$b_1$ and the Angular Momentum Sum Rule in the Deuteron

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Abstract. We formulate the generalized deep inelastic tensor spin structure of the deuteron which can be obtained from deeply virtual Compton scattering and meson production experiments. We discuss its connection to the total quark angular momentum sum rule for a spin-one hadronic system within a gauge invariant decomposition of hadronic spin.

1. Introduction
In Ref.[1] a sum rule for the deuteron angular momentum was derived. The sum rule involves the vector polarization aspects of the deuteron, namely

$$J_q = \frac{1}{2} \int dx \, x H_5^q(x, 0, 0),$$

(1)

where $H_5^q(x, \xi, t)$ is one of the five Generalized Parton Distributions (GPDs), $H_i$, $i = 1, 5$, in terms of which the vector component of the generalized deuteron quark-quark correlator is parametrized ($x$ is the struck quark’s momentum fraction, $\xi$ is the skewness parameter, $t$ is the four-momentum transfer squared between the initial and final protons) [2]. The first Mellin moment of $H_5^q$ gives the quarks contribution to the deuteron magnetic form factor $G_2(t) \equiv G_M(t)$. Eq.(1) can be compared to the corresponding relation for the nucleon [3],

$$J_q = \frac{1}{2} \int dx \, x [H_q(x, 0, 0) + E_q(x, 0, 0)],$$

(2)

where the first moment of the GPD sum, $H_q(x, \xi, t) + E_q(x, \xi, t)$, is the quark contribution to the nucleon magnetic form factor, $F_1(t) + F_2(t) \equiv G_M(t)$. Similarly to the proton GPD $E_q$, $H_5^q$ does not have a forward partonic limit.

The deuteron spin structure provides a much richer context to study hadronic spin. Its deep inelastic tensor polarized structure is described at leading twist by the structure function $b_1 \equiv H_5(x, 0, 0)$ [4]. $b_1$ is equal to zero for two nucleons moving independently from one another. Therefore it can be seen as a more straightforward measure of nuclear modifications in hard processes. An initial experimental measurement from HERMES [5] in the kinematic region $0.01 < x < 0.45$ shows that both $b_1$ and the integral $\int b_1(x)dx$ are sensibly different from zero, thus suggesting both the presence of non trivial nuclear effects and of tensor polarized sea quarks in the deuteron [6]. In this contribution, we explore the off-forward deuteron tensor polarized
structure with the aim of making a connection to the puzzles associated with the partonic interpretation of hadronic spin. We give a brief overview of the angular momentum sum rule in Section 2, in Section 3 we describe the angular momentum sum rule for the deuteron, in Section 4 we give a formulation of the generalized tensor structure function $H_5$ in terms of the nucleon components. In particular, our formulation uncovers new contributions from the nucleon helicity flip GPDs in the deuteron. In Section 5 we draw our conclusions.

2. Generalized Parton Distributions and Orbital Angular Momentum

One of the issues driving the research program at Jefferson Labs 12 GeV upgrade is the investigation of the various components of the proton’s angular momentum sum rule. Experiments performed since the ’80s, have, in fact, confirmed that only about 30% of the proton spin is accounted by quarks spin, and that the quark contribution is dominated by the valence component (see review in [7]). Current efforts, both in theory and experiment, are therefore directed towards determining the contributions of the orbital angular momentum (OAM) of the quarks, as well as of the spin and OAM of the gluons.

The sum rule relating the energy momentum tensor’s (EMT) form factors to the nucleon angular momentum components reads [3]

$$\frac{1}{2} = \sum q J_q + J_g, \quad (3)$$

where $J_q$ and $J_g$ are the quarks and gluons total angular momenta. One can subsequently consider the decomposition

$$J_q = L_q + S_q, \quad (4)$$

where the longitudinal component, $S_q = \frac{1}{2} \Delta \Sigma q$, is the quark’s helicity, while the gluon angular momentum cannot be separated into its spin and orbital components gauge invariantly. In an alternative, precedent formulation [8], it was pointed out that on the light front one can write the longitudinal component as

$$\frac{1}{2} = \frac{1}{2} \Delta \Sigma q + L_q + \Delta G + L_g, \quad (5)$$

where $\Delta \Sigma q$ is the same quark helicity term appearing in Eq.(4), and the other terms are identified with the quarks orbital angular momentum and the gluons spin and orbital angular momentum, respectively. $J_q$ in Eqs.(3) and (4) is connected to the GPDs which can be observed in deeply virtual compton scattering (DVCS) experiments through Eq.(2). Eq.(1) was derived similarly by working out the connection between the moments of the GPDs appearing in the parametrization of the hadronic tensor for a spin 1 system [2], and the form factors of the EMT [1], which is further discussed in Section 3.

An important issue is how to observe the quarks and gluons orbital angular momentum components in either decomposition from an independent measurement. In [9] $L_q$ was shown to be associated to a twist three GPD,

$$L_q = - \int_{-1}^{1} dx G_2(x, 0, 0). \quad (6)$$

On the other hand, several issues remain to be solved before $L_{q(g)}$ can be uniquely identified with observables.

In this context, exploring the various angular momentum components in the deuteron provides yet another handle on the hadronic spin puzzle. Complementary information can, in fact, be obtained both by accessing different polarization states that are not allowed in a spin 1/2 target, and because of the isoscalar property leading to specific cancellations of the $u$ and $d$ quarks contributions.
3. **Angular momentum in a Spin 1 target**

In Ref.[1] we connected the deuteron GPDs to angular momentum by considering the parametrization of the matrix element of the deuteron EMT, $T_{q,g}^{\mu\nu}$, in terms of a new set of gravitational form factors,

$$
\langle p'|T_{\mu\nu}|p \rangle = -\frac{1}{2} P^\mu P^\nu (\epsilon^* \epsilon) G_1(t) - \frac{1}{4} P^\mu P^\nu (\epsilon^* \epsilon) \frac{(e^* P)}{M^2} G_2(t) 
- \frac{1}{2} \left[ (\epsilon^* \epsilon) G_3(t) - \frac{1}{4} \left[ (\epsilon^* \epsilon) \frac{(e^* P)}{M^2} G_4(t) \right] \right] 
+ \frac{1}{4} \left[ \left( (\epsilon^* \epsilon) - \frac{\epsilon^* (e^* P)}{P^2} \right) \frac{(e^* P)}{M^2} G_5(t) \right]
+ \frac{1}{4} \left[ \left( (\epsilon^* \epsilon) - \frac{\epsilon^* (e^* P)}{P^2} \right) \frac{(e^* P)}{M^2} G_6(t) \right]
+ \frac{1}{2} \left[ (\epsilon^* \epsilon) \frac{(e^* P)}{P^2} \right] G_7(t) + g_{\mu\nu} (\epsilon^* \epsilon) M^2 G_8(t)
$$

where $P = p + p'$ and $\Delta = p' - p$, and $t = \Delta^2$. $\epsilon^\mu(\Lambda), \epsilon^\mu(\Lambda')$ are the polarization vectors of the deuteron in the initial and final helicity states, respectively.

The various EMT form factors can be then expressed as combinations of second Mellin moments of the deuteron GPDs, proceeding analogously to the spin 1/2 case. The deuteron GPDs were originally identified in [2]. Following Ref.[2], we write the vector component of the deuteron correlator as

$$W_{\Lambda \Lambda'}^{\gamma^+} = \int \frac{dx}{2\pi} e^{ixP^+z^-} \langle p', \Lambda' | \psi\left(-\frac{z^+}{2}\right) \gamma^+ \frac{(z^-)}{2} | p, \Lambda \rangle |_{z^+ = 0, x_T = 0}$$

$$= -(\epsilon^* \epsilon) H_1 + \frac{(en)(e^* P) + (\epsilon^* n)(eP)}{(Pn)} H_2 - \frac{(eP)(e^* P)}{2M^2} H_3$$

$$+ \frac{(en)(e^* P) - (\epsilon^* n)(eP)}{(Pn)} H_4 + \left\{ 4M^2 \frac{(en)(e^* n)}{(Pn)^2} + \frac{1}{3} (\epsilon^* \epsilon) \right\} H_5,$$

where $\Lambda, \Lambda'$ are the deuteron helicities, and $n$ is a light-like vector. The first Mellin moments of the deuteron GPDs satisfy the following relations:

$$\int dx H_1(x, \xi, t) = G_C(t) - \frac{2}{3} \frac{-t}{4M_D^2} G_Q,$$

$$\int dx H_2(x, \xi, t) = G_M(t),$$

$$\int dx H_3(x, \xi, t) = \frac{1}{1-\frac{t}{4M_D^2}} \left[ G_M(t) - G_C(t) + \left(1 + \frac{t}{3} \frac{-t}{4M_D^2} \right) G_Q(t) \right] \right),$$

$$\int dx H_4(x, \xi, t) = 0,$$

$$\int dx H_5(x, \xi, t) = 0,$$

where $G_C, G_M$, and $G_Q$ are the deuteron charge, magnetic and quadrupole form factors. Notice that Eq.(9e) is valid assuming a zero tensor polarization of the antiquarks.
For the second moments, one finds [1]

\[
\begin{align}
2 & \int_{-1}^{1} dx x [H_1(x, \xi, t) \pm \frac{1}{3} H_5(x, \xi, t)] = G_1(t) + \xi^2 G_3(t), \\
2 & \int_{-1}^{1} dx x H_2(x, \xi, t) = G_5(t), \\
2 & \int_{-1}^{1} dx x H_3(x, \xi, t) = G_2(t) + \xi^2 G_4(t), \\
-4 & \int_{-1}^{1} dx x H_4(x, \xi, t) = \xi G_6(t), \\
2 & \int_{-1}^{1} dx x H_5(x, \xi, t) = \frac{t}{4M_D^2} G_6(t) + G_7(t).
\end{align}
\]

Several relations and observations follow. Eq.(10b) leads to the angular momentum relation, Eq.(1). The total quarks momentum is described in the forward limit by the form factor $G_1$, which is analogous to the EMT form factor, $A_q$, in the spin 1/2 case [3]. However, differently from the spin 1/2 case, the form factor $G_1(t)$ does not simultaneously enter the angular momentum sum rule. Because of parity and time reversal symmetry properties, $G_7(t)$ is zero at $t = 0$. Therefore we obtain the result

\[
\int_{-1}^{1} dx x H_5(x, 0, 0) = \int dx x b_1(x) = 0
\]

for the second moment of $b_1$. Although initial experimental results from HERMES are consistent with the presence of a node in $xb_1(x)$ [5], this sum rule could be verified with future dedicated experiments in the large $x$ region [10].

4. Off-forward tensor polarized deuteron structure

The correlator in Eq.(8) is expressed in terms of kinematical variables defined in the “symmetric frame”, where we define: $P = (p + p')/2$, the average proton momentum, and $\Delta = p - p'$. $P$ is along the z-axis with momentum, $P_3 \approx P_z$. In the symmetric frame, the deuteron polarization vectors are obtained by rotating $\epsilon_{p'}^n \equiv 1/\sqrt{2}(0; \pm 1, -i, 0)$ to the planes transverse to $p$ and $p'$, respectively. Connecting the correlator in Eq.(8) to the helicity amplitudes for DVCS from a spin-1 hadron allows one to identify observables that would give access to the various deuteron GPDs, and, in particular, the tensor polarized structure function.

The amplitudes for photon-deuteron scattering used in DVCS are given by,

\[
f_{\Lambda, \Lambda', \Lambda_\gamma, \Lambda'} = \sum_{\lambda, \lambda'} g^{\Lambda, \Lambda_\gamma}_{\lambda, \lambda'}(x, \xi, t, Q^2) \otimes A_{\Lambda, \Lambda', \Lambda_\lambda}(x, \xi, t),
\]

where $\Lambda_\gamma$ denotes the helicity of the initial (virtual) photon, $\Lambda_{\lambda'}$ denotes the helicity of the outgoing photon; $\Lambda(\Lambda')$ the helicity of the incoming (outgoing) deuteron; $\lambda(\lambda')$ the helicity of the initial (final) quark. In Eq.(12) we describe the factorization into the hard scattering part, $g^{\Lambda, \Lambda_\gamma}_{\lambda, \lambda'}$ for the partonic subprocess $\gamma^* + q \rightarrow \gamma + q$, and the quark-proton helicity amplitudes, $A_{\Lambda, \Lambda', \lambda}$ that contain the GPDs.

The subprocess amplitude is given by

\[
g^{\Lambda, \Lambda_\gamma}_{\lambda, \lambda'}(x, \xi) = \left[ \bar{u}(k', \lambda') \gamma^\mu \gamma^\nu u(k, \lambda) \right] \left( \frac{\epsilon_{\lambda'} \cdot \epsilon_\mu}{s - i\epsilon} + \frac{\epsilon_{\mu} \cdot \epsilon_{\nu}}{\bar{u} - i\epsilon} \right) q^-,
\]

where $s$ is the invariant mass of the quark-proton pair, $q^-$ is the momentum of the outgoing quark. The factors $\epsilon_{\lambda'} \cdot \epsilon_\mu / (s - i\epsilon)$ and $\epsilon_{\mu} \cdot \epsilon_{\nu} / (\bar{u} - i\epsilon)$ represent the quark helicity amplitudes for DVCS from a spin-1 hadron and contain the GPDs.
where, in the Bjorken limit, \((k + q) \approx \gamma^+ q^\prime\), \(s = (k + q)^2 \approx Q^2(X - \zeta)/\zeta\) and \(\hat{u} = (k' - q)^2 \approx Q^2X/\zeta\), and \(q^- \approx (Pq)/P^+ = Q^2/(2\zeta P^+)\), with \(\zeta = 2\xi/(1 + \xi)\), and \(X = (x + \xi)/(1 + \xi)\).

For DVCS one can consider either the sum over the transverse helicities of the initial photon (unpolarized), or the difference of the helicities in which case the transverse photon is polarized. The outgoing photon is on-shell, thus purely transverse or helicity \pm 1. So the leading incoming virtual photon will have the same helicity in the collinear limit. Only \(g^+_{+\gamma} (= g^-_{-\gamma}\) via Parity conservation) for the direct \(\hat{s}\) pole term or \(g^+_{+\gamma} (= g^-_{-\gamma})\) for the crossed \(\hat{u}\) pole term will survive.

For either allowed combination, \(g^{S(A)}_{++} = g^{++}_{++} + g^{--}_{++}\) and \(g^{A}_{++} = g^{++}_{++} - g^{--}_{++}\) such that

\[
g^{S(A)} = g^{++}_{++} \pm g^{--}_{++} = \sqrt{X(X - \zeta)} \left( \frac{1}{X - \zeta + i\epsilon} \pm \frac{1}{X - i\epsilon} \right) , \tag{14}\]

From the convolution in Eq.(12) we obtain the following decomposition of the transverse photon helicity amplitudes, that we divide into symmetric (S) and antisymmetric (A) components:

\[
f^{S}_{++} = g^{S}_{++} \otimes (A_{++,++} + A_{+-,-+}) \tag{15a} \]
\[
f^{S}_{+-} = g^{S}_{++} \otimes (A_{--,++} + A_{+-,-+}) \tag{15b} \]
\[
f^{S}_{00} = g^{S}_{++} \otimes 2A_{0+,0+} \tag{15c} \]
\[
f^{S}_{0+} = g^{S}_{++} \otimes (A_{0+,0+} + A_{+-,0-}) \tag{15d} \]
\[
f^{S}_{0-} = g^{S}_{++} \otimes (A_{0+,0+} + A_{+-,0-}) \tag{15e} \]

and

\[
f^{A}_{++} = g^{A}_{++} \otimes (A_{++,++} - A_{+-,-+}) \tag{16a} \]
\[
f^{A}_{+-} = g^{A}_{++} \otimes (A_{--,++} - A_{+-,-+}) \tag{16b} \]
\[
f^{A}_{0+} = g^{A}_{++} \otimes (A_{0+,0+} - A_{+-,0-}) \tag{16c} \]
\[
f^{A}_{0-} = g^{A}_{++} \otimes (A_{0+,0+} - A_{+-,0-}) . \tag{16d} \]

The amplitudes \(A_{\Lambda',\Lambda,\Lambda'}\) implicitly contain an integration over the unobserved quark’s transverse momentum, \(k_T\), and are functions of \(x_{ Bj} = Q^2/2M\nu \approx \zeta\), \(t\) and \(Q^2\). The convolution integral in Eq.(12) is given by \(\otimes \rightarrow \int^1_0 dx\).

The connection with the correlator is carried out by considering,

\[
A_{\Lambda',\Lambda,\Lambda'} = \int \frac{dz^-}{2\pi} e^{ixP^+ z^-} \langle p', \Lambda' | \mathcal{O}_{\Lambda,\Lambda}(z) | p, \Lambda \rangle \big|_{z^+ = 0, \zeta = 0} , \tag{17}\]

\[
\mathcal{O}_{\pm \pm}(z) = -i\bar{\psi} \left( -\frac{z}{2} \right) \gamma^+(1 \pm \gamma_5) \psi \left( \frac{z}{2} \right) \tag{18}\]

The vector correlator, \(W^{++}_{\Lambda',\Lambda'} = V_{\Lambda,\Lambda'}\), is given by the symmetric combinations \(A_{\Lambda',\Lambda,\Lambda'} + A_{\Lambda,\Lambda',\Lambda} - A_{\Lambda',\Lambda',\Lambda} + A_{\Lambda,\Lambda,\Lambda'}\).
corresponding to Eqs.(15), namely,

\begin{align}
V_{++} &\equiv f_{++}^S = H_1 - \frac{\Delta_T^2}{2M^2} \frac{1}{1 - \xi^2} H_3 - \frac{1}{3} H_5 \\
V_{+-} &\equiv f_{+-}^S = -\frac{\Delta_T^2}{2M^2} \frac{1}{1 - \xi^2} H_3 \\
V_{00} &\equiv f_{00}^S = -\frac{1 + \xi^2}{1 - \xi^2} \left( -1 + \frac{\Delta_T^2}{4M^2} \right) H_1 + \left[ 2(1 - \xi^2) + 2 \left( -1 + \frac{\Delta_T^2}{4M^2} \right) \frac{1}{1 - \xi^2} \right] \frac{\Delta_T^2}{4M^2} \right) \frac{1}{1 - \xi^2} \right] H_3 \\
V_{0+} &\equiv f_{0+}^S = -\frac{\Delta_T}{2M} \frac{1 + \xi^2}{1 - \xi^2} H_1 + \left[ \frac{\Delta_T}{2M} \frac{4\xi}{1 - \xi^2} H_2 - \frac{1}{3} \frac{\Delta_T}{2M} \frac{1 + \xi^2}{1 - \xi^2} \right] \frac{\Delta_T^2}{4M^2} \right) \frac{1}{1 - \xi^2} \right] H_3 \\
V_{0+} &\equiv f_{0+}^S = \frac{\Delta_T}{2M} \frac{1 + \xi^2}{1 + \xi} \left[ M(1 - \xi) + M \left( 1 + \frac{\Delta_T^2}{4M^2} \right) \frac{1}{1 - \xi} \right] H_3 + \frac{1}{3} \frac{\Delta_T}{2M} \frac{1 + \xi^2}{1 + \xi} H_5.
\end{align}

In the forward limit (\(\xi = 0, t = 0\)), these expressions reduce to

\begin{align}
V_{++} &\equiv f_{++}^S = H_1 - \frac{1}{3} H_5 \\
V_{+-} &\equiv f_{+-}^S = 0 \\
V_{00} &\equiv f_{00}^S = H_1 + \frac{2}{3} H_5 \\
V_{0+} &\equiv f_{0+}^S = 0 \\
V_{0+} &\equiv f_{0+}^S = 0.
\end{align}

From the above relations, we can solve for \(H_5\), in terms of the helicity structures,

\[ H_5(x, 0, 0) = V_{00} - V_{++}. \]

To explain the partonic structure of the deuteron helicity amplitudes we start from a picture in terms of bound nucleons. The quark-deuteron helicity amplitudes appearing in Eq.(15) are now labeled as \(A_D^{\Lambda, \lambda_q, \lambda_N, \lambda_q}\). We define them in terms of a convolution of quark-nucleon amplitudes, \(A_N^{\lambda_q, \lambda_q, \lambda_N, \lambda_q}\), and the nucleon-deuteron amplitudes, \(B^{\Lambda, \lambda_q, \lambda_N, \lambda_N}\), where the integral is taken over the unobserved nucleon momenta as

\[ A_D^{\Lambda, \lambda_q, \lambda_N, \lambda_q} = \sum_{N=n, p, \lambda_N, \lambda_N} B^{\Lambda, \lambda_q, \lambda_N, \lambda_N} \otimes A_N^{\lambda_q, \lambda_q, \lambda_N, \lambda_q}, \]

with \(\Lambda, \lambda_N, \lambda_q\), being the deuteron, nucleon, and quark helicities, respectively (see Figure 1).
The deuteron helicity amplitudes are written as,

\[ V_{++} = A_{++}^D + A_{++}^D = \sum_{N=n,p} \sum_{\lambda_N} \sum_{\lambda_N'} B_{+\lambda_N,+\lambda_N} \otimes (A_{++}^N + A_{++}^N) \]

\[ = \sum_{N=n,p} [(B_{++,++} + B_{+-,+-}) \otimes (A_{++}^N + A_{++}^N)] \]

\[ V_{00} = A_{0+0+}^D = \sum_{N=n,p} \sum_{\lambda_N} \sum_{\lambda_N'} B_{0\lambda_N,0\lambda_N} \otimes (A_{++}^N + A_{--}^N) \]

\[ = \sum_{N=n,p} B_{0+0+} \otimes (A_{++}^N + A_{--}^N) \]

The deuteron helicity amplitudes are written as,

\[ B_{\lambda_N^N',\lambda_N^N} = \sum_{\lambda_N^2} \chi_{\lambda_N^N} \chi_{\lambda_N^N}^{\lambda_N^2} (z', p') \chi_{\lambda_N^N} \chi_{\lambda_N^N}^{\lambda_N^2} (z, p), \]  

where \( \lambda_N \) (\( \lambda_N' \)) are the initial (returning) nucleons’ helicities, \( \lambda_{N2} \) is the spectator nucleon one, and \( \chi_{\lambda_N^N} \chi_{\lambda_N^N}^{\lambda_N^2} (z, p) \) is the deuteron wave function [11, 12],

\[ \chi_{\lambda_N^N} \chi_{\lambda_N^N}^{\lambda_N^2} (z, p) = \mathcal{N} \sum_{L,m_L,m_S} \left( \begin{array}{ccc} j_1 & j_2 & 1 \\ \lambda_{N1} & \lambda_{N2} & m_S \end{array} \right) \left( \begin{array}{ccc} L & S & J \\ m_L & m_S & \Lambda \end{array} \right) Y_{Lm_L} \left( \frac{p}{pl} \right) u_L(p). \]  

In Eq.(34), \( j_1 = j_2 = 1/2, S = J = 1 \). \( Y_{Lm_L} \), where \( L = 0, 2 \), can be written in terms of the nucleon’s light cone variables, consistently with the formalism for describing deep inelastic processes from nuclear targets [13] in the approximation where the quarks’ \( k_\perp \) dependence is trivially integrated over, and no off-shell effects are considered [14].

By evaluating the various expressions in Eqs.(31,32) one has that:

- \((A_{++}^N + A_{++}^N) \propto H^N \) and \((A_{++}^N + A_{++}^N) \propto E^N \).
- \((B_{++} + B_{++}) \) contains both the \( S \) and \( D \) wave contributions, while \( B_{0+0+} \) contains only \( S \) wave contributions. Both terms multiply the same GPD, \( H^N \).
- A new term \((B_{++} + B_{++}) \), describing the bound nucleon’s helicity flip, and multiplying the GPD \( E^N \), appears in the off-forward configuration.
Therefore in the forward limit, due to the cancellation of the $S$ wave contributions between the $V_{++}$ and $V_{0+}$ terms, $b_1 \equiv H_5(x,0,0)$ is solely proportional to the $D$ wave contribution in the deuteron light cone wave function.

However, in off-forward kinematics we predict a modification and possible enhancement of the tensor structure function, $H_5(x,\xi,t)$ because of the appearance of the nucleon spin flip term $\propto E^N$ in $V_{++}$.

5. Conclusions

In conclusion, the study of deuteron polarization observables in the exclusive channels can offer unique insights in the context of the proton spin puzzle where observables are currently being sought for possible direct measurements of the quarks and gluons orbital angular momentum components.

In this contribution we discussed the relations between the deuteron energy momentum tensor form factors and GPDs. The angular momentum involves the GPD $H_2$ that is related to the magnetic properties of the deuteron. Therefore, measurements of the angular momentum sum rule in the deuteron involve DVCS type observables sensitive to the vector polarization.

We also derived a relation for the second moment of the tensor structure function: $xb_1(x)$ integrates to zero at $t = 0$, thus suggesting the presence of a node in this function. Furthermore, we analyzed the helicity structure of $H_5(x,\xi,t)$ and explained that while the dominance of the $D$ wave contribution for this function persists in the off-forward limit, we also uncover a new term that would not be present in inclusive scattering, and is proportional to the bound nucleon spin flip amplitude. Different theoretical scenarios can be envisaged. If this term turns out to be large, it will open investigations of the off-forward tensor structure of the deuteron beyond the $D$ wave component of its wave function. If instead it is small as a consequence of the deuteron isoscalarity ($E_N = E_p + E_n$ is also expected to be small), this would trigger a different viewpoint from where to investigate the off-forward deuteron tensor structure, which would be directly sensitive to the gluon contributions (see also e.g. [15]).

A future program of measurements of polarized observables in DVCS from the deuteron stands therefore as an exciting possibility.

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