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Medium modifications of the bound nucleon GPDs and incoherent DVCS on nuclear targets

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ABSTRACT

We study incoherent DVCS on ^{4}He in the $^4\mathrm{He}(e, e'\gamma p)X$ reaction, which probes possible medium-modifications of the bound nucleon GPDs and elastic form factors. Assuming that the bound nucleon GPDs are modified in proportion to the corresponding bound nucleon elastic form factors, as predicted in the quark–meson coupling model, we develop an approach to calculate various incoherent nuclear DVCS observables. As an example, we compute the beam-spin DVCS asymmetry, and predict the x_B- and t-dependence of the ratio of the bound to free proton asymmetries, $A_{\mathrm{LU}}^{p^*}(\phi)/A_{\mathrm{LU}}^{p}(\phi)$. We find that the deviation of $A_{\mathrm{LU}}^{p^*}(\phi)/A_{\mathrm{LU}}^{p}(\phi)$ from unity is as much as $\sim 6\%$.

Properties of hadrons in a nuclear medium are expected to be modified compared to those in the vacuum. As indicated by measurements of unpolarized deeply inelastic scattering (DIS) of leptons off nuclear targets, unpolarized parton (quarks and gluons) distributions are appreciably modified by the nuclear medium over the entire range of values of Bjorken x_B [1–5]. Even stronger medium modifications for the polarized parton distributions have been predicted for polarized DIS off nuclear targets [6–8].

The pattern of nuclear modifications emerging from DIS off nuclear targets (and other processes, such as e.g. proton–nucleus Drell–Yan scattering) can be briefly summarized as follows. At small values of Bjorken x_B, $x_B < 0.05$, the ratio $F_2^A(x,Q^2)/[AF_2^N(x,Q^2)] < 1$, where $F_2^A(x,Q^2)$ and $F_2^N(x,Q^2)$ are the inclusive nuclear and nucleon structure functions, respectively. This suppression is called nuclear shadowing and is explained as the effect of the attenuation due to multiple coherent interactions with the target nucleons [5]. The effect of nuclear shadowing increases with the target atomic number A as $A^{1/3}$ and is as large as 30% for heavy targets. As one increases x_B, $0.05 < x_B < 0.2$, the ratio $F_2^A(x,Q^2)/[AF_2^N(x,Q^2)]$ increases above unity by a few percent. This enhancement is called antishadowing. While no widely accepted explanation of antishadowing exists, it can be dynamically generated by taking into account both the Pomeron and Reggeon exchanges in the interaction of the virtual photon with the target nucleons [9] as well as by the excess of pions in nuclei [10]. For intermediate values of x_B, $0.2 < x_B \leqslant 0.8$, the ratio $F_2^A(x,Q^2)/[AF_2^N(x,Q^2)]$ is again less than unity and this is usually what is called the EMC effect [1]. It is important to point out that, while there is no universal and generally accepted explanation of the EMC effect, it cannot be explained by traditional nuclear physics, where the nucleus consists of nucleons whose properties are not modified by the nuclear environment [2,11]. The large number of approaches and models for the explanation of the EMC effect can be grouped into two large classes [3]: the models introducing non-nucleon degrees of freedom (such as e.g. the pion cloud [12,13]) and the models assuming some kind of modifications of the nucleons themselves in the nuclear medium which mentioned earlier [14–20]. Our analysis falls into the latter category. Finally, in the large x limit ($x_B > 0.8$), the ratio $F_2^A(x,Q^2)/[AF_2^N(x,Q^2)] > 1$ as a consequence of Fermi motion and the fact $F_2^N(x,Q^2)$ vanishes in the $x_B \to 1$ limit.

It should be noted that pion excess models (some of them are mentioned above) automatically lead to the enhancement of sea quarks in nuclei, which seems to contradict the nuclear Drell–Yan data from FNAL [21]. However, a number of recent theoretical papers challenges the “naive” relation of the nuclear Drell–Yan rates to the nuclear sea quark parton distributions by discussing initial-state interactions of the quarks going into the nucleus that lower the effective momentum of the quark at the point where it annihilates. This can give very big corrections, see e.g. [22–25].

There has also been considerable interest in the possible modification of the bound nucleon elastic form factors. The polarization transfer measurement in the $^4\mathrm{He}(\vec{e}, e'\vec{p})^3\mathrm{H}$ reaction at the Hall

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A Jefferson Lab experiment [26,27] probes the possible medium modifications of the bound-nucleon form factors and can be described either by the inclusion of the modified elastic form factors as predicted by the quark–meson coupling (QMC) model [28] or by the inclusion of the strong charge-exchange final-state interaction (FSI) [29]. However, such a strong FSI may not be consistent with the induced polarization data – see Ref. [27] for details. In addition to the modifications of inclusive structure functions (parton distributions) and elastic form factors of the bound nucleon, the QMC model [14,15] predicts modifications of various hadron properties in a nuclear medium [16].

Deeply virtual Compton scattering (DVCS) interpolates between the inclusive DIS and elastic scattering reactions (see Refs. [30–33] for reviews). Therefore, it is natural to expect that generalized parton distributions (GPDs) of the bound nucleon, which are probed by various observables measured in DVCS, should also be modified in the nuclear medium. An early investigation [34,35] of such modifications in DVCS on ^{4}He assumed that in-medium nucleon GPDs are modified through the kinematic off-shell effects associated with the modification of the relation between the struck quark's transverse momentum and its virtuality.

On the experimental side, DVCS on ^{4}He in the coherent (the target nucleus remains intact) and incoherent (the target nucleus breaks up) regimes will be measured at Jefferson Lab [36]. The expected experimental accuracy will be sufficiently high to distinguish between different theoretical predictions and to extract the effects of the medium modifications of the bound nucleon GPDs. Note that the first data on coherent and incoherent DVCS on a wide range of nuclear targets was taken and analyzed by the HERMES Collaboration [37]. However, the accuracy of the data was not sufficient to extract the relatively small effects associated with medium modifications.

In this work, we compute medium modifications of the bound proton GPDs and their influence on *incoherent* DVCS on nuclear targets, $eA \to e'\gamma p X$, where A denotes any nuclear target. As a practical application, we consider the beam-spin DVCS asymmetry, A_{LU}, for a ^{4}He nucleus, since this DVCS observable will soon be measured at Jefferson Lab [36]. We find the following trend for the ratio of the bound to free proton beam-spin asymmetries: $A_{\mathrm{LU}}^{p^*}/A_{\mathrm{LU}}^{p} < 1$ for small t and x_B, and $A_{\mathrm{LU}}^{p^*}/A_{\mathrm{LU}}^{p} > 1$ as t and x_B are increased. The deviation of $A_{\mathrm{LU}}^{p^*}/A_{\mathrm{LU}}^{p}$ from unity arises mainly from the medium modification of the bound proton elastic form factor, $F_2^{p^*}(t)$.

The kinematics of DVCS on a hadronic (nuclear) target, $e(k)A(P_A) \to e(k')\gamma A'(P'_A)$, is presented in Fig. 1. The corresponding scattering amplitude reads

$$\mathcal{T}^A_{\mathrm{DVCS}} = -\bar{u}(k')\gamma_\mu u(k)\frac{1}{Q^2}H^{\mu\nu}\epsilon^*_\nu, \tag{1}$$

where the spinor $u(k)$ [$\bar{u}(k')$] corresponds to the initial [final] lepton. Q^2 is the virtuality of the exchanged photon, and ϵ^*_ν is the polarization vector of the final real photon. Note that the final nuclear state A' could be both elastic (coherent DVCS) and inelastic (incoherent DVCS).

Information on the target response is contained in the DVCS hadronic tensor, $H^{\mu\nu}$, which is defined as a matrix element of the T-product of two electromagnetic currents,

$$H^{\mu\nu} = -i\int d^4x\, e^{-iq\cdot x}\langle P'_A|T\{J^\mu(x)J^\nu(0)\}|P_A\rangle, \tag{2}$$

where q ($-q^2 = Q^2$) is the momentum of the virtual photon. To the leading twist accuracy, $H^{\mu\nu}$ of a spinless nucleus is expressed in terms of a single generalized parton distribution, H^A, convoluted with the hard scattering coefficient function $C^+(x,\xi_A)$, see e.g. Ref. [31],

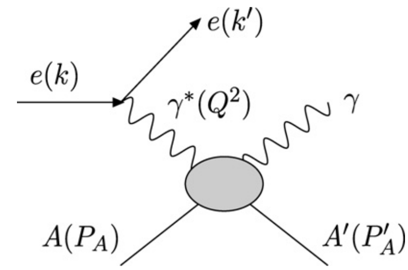

Fig. 1. Kinematics of DVCS on a nuclear target.

$$H^{\mu\nu} = -g_\perp^{\mu\nu}\int_{-1}^{1} dx\, C^+(x,\xi_A)H^A(x,\xi_A,t,Q^2)$$
$$\equiv -g_\perp^{\mu\nu}\mathcal{H}^A(\xi_A,t,Q^2), \tag{3}$$

where $g_\perp^{\mu\nu} = g^{\mu\nu} - \tilde{p}^\mu n^\nu - \tilde{p}^\nu n^\mu$ is defined by the two light-like vectors $\tilde{p} = 1/\sqrt{2}(1,0,0,1)$ and $n = 1/\sqrt{2}(1,0,0,-1)$, and $C^+(x,\xi_A) = 1/(x-\xi_A+i\epsilon) + 1/(x+\xi_A-i\epsilon)$. The function $\mathcal{H}^A$ is often called the Compton form factor (CFF). It depends on the momentum transfer, $t = (P'_A - P_A)^2$, the longitudinal momentum transfer (skewedness), $\xi_A = -(P'_A - P_A)\cdot n/(P'_A + P_A)\cdot n \approx x_A/(2-x_A)$, where $x_A = Q^2/(2P_A\cdot q)$ is the Bjorken variable, and the virtuality, Q^2.

Incoherent DVCS on a nuclear target occurs on a single nucleon, $eA \to e'\gamma NX$. Therefore, the squared amplitude of DVCS on a nuclear target, $|\mathcal{T}^A_{\mathrm{DVCS}}|^2$, can be expressed in terms of the squared amplitude of DVCS on the bound nucleons, $|\mathcal{T}^{N^*}_{\mathrm{DVCS}}|^2$. This is graphically presented in Fig. 2. Below we give the derivation and explain the notation in Fig. 2. In our work, we follow the example of the derivation of GPDs of deuterium by Cano and Pire [38]. For a similar formalism, see also [34,35].

The connection between $|\mathcal{T}^A_{\mathrm{DVCS}}|^2$ and $|\mathcal{T}^{N^*}_{\mathrm{DVCS}}|^2$ can be derived straightforwardly using the notion of the nuclear light-cone (LC) wave function. In the formalism of LC quantization, each state is characterized by its plus-momentum, $p^+ = p\cdot n = (p^0+p^3)/\sqrt{2}$, the transverse momentum, $\vec{p}_\perp$, and the helicity, λ. The nuclear state $|P_A\rangle$ is expressed in terms of the nuclear LC wave function, ϕ_A, and the product of nucleon states, $|p_i^+, \vec{p}_{\perp i}, \lambda_i\rangle$ [39],

$$|P_A^+, \vec{P}_{\perp A}\rangle$$
$$= \sum_{\lambda_i}\int\prod_{i=1}^{A}\frac{d\alpha_i}{\sqrt{\alpha_i}}\frac{d^2\vec{k}_{\perp i}}{16\pi^3}16\pi^3\delta\left(\sum_{j=1}^{A}\alpha_j - 1\right)\delta\left(\sum_{j=1}^{A}\vec{k}_{\perp j}\right)$$
$$\times\,\phi_A(\alpha_i,\vec{k}_{\perp i},\lambda_i)|\alpha_i P_A^+, \vec{k}_{\perp i} + \alpha_i\vec{P}_{\perp A}, \lambda_i\rangle, \tag{4}$$

where $\alpha_i = p_i^+/P_A^+$ is the fraction of the nucleus plus-momentum carried by nucleon i.

Substituting Eq. (4) for the initial nuclear state $|P_A\rangle$ in the nuclear hadronic tensor defined by Eq. (2), and using the assumption that the final nuclear state $|P'_A\rangle$ consists of an active nucleon N^* and $A-1$ spectators (see Fig. 2), one obtains the relation between $|\mathcal{T}^A_{\mathrm{DVCS}}|^2$ and $|\mathcal{T}^{N^*}_{\mathrm{DVCS}}|^2$,

$$\left|\mathcal{T}^A_{\mathrm{DVCS}}(\xi_A,t)\right|^2 = \int_{\alpha_{\min}}^{1}\frac{d\alpha}{\alpha}\rho_A^N(\alpha,\lambda)\sum_\lambda\left|\mathcal{T}^{N^*}_{\mathrm{DVCS}}(\xi_N,t)\right|^2, \tag{5}$$

where the nucleon light-cone distribution $\rho_A^N(\alpha,\lambda)$ is defined in terms of the nuclear LC wave function [40]:

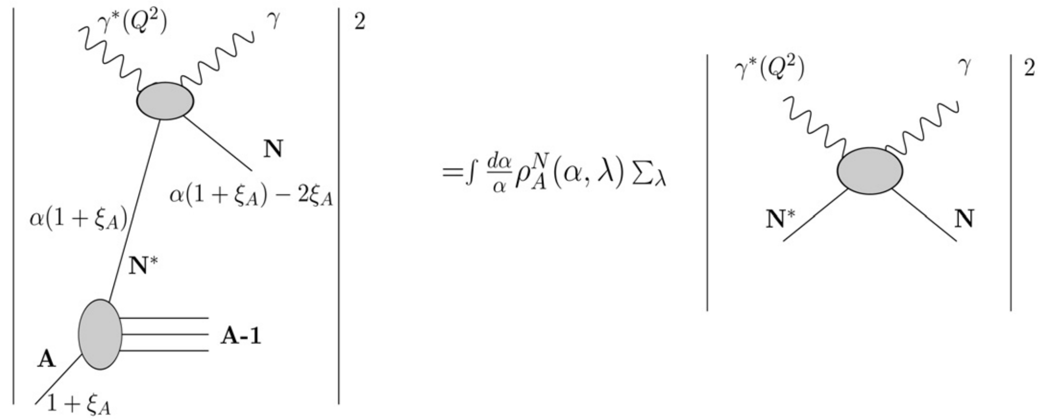

Fig. 2. Illustration of the connection between incoherent DVCS on a nuclear target and DVCS on a bound nucleon (see Eq. (5)). The relevant light-cone fractions are also indicated.

$$\rho_A^N(\alpha,\lambda)=\int \frac{d^2\vec{k}_\perp}{16\pi^3}\sum_{\lambda_i}\prod_{i=2}^{A}\frac{d\alpha_i\, d^2\vec{k}_{\perp i}}{16\pi^3}\delta\left(\alpha+\sum_{j=2}^{A}\alpha_j-1\right)$$
$$\times 16\pi^3\delta\left(\vec{k}_\perp+\sum_{j=2}^{A}\vec{k}_{\perp j}\right)\left|\phi_A(\alpha,\vec{k}_\perp,\lambda,\alpha_i,\vec{k}_{\perp i},\lambda_i)\right|^2. \tag{6}$$

Here the light-cone fraction α and the transverse momentum $\vec{k}_\perp$ refer to the interacting nucleon, while α_i and $\vec{k}_{\perp i}$ refer to the spectator nucleons. The distribution $\rho_A^N(\alpha,\lambda)$ introduced above is normalized to unity,

$$\sum_\lambda \int_0^1 d\alpha\, \rho_A^N(\alpha,\lambda)=1. \tag{7}$$

In Fig. 2, we also show the relevant light-cone momentum fractions (we use the standard symmetric frame [30]): the nucleus carries the plus-momentum $P_A^+=(1+\xi_A)\bar{P}_A^+$, where $\bar{P}_A^+=(P_A^+ + P_A^{\prime +})/2$; the active nucleon has $p^+=\alpha(1+\xi_A)\bar{P}_A^+$ in the initial state and $p^{\prime +}=(\alpha(1+\xi_A)-2\xi_A)\bar{P}_A^+$ in the final state. The requirement $p^{\prime +}\geqslant 0$ determines the minimal value of α, $\alpha_{\min}=2\xi_A/(1+\xi_A)$.

In Eq. (5), the skewedness, ξ_N, is defined with respect to the active nucleon in the symmetric frame [38]:

$$\xi_N=\frac{\xi_A}{\alpha(1+\xi_A)-\xi_A}. \tag{8}$$

The light-cone distribution $\rho_A^N(\alpha)$ is peaked around $\alpha\approx 1/A$. From our experience with the EMC effect [1], it is well known that, except for large x_B, the effect of Fermi motion is small [2–4]. Therefore, we neglect Fermi motion of the bound nucleon and evaluate $|\mathcal{T}_{\mathrm{DVCS}}^{N^*}(\xi_N,t)|^2$ at $\alpha=1/A$. Using the normalization condition of Eq. (7), we obtain an approximate expression for $|\mathcal{T}_{\mathrm{DVCS}}^{A}|^2$,

$$\left|\mathcal{T}_{\mathrm{DVCS}}^{A}(\xi_A,t)\right|^2\simeq\sum_N\frac{1}{2}\sum_\lambda\left|\mathcal{T}_{\mathrm{DVCS}}^{N^*}\big(\langle\xi_N\rangle,t\big)\right|^2, \tag{9}$$

where the factor 1/2 comes from the normalization condition of Eq. (7), and the average nucleon skewedness, $\langle\xi_N\rangle$, is defined as

$$\langle\xi_N\rangle\equiv\frac{\xi_A}{\frac{1}{A}(1+\xi_A)-\xi_A}. \tag{10}$$

To compare with the experiment, it is convenient to rescale x_A and to define the Bjorken variable, x_B, with respect to the nucleon:

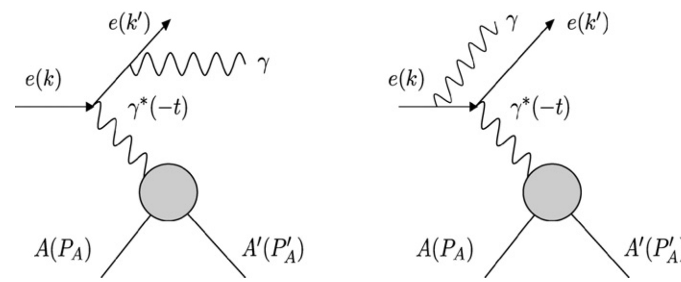

Fig. 3. The Bethe–Heitler process.

$$x_B\equiv A\frac{Q^2}{2P_A\cdot q}\equiv A x_A. \tag{11}$$

The corresponding skewedness ξ, $\xi=x_B/(2-x_B)$, coincides with that given by Eq. (10). Using the Bjorken x_B of Eq. (11), and the fact that both sides of Eq. (9) depend on the same skewedness ξ, we obtain:

$$\left|\mathcal{T}_{\mathrm{DVCS}}^{A}(\xi,t)\right|^2=\sum_N\frac{1}{2}\sum_\lambda\left|\mathcal{T}_{\mathrm{DVCS}}^{N^*}(\xi,t)\right|^2. \tag{12}$$

The interpretation of Eq. (12) is intuitive: the probability of incoherent DVCS on a nuclear target is a sum of the probabilities of DVCS on individual nucleons.

Since Eq. (12) is based on the decomposition of Eq. (4) and does not depend on the type of the elementary interaction with the active nucleon, similar relations also hold for the Bethe–Heitler (BH) amplitude (see Fig. 3) and for the interference between the DVCS and BH amplitudes (see Ref. [41] for details of the definitions of the BH and interference amplitudes):

$$\left|\mathcal{T}_{\mathrm{BH}}^{A}(\xi,t)\right|^2=\sum_N\frac{1}{2}\sum_\lambda\left|\mathcal{T}_{\mathrm{BH}}^{N^*}(\xi,t)\right|^2,$$
$$\left|\mathcal{I}^{A}(\xi,t)\right|^2=\sum_N\frac{1}{2}\sum_\lambda\left|\mathcal{I}^{N^*}(\xi,t)\right|^2. \tag{13}$$

The practical corollary of Eqs. (12) and (13) is the following: expressions for DVCS observables (cross section asymmetries) in incoherent nuclear DVCS on a spinless nuclear target are exactly the same as those for the sum of individual bound nucleons.

In this work, we apply Eqs. (12) and (13) to incoherent DVCS on ^{4}He in the situation when a proton in the final state is detected, $e\,^4\mathrm{He}\to e'\gamma p X$. In this case, the neutrons do not contribute and Eqs. (12) and (13) become

$$\begin{aligned}
\left|\mathcal{T}^{A}_{\rm DVCS}(\xi,t)\right|^2 &= \sum_{\lambda}\left|\mathcal{T}^{p^*}_{\rm DVCS}(\xi,t)\right|^2,\\
\left|\mathcal{T}^{A}_{\rm BH}(\xi,t)\right|^2 &= \sum_{\lambda}\left|\mathcal{T}^{p^*}_{\rm BH}(\xi,t)\right|^2,\\
\left|\mathcal{I}^{A}(\xi,t)\right|^2 &= \sum_{\lambda}\left|\mathcal{I}^{p^*}(\xi,t)\right|^2.
\end{aligned} \tag{14}$$

Note that although Eq. (14) does not contain an explicit reference to the Fermi motion of the bound nucleon, it does implicitly contain some effects of the Fermi motion through the self-consistent change of the internal structure of the bound nucleon via the medium-modified proton elastic form factors (see below).

As we mentioned in the Introduction, the GPDs of the bound nucleon may generally differ from the GPDs of the free nucleon. In this work, we assume that the GPDs of the bound proton are modified in proportion to the corresponding bound proton elastic form factors,

$$\begin{aligned}
H^{q/p^*}(x,\xi,t,Q^2) &= \frac{F_1^{p^*}(t)}{F_1^p(t)} H^q(x,\xi,t,Q^2),\\
E^{q/p^*}(x,\xi,t,Q^2) &= \frac{F_2^{p^*}(t)}{F_2^p(t)} E^q(x,\xi,t,Q^2),\\
\tilde{H}^{q/p^*}(x,\xi,t,Q^2) &= \frac{G_1^*(t)}{G_1(t)} \tilde{H}^q(x,\xi,t,Q^2),
\end{aligned} \tag{15}$$

where the GPDs H^{q/p^*}, E^{q/p^*} and $\tilde{H}^{q/p^*}$ and the elastic form factors $F_1^{p^*}$ (Dirac form factor), $F_2^{p^*}$ (Pauli form factor) and G_1^* (axial form factor) refer to the bound proton, while H^q, E^q, and $\tilde{H}^q$ and F_1^p, F_2^p and G_1 refer to those of the free proton. The assumption of Eq. (15) is rather simple, since the medium-modifications result only in the t-dependent renormalization and do not change the shape of the in-medium GPDs. The GPDs $H^{q/p^*}(x,\xi,t,Q^2)$ and $E^{q/p^*}(x,\xi,t,Q^2)$ in a ^{4}He nucleus are constrained to reproduce the extracted bound proton elastic electromagnetic form factors after integration over x, as the QMC model predicted [26] (see below). Note also that we have ignored the insignificant kinematically-suppressed contribution of the GPD $\tilde{E}$ to the DVCS beam-spin asymmetry [41].

The bound proton form factors have been calculated in the QMC model [28,42,43]. Since these form factors depend on the nuclear density, the in-medium form factors in Eq. (15) must be averaged over the nuclear density distribution in ^{4}He ($A = {}^4$He below),

$$\begin{aligned}
F_1^{p^*}(t) &= \int d^3\vec{r}\,\rho_A(r) F_1^{p^*}(t,\rho_A(r)),\\
F_2^{p^*}(t) &= \int d^3\vec{r}\,\rho_A(r) F_2^{p^*}(t,\rho_A(r)),\\
G_1^*(t) &= \int d^3\vec{r}\,\rho_A(r) G_1^*(t,\rho_A(r)),
\end{aligned} \tag{16}$$

where $F_1^{p^*}(t,\rho_A(r))$, $F_2^{p^*}(t,\rho_A(r))$ and $G_1^*(t,\rho_A(r))$ are the nuclear density-dependent bound proton form factors, and $\rho_A(r)$ ($\equiv \rho_{^4{\rm He}}(r)$) is the nuclear density distribution in ^{4}He calculated in Ref. [44]. In Fig. 4, we show the resulting ratios $F_1^{p^*}(t)/F_1^p(t)$, $F_2^{p^*}(t)/F_2^p(t)$ and $G_1^*(t)/G_1(t)$ as functions of $-t$ [28,42,43].

For the free proton GPDs, we use the double distribution model [45] based on valence quark PDFs. In particular, we use

$$\begin{aligned}
&H^q(x,\xi,t,Q^2)\\
&\quad = \int_0^1 d\beta \int_{-1+|\beta|}^{1-|\beta|} d\alpha\, \delta(\beta+\alpha\xi-x)\pi(\beta,\alpha)\beta^{-\alpha'(1-\beta)t} q_v(\beta,Q^2),
\end{aligned}$$

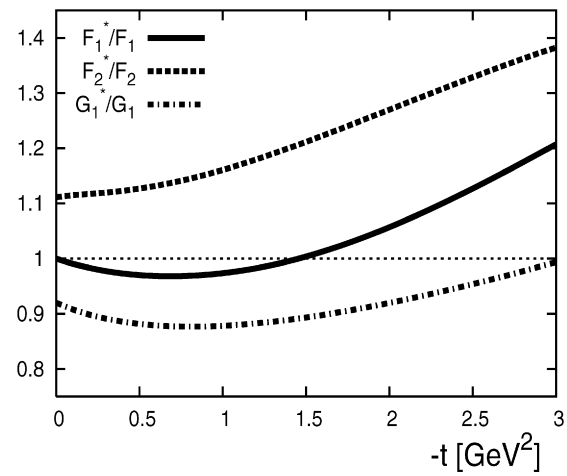

Fig. 4. The bound (^{4}He) to free proton ratios of elastic form factors $F_1^{p^*}(t)/F_1^p(t)$, $F_2^{p^*}(t)/F_2^p(t)$ and $G_1^*(t)/G_1(t)$ as functions of the momentum transfer t, see Eq. (16).

$$\begin{aligned}
&E^q(x,\xi,t,Q^2)\\
&\quad = \int_0^1 d\beta \int_{-1+|\beta|}^{1-|\beta|} d\alpha\, \delta(\beta+\alpha\xi-x)\pi(\beta,\alpha)\beta^{-\alpha'(1-\beta)t} e_v^q(\beta,Q^2),\\
&\tilde{H}^q(x,\xi,t,Q^2)\\
&\quad = \int_0^1 d\beta \int_{-1+|\beta|}^{1-|\beta|} d\alpha\, \delta(\beta+\alpha\xi-x)\pi(\beta,\alpha)\beta^{-\alpha'(1-\beta)t} \Delta q_v(\beta,Q^2),
\end{aligned} \tag{17}$$

where q_v and Δq_v are the valence unpolarized and polarized quark PDFs, respectively, while $e_v^q(\beta)$ is the valence part of the forward limit of the GPD E^q. The profile function $\pi(\beta,\alpha)$ is taken in a standard form [31]:

$$\pi(\beta,\alpha) = \frac{3}{4}\frac{(1-\beta)^2-\alpha^2}{(1-\beta)^3}. \tag{18}$$

The t-dependence of GPDs is introduced through the Regge theory-motivated factor with the slope parameter $\alpha' = 1.105$ GeV^{-2}, which leads to a good description of the proton and neutron elastic form factors [46].

For the unpolarized quark PDFs, we use the leading-order (LO) CTEQ5L parameterization [47], while for the polarized quark PDFs, we use the LO GRSV 2000 parameterization [48]. The model for the forward limit of the GPD E^q is taken from Ref. [46]. Explicitly, it is given by

$$\begin{aligned}
e_v^u(x,Q^2) &= \frac{k_u}{N_u}(1-x)^{\eta_u} u_v(x,Q^2),\\
e_v^d(x,Q^2) &= \frac{k_d}{N_d}(1-x)^{\eta_d} d_v(x,Q^2),
\end{aligned} \tag{19}$$

where $k_u = 1.673$ and $k_d = -2.033$ are the anomalous magnetic moments; $\eta_u = 1.713$ and $\eta_d = 0.566$ are determined from fits to the nucleon elastic form factors; N_u and N_d are the normalization factors,

$$\begin{aligned}
N_u &= \int_0^1 dx\,(1-x)^{\eta_u} u_v(x,Q^2),\\
N_d &= \int_0^1 dx\,(1-x)^{\eta_d} d_v(x,Q^2).
\end{aligned} \tag{20}$$

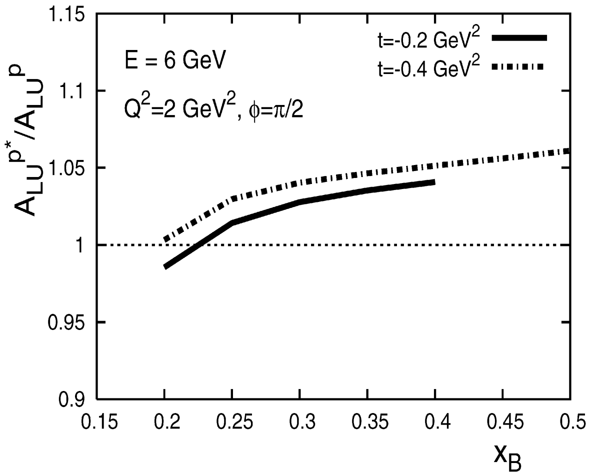

Fig. 5. The ratio of the bound (incoherent ^{4}He) to free proton beam-spin DVCS asymmetries, $A_{\mathrm{LU}}^{p^*}(\phi)/A_{\mathrm{LU}}^{p}(\phi)$, as a function of Bjorken x_B at $E = 6$ GeV, $Q^2 = 2$ GeV^2, $\phi = \pi/2$ and two values of t.

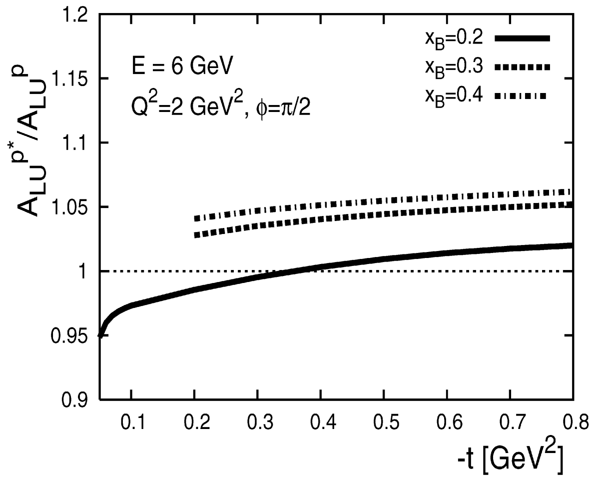

Fig. 6. The ratio of the bound (incoherent ^{4}He) to free proton beam-spin DVCS asymmetries, $A_{\mathrm{LU}}^{p^*}(\phi)/A_{\mathrm{LU}}^{p}(\phi)$, as a function of the momentum transfer t at $E = 6$ GeV, $Q^2 = 2$ GeV^2, $\phi = \pi/2$ and three values of x_B.

In summary, the bound proton GPDs are given by Eqs. (17)–(20). Since for the case of incoherent DVCS on ^{4}He, $e\,^4\mathrm{He} \to e'\gamma p X$, the scattering amplitudes squared are the same as those for the bound proton (see Eq. (14)), we may use the standard formalism developed for the free nucleon [41] to calculate various DVCS observables (cross section asymmetries). Our results are presented in Figs. 5 and 6.

In Fig. 5 we present the ratio of the bound (incoherent ^{4}He) to free proton beam-spin DVCS asymmetries, $A_{\mathrm{LU}}^{p^*}(\phi)/A_{\mathrm{LU}}^{p}(\phi)$, as a function of Bjorken x_B at the fixed energy of the lepton beam, $E = 6$ GeV, and virtuality $Q^2 = 2$ GeV^2. This asymmetry is measured with a linearly polarized lepton beam and an unpolarized target. The $A_{\mathrm{LU}}(\phi)$ asymmetry is mostly sensitive to the imaginary part of the Compton form factor, $\mathrm{Im}\,\mathcal{H}^A$ (see Eq. (3)), and behaves as $A_{\mathrm{LU}} \propto \mathrm{Im}\,\mathcal{H}^A \sin\phi$, where ϕ is the angle between the leptonic and hadronic (production) planes. (See Ref. [41] for the details.) Note that in Fig. 5, $A_{\mathrm{LU}}(\phi)$ is evaluated at $\phi = \pi/2$.

As seen from Fig. 5, effects of the medium-modifications in the kinematic region under study do not exceed $\sim 6\%$. Their trend can be understood by analyzing the approximate expression for $A_{\mathrm{LU}}(\phi)$ [41],

$$A_{\mathrm{LU}}(\phi) \propto \mathrm{Im}\left(F_1^{p^*}\mathcal{H}^{p^*} + \frac{x_B}{2-x_B}\left(F_1^{p^*} + F_2^{p^*}\right)\tilde{\mathcal{H}}^{p^*} - \frac{t}{4m_N^2}F_2^{p^*}\mathcal{E}^{p^*}\right)\Big/ f\left(F_1^{p^*}, F_2^{p^*}\right)\sin\phi, \tag{21}$$

where $\mathcal{H}^{p^*}$, $\mathcal{E}^{p^*}$ and $\tilde{\mathcal{H}}^{p^*}$ are the Compton form factors of the respective bound nucleon GPDs; $f(F_1^{p^*}, F_2^{p^*})$ is a certain function (dominated by the Bethe–Heitler amplitude squared) of the elastic form factors. Note that the argument of the elastic form factors is the invariant momentum transfer t (see Fig. 3).

At small x_B and t, the contributions of $\tilde{\mathcal{H}}^{p^*}$ and $\mathcal{E}^{p^*}$ in Eq. (21) are unimportant and $A_{\mathrm{LU}}^{p^*}(\phi)/A_{\mathrm{LU}}^{p}(\phi) < 1$ because of the increase of $f(F_1^{p^*}, F_2^{p^*})$ for the bound nucleon. This comes mainly from the enhancement, $F_2^{p^*} > F_2^{p}$, in ^{4}He. (See Fig. 4.)

As x_B and t are increased, $\tilde{\mathcal{H}}^{p^*}$ and $\mathcal{E}^{p^*}$ in Eq. (21) start to play a progressively more important role (the contribution of $\tilde{\mathcal{H}}^{p^*}$ is more important). Thus, the medium-enhancement of the term proportional to $(F_1^{p^*} + F_2^{p^*})\tilde{\mathcal{H}}^{p^*}$ wins over the enhancement of the denominator in Eq. (21), and makes $A_{\mathrm{LU}}^{p^*}(\phi)/A_{\mathrm{LU}}^{p}(\phi) > 1$.

In Fig. 6 we present the ratio $A_{\mathrm{LU}}^{p^*}(\phi)/A_{\mathrm{LU}}^{p}(\phi)$ as a function of the invariant momentum transfer t, in the same kinematics as in Fig. 5. The size of the medium-modification is similar to that shown in Fig. 5, and the trend of the medium modifications of the ratio $A_{\mathrm{LU}}^{p^*}(\phi)/A_{\mathrm{LU}}^{p}(\phi)$ has a similar interpretation.

Our numerical predictions are based on the particular model of the nucleon GPDs summarized by Eqs. (17)–(20). We have explicitly checked that taking a different profile function $\pi(\beta,\alpha)$ in Eq. (18), e.g. $\pi(\beta,\alpha) = 15/16[(1-\beta)^2 - \alpha^2]^2/(1-\beta)^5$, does not change our numerical prediction. Since we present our results in the form of the ratio of the nuclear to nucleon DVCS asymmetries, we expect that details and subtleties of the nucleon GPDs should mostly cancel in the ratio and, thus, our predictions summarized in Figs. 5 and 6 should be stable against variations of the parameterization of the nucleon GPDs.

In our analysis, we did not address the issue of possible final state interactions (FSI) between the produced proton (nucleon) and the remaining $A = 3$ system. In principle, this is a separate, rather involved analysis. However, based on the observation that the non-charge-exchange FSI for the $^4\mathrm{He}(\vec{e}, e'\vec{p})^3\mathrm{H}$ reaction are rather small [29] and on the observation that the large charge-exchange final-state interaction (FSI) for the same reaction are inconsistent with the polarization transfer data [27], one should not expect FSI for our case of incoherent DVCS, $^4\mathrm{He}(e, e'\gamma p)X$, that are larger than a few percent. Therefore, the theoretical uncertainty associated with the FSI is not large and should not affect our conclusions. One should emphasize that the medium modifications of the bound nucleon GPDs and FSI are two separate effects. Once the effect of FSI for incoherent DVCS on ^{4}He is estimated, it should be added on the top of the medium modification effects discussed in the present Letter.

Finally, we would like to compare our results in Figs. 5 and 6 with the predictions of Liuti and Taneja [34]. While in our model of the bound proton GPDs in ^{4}He, the effects of Fermi motion, off-shellness, and the internal structure change of the bound nucleon are encoded in the medium-modified proton elastic form factors, the approach of Ref. [34] explicitly takes into account such effects in the bound nucleon GPD. Furthermore, the bound nucleon GPDs in the approach of Ref. [34] are modified through the kinematic off-shell effects associated with the modification of the relation between the struck quark's transverse momentum and its virtuality.

First we discuss the t-dependence. While our prediction for the t-dependence of $A_{\mathrm{LU}}^{p^*}(\phi)/A_{\mathrm{LU}}^{p}(\phi)$ is similar to that of Ref. [34], the size of the nuclear modifications is significantly smaller in our case. Although the x_B-dependence of incoherent DVCS was not presented in Ref. [34], the x_B-dependence of $A_{\mathrm{LU}}^{p^*}(\phi)/A_{\mathrm{LU}}^{p}(\phi)$, which is based on the same model as presented in Ref. [34], was given in the proposal of the future Jefferson Lab experiment [36]. Our predictions for the x_B-dependence of $A_{\mathrm{LU}}^{p^*}(\phi)/A_{\mathrm{LU}}^{p}(\phi)$ are very different both in shape and in size from those presented in Ref. [36]. In particular, our prediction for the in-medium modification is much smaller in magnitude. The future Jefferson Lab experiment on DVCS on ^{4}He will be able to distinguish between our predictions and those of Ref. [34].

In conclusion, we have studied incoherent DVCS on ^{4}He in the $^4\mathrm{He}(e, e'\gamma p)X$ reaction, which probes medium-modifications of the bound proton GPDs and elastic form factors. Assuming that the proton GPDs are modified in proportion to the corresponding bound proton elastic form factors, as predicted in the quark–meson coupling model, we have developed an approach to calculate various incoherent nuclear DVCS observables. As an example, we have computed the beam-spin DVCS asymmetry and made predictions for the x_B- and t-dependence of the ratio of the bound to free proton asymmetries, $A_{\mathrm{LU}}^{p^*}(\phi)/A_{\mathrm{LU}}^{p}(\phi)$. We have found that the deviation of $A_{\mathrm{LU}}^{p^*}(\phi)/A_{\mathrm{LU}}^{p}(\phi)$ from unity is as much as $\sim 6\%$. We checked that our predictions are stable against the variation of the model of the nucleon GPDs. Also, based on the studies of final state interactions in $^4\mathrm{He}(\vec{e}, e'\vec{p})^3\mathrm{H}$ quasi-elastic scattering, we argue that the effect of the FSI should not exceed a few percent in our case of incoherent DVCS on ^{4}He.

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