp})}{(1 - x_1)} \right\}_1 \\
&\quad \times \left\{ \frac{\psi_N^{\dagger\lambda_A,\eta_1}(p_A, x_1', k_{1\perp})}{1 - x_1'} \bar{u}_{\eta_1}(p_A - k_1)\left[-igT_c^F \gamma^\mu\right] u_{\xi_2}(p_2 - k_2) \frac{\psi_N^{\lambda_2,\xi_2}(p_2, x_2, k_2)}{(1 - x_2)} \right\}_2 \\
&\quad \times G^{\mu,\nu}(r) \frac{dx_1}{x_1} \frac{d^2k_{1\perp}}{2(2\pi)^3} \frac{dx_2}{x_2} \frac{d^2k_{2\perp}}{2(2\pi)^3} \frac{\Psi_D^{\lambda_D,\lambda_1,\lambda_2}(\alpha, p_\perp)}{(1 - \alpha)} \frac{d\alpha}{\alpha} \frac{d^2p_\perp}{2(2\pi)^3},
\end{aligned} \tag{2}$$

where the four-momenta: p_1, p_2, k_1, k_2, r, p_A and p_B are defined in Fig. 1. Note that k_1 and k_2 define the four-momenta of residual quark–gluon system of the nucleons without specifying their actual composition. $s' = s - M_d^2$, where $s = (q + p_d)^2$. x_1, x_1', x_2 and x_2' are the light-cone momentum fractions of initial and final nucleons carried out by spectator system in the nucleons $x_{1(2)} = k_{1(2)+}/p_{1(2)+}$, $x_{1(2)}' = k_{1(2)+}/p_{A(B)+}$.[2] $\alpha = p_{2+}/p_{d+}$ is the light cone momentum fraction of the deuteron carried by one of the nucleons and $p_\perp$ is the relative transverse momentum of the nucleons in the deuteron. The denominator $(1 - x_1)s'(\alpha_c - \alpha + i\varepsilon)$ is obtained from the denominator of

[2] The light cone four-momentum is defined as $(p_+, p_-, p_\perp)$, where $p_\pm = E \pm p_z$. Here the z axis is defined in the direction opposite to the incoming photon momentum.

knocked-out quark propagator, $(p_1 - k_1 + q)^2 - m_q^2 + i\varepsilon$ by expressing it through α and

$$\alpha_c = 1 + \frac{1}{s'}\left[\tilde{m}_N^2 - \frac{\tilde{m}_R^2(1-x_1) + m_q^2 x + (k_1 - xp_1)_\perp^2}{x_1(1-x_1)}\right], \tag{3}$$

where $\tilde{m}_N^2 = p_{1-}p_{d+}(1-\alpha) - p_\perp^2$ and $\tilde{m}_R^2 = k_{1-}p_{d+}(1-\alpha)x_1 - k_{1\perp}^2$ are an effective masses of the off-shell nucleon and its residual system, respectively. m_q represents the current quark mass of the knocked-out quark. The scattering process in Eq. (2) can be described through the combination of the following blocks:

(a) $\Psi_D^{\lambda_D,\lambda_1,\lambda_2}(\alpha, p_\perp)$, is the light-cone deuteron wave function which describes the transition of the deuteron with helicity λ_D into two nucleons with λ_1 and λ_2 helicities, respectively.
(b) The term in $\{\ldots\}_1$ describes the "knocking out" a ξ_1-helicity quark from the λ_1-helicity nucleon by an incoming photon with helicity λ_γ. Subsequently, the "knocked-out" ζ_1-helicity quark exchanges gluon, $([-igT_c^F\gamma^\nu])$, with a quark from second nucleon producing a final η_2-helicity quark which enters the nucleon B with helicity λ_B.
(c) The term in $\{\ldots\}_2$ describes the emerging ξ_2-helicity quark from λ_2-helicity nucleon which then exchanges a gluon, $([-igT_c^F\gamma^\mu])$, with the knocked-out quark and produces a final η_1-helicity quark which enters the nucleon with helicity λ_A.
(d) The propagator of the exchanged gluon is $G^{\mu\nu}(r) = d^{\mu\nu}/(r^2 + i\varepsilon)$ with polarization matrix, $d^{\mu\nu}$ (fixed by light-cone gauge), and $r = (p_2 - k_2 + l) - (p_1 - k_1 + q)$, with $l = (p_B - p_2)$.

In Eq. (2) the $\psi_N^{\lambda,\tau}$ represents everywhere a τ-helicity single quark wave function of λ-helicity nucleon and u_τ is the quark spinor defined in the helicity basis. We keep only the $u_\zeta \bar{u}_\zeta$ term in the numerator of the knocked-out quark propagator, since this is the only term that contributes through the soft (dominant) component of the deuteron wave function.

Next, we integrate Eq. (2) by α, taking into account only on-mass shell contribution of struck quark propagator, i.e., the second term in the decomposition: $(\alpha_c - \alpha + i\varepsilon)^{-1} = \mathcal{P}(\alpha_c - \alpha)^{-1} - i\pi\delta(\alpha_c - \alpha)$. The on-mass shell approximation allows us to evaluate the photon–quark interaction vertex, for which, in vanishing current quark mass approximation one obtains:

$$\bar{u}_\zeta(p_1 - k_1 + q)\left[-ie_q\epsilon_\mu^{\lambda_\gamma}\gamma^\mu\right]u_{\xi_1}(p_1 - k_1) = e_q\sqrt{2s'}\sqrt{[1-(1-\alpha)(1-x_1)]\,(1-\alpha)(1-x_1)}\,\delta^{\zeta,\lambda_\gamma}\delta^{\lambda_\gamma,\xi_1}. \tag{4}$$

Two important features of the above equation should be emphasized: (i) an energetic photon selects only those quarks from a nucleon that have the same helicity that the photon has ($\xi_1 = \lambda_\gamma$); (ii) the helicity of the initial quark is conserved after it was struck by incoming photon ($\zeta = \xi_1$). Inserting Eq. (4) into Eq. (2) and taking the $d\alpha$ integral by estimating it through the residue at the pole $\alpha = \alpha_c$ one obtains:

$$\begin{aligned}
&\langle\lambda_A, \lambda_B|A|\lambda_\gamma, \lambda_D\rangle \\
&\quad = \sum_{\substack{(\eta_1,\eta_2),(\xi_2),\\(\lambda_1,\lambda_2)}} \int \frac{e_q\sqrt{2}}{(1-x_1)\sqrt{s'}}\sqrt{[1-(1-\alpha_c)(1-x_1)](1-\alpha_c)(1-x_1)} \\
&\qquad \times \left\{\frac{\psi_N^{\dagger\lambda_B,\eta_2}(p_B, x_2', k_{2\perp})}{1-x_2'}\bar{u}_{\eta_2}(p_B - k_2)\left[-igT_c^F\gamma^\nu\right]u_{\lambda_\gamma}(p_1 - k_1 + q)\right. \\
&\qquad \times \frac{\psi_N^{\lambda_1,\lambda_\gamma}(p_1, x_1, k_{1\perp})}{(1-x_1)}\frac{\psi_N^{\dagger\lambda_A,\eta_1}(p_B, x_1', k_{1\perp})}{1-x_1'}
\end{aligned}$$

$$\times\, \bar{u}_{\eta_1}(p_A-k_1)\big[-igT_c^F\gamma^\mu\big]u_{\xi_2}(p_2-k_2)\frac{\psi_N^{\lambda_2,\xi_2}(p_2,x_2,k_2)}{(1-x_2)}G^{\mu,\nu}(r)\frac{dx_1}{x_1}\frac{d^2k_{1\perp}}{2(2\pi)^3}\frac{dx_2}{x_2}\frac{d^2k_{2\perp}}{2(2\pi)^3}\Bigg\}$$
$$\times\, \frac{\Psi^{\lambda_D,\lambda_1,\lambda_2}(\alpha,p_\perp)}{(1-\alpha)\alpha}\frac{d^2p_\perp}{4(2\pi)^2}. \tag{5}$$

One can relate the expression in $\{\ldots\}$ to the quark-interchange kernel of NN interaction. Taking into account the fact that the deuteron wave function peaks strongly at $\alpha_c=1/2$ we approximate Eq. (5), choosing $\alpha_c=1/2$. In this case in $x_1\to 0$ limit, which corresponds to the Feynman picture of hard scattering [25], Eq. (5) factorizes into the product of γ–quark scattering vertex and quark-exchange amplitude of NN scattering [24]. In the case of the minimal Fock component approximation, in which x_1, $(1-x_1)\sim 1$, the factorization is correct up to the scaling function $f(\theta_{\mathrm{cm}})$, with $f(\theta_{\mathrm{cm}}=90^\circ)\approx 1$ [24]. Using this factorization, for Eq. (5) one obtains:

$$\langle\lambda_A,\lambda_B|A_{Q_i}|\lambda_\gamma,\lambda_D\rangle$$
$$=\sum_{\substack{(\eta_1,\eta_2),(\xi_2),\\(\lambda_1,\lambda_2)}}\int\frac{eQ_if(\theta_{\mathrm{cm}})}{\sqrt{2s'}}\langle\eta_2,\lambda_B|\langle\eta_1,\lambda_A|A^i_{\mathrm{QIM}}\big(s,l^2\big)|\lambda_1,\lambda_\gamma\rangle|\lambda_2\xi_2\rangle\Psi^{\lambda_D,\lambda_1,\lambda_2}(\alpha_c,p_\perp)\frac{d^2p_\perp}{(2\pi)^2}, \tag{6}$$

where $\langle\eta_2,\lambda_B|\langle\eta_1,\lambda_A|A^i_{\mathrm{QIM}}(s,l^2)|\lambda_1,\lambda_\gamma\rangle|\lambda_2,\xi_2\rangle$ is the quark-interchange kernel (with quark-i interacting with the photon) corresponding to the expression in $\{\ldots\}$ in Eq. (5). Here $|\lambda,\eta\rangle$ represents η-helicity quark wave function of λ-helicity nucleon. Since the momenta of interacting quarks are large $(1-x_1\sim 1)$ one can assume that the interchanging quarks carry the helicities of a parent nucleons (i.e., $\eta=\lambda$). This allows us to express the scattering amplitude in Eq. (6) through the helicities of the photon, deuteron and nucleons as follows:

$$\langle\lambda_A,\lambda_B|A_{Q_i}|\lambda_\gamma,\lambda_D\rangle$$
$$=\sum_{\lambda_2}\int\frac{ef(\theta_{\mathrm{cm}})}{\sqrt{2s'}}Q_i\langle\lambda_A,\lambda_B|A^i_{\mathrm{QIM}}\big(s,l^2\big)|\lambda_\gamma,\lambda_2\rangle\Psi^{\lambda_D,\lambda_\gamma,\lambda_2}(\alpha_c,p_\perp)\frac{d^2p_\perp}{(2\pi)^2}, \tag{7}$$

where $|\lambda_1,\lambda_2\rangle$ represents two nucleons having λ_1 and λ_2 helicities, respectively. Note that A^i_{QIM} in the above equation is weighted with the charge of the knocked-out quark Q_i, thus it cannot be directly related to the quark interchange amplitude of $pn\to pn$ scattering.

To calculate the total scattering amplitude within HRM we sum all amplitudes of topologies of Fig. 1. Identifying λ_A and λ_B with the helicities of proton and neutron, respectively, one obtains:

$$\langle p_{\lambda_A},n_{\lambda_B}|A|\lambda_\gamma,\lambda_D\rangle=\sum_{i\in D}\big[\langle p_{\lambda_A},n_{\lambda_B}|A_{Q_i}|\lambda_\gamma,\lambda_D\rangle-\langle n_{\lambda_B},p_{\lambda_A}|A_{Q_i}|\lambda_\gamma,\lambda_D\rangle\big], \tag{8}$$

where one sums valence quarks of the deuteron. This sum can be performed within the quark-interchange model of hadronic interactions, which allows us to represent the NN scattering amplitude as follows [18]:

$$\langle a'b'|A|ab\rangle=\frac{1}{2}\langle a'b'|\sum_{i\in a,j\in b}[I_iI_j+\vec{\tau}_i\cdot\vec{\tau}_j]F_{i,j}(s,t)|ab\rangle, \tag{9}$$

where I_i and τ_i are identity and Pauli matrices defined in $SU(2)$ flavor (isospin) space of the interchanged quarks. The kernel, $F_{i,j}(s,t)$ describes an interchange of i and j quarks.[3]

[3] The additional assumption of helicity conservation allows us to express the kernel in the form [18]: $F_{i,j}(s,t)=\frac{1}{2}[I_iI_j+\vec{\sigma}_i\cdot\vec{\sigma}_j]\tilde{F}_{i,j}(s,t)$, where I_i and σ_i operate in $SU(2)$ helicity (H-spin) space of exchanged (i,j) quarks [18]. However, for our discussion the assumption of helicity conservation is not required.

One can use Eq. (9) to calculate the quark-charge weighted QIM amplitude, $\langle a'b'|A^Q|ab\rangle$, to obtain:

$$\langle a'b'|A^Q|ab\rangle\big|_{a,b\in D} = \frac{1}{2}\langle a'b'|\sum_{i\in a,\, j\in b}[I_iI_j + \vec{\tau}_i\cdot\vec{\tau}_j](Q_i+Q_j)F_{i,j}(s,t)|ab\rangle = (Q_u+Q_d)\langle a'b'|A|ab\rangle$$
$$= \frac{1}{3}\langle a'b'|A|ab\rangle. \tag{10}$$

The above result can be understood qualitatively: since the number of u and d quarks in the deuteron are equal one has the same number of diagrams with knocked out u and d quarks. Using Eqs. (7), (8) and (10) for $\gamma d\to pn$ amplitude one obtains:

$$\begin{aligned}&\langle p_{\lambda_A}, n_{\lambda_B}|A|\lambda_\gamma,\lambda_D\rangle\\ &\quad= \sum_{\lambda_2}\frac{f(\theta_{\mathrm{cm}})}{3\sqrt{2s'}}\big(\langle p_{\lambda_A}, n_{\lambda_B}|A_{pn}(s,t_n)|p_{\lambda_\gamma}, n_{\lambda_2}\rangle + \langle p_{\lambda_A}, n_{\lambda_B}|A_{pn}(s,u_n)|n_{\lambda_\gamma}p_{\lambda_2}\rangle\big)\\ &\qquad\times\int \Psi^{\lambda_D,\lambda_\gamma,\lambda_2}(\alpha_c, p_\perp)\frac{d^2p_\perp}{(2\pi)^2},\end{aligned} \tag{11}$$

where $t_n = (p_B - \frac{1}{2}p_D)^2$, $u_n = (p_A - \frac{1}{2}p_D)^2$ and A_{pn} is the helicity amplitude of pn scattering, which is factorized from the integral. In the factorization we take into account also the antisymmetry of the deuteron wave function with respect to $p\leftrightarrow n$. This factorization is justified due to the fact that at $\alpha_c = \frac{1}{2}$ the momenta involved in the integration, $p_\perp \leqslant 300\ \mathrm{MeV}/c$ are much smaller than the transferred momenta in the A_{pn} amplitude. For the same reason one can approximate the light-cone deuteron wave function that enters in Eq. (11) through rather well-known nonrelativistic deuteron wave function [24,26]: $\Psi^{\lambda_D,\lambda_1,\lambda_2} = (2\pi)^{3/2}\Psi_{NR}^{J_D,\lambda_1,\lambda_2}\sqrt{m}$, where $\Psi_{NR}^{\lambda_D,\lambda_1,\lambda_2} = [u(k) + w(k)\sqrt{1/8}\,S_{12}]\xi_1^{\lambda_D,\lambda_1,\lambda_2}$, with $u(k)$ and $w(k)$ corresponding to the s- and d-waves normalized as $\int |u(k)|^2(|w(k)|^2)\,d^3k = 1$ and $\xi_1^{\lambda_D,\lambda_1,\lambda_2}$ represents the spin component of the wave function.

3. Predictions for polarization observables

3.1. Definition of observables

We will discuss several polarization observables of reaction (1) for which there are ongoing experimental investigations [6,16]. These are:

- Recoil-proton polarization P_y which corresponds to the measurement of asymmetry in the spin component of the protons parallel/antiparallel to the direction of $y = \hat{q}\times\hat{p}_p$ for the reaction with unpolarized photon and deuteron.
- Transfered polarizations $C_{x'}$ and $C_{z'}$, which correspond to the measurement of asymmetry in the spin component of the protons parallel/antiparallel to the directions of $\hat{x}' = \hat{p}_p\times\hat{y}$ and $\hat{p}_p$, respectively, for the reaction with circularly polarized photons and unpolarized deuteron.
- Cross section asymmetry Σ for the reaction with linearly polarized photons.

These observables are expressed through the helicity amplitudes $\langle\lambda_p\lambda_n|A|\lambda_\gamma,\lambda_d\rangle$ as follows [22,27]:

$$f(\theta)P_y = 2\,\mathrm{Im}\sum_{i=1}^{3}\big[F_{i+}^\dagger F_{[i+3]-} + F_{i-}F_{[i+3]+}^\dagger\big],\qquad f(\theta)C_{x'} = 2\,\mathrm{Re}\sum_{i=1}^{3}\big[F_{i+}^\dagger F_{[i+3]-} + F_{i-}F_{[i+3]+}^\dagger\big],$$

$$f(\theta)C_{z'} = \sum_{i=1}^{6}\big[|F_{i+}|^2 - |F_{i-}|^2\big], \qquad f(\theta)\Sigma = -2\,\mathrm{Re}\Big[\sum_{\pm}\big(F^{\dagger}_{1\pm}F_{3\mp} - F^{\dagger}_{4\pm}F_{6\mp}\big) - F^{\dagger}_{2+}F_{2-} + F^{\dagger}_{5+}F_{5-}\Big],$$

$$f(\theta) = \sum_{i=1}^{6}\sum_{\pm}|F_{i\pm}|^2, \tag{12}$$

where $F_{i\pm} = \langle \pm, \pm|A|1, 2-i\rangle$, for $i = 1, 2, 3$ and $F_{i\pm} = \langle \pm, \mp|A|1, 5-i\rangle$, for $i = 4, 5, 6$.

3.2. HRM predictions

Based on Eq. (11) one calculates the observables defined in Eq. (12) expressing them through the helicity amplitudes of pn scattering. Derivations are simplified further by using the fact that the momenta relevant in the deuteron wave function are $\leqslant 300\ \mathrm{MeV}/c$. As a result one can restrict by s wave contribution in the deuteron wave function only. In this case the radial part of the deuteron wave function in Eq. (12) will cancel out and one obtains:

$$P_y = -\frac{2\,\mathrm{Im}\{\phi_5^{\dagger}[2(\phi_1+\phi_2)+\phi_3-\phi_4]\}}{2|\phi_1|^2+2|\phi_2|^2+|\phi_3|^2+|\phi_4|^2+6|\phi_5|^2}, \qquad C_{x'} = \frac{2\,\mathrm{Re}\{\phi_5^{\dagger}[2(\phi_1-\phi_2)+\phi_3+\phi_4]\}}{2|\phi_1|^2+2|\phi_2|^2+|\phi_3|^2+|\phi_4|^2+6|\phi_5|^2},$$

$$C_{z'} = \frac{2|\phi_1|^2-2|\phi_2|^2+|\phi_3|^2-|\phi_4|^2}{2|\phi_1|^2+2|\phi_2|^2+|\phi_3|^2+|\phi_4|^2+6|\phi_5|^2}, \qquad \Sigma = \frac{2\,\mathrm{Re}[|\phi_5|^2-\phi_3^{\dagger}\phi_4]}{2|\phi_1|^2+2|\phi_2|^2+|\phi_3|^2+|\phi_4|^2+6|\phi_5|^2}, \tag{13}$$

where off-shell helicity amplitudes of pn scattering are:

$$\phi_1(s,t_n,u_n) = \langle +,+|A_{pn\to pn}+A_{pn\to np}|+,+\rangle,$$
$$\phi_2(s,t_n,u_n) = \langle +,+|A_{pn\to pn}+A_{pn\to np}|-,-\rangle,$$
$$\phi_3(s,t_n,u_n) = \langle +,-|A_{pn\to pn}+A_{pn\to np}|+,-\rangle,$$
$$\phi_4(s,t_n,u_n) = \langle +,-|A_{pn\to pn}+A_{pn\to np}|-,+\rangle,$$
$$\phi_5(s,t_n,u_n) = \langle +,+|A_{pn\to pn}+A_{pn\to np}|+,-\rangle. \tag{14}$$

Due to the relation: $A_{pn\to pn/np} = \frac{A^{I=1}}{2}(+/-)\frac{A^{I=0}}{2}$, in which I is the isospin of the pn system one observes that in the on-shell limit HRM predicts a dominance of the isovector channel in pn rescattering at $\theta_{\mathrm{cm}} = 90^\circ$. In this case one has the following features of on-shell ϕ-amplitudes at $\theta_{\mathrm{cm}} = 90^\circ$: (i) $\phi_5 = 0$ and (ii) $\phi_3 = -\phi_4$. Furthermore, for any given isospin state and θ_{cm} there is a hierarchy in helicity amplitudes in the hard regime of the scattering (see, e.g., [20,21]):[4]

$$|\phi_1| \geqslant |\phi_3|, |\phi_4| > |\phi_5| > |\phi_2|. \tag{15}$$

Based on the above features one can do following rather general observations for polarization observables of Eq. (13):

– P_y and $C_{x'}$ should be small at large θ_{cm}, due to the fact that on-shell ϕ_5 approaches zero at $\theta_{\mathrm{cm}} \to 90^\circ$. Thus the smallness of P_y and $C_{x'}$ at 90° will not necessarily indicate an onset of helicity-conserving regime in the scattering amplitude. This observation can be checked by looking at θ_{cm} dependence of P_y and $C_{x'}$. Their increase with θ_{cm} going away from 90° will confirm the present conjecture.[5]

[4] This hierarchy is well founded phenomenologically, even with observed finite effects of helicity nonconservation (see, e.g., [20]).

[5] Inclusion of the d wave in the deuteron wave function will not change the result, since the additional terms associated with the d wave are proportional to ϕ_5, too.

- Using relations of Eq. (15), from Eq. (13) one can conclude that the relative sign of P_y and $C_{x'}$ is related predominantly to the relative phase of ϕ_5 and ϕ_1. For example, if real and imaginary parts of both ϕ_5 and ϕ_1 have same signs then P_y and $C_{x'}$ will have an opposite signs.
- Based on Eq. (15) on expects $C_{z'}$ to have a positive values ≈ 0.5–1.
- The relative sign of ϕ_3 and ϕ_4 defines the sign of Σ. If isovector channel is dominant in the hard pn rescattering then one expects $\Sigma > 0$ at $\theta_{\rm cm} = 90°$.

3.3. Numerical estimates

We discuss the numerical estimates for illustration purposes only. Since there are practically no available data on helicity pn amplitudes for hard scattering kinematics, we model them based on quark-interchange framework of the scattering and the fact that HRM predicts the dominance of isovector state NN rescattering at $\theta_{\rm cm} = 90°$. These two features are reflected in the following parameterization (see, e.g., [17,18,21]):

$$\begin{aligned}
\phi_1 &= \phi_1(0)\left[\frac{17}{62}\big(F(z_t)+F(z_u)\big)+\frac{14}{62}\big(F(-z_t)+F(-z_u)\big)\right],\\
\phi_3 &= \phi_3(0)\left[\frac{25}{94}\big(F(z_t)+F(z_u)\big)+\frac{22}{94}\big(F(-z_t)+F(-z_u)\big)\right],\\
\phi_4 &= \phi_4(0)\left[\frac{1}{4}\big(F(z_t)+F(z_u)\big)+\frac{1}{4}\big(F(-z_t)+F(-z_u)\big)\right],
\end{aligned}\tag{16}$$

where $\phi_i(0) \equiv \phi_i^{I=1}(0) \approx \phi_i^{pp}(\theta_{\rm cm} = 90°)$ $(i = 1,2,3,4)$ and the angular function is defined according to Ref. [21]: $F(z) = 1/[(1+z)(1-z)^3]$, with $z_t = 1 + \frac{2t_n}{s-4m^2}$, and $z_u = -1 - \frac{2u_n}{s-4m^2}$. We define ϕ_2 as:

$$\phi_2 = \frac{\phi_2(0)}{\phi_1(0)}\phi_1. \tag{17}$$

Because of (15) the observables of Eq. (13) depend weakly on the particular choice of ϕ_2. To asses the values of $\phi_{1,2,3,4}(0)$ we use the phenomenological parameterizations of [20], which successfully describe the available polarization and cross section data on hard pp scattering:

$$\begin{aligned}
&\phi_1(0) = \frac{\phi_+ + \phi_-}{\sqrt{2}}, \qquad \phi_2(0) = \frac{\phi_+ - \phi_-}{\sqrt{2}}, \qquad \phi_4(0) = -\phi_3(0),\\
&\phi_{\pm,3}(0) = \frac{N}{(s/Gev^2)^4}\big(B_{\pm,3} + C_{\pm,3}e^{i[\Psi_{\pm,3}(s)+\delta_{\pm,3}]}\big),
\end{aligned}\tag{18}$$

where $\phi_{\pm,3} = a\ln(s/\Lambda^2)/\ln(s/\Lambda_i^2)$ with $\Lambda \equiv \Lambda_{\rm QCD} = 0.2$ and all remaining parameters: B_i, C_i, a, Λ_i are defined in Ref. [20] (see Table 1 of [20]).

For ϕ_5 we use relation that ensures a vanishing value at $\theta_{\rm cm} = 90$ in the on-shell limit [21]

$$\phi_5 = R_{5-}\phi_1 + R_{5+}(\phi_3 + \phi_4), \tag{19}$$

where $R_{5\pm}$ is an angular factor defined similar to [21]:

$$R_{5\pm}(\hat t, \hat u) = \epsilon\left[\frac{1}{\sqrt{-\hat t}} \pm \frac{1}{\sqrt{-\hat u}}\right]. \tag{20}$$

We consider two values for ϵ: $\epsilon = \sqrt{(s-4m^2)/2}$ corresponding to the assumption that the smallness of ϕ_5 at large angles is related only to the condition: $\phi_5(\theta_{\rm cm} = 90°) = 0$, and $\epsilon \approx 0.1$—characteristic value obtained from the analysis of ϕ_5 for pp scattering which takes into account an additional suppression due to helicity conservation

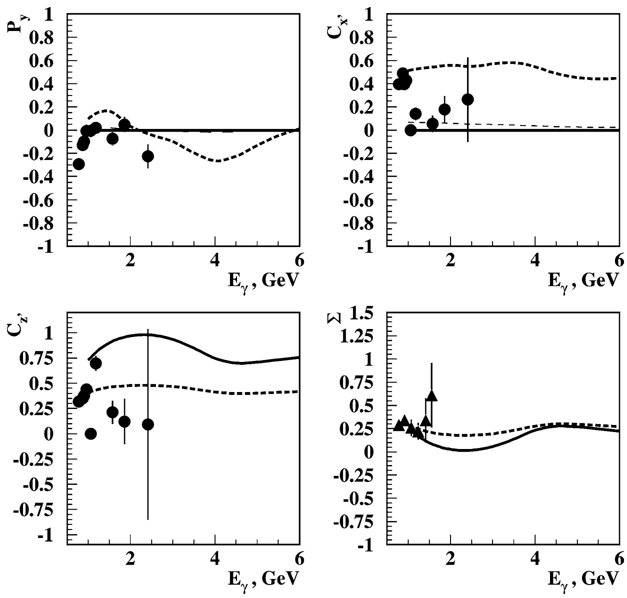

Fig. 2. The photon energy dependence of P_y, $C_{x'}$ $C_{z'}$ and Σ at $\theta_{\rm cm} = 90°$ photodisintegration of the deuteron. The curves are described in the text. The P_y, $C_{x'}$ and $C_{z'}$ data are from Ref. [5]. The Σ data are from Ref. [6].

[21]. Note that because of the overall smallness of ϕ_5 at large $\theta_{\rm cm}$ the unpolarized cross section is practically insensitive to the particular choice of ϵ.

In the hard regime when helicities are conserved ϕ_5 vanishes and its nonzero value is related mainly to the soft component of NN scattering (see, e.g., Ref. [19]). Therefore, the fact that one can identify the kernel of hard rescattering in Eq. (5) with the hard kernel of NN scattering does not justify the replacement of $\hat{t}$ and $\hat{u}$ in $R_{5\pm}$ by t_n and u_n. Furthermore, we will refer such a replacement as an "on shell" approximation for ϕ_5. Additionally, we consider an "off-shell" approximations in which in the first case ("off-shell I") we identify

$$\hat{t} = -\frac{s-4m^2}{2}(1-z_t) \quad \text{and} \quad \hat{u} = -\frac{s-4m^2}{2}(1+z_t)$$

and in the second case ("off-shell II")

$$\hat{t} = -\frac{s-4m^2}{2}(1+z_u) \quad \text{and} \quad \hat{u} = -\frac{s-4m^2}{2}(1-z_u).$$

Note that these are only choices which satisfies the condition, $|\hat{t}| < |\hat{u}|$ at $\theta_{\rm cm} < 90$ (forward angles). The above ambiguity naturally disappears in the on-shell limit.

Fig. 2 demonstrates the HRM predictions for energy dependences of P_y, $C_{x'}$ $C_{z'}$ and Σ at $\theta_{\rm cm} = 90°$. Thick and thin curves represent the calculations with parameter ϵ in Eq. (20) chosen $\sqrt{(s-4m^2)/2}$ and 0.1, respectively. Solid and dashed curves correspond to the "on-shell" and "off-shell" approximation for ϕ_5. Note that at $\theta_{\rm cm} = 90°$ both off-shell approximations give an identical results. According to Eq. (13) the "on-shell" approximation predicts P_y and $C_{x'}$ to be exactly zero at $\theta_{\rm cm} = 90°$. Thus vanishing P_y and $C_{x'}$ do not indicate unambiguously the onset of helicity conservation regime. The existing data do not rule out the large values for helicity flip amplitudes (thick curves). It is interesting to note that within HRM the small value of $C_{z'}$ favors a nonvanishing contribution from ϕ_2 and ϕ_5. Thus the accurate measurement of $C_{z'}$ will have an utmost importance.

The "on-shell" and "off-shell" approximations can be discriminated unambiguously through the study of angular dependences of the observables of Fig. 2. Fig. 3 demonstrates such a dependence for the reaction with $E_\gamma = 4$ GeV.

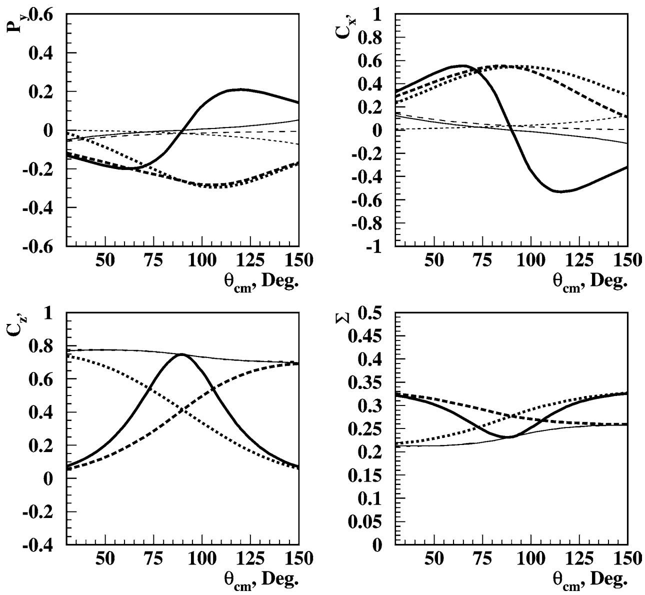

Fig. 3. The prediction of $\theta_{\rm cm}$ dependence of P_y, $C_{x'}$ $C_{z'}$ and Σ at $E_\gamma = 4$ GeV photodisintegration of the deuteron. The definition of the curves are the same as in Fig. 2.

The definition of the curves are the same as for Fig. 1, with dashed and doted curves representing "off-shell I" and "off-shell II" approximations. HRM predicts a qualitatively different dependences for P_y, $C_{x'}$ and $C_{z'}$ for "on-shell" and "off-shell" approximations of ϕ_5, when no additional suppression due to helicity conservation is assumed ($\epsilon = (s-4m^2)/2$) (thick curves). If the regime of helicity conservation is established then the difference between "on-shell" and "off-shell" approximations become unimportant (thin curves) and in both cases HRM predicts a vanishing values for P_y and $C_{x'}$. The dominance of the isovector channel in hard NN rescattering is reflected in the positive asymmetry of Σ.

4. Summary

Polarization observables in $\gamma D \to pn$ have been studied within the hard rescattering mechanism of deuteron photodisintegration. Within this model P_y, $C_{x'}$ $C_{z'}$ and Σ asymmetries are expressed through the helicity amplitudes of hard pn scattering. At $\theta_{\rm cm} = 90°$ HRM predicts a dominance of the isovector channel in the hard pn reinteraction.

Based on the general constraints on NN helicity amplitudes we predict several qualitative features of P_y, $C_{x'}$ $C_{z'}$ and Σ. These are the vanishing values of P_y, $C_{x'}$ at $\theta_{\rm cm} = 90°$ due to $\phi_5^{I=1}(\theta = 90°)$, positive large value for $C_{z'}$ if helicity-conserving regime is established, as well as a positive sign for Σ.

Within the quark-interchange framework we model the pn helicity amplitudes expressing unknown parameters through the existing parameterization of pp amplitudes. Our numerical predictions are in reasonable agreement with the existing data, indicating that the available data are not sufficient to relate unambiguously the observed smallness of P_y, $C_{x'}$ to the onset of the helicity-conserving regime. Within HRM this smallness can be explained rather by the vanishing ϕ_5 amplitude for NN scattering at 90° in isovector channel. On the other hand, the vanishing helicity-nonconserving amplitudes within HRM predict a sizable asymmetry for $C_{z'}$. Thus it is very important to have an accurate measurement of $C_{z'}$. In addition, the study of the angular dependences of P_y, $C_{x'}$ and $C_{z'}$ will clarify unambiguously the question whether the smallness of P_y, $C_{x'}$ is related to the vanishing ϕ_5 at $\theta_{\rm cm} = 90°$

or the onset of helicity-conserving regime of high energy scattering. The experimental verification of the sign of Σ will check HRM observation that $\theta_{\rm cm} = 90^\circ$ scattering is dominated by hard pn rescattering in the isovector channel.

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