New Structure in the Deuteron

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We demonstrate that a paradigm shift from considering the deuteron as a system of a bound proton and neutron to that of a pseudovector system in which we observe a proton and neutron results in the possibility of probing a new “incomplete” P-statelike structure on the light front (LF). This occurs at extremely large internal momenta, which can be achieved in a high energy transfer electrodisintegration of the deuteron. Investigating the deuteron on the light front, where the vacuum fluctuations are suppressed, we found that this new structure, together with the conventional S and D states, is leading order in transferred energy of the reaction and thus not suppressed on the light front. The incompleteness of the observed P state results in a violation of the angular condition that can happen only if the deuteron contains non-nucleonic structures, such as $\Delta\,\Delta$, N^*N or hidden color components. We demonstrate that experimentally verifiable signatures of incomplete P states are angular anisotropy of the light front momentum distribution of the nucleon in the deuteron, as well as an enhancement of the tensor polarization strength beyond the S- and D-wave predictions at large internal momenta in the deuteron.

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One of the outstanding issues of strong interaction physics is understanding the dynamics of the transition between hadronic to quark-gluon phases of matter. Such transitions at high temperature are relevant to the evolution of the Universe after the big bang and can be studied experimentally in heavy ion collisions. Transitions at low (near zero) temperatures and high densities (“cold-dense” transitions) are relevant for superdense nuclear matter that can exist at the cores of neutron stars and can set the limits of matter density before it collapses to a black hole. However, direct exploration of cold-dense transitions is severely restricted.

Currently the accepted ways of investigating such transitions are (i) Studying the nuclear medium modification of quark-gluon structure of bound nucleons. Such a modification was discovered in 1983 by the European Muon Collaboration [1], commonly referred to as EMC effect. Little progress has been made in understanding this effect for past 40 years (for reviews, see Refs. [2,3]), including the observation of the dependence of the effect on the local nuclear density [4] and the important role of short range nucleonic correlations in the EMC effect for medium to large nuclei [5,6]. In all these cases, the role of the hadronic to quark-gluon transition is not clearly understood. (ii) Studying the implications of the transition of baryonic matter to quark matter in the cores of neutron stars. The situation with the existence of quark matter in the cores of neutron stars is even more unclear than with the EMC effect. With the observation of unexpectedly large neutron star masses [7] ($\approx 2.08 M_{\odot}$) it was expected that if such stars would have radii of $R < 10$ km it will be indicative of a large quark matter component in their cores. However, the observed radii for large mass neutron stars are above $R \geq 12$ km (e.g., Ref. [8]).

Currently, the progress in advancing the studies of the EMC effect is seen in performing a new generation of experiments in which the density of nuclear medium is controlled by tagging a recoil nucleon which is in a short range correlation with the probed nucleon (e.g., Ref. [9]). The neutron star studies rely on improving the detection techniques that will allow the identification of anomalously small-sized neutron stars.

In the present Letter, we are suggesting a new method of studying the baryon-quark transition using the simplest known atomic nucleus, the deuteron.

Deuteron on the light front (LF).—Our current mindset about the deuteron is fully nonrelativistic, within which, the observation that it has total spin, $J = 1$ and positive parity, P, together with the relation that for the nonrelativistic wave function, $P = (-1)^l$, one concludes that the deuteron consists of S- and D-partial waves for the proton-neutron system.

However, if we are interested in the deuteron structure at internal momenta comparable with the nucleon rest mass then a nonrelativistic framework is not valid and the problem is more fundamental, related to the description of a relativistic bound system. There were many important works to account for relativistic effects in the deuteron wave function (see, e.g., Refs. [10–12] and the reviews [13,14]). Our approach is similar to the one used in QCD for calculation of quark distributions in hadrons, in which the light-front description of the scattering process allows us to suppress vacuum fluctuations that overshadow the composite structure of the hadron [15].

To discuss the relativistic structure of the deuteron one needs to identify the process in which the deuteron structure is probed. In our case we consider the high-momentum transfer electrodisintegration process:

$$e + d \to e' + p + n, \tag{1}$$

in which one of the nucleons are struck by the incoming probe and the spectator nucleon is probed with momenta comparable to nucleon masses. If one can neglect (or remove) the effects related to final state interactions of two outgoing nucleons, then the above reaction at high Q^2 measures the probability of probing a proton and neutron in the deuteron at very large relative momenta. In such a formulation the deuteron is not a composite system consisting of a proton and neutron, but it is a composite pseudovector ($J = 1$, $P = +$) "particle" from which one extracts a proton and neutron. Thus we formulate the question not as how to describe relativistic motion of proton and neutron in the deuteron, but how such a proton and neutron are produced at such extreme conditions relating it to the dynamical structure of the LF deuteron wave function. In such formulation the latter may include internal elastic $pn \to pn$ as well as inelastic $\Delta\Delta \to pn$, $N^*N \to pn$ or $N_cN_c \to pn$ transitions. Here, Δ and N^* denote Δ isobar and N^* resonances, while N_c is a color octet baryonic state contributing to the hidden-color component in the deuteron. The framework for calculation of reaction (1) in the relativistic domain is the LF approach (e.g., Refs. [16–21]) in which one introduces the LF deuteron wave function

$$\psi_d^{\lambda_d}(\alpha_i, p_\perp, \lambda_1\lambda_2) = -\frac{\bar{u}(p_2,\lambda_2)\bar{u}(p_1,\lambda_1)\Gamma_d^\mu \chi_\mu^{\lambda_d}}{\frac{1}{2}(m_d^2 - 4\frac{m_N^2+p_\perp^2}{\alpha_i(2-\alpha_i)})\sqrt{2(2\pi)^3}}, \tag{2}$$

where $\alpha_i = 2(p_{i+}/p_{d+})$, $(i = 1, 2)$ and $\alpha_1 + \alpha_2 = 2$ are LF momentum fractions of two nucleons coming out from the deuteron that has four-momentum p_d^μ. Absorbing the energy denominator into the vertex function and using crossing symmetry one obtains

$$\begin{aligned}\psi_d^\mu(\alpha_i, p_\perp, \lambda_1, \lambda_2) &= -\bar{u}(p_2,\lambda_2)\Gamma_d^\mu(k)\frac{(i\gamma_2\gamma_0)}{\sqrt{2}}\bar{u}(p_1,\lambda_1)^T \\ &= -\sum_{\lambda_1'}\bar{u}(p_1,\lambda_1)\Gamma_d^\mu\gamma_5\frac{\epsilon_{\lambda_1,\lambda_1'}}{\sqrt{2}}u(p_1,\lambda_1'),\end{aligned} \tag{3}$$

where $u(p,\lambda)$'s are the LF bispinors of the proton and neutron [22] and $\epsilon_{i,j}$ is the two-dimensional Levi-Civita tensor, with $i, j = \pm 1$ helicity of the nucleon. Since the deuteron is a pseudovector particle, due to γ_5 in Eq. (3), the vertex Γ_d^μ is a four-vector which we can construct in a general form that explicitly satisfies time reversal, parity, and charge conjugate symmetries. Noticing that at the $d \to pn$ vertex on the light front the "−" ($p^- = E - p_z$) components of the four-momenta of the particles are not conserved, in addition to the four-momenta of two nucleons, p_1^μ and p_2^ν, one has the additional four-momentum

$$\Delta^\mu \equiv p_1^\mu + p_2^\mu - p_d^\mu \equiv (\Delta^-, \Delta^+, \Delta_\perp) = (\Delta^-, 0, 0), \tag{4}$$

where

$$\begin{aligned}\Delta^- = p_1^- + p_2^- - p_d^- &= \frac{m_N^2+k_\perp^2}{p_1^+} + \frac{m_N^2+k_\perp^2}{p_2^+} - \frac{M_d^2}{p_d^+} \\ &= \frac{1}{p_d^+}\left[\frac{4(m_N^2+k_\perp^2)}{\alpha_1(2-\alpha_1)} - M_d^2\right] = \frac{4}{p_d^+}\left[m_N^2 - \frac{M_d^2}{4} + k^2\right].\end{aligned} \tag{5}$$

Here k is the relative momentum in the pn center of mass (c.o.m.) system defined as

$$k = \sqrt{\frac{m_N^2 + k_\perp^2}{\alpha_1(2-\alpha_1)} - m_N^2} \quad \text{and} \quad \alpha_1 = \frac{E_k + k_z}{E_k}, \tag{6}$$

where $E_k = m^2 + k^2$. With p_1^μ, p_2^μ, and Δ^μ four-vectors the Γ_d^μ four-vector function is constructed in the following form:

$$\begin{aligned}\Gamma_d^\mu &= \Gamma_1\gamma^\mu + \Gamma_2\frac{(p_1-p_2)^\mu}{2m_N} + \Gamma_3\frac{\Delta^\mu}{2m_N} + \Gamma_4\frac{(p_1-p_2)^\mu \not{\Delta}}{4m_N^2} \\ &\quad + i\Gamma_5\frac{1}{4m_N^3}\gamma_5\epsilon^{\mu\nu\rho\gamma}(p_d)_\nu(p_1-p_2)_\rho(\Delta)_\gamma + \Gamma_6\frac{\Delta^\mu\not{\Delta}}{4m_N^2},\end{aligned} \tag{7}$$

where Γ_i, $(i = 1, 6)$ are scalar functions describing dynamics of the pn component being observed in the deuteron.

High energy approximation.—For the large Q^2 limit, the LF momenta for reaction (1) are chosen as follows:

$$\begin{aligned}p_d^\mu &\equiv (p_d^-, p_d^+, p_{d\perp}) = \left(\frac{Q^2}{x\sqrt{s}}\left[1 + \frac{x}{\tau} - \sqrt{1 + \frac{x^2}{\tau}}\right],\right. \\ &\qquad \left.\frac{Q^2}{x\sqrt{s}}\left[1 + \frac{x}{\tau} + \sqrt{1 + \frac{x^2}{\tau}}\right], 0_\perp\right) \\ q^\mu &\equiv (q^-, q^+, q_\perp) = \left(\frac{Q^2}{x\sqrt{s}}\left[1 - x + \sqrt{1 + \frac{x^2}{\tau}}\right],\right. \\ &\qquad \left.\frac{Q^2}{x\sqrt{s}}\left[1 - x - \sqrt{1 + \frac{x^2}{\tau}}\right], 0_\perp\right),\end{aligned} \tag{8}$$

where $s = (q + p_d)^2$, $\tau = (Q^2/M_d^2)$, and $x = (Q^2/M_d q_0)$, with q_0 being the virtual photon energy in the deuteron rest frame. The high energy nature of this process results in, $p_d^+ \sim \sqrt{Q^2} \gg m_N$. Then one observes in Eq. (5) that the Δ^- term is suppressed by the large p_d^+ factor.

Analyzing now the vertex function (7) one observes that $(\Delta^-/2m_N)$ is a small parameter. Here the Γ_3 and Γ_4 terms enter with order $\mathcal{O}^1(\Delta^-/2m_N)$, while the Γ_6 term enters as $\mathcal{O}^2(\Delta^-/2m_N)$. The situation with the Γ_5 term is, however, different; since the covariant components $\Delta_+ = \frac{1}{2}\Delta^-$ and $p_{d,-} = \frac{1}{2}p_d^+$, the term with $\epsilon^{\mu+\perp-}$ is leading order $[\mathcal{O}^0(\Delta^-/2m_N)]$ due to the fact that the large p_d^+ factor is canceled in the $p_{d,-}\Delta_+ = \frac{1}{4}p_d^+\Delta^-$ combination.

Keeping the leading, $\mathcal{O}^0(\Delta^-)$, terms in Eq. (7) the LF deuteron wave function reduces to [20,21]

$$\psi_d^{\lambda_d}(\alpha_i, p_\perp) = -\sum_{\lambda_2,\lambda_1,\lambda_1'} \bar{u}(p_2,\lambda_2)\left\{\Gamma_1\gamma^\mu + \Gamma_2\frac{(p_1-p_2)^\mu}{2m_N} + \sum_{i=1}^{2} i\Gamma_5 \frac{1}{8m_N^3}\epsilon^{\mu+i-}p_d^+ k_i \Delta^-\right\} \times \gamma_5 \frac{\epsilon_{\lambda_1,\lambda_i'}}{\sqrt{2}} u(p_1,\lambda_1')\chi_\mu^{\lambda_d}, \tag{9}$$

where $k_i = [(p_{1,i} - p_{2,i})/2]$, for $i = 1$, 2. The deuteron's polarization four-vector is chosen as

$$\chi_\mu^{\lambda_d} = (\chi_0^{\lambda_d}, \chi_\perp^{\lambda_d}, \chi_z^{\lambda_d}) = \left(\frac{p_{12}s_{d,z}}{M_{12}}, s_{s,\perp}, \frac{E_{12}s_{d,z}}{M_{12}}\right), \tag{10}$$

where $\mathbf{p_{12}} = (p_{1_z} + p_{2,z}, 0_\perp)$, $E_{12} = \sqrt{M_{12}^2 + p_{12}^2}$, and $M_{12}^2 = s_{NN} = 4[(m_N^2 + k_\perp^2)/\alpha_1(2-\alpha_1)]$.

Since the wave function in Eq. (9) is Lorentz boost invariant along the z axis, it is convenient to calculate it in the deuteron c.o.m. frame obtained by boosting with velocity $v = (\mathbf{p_{12}}/E_{12})$. Such a transformation results in [21]

$$\psi_d^{\lambda_d}(\alpha_i, k_\perp) = -\sum_{\lambda_2,\lambda_1,\lambda_1'} \bar{u}(-k,\lambda_2)\left\{\Gamma_1\gamma^\mu + \Gamma_2\frac{\tilde{k}^\mu}{m_N} + \sum_{i=1}^{2} i\Gamma_5 \frac{1}{8m_N^3}\epsilon^{\mu+i-}p_d'^+ k_i \Delta'^-\right\} \times \gamma_5 \frac{\epsilon_{\lambda_1,\lambda_i'}}{\sqrt{2}} u(k,\lambda_1')s_\mu^{\lambda_d}, \tag{11}$$

where $\tilde{k}^\mu = (0, k_z, k_\perp)$ with $k_\perp = p_{1\perp}$, $k^2 = k_z^2 + k_\perp^2$ and $E_k = (\sqrt{S_{NN}}/2)$ and $s_\mu^{\lambda_d} = (0, \mathbf{s_d^\lambda})$ in which

$$s_d^1 = -\frac{1}{\sqrt{2}}(1,i,0), \qquad s_d^1 = \frac{1}{\sqrt{2}}(1,-i,0)s_d^0 = (0,0,1). \tag{12}$$

In Eq. (11) "primed" variables correspond to the Lorentz boosts of the respective unprimed quantities:

$$p_d'^+ = \sqrt{s_{NN}}, \qquad \Delta'^- = \frac{1}{\sqrt{s_{NN}}}\left[\frac{4(m_N^2+k_\perp^2)}{\alpha_1(2-\alpha_1)} - M_d^2\right]. \tag{13}$$

Since the term related to Γ_5 is proportional to $[4(m_N^2 + k_\perp^2)/\alpha_1(2-\alpha_1)] - M_d^2$, which diminishes at small momenta, only the Γ_1 and Γ_2 terms will contribute in the nonrelativistic limit defining the S and D components of the deuteron. Thus, the LF wave function in Eq. (11) provides a smooth transition to the nonrelativistic deuteron wave function. This can be seen by expressing Eq. (11) through two-component spinors:

$$\psi_d^{\lambda_d}(\alpha_1, k_t, \lambda_1, \lambda_2) = \sum_{\lambda_1'} \phi_{\lambda_2}^\dagger \sqrt{E_k}\left[\frac{U(k)}{\sqrt{4\pi}}\sigma\mathbf{s_d^{\lambda_d}} - \frac{W(k)}{\sqrt{4\pi}\sqrt{2}}\left(\frac{3(\sigma\mathbf{k})(\mathbf{k}s_d^\lambda)}{k^2} - \sigma\mathbf{s_d^\lambda}\right) + (-1)^{\frac{1+\lambda_d}{2}}P(k)Y_1^{\lambda_d}(\theta,\phi)\delta^{1,|\lambda_d|}\right]\frac{\epsilon_{\lambda_1,\lambda_1'}}{\sqrt{2}}\phi_{\lambda_1'}. \tag{14}$$

Here the first two terms have explicit S and D structures where the radial functions are defined as

$$U(k) = \frac{2\sqrt{4\pi}\sqrt{E_k}}{3}\left[\Gamma_1\left(2 + \frac{m_N}{E_k}\right) + \Gamma_2\frac{k^2}{m_N E_k}\right],$$
$$W(k) = \frac{2\sqrt{4\pi}\sqrt{2E_k}}{3}\left[\Gamma_1\left(1 - \frac{m_N}{E_k}\right) - \Gamma_2\frac{k^2}{m_N E_k}\right]. \tag{15}$$

This relation is known for pn-component deuteron wave function [16,23], which allows us to model the LF wave function through known radial S and D wave functions evaluated at the LF relative momentum k defined in Eq. (6).

However, in addition to the S and D terms, our observation is that due to the Γ_5 term there is an additional leading contribution, which because of the relation $Y_1^\pm(\theta,\phi) = \mp i\sqrt{(3/4\pi)}\sum_{i=1}^2[(k\times s_d^{\pm 1})_z/k]$, has a P-wavelike structure, where the P- radial function is defined as

$$P(k) = \sqrt{4\pi}\frac{\Gamma_5(k)\sqrt{E_k}}{\sqrt{3}}\frac{k^3}{m_N^3}. \tag{16}$$

It is worth emphasizing that this term is purely relativistic in origin: as it follows from Eq. (16) it has an extra (k^2/m_N^2) factor in addition to the $(k^{l=1}/m_N)$ term characteristic to the radial P wave. As a result, one has a smooth transition to S and D states in the nonrelativistic limit.

The interesting feature of the above result, which we will discuss in the next section, is that the P wave is "incomplete," that is it contributes only for $\lambda_d = \pm 1$ polarizations of the deuteron.

Closing this section we would like to mention that the consideration of six invariant vertex functions and the contribution of P-radial waves in the relativistic description of the deuteron were discussed earlier in the literature; see, e.g., Refs. [10,24]. However, to the best of our knowledge, the observation that Γ_5 is a leading term on the light front (while $\Gamma_{3,4,6}$ terms are suppressed) in the high energy limit and it results in a noncomplete P-wave contribution are original results of the present Letter.

Light front density matrix of the deuteron.—Using Eq. (14) one defines the unpolarized deuteron light-front density matrix in the form [2,16]

$$\rho_d(\alpha, k_\perp) = \frac{n_d(k, k_\perp)}{2-\alpha}, \tag{17}$$

where the LF momentum distribution is expressed through the radial wave functions as follows:

$$n_d(k,k_\perp) = \frac{1}{3}\sum_{\lambda_d=-1}^{1}|\psi_d^{\lambda_d}(\alpha,k_\perp)|^2 = \frac{1}{4\pi}\left(U(k)^2 + W(k)^2 + \frac{k_\perp^2}{k^2}P^2(k)\right). \quad (18)$$

The LF density matrix satisfies the baryonic and momentum sum rules as follows:

$$\int \rho_d(\alpha,k_\perp)\frac{d\alpha}{\alpha} = 1 \quad \text{and} \quad \int \alpha\rho_d(\alpha,k_\perp)\frac{d\alpha}{\alpha} = 1. \quad (19)$$

From the above, the normalization condition for the radial wave functions is

$$\int \left(U(k)^2 + W(k)^2 + \frac{2}{3}P^2(k)\right)k^2 dk = 1. \quad (20)$$

The Γ_5-term and non-nucleonic component in the deuteron.—The unusual result of Eq. (14) is that the P-wavelike term enters only for deuteron polarizations $\lambda_d = \pm 1$. The later is the reason that momentum distribution in Eq. (18) depends explicitly on the transverse component of the relative momentum on the light front. Such behavior is impossible for nonrelativistic quantum mechanics of the deuteron since in this case the potential of the interaction is real (no inelasticities) and the solution of Lippmann-Schwinger equation for partial S and D waves satisfies the "angular condition," according to which the momentum distribution in the unpolarized deuteron depends on the magnitude of the relative momentum only. Our result does not contradict the properties of the nonrelativistic deuteron wave function since, as discussed earlier, according to Eq. (16) the P wave is purely relativistic in nature. On the other hand, in the relativistic domain the definition of the interaction potential is not straightforward to allow the use of quantum-mechanical arguments in claiming that the momentum distribution in Eq. (18) should satisfy the angular condition (i.e., depends only on the magnitude of k).

For the relativistic domain, on the light front, the analog of the Lippmann-Schwinger equation is the Weinberg type equation [25], using which for the Nucleon-Nucleon (NN) scattering amplitude, in which only nucleonic degrees are considered, in the c.o.m. of the NN system, one obtains [26]

$$T_{NN}(\alpha_i,k_{i\perp},\alpha_f,k_{f,\perp}) \equiv T_{NN}(k_{i,z},k_{i\perp},k_{f,z},k_{f,\perp}) = V(k_{i,z},k_{i\perp},k_{f,z},k_{f,\perp}) + \int V(k_{i,z},k_{i\perp},k_{m,z},k_{m,\perp})\frac{d^3k_m}{(2\pi)^3\sqrt{m^2+k_m^2}} \times \frac{T_{NN}(k_{m,z},k_{m\perp},k_{f,z},k_{f,\perp})}{4(k_m^2-k_f^2)}, \quad (21)$$

where the "i," "m," and "f" subscripts correspond to initial, intermediate, and final NN states, respectively, and momenta $k_{i,m,f}$ are defined similar to Eq. (6). The realization of the angular condition for the relativistic case will require that the light-front potential satisfy the condition

$$V(k_{i,z},k_{i\perp},k_{m,z},k_{m,\perp}) = V[k_i^2,(\vec{k}_m-\vec{k}_i)^2]. \quad (22)$$

Such a condition is obvious for the on-shell limit, since the Lorentz invariance of the T_{NN} amplitude requires

$$T_{NN}^{\text{on shell}}(k_{i,z},k_{i\perp},k_{m,z},k_{m,\perp}) = T_{NN}^{\text{on shell}}[k_i^2,(\vec{k}_m-\vec{k}_i)^2] \quad (23)$$

and the existence of the Born term in Eq. (21) indicates that the potential V satisfies the same condition in the on-shell limit.

For the off-shell potential the angular condition is not obvious. In Refs. [2,26,27] it was shown that requirements of the potential V satisfying angular condition in the on-shell limit and that it can be constructed through the series of elastic pn scatterings result in a potential which is an analytic function of angular momentum. With the assumption that the potential, analytically continued to the complex angular momentum space, does not diverge exponentially, it was shown that the V and T_{NN} functions satisfy the angular condition [Eqs. (22), (23)] in general. Using the same potential to calculate the LF deuteron wave function will result in a momentum distribution dependent only on the magnitude of the relative pn momentum. This observation requires a consideration of the pn component only in the deuteron.

Inclusion of the inelastic transitions will completely change the LF equation for the pn scattering. For example, the contribution of N^*N transition to the elastic NN scattering:

$$T_{NN}(k_{i,z},k_{i\perp},k_{f,z},k_{f,\perp}) = \int V_{NN^*}(k_{i,z},k_{i\perp},k_{m,z},k_{m,\perp}) \times \frac{d^3k_m}{(2\pi)^3\sqrt{m^2+k_m^2}} \times \frac{T_{N^*N}(k_{m,z},k_{m\perp},k_{f,z},k_{f,\perp})}{4(k_m^2-k_f^2+m_{N*}^2-m_N^2)}, \quad (24)$$

will not require the condition of Eq. (22) with the transition potential having also an imaginary component. Equation (24) cannot be described with any combination of elastic NN interaction potentials that satisfies the angular condition. The same will be true also for $\Delta\Delta \to NN$ and

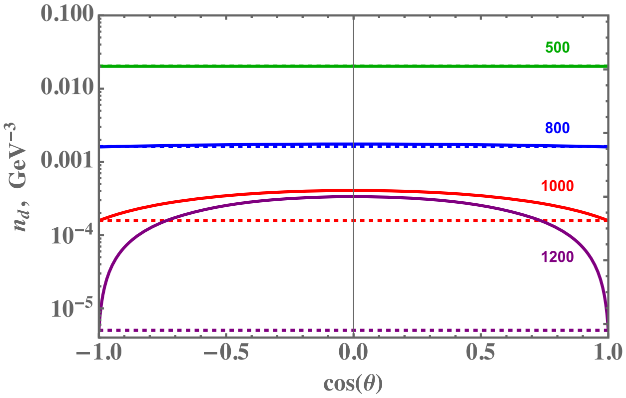

FIG. 1. LF momentum distribution of the deuteron as a function of $\cos\theta$, for different values of k. Dashed lines, deuteron with pn component only; solid lines, with P-wavelike component included.

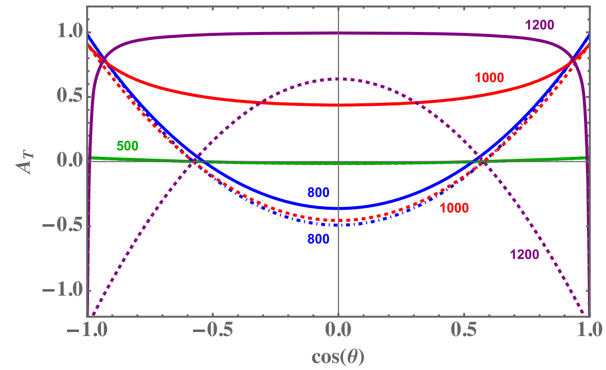

FIG. 2. Tensor asymmetry as a function of $\cos\theta$ for different k. Dashed lines, deuteron with pn component only; solid lines, with P component included.

$N_c, N_c \to NN$ transitions. This indicates that if the Γ_5 term is not zero then it should originate from a non-nucleonic component in the deuteron.

Estimate of the possible effects.—Our prediction is that the observation of anisotropic LF momentum distribution depending on the center of mass k and $k_\perp$ separately will indicate the presence of a nonnucleonic component in the deuteron. Since this effect is due to the P-wavelike structure, (originating from the Γ_5 term) which has an extra (k^2/m_N^2) factor [Eq. (15)] compared to the S and D radial waves, one expects it to become important at $k > m_N$.

To give qualitative estimates of the possible effects we evaluate the Γ_5 vertex function assuming two color-octet baryon transition to the pn system ($N_c N_c \to pn$) through one-gluon exchange, parametrizing it in the dipole form $\{A/[1+(k^2/0.71)]^2\}$. The parameter A is estimated by assuming 1% contribution to the total normalization from the P wave in Eq. (20). The latter is consistent with the experimental estimation in Ref. [28] of 0.7%. In Fig. 1 we consider the dependence of the momentum distribution of Eq. (18) as a function of $\cos\theta = [(\alpha-1)E_k/k]$ for different values of k. Notice that if the momentum distribution is generated by a pn component only, the angular condition is satisfied, and no dependence should be observed.

As the figure shows one may expect measurable angular dependence at $k \gtrsim 1$ GeV/c, which is consistent with the expectation that the inelastic transition in the deuteron corresponding to the non-nucleonic components takes place at $k \gtrsim 800$ MeV/c. Additionally, due to the fact that the P component contributes only for $\lambda_d = \pm 1$ polarizations of the deuteron [Eq. (14)] one expects an enhanced effect in the asymmetry from scattering off the tensor polarized deuteron:

$$A_T = \frac{n_d^{\lambda_d=1}(k,k_\perp) + n_d^{\lambda_d=-1}(k,k_\perp) - 2n_d^{\lambda_d=0}(k,k_\perp)}{n_d(k,k_\perp)}. \quad (25)$$

As Fig. 2 shows the presence of a non-nucleonic component will be visible already at $k \approx 800$ MeV/c, resulting in a qualitative difference in asymmetry at larger momenta as compared with the asymmetry predicted by the deuteron wave function with a pn component only.

The reason why a small, 1% effect in overall normalization gives large measurable effect in LF momentum distribution and asymmetry at $k \geq 1$ GeV/c is due to the fact that the observed "incomplete P wave" structure enters with the (p^2/m_N^2) prefactor [see Eq. (16)], which significantly amplifies the effect at very large internal momenta.

Outlook on experimental verification of the effect.—The prediction that a non-nucleonic component in the deuteron wave function may result in angular dependence of the LF-momentum distribution can be verified at c.o.m. momenta $k \gtrsim 1$ GeV/c. This seems incredibly large momenta to be measured in experiment. However, the first such measurement at high Q^2 disintegration of the deuteron has already been performed at Jefferson Lab [29] reaching $k\sim$ 1 GeV/c. It is intriguing that the results of this measurement qualitatively disagree with predictions based on conventional deuteron wave functions once $k \gtrsim 800$ MeV/c. The planned new measurements [30] will significantly improve the quality of the data, allowing possible verification of the effects discussed in this Letter. It is worth mentioning that the analysis of the experiment will require a careful account for competing nuclear effects such as final state interactions, (FSI) for which there has been significant theoretical and experimental progress during the last decade [31,32]. The advantage of high energy scattering is that the eikonal regime is established, which makes FSI to be strongly isolated in transverse kinematics and be suppressed in near collinear directions. Additionally, the comparison with the first high Q^2 experimental data [32] indicates that the accuracy of FSI calculations increases with Q^2 which will allow a meaningful analysis of new high-Q^2 data.

If the experiment will not find the angular dependence in the momentum distribution this will allow us to set a new limit on the dominance of the pn component at instantaneous high nuclear densities that correspond to $\sim$1 GeV/c internal momentum in the deuteron. If, however, the angular dependence is found, it will motivate theoretical modeling of non-nucleonic components in the deuteron, such as $\Delta\Delta$, N^*N, or hidden-color N_cN_c that can reproduce the observed result. In both cases the results of such studies will advance the understanding of the dynamics of high density nuclear matter and the relevance of the quark-hadron transition. Measuring of tensor asymmetries will significantly complement the above studies. Another venue of studies will be the extension to deep-inelastic processes on tensor polarized target aiming at the measurement of the b_1 structure function [33] as well as systematic studies of generalized parton distributions in the deuteron at high Bjorken x [34].

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[1] J. J. Aubert *et al.* (European Muon Collaboration), Phys. Lett. **123B**, 275 (1983).
[2] L. L. Frankfurt and M. I. Strikman, Phys. Rep. **160**, 235 (1988).
[3] O. Hen, G. A. Miller, E. Piasetzky, and L. B. Weinstein, Rev. Mod. Phys. **89**, 045002 (2017).
[4] J. Seely*et al.*, Phys. Rev. Lett. **103**, 202301 (2009).
[5] L. B. Weinstein, E. Piasetzky, D. W. Higinbotham, J. Gomez, O. Hen, and R. Shneor, Phys. Rev. Lett. **106**, 052301 (2011).
[6] B. Schmookler *et al.* (CLAS Collaboration), Nature (London) **566**, 354 (2019).
[7] E. Fonseca*et al.*, Astrophys. J. Lett. **915**, L12 (2021).
[8] M. C. Miller*et al.*, Astrophys. J. Lett. **918**, L28 (2021).
[9] W. Melnitchouk, M. Sargsian, and M. I. Strikman, Z. Phys. A **359**, 99 (1997).
[10] W. W. Buck and F. Gross, Phys. Rev. D **20**, 2361 (1979).
[11] R. G. Arnold, C. E. Carlson, and F. Gross, Phys. Rev. C **23**, 363 (1981).
[12] R. Dymarz and F. C. Khanna, Phys. Rev. Lett. **56**, 1448 (1986).
[13] M. Garcon and J. W. Van Orden, Adv. Nucl. Phys. **26**, 293 (2001).
[14] R. A. Gilman and F. Gross, J. Phys. G **28**, R37 (2002).
[15] R. P. Feynman, *Photon-Hadron Interactions* (CRC Press, Boca Raton, 1972).
[16] L. Frankfurt and M. Strikman, Phys. Rep. **76**, 215 (1981).
[17] G. A. Miller, Prog. Part. Nucl. Phys. **45**, 83 (2000).
[18] S. J. Brodsky, H. C. Pauli, and S. S. Pinsky, Phys. Rep. **301**, 299 (1998).
[19] T. Frederico, E. M. Henley, and G. A. Miller, Nucl. Phys. **A533**, 617 (1991).
[20] F. Vera, arXiv:2108.11502.
[21] M. M. Sargsian and F. Vera (to be published).
[22] G. P. Lepage and S. J. Brodsky, Phys. Rev. D **22**, 2157 (1980).
[23] J. Carbonell and V. A. Karmanov, Nucl. Phys. **A581**, 625 (1995).
[24] J. Carbonell, B. Desplanques, V. A. Karmanov, and J. F. Mathiot, Phys. Rep. **300**, 215 (1998).
[25] S. Weinberg, Phys. Rev. **150**, 1313 (1966).
[26] L. L. Frankfurt and M. I. Strikman, in *Modern Topics in Electron Scattering*, edited by B. Frois and I. Sick (World Scientific Publishing Company, Singapore, 1991).
[27] L. L. Frankfurt, M. I. Strikman, L. Mankiewicz, and M. Sawicki, Few Body Syst. **8**, 37 (1990).
[28] P. V. Degtyarenko, Y. V. Efremenko, V. B. Gavrilov, and G. A. Leksin, Z. Phys. A **335**, 231 (1990).
[29] C. Yero *et al.* (Hall Collaboration), Phys. Rev. Lett. **125**, 262501 (2020).
[30] W. U. Boeglin*et al.*, arXiv:1410.6770.
[31] M. M. Sargsian, Phys. Rev. C **82**, 014612 (2010).
[32] W. Boeglin and M. Sargsian, Int. J. Mod. Phys. E **24**, 1530003 (2015).
[33] W. Cosyn, Y. B. Dong, S. Kumano, and M. Sargsian, Phys. Rev. D **95**, 074036 (2017).
[34] E. R. Berger, F. Cano, M. Diehl, and B. Pire, Phys. Rev. Lett. **87**, 142302 (2001).