

Coherent deeply virtual Compton scattering on ^4He beyond leading power

V. Martínez-Fernández^{1,2}, B. Pire³, P. Sznajder⁴, and J. Wagner⁴

¹*Université Paris-Saclay - CEA - IRFU, 91191 Gif-sur-Yvette, France*

²*Center for Frontiers in Nuclear Science, Stony Brook University, Stony Brook, New York 11794, USA*

³*Centre de Physique Théorique, CNRS, École polytechnique, I.P. Paris, 91128 Palaiseau, France*

⁴*National Centre for Nuclear Research (NCBJ), Pasteura 7, 02-093 Warsaw, Poland*



(Received 19 May 2026; accepted 28 July 2026; published 9 September 2026)

Coherent hard exclusive reactions on light nuclei provide access to their quark and gluon structure and enable three-dimensional tomography of these complex systems. We study deeply virtual Compton scattering on a helium-4 target, including both kinematic twist-3 and twist-4 corrections, as well as next-to-leading-order corrections to the twist-2 amplitude in the strong coupling α_s . We show that these contributions are crucial for achieving a precise description of the data and, as a result, obtain the first tomographic image of the helium-4 nucleus at the quark-gluon level.

DOI: [10.1103/h7kn-ql7j](https://doi.org/10.1103/h7kn-ql7j)

I. INTRODUCTION

Studying coherent exclusive scattering has long been recognized [1–5] as the best way to progress in our understanding of the quark and gluon structure of light nuclei, and therefore as a key step toward the understanding of the confinement mechanism beyond the nucleon or meson case. The golden channel, as in the nucleon case, is the *electroproduction of a real photon* [6]

$$e^-(k) + A(p) \rightarrow e^-(k') + A(p') + \gamma(q'), \quad q^2 = 0, \quad (1)$$

where e^- stands for the electron, A represents the nucleus, and γ is the produced real photon, while the momenta of the corresponding particles are given in parentheses. This process can be understood as the superposition of two types of scatterings depicted in Fig. 1: a purely electromagnetic Bethe-Heitler (BH) process where the final-state photon is the result of leptonic bremsstrahlung and it is characterized by the electromagnetic form factors (EFFs) of the target; and the deeply virtual Compton scattering (DVCS) consisting of an absorption of a virtual photon and subsequent emission of a real one.

The DVCS amplitude is conveniently described in terms of Compton form factors (CFFs), which are complex-valued convolutions of perturbatively calculable coefficient functions, describing the hard parton-photon interaction, with generalized parton distributions (GPDs) [7–9]. This factorized form of the scattering amplitude applies in the generalized Bjorken regime, where the momentum transfer to the target

remains small in comparison with the photon virtuality. Denoting the latter by $-q^2 = -(k - k')^2 = Q^2 > 0$ and the former by the Mandelstam variable $t = (p' - p)^2 < 0$, one finds that the leading contribution, referred to as the leading-twist (LT) approximation, is obtained in the limit $|t|/Q^2 \rightarrow 0$.

The real experiments obviously do not realize the leading-twist (LT) limit. Consequently, deviations from this limit manifest as corrections that enter the coefficient functions of DVCS in powers of the form $(|t|/Q^2)^{(\tau-2)/2}$, commonly referred to as kinematic-twist (τ) corrections [10,11].¹ The importance of systematically accounting for these corrections is multifaceted. Including them improves theoretical accuracy and enables the use of a larger portion of experimental data for the extraction of CFFs, thereby relaxing kinematic cuts that remove points affected by higher-twist contributions. Besides these kinematic corrections, there are also genuine higher-twist contributions associated with matrix elements of higher-twist operators. These are beyond the scope of the present study, in which we focus exclusively on the kinematic-twist sector.

Furthermore, hadron and nuclear tomography [12–14], which provides transverse imaging of their internal structure, requires integration over all possible values of t . The inclusion of kinematic t -dependent corrections is therefore essential for a consistent and complete tomographic analysis. Similarly, since GPDs depend explicitly on the momentum transfer t , incorporating recoil effects in data analyses is crucial for achieving a more reliable modeling of these distributions, done thus far under a leading twist formalism.

In addition, moments of GPDs are related to the form factors parametrizing the quantum chromodynamics (QCD)

Published by the American Physical Society under the terms of the [Creative Commons Attribution 4.0 International](https://creativecommons.org/licenses/by/4.0/) license. Further distribution of this work must maintain attribution to the author(s) and the published article's title, journal citation, and DOI. Funded by SCOAP³.

¹Although they are sometimes called target mass corrections, we explicitly show below that their effect is not proportional to the nucleus mass, but rather to the momentum transfer.

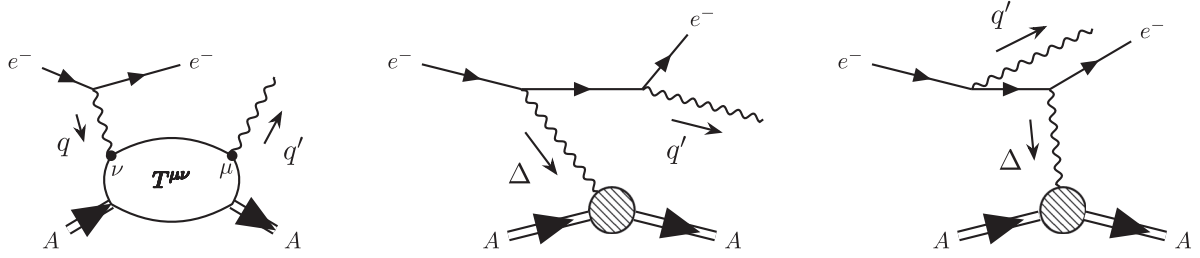


FIG. 1. Deeply virtual Compton scattering (DVCS, on the left) and the two cases of Bethe-Heitler background, BH (center) and BHx (right). The tensor $T^{\mu\nu}$ is the so-called Compton tensor, see main text.

energy-momentum tensor, known as gravitational form factors, which provide access to the mechanical properties of hadrons and nuclei [15]. Recent works [16,17] extended the relation between Compton form factors and gravitational form factors beyond leading twist by incorporating kinematic-twist corrections, thus allowing Compton scattering observables to probe a broader class of gravitational form factors.

In summary, twist effects are not only relevant for precision analyses, but they also provide unique insights into the mechanical structure of hadronic matter. As a result, recovering these recoil and mass effects must become an integral part of the standard for theoretical calculations, alongside with next-to-leading order (NLO) corrections in the strong coupling constant α_s whose relevance have already been highlighted [18]. Such corrections also include contributions of the gluon-transversity GPD responsible for a two-unit photon-helicity flip at LT [19], making it a competing effect to the LO twist-4 contribution by quarks to said helicity flip.

This motivates revisiting DVCS with modern theoretical tools, combining higher-twist and NLO accuracy with phenomenological studies. As a first step, in this manuscript and in the companion letter [20] we focus on DVCS off a spinless target, and in particular on helium-4. Compared to the nucleon, spinless targets are described by a reduced set of GPDs, which makes them a natural starting point for such an analysis. Among them, helium-4 is especially interesting, as it provides a clean setting to explore nuclear effects in hadron structure. Moreover, existing JLab measurements [21,22] and ongoing experimental activity make this study both timely and phenomenologically relevant.

Our line of reasoning for building a model of helium-4 GPDs is very different from the ones used in [23] where the nucleus wave function is convoluted with nucleon GPDs. In contrast to this approach, ours respects Lorentz invariance and treats the nuclear state as a QCD object without any prejudice on its nucleon decomposition. A comparison of the tomographic picture of the nucleus based on the same experimental data but with these different approaches would be very welcome, but is outside the scope of our study.

This paper is organized as follows. Section II presents the calculation of the cross section and the corresponding amplitudes at twist-4 and NLO accuracy. Section III is devoted to the modeling of helium-4 GPDs. The results of the fit to DVCS data are presented in Sec. IV, including the first tomographic image of the helium-4 nucleus in terms of

quarks and gluons, as well as numerical estimates of cross sections and amplitudes for JLab experiments. Section V concludes with a summary of our main results and final remarks. For completeness, Appendix A collects the relevant kinematics, Appendix B provides details of the amplitude calculations based on the Kleiss-Stirling spinor techniques [24,25] that we use in Sec. II, and we collect in Appendix C the expressions for helicity amplitudes.

II. AMPLITUDES AND CROSS SECTION

The differential cross section for a beam of charge sign χ and helicity s , producing a real photon with polarization λ , can be written as

$$\frac{d^4\sigma_{\chi}^{s\lambda}}{dx_A dQ^2 dt d\phi} = \frac{\alpha_{\text{EM}}^3}{8\pi} \frac{x_A y^2}{Q^4 \sqrt{1+\omega^2}} \left| \frac{\mathcal{M}_{\chi}^{s\lambda}}{e^3} \right|^2, \quad (2)$$

where the inelasticity variable is given by $y = pq/(pk)$, $\omega = 2x_A M/Q$, M is the nucleus mass, the nuclear Bjorken variable x_A is defined as

$$x_A = \frac{Q^2}{2pq} \approx \frac{x_B}{A}, \quad (3)$$

and the approximate relation to the proton Bjorken variable x_B holds for a nucleus with mass number A (up to binding energy corrections). The angle ϕ is defined according to the Trento convention [26] as presented in Fig. 2. The detailed parametrization of all four-momenta in the target rest frame (TRF) of Ref. [6], which is related to the standard Trento frame by a 180° rotation about the y axis, is briefly recalled in Appendix A.

The total amplitude is given by the sum of the three contributions corresponding to the diagrams shown in Fig. 1:

$$i\mathcal{M}_{\chi}^{s\lambda} = i\mathcal{M}_{\chi, \text{DVCS}}^{s\lambda} + i\mathcal{M}_{\text{BH}}^{s\lambda} + i\mathcal{M}_{\text{BHx}}^{s\lambda}. \quad (4)$$

In the following, we describe the calculation of each of these contributions.

A. DVCS amplitude

The amplitude for the DVCS process, shown in the left panel of Fig. 1, is given by

$$i\mathcal{M}_{\chi, \text{DVCS}}^{s\lambda} = \frac{ie^3\chi}{q^2} \bar{u}(k', s) \gamma_\nu u(k, s) T^{\mu\nu} \varepsilon'_\mu(-\lambda), \quad (5)$$

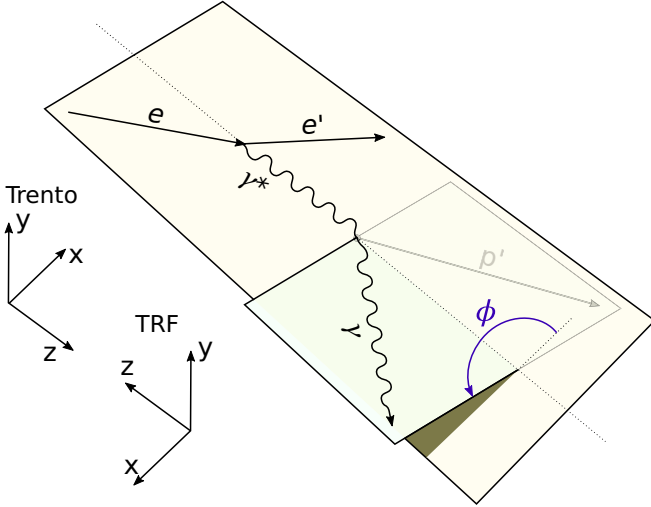


FIG. 2. Target rest frame (TRF) and Trento frame.

where the Compton tensor $T^{\mu\nu}$ ($\bar{q} = (q + q')/2$ and $\mathcal{T}$ represents the time ordering),

$$\begin{aligned} T^{\mu\nu} &= i \int d^4z e^{i\bar{q}z} \langle p' | \mathcal{T} \{ j^\mu(z/2) j^\nu(-z/2) \} | p \rangle \\ &= i \int d^4z e^{iq'z} \langle p' | \mathcal{T} \{ j^\mu(z) j^\nu(0) \} | p \rangle \\ &= \sum_{\Lambda, \Lambda'} T_{(\Lambda\Lambda')}^{\mu\nu} \mathcal{A}^{\Lambda\Lambda'}, \end{aligned} \quad (6)$$

can be parametrized by means of photon-helicity-dependent amplitudes $\mathcal{A}^{\Lambda\Lambda'}$ [11,27]. For an incoming photon with polarization Λ and an outgoing one with polarization Λ' , these take

the form

$$\mathcal{A}^{\Lambda\Lambda'} = (-1)^{\Lambda-1} (\epsilon'_\mu(\Lambda'))^* T^{\mu\nu} \epsilon_\nu(\Lambda), \quad (7)$$

where the polarization vectors are defined in Eqs. (A25) and (A26). After accounting for parity symmetry, the sum over photon polarizations is limited to $\Lambda\Lambda' \in \{++, +-, 0+\}$, where $\pm$ stand for transverse polarizations while 0 for the longitudinal one. The corresponding Lorentz structures, expressed by means of the polarization vectors

$$T_{(++)}^{\mu\nu} = \epsilon^\mu(+) \epsilon^\nu(-) + \epsilon^\mu(-) \epsilon^\nu(+), \quad (8)$$

$$T_{(+-)}^{\mu\nu} = \epsilon^\mu(+) \epsilon^\nu(+) + \epsilon^\mu(-) \epsilon^\nu(-), \quad (9)$$

$$T_{(0+)}^{\mu\nu} = \epsilon^\mu(+) \epsilon^\nu(0) + \epsilon^\mu(-) \epsilon^\nu(0), \quad (10)$$

take the form

$$T_{(++)}^{\mu\nu} = -g_{\perp}^{\mu\nu} = -\left(g^{\mu\nu} - \frac{n^\mu n'^\nu + n^\nu n'^\mu}{nn'} \right), \quad (11)$$

$$T_{(+-)}^{\mu\nu} = \frac{1}{|\bar{p}_\perp|^2} [p_\perp^\mu p_\perp^\nu - \tilde{p}_\perp^\mu \tilde{p}_\perp^\nu], \quad (12)$$

$$T_{(0+)}^{\mu\nu} = \frac{\sqrt{2}}{|\bar{p}_\perp|} p_\perp^\mu \left[\frac{1}{Q} q^\nu - \frac{2Q}{Q^2} q'^\nu \right], \quad (13)$$

where our choice of longitudinal vectors n and n' is described in Appendix A, as well as the exact parametrization of the transverse momenta $p_\perp$ and $\tilde{p}_\perp$. We also introduced the hard scale of the process given by $Q^2 = Q^2 + t$ and the dual transverse four-vector $\tilde{p}_\perp^\mu = \epsilon_{\perp}^{\mu\nu} p_\nu$ by means of the transverse Levi-Civita tensor: $\epsilon_{\perp}^{\mu\nu} = \epsilon^{\mu\nu\alpha\beta} n_\alpha n'_\beta / (nn')$.

We can express our final result for the DVCS amplitude in the following way:

$$\begin{aligned} i\mathcal{M}_{\chi, \text{DVCS}}^{s\lambda} &= \frac{-ie^3 \chi}{Q^2} \left[\mathcal{A}^{++} \left(\frac{\sqrt{2}}{2|\bar{p}_\perp|} \right) \left(\sum_{j=1}^3 B_j g(s, k', L_j, k) - i\lambda p_\mu \epsilon_{\perp}^{\mu\nu} j_\nu(s, k', k) \right) \right. \\ &\quad \left. + \mathcal{A}^{+-} \left(\frac{\sqrt{2}}{2|\bar{p}_\perp|} \right) \left(\sum_{j=1}^3 B_j g(s, k', L_j, k) + i\lambda p_\mu \epsilon_{\perp}^{\mu\nu} j_\nu(s, k', k) \right) - \mathcal{A}^{0+} \frac{2Q}{Q^2} g(s, k', q', k) \right], \end{aligned} \quad (14)$$

where $B_j \in \{1, -p^- Q^2, -\alpha\}$, $L_j \in \{r_2, q', n'\}$, and the functions g and j_μ are defined by Eqs. (B11) and (B12), respectively. If the CFFs are kept to twist-4 accuracy, as it is the case of Eqs. (15)–(17), then in the last line we can approximate $2Q/Q^2 \approx 2/Q$ since $Q^2 = Q^2 + t$. Otherwise, the expression (14) is exact to any order in kinematic twist and power of α_s .

In our analysis, we take into account twist-3 and twist-4 corrections, as well as NLO corrections to the leading-twist contribution. We write them schematically as

$$\mathcal{A}^{++}(\xi, t) = \int_{-1}^1 dx \frac{1}{\xi} \left\{ \left[C_0(x/\xi) + \frac{\alpha_S C_F}{4\pi} C_{\text{NLO}}^q(x/\xi) + \frac{t}{Q^2} C_{\text{HT}}(x, \xi) \right] H^\Sigma(x, \xi, t) + \frac{\alpha_S T_F}{4\pi} \frac{1}{\xi} C_{\text{NLO}}^g(x/\xi) H^g(x, \xi, t) \right\}, \quad (15)$$

$$\mathcal{A}^{+-}(\xi, t) = \int_{-1}^1 dx \frac{1}{\xi} \left\{ \frac{t}{Q^2} C_{\text{HT}}^{+-}(x, \xi) H^\Sigma(x, \xi, t) + \frac{\alpha_S T_F}{4\pi} \frac{1}{\xi} C_{\text{NLO}}^{+-}(x/\xi) H_T^g(x, \xi, t) \right\}, \quad (16)$$

$$\mathcal{A}^{0+}(\xi, t) = \int_{-1}^1 dx \frac{1}{\xi} \left\{ \frac{\sqrt{-t}}{Q} C_{\text{HT}}^{0+}(x, \xi) H^\Sigma(x, \xi, t) \right\}, \quad (17)$$

where

$$H^\Sigma(x, \xi, t) = \sum_{q \in \{u, d, s\}} e_q^2 H^q(x, \xi, t). \quad (18)$$

Our definitions of gluon H^g and quark H^q GPDs follow the review [9] with the skewness variable as in Eq. (A28), and e_q the electric charge of quark q relative to the positron. The NLO coefficient functions, as well as those at the HT are known [11,18,19,27–29], and were, for completeness, gathered and explained in detail in Appendix C.

B. Bethe-Heitler amplitudes

The detailed derivation of the expressions for BH amplitudes is presented in Appendix B. Here, we shortly summarize the results. For the Feynman diagram on the middle of Fig. 1, the corresponding amplitude reads

$$i\mathcal{M}_{\text{BH}}^{\text{sl}} = \frac{ie^3 F(t)}{(k' + q')^2 t} \left(\frac{\sqrt{2}}{2|\bar{p}_\perp|} \right) \sum_{i=1}^2 \sum_{\ell=1}^4 g(s, R_i, \tilde{r}_\ell, k) \times \left[\sum_{j=1}^3 B_j g(s, k', L_j, R_i) - i\lambda p_v \epsilon_\perp^{\nu\rho} j_\rho(s, k', R_i) \right], \quad (19)$$

where $B_j \in \{1, -p^- Q^2, -\alpha\}$, $L_j \in \{r_2, q', n'\}$, $R_i \in \{k', q'\}$, and $\tilde{r}_\ell \in \{r_1, r_2, r'_1, r'_2\}$.

Conversely, for the right plot in Fig. 1,

$$i\mathcal{M}_{\text{BHX}}^{\text{sl}} = \frac{ie^3 F(t)}{(k - q')^2 t} \left(\frac{\sqrt{2}}{2|\bar{p}_\perp|} \right) \sum_{i=1}^2 \sum_{\ell=1}^4 \sigma(\tilde{R}_i) g(s, k', \tilde{r}_\ell, \tilde{R}_i) \times \left[\sum_{j=1}^3 B_j g(s, \tilde{R}_i, L_j, k) - i\lambda p_v \epsilon_\perp^{\nu\rho} j_\rho(s, \tilde{R}_i, k) \right], \quad (20)$$

where, again, $B_j \in \{1, -p^- Q^2, -\alpha\}$, $L_j \in \{r_2, q', n'\}$, and $\tilde{r}_\ell \in \{r_1, r_2, r'_1, r'_2\}$, while $\tilde{R}_i \in \{k, q'\}$ carry the signature function $\sigma(\tilde{R}_i) \in \{+1, -1\}$, respectively.

III. MODEL FOR HELIUM-4 GPDS

Let us now model the helium-4 GPDs as rigorously as possible, while maintaining a sufficiently flexible parametrization so that it reproduces existing data, has the potential to describe future data, and allows for the reliable extraction of tomographic features. As in the case of modeling nucleon GPDs (see, for instance, Ref. [30]), the fundamental property of polynomiality is preserved by starting from double distributions. The limits $(\xi, t) \rightarrow (0, 0)$ of GPDs $H^{q,g}$ are readily enforced by using a Radyushkin-type ansatz [31] with nuclear PDFs extracted from experimental data [32]. For valence quarks, we require that the measured dependence on the t variable of the electromagnetic form factor is recovered as the first x moment of the GPD. For sea quarks and gluons, the dependence on t will be constrained by sparse coherent DVCS experimental data. In all these cases, we rely on isospin

symmetry, assuming an exact symmetry between up and down quarks. A somewhat different approach is adopted for modeling the GPD H_T^g . Given the limited empirical constraints currently available for this distribution, we primarily rely on theory-driven constraints in this case.

Our models for helium-4 GPDs, $H^i(x, \xi, t; \mu^2)$, where $i = \{u_{\text{val}}, u_{\text{sea}}, d_{\text{val}}, d_{\text{sea}}, s, g\}$ denotes the parton type explicitly distinguishing between valence and sea components, are based on double distributions, $F^i(\beta, \alpha, t; \mu^2)$. Here, μ is the energy scale at which these distributions are fitted. Both types of objects are related by

$$H^i(x, \xi, t; \mu^2) = \int_{-1}^1 d\beta \int_{-1+|\beta|}^{1-|\beta|} d\alpha \delta(\beta + \xi\alpha - x) \times F^i(\beta, \alpha, t; \mu^2). \quad (21)$$

In the current analysis, we do not include any D term [33], as the available data lack sensitivity to the real parts of the amplitudes. In the following equations, for brevity, we will suppress the scale dependence.

The double distributions are expressed as products of generalized t -dependent PDFs, $f_i^A(\beta, t)$, and a profile function, $h_i(\beta, \alpha)$, which generates the dependence on the skewness variable ξ :

$$F^i(\beta, \alpha, t) = f_i^A(\beta, t) h_i(\beta, \alpha). \quad (22)$$

We stress that $H^i(x, \xi, t)$ denotes the total contribution of i parton to helium-4. In particular, for $x > 0$ the so-called forward limits of our models for quarks, $q = \{u_{\text{val}}, u_{\text{sea}}, d_{\text{val}}, d_{\text{sea}}, s\}$, are

$$H^q(x, 0, 0) = f_q^A(x), \quad (23)$$

which represents the parton distribution functions describing the probability of finding a quark carrying a given fraction of the helium-4 momentum (rather than a given fraction of the proton-in-helium-4 momentum). The corresponding forward limit for gluons is

$$H^g(x, 0, 0) = x f_g^A(x). \quad (24)$$

We also note that this work employs the full evolution equations for GPDs, rather than the so-called forward evolution governed by PDFs that is often customary in phenomenological analyses. For this purpose, we utilize the APFEL++ package [34,35], which implements the evolution equations for GPDs at LO.

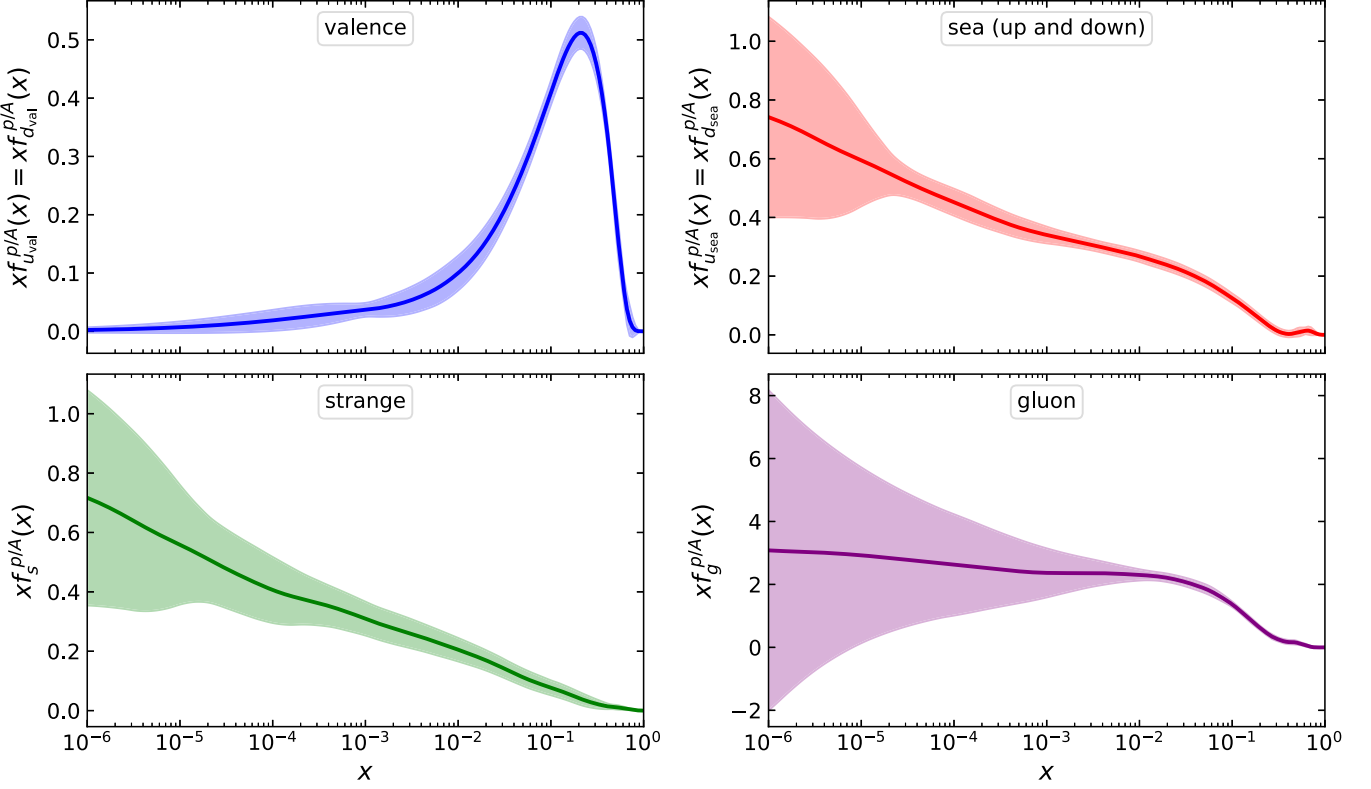
The generalized PDF is

$$f_i^A(\beta, t) = f_i^A(\beta) \frac{k_i(|\beta|, t)}{k_i(|\beta|, 0)}, \quad (25)$$

where $f_i(\beta)$ accounts for both partons ($\beta > 0$) and antipartons ($\beta < 0$), and we have

(1) for $A|\beta| \leq 1$

$$f_i^A(\beta) = A^2 \begin{cases} \Theta(\beta) f_i^{p/A}(A|\beta|), & \text{for } i = u_{\text{val}}, d_{\text{val}}, \\ \text{sgn}(\beta) f_i^{p/A}(A|\beta|), & \text{for } i = u_{\text{sea}}, d_{\text{sea}}, s, \\ |\beta| f_i^{p/A}(A|\beta|), & \text{for } i = g, \end{cases} \quad (26)$$

FIG. 3. Parametrizations of PDFs for helium-4 obtained by NNPDF group [32] for $\mu^2 = 2 \text{ GeV}^2$.

(2) otherwise

$$f_i^A(\beta) = 0. \quad (27)$$

Here, $\Theta(\cdot)$ and $\text{sgn}(\cdot)$ are the step and signum functions, respectively, while $f_i^{p/A}(A|\beta|)$ are PDFs of valence and sea quarks, and of gluons in the nucleon in helium-4. In this analysis we have chosen to use nNNPDF30 parametrizations [32]. Since these parametrizations are defined with respect to proton momentum, we rescale $|\beta|$ by A . Furthermore, since they give parton densities per nucleon, in order to ensure the correct number of valence quarks per helium-4 nucleus, i.e.,

$$\int_{-1}^1 d\beta f_{u_{val}}^A(\beta) = 2Z + (A - Z) = 6 \quad (28)$$

and

$$\int_{-1}^1 d\beta f_{d_{val}}^A(\beta) = Z + 2(A - Z) = 6, \quad (29)$$

in Eq. (26) we include the prefactor A (another power of A due to the rescaling of the arguments). To avoid probing the PDFs in unknown domains, we explicitly set $f_i^A(\beta) = 0$ for $\beta > 1/A$.

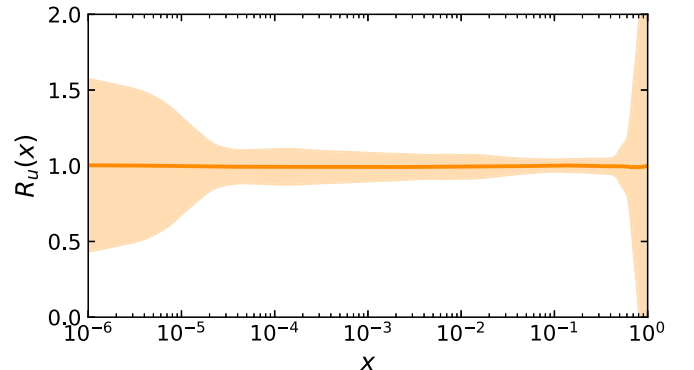
Figure 3 shows the nNNPDF30 parametrizations of the helium-4 PDFs, while Fig. 4 shows the nuclear modification ratio, defined as

$$R_i(x) = \frac{A f_i^{p/A}(x)}{Z f_i^p(x) + (A - Z) f_i^n(x)}. \quad (30)$$

Here, $f_i^p(x)$ and $f_i^n(x)$ denote the PDFs of the proton and neutron, respectively, coming from the same set of phenomenological solutions, and no distinction is made between valence and sea contributions.

In this analysis, we refrain from using fully factorized dependencies in x and t . Instead, for valence quarks, $i = \{u_{val}, d_{val}\}$, we have

$$k_i(|\beta|, t) = \left(\frac{1}{1 - p_0(1 - |\beta|)^2 t} \right)^{p_1} \times \prod_{j=1}^n \left(\frac{|p_{2,j} + t|}{|p_{2,j} + t| - p_{3,j}(1 - |\beta|)^2 t} \right)^{p_{4,j}}. \quad (31)$$

FIG. 4. The nuclear modification ratio $R_q(x)$, see Eq. (30), for up quarks and $\mu^2 = 2 \text{ GeV}^2$.

This ansatz consists of two distinct components. The first, a dipole-like ansatz, captures the main trend in t for a given x . The second, represented by the product, accounts for n diffractive minima occurring at $-t = \{p_{2,1}, p_{2,2}, \dots\}$. Our ansatz ensures that

$$\lim_{|\beta| \rightarrow 1} k_i(|\beta|, t) = \text{const.}, \quad (32)$$

which reflects a theory-driven constraint arising from hadron tomography. Specifically, in the limit $x = 1$, where the entire momentum of the hadron is carried by a single parton, the position of this parton coincides with the origin of the coordinate system in which the hadronic tomography is defined.

The free parameters of Eq. (31) are fixed by a fit to the helium-4 elastic form factor data, as detailed in the Sec. IV, using the relation

$$F(t) = \frac{1}{Z} \int_{-1}^1 dx (e_u f_{u,\text{val}}^A(x, t) + e_d f_{d,\text{val}}^A(x, t)). \quad (33)$$

For sea quarks and gluons, $i = \{u_{\text{sea}}, d_{\text{sea}}, s, g\}$, we use a different ansatz, namely,

$$k_i(|\beta|, t) = \exp(p(1 - |\beta|^2)t). \quad (34)$$

The slope p will be fitted to CLAS data [21], as described in the Sec. IV.

For the profile function in Eq. (22), we use the classic Radyushkin's ansatz,

$$h_i(\beta, \alpha) = \frac{\Gamma(2b_i + 2)}{2^{2b_i+1} \Gamma^2(b_i + 1)} \frac{((1 - |\beta|)^2 - \alpha^2)^{b_i}}{(1 - |\beta|)^{2b_i+1}}, \quad (35)$$

whose normalization

$$\int_{-1+|\beta|}^{1-|\beta|} d\alpha h_i(\beta, \alpha) = 1 \quad (36)$$

ensures the desired reduction at $\xi = 0$, already expressed in Eqs. (23) and (24).

The coefficients controlling the strength of the skewness effect are $b_i = 1$ for valence quarks and $b_i = 2$ for sea quarks and gluons, which are typical values adopted in GPD models for the proton, for instance in Ref. [30].

For the purpose of our analysis, we also need a model for H_T^g , which contributes at NLO to the $\mathcal{A}^{+-}$ amplitude. Since little is known about this GPD, we bound it by $H^g(x, \xi, 0)$ as $t \rightarrow 0$, while for $t \ll 0$ we use theory-driven constraints, namely, the positivity inequality [36] derived in Ref. [37] for $x > \xi$:

$$|H_T^g(x, \xi, t)| < \frac{1}{\xi} \sqrt{\frac{x^2 - \xi^2}{1 - \xi^2}} \sqrt{\frac{t}{t - t_0}} \sqrt{\frac{t_0}{t - t_0}} \times \sqrt{g\left(\frac{x + \xi}{1 + \xi}\right) g\left(\frac{x - \xi}{1 - \xi}\right)}, \quad (37)$$

where

$$t_0 = -4M^2 \frac{\xi^2}{1 - \xi^2}. \quad (38)$$

We note that

$$\sqrt{\frac{t}{t - t_0}} \sqrt{\frac{t_0}{t - t_0}} = 2M\xi \sqrt{-1/t} + O(\xi^3), \quad (39)$$

so for $\xi \rightarrow 0$

$$|H_T^g(x, \xi \rightarrow 0, t)| < x g(x) \sqrt{\frac{1}{-t/(4M^2)}}. \quad (40)$$

This suggests the following ansatz, which we adopt in this analysis

$$H_T^g(x, \xi, t) = H^g(x, \xi, 0) \sqrt{\frac{1}{1 - t/(4M^2)}}. \quad (41)$$

The main feature of this solution is its conservativeness: at $t = 0$, we find $H_T^g(x, \xi, 0) = H^g(x, \xi, 0)$, whereas for $|t| > 0$, $H_T^g(x, \xi, t)$ exhibits a much milder dependence on t than $H^g(x, \xi, t)$.

It remains instructive to verify whether our model for H_T^g , now based on the double distribution representation, fulfills the positivity constraint for $\xi > 0$.

The results are illustrated in Fig. 5, which displays our model against the positivity bound. We observe that the constraint is only violated in the large- x tail of the distribution. This is expected, as in our analysis the gluon PDF for helium-4 vanishes identically for $x > 1/A$, see Eq. (27). Consequently, the right-hand side of Eq. (37) always vanishes for $x > (1 + \xi(1 - A))/A$, whereas the left-hand side retains a nonzero dependence due to the inherent properties of the radon transform. We note that since our phenomenological analysis focuses on the small- ξ region where these effects are suppressed, this localized violation at large x does not undermine the overall validity of our results.

As the positivity constraint does not fix the sign of the GPD H_T^g , it would seem reasonable to vary it in the phenomenological analysis. We, however, refrain from doing so, as we find the contribution from H_T^g , despite taking a conservative estimate, to be negligible (see the next section).

IV. RESULTS

In this section, we collect the main results of our analysis of helium-4 GPDs. We first determine the model parameters from elastic form factor and beam-spin asymmetry data, and then discuss the resulting tomographic picture. We also examine the impact of NLO and kinematic higher-twist effects on the description of present data and on predictions for future JLab12 observables.

Before turning to the results in detail, we briefly comment on how the extraction framework used in this work is related to the kinematics employed for tomography. In the present work, the GPDs are extracted using the approach of Refs. [11,38], which is particularly well suited for the systematic organization of kinematic higher-twist corrections. When discussing tomography, however, we employ the usual symmetric frame [9], in which the $\xi = 0$ limit corresponds to a purely transverse momentum transfer, $t = -\Delta_T^2$, thus recovering the conventional impact-parameter representation

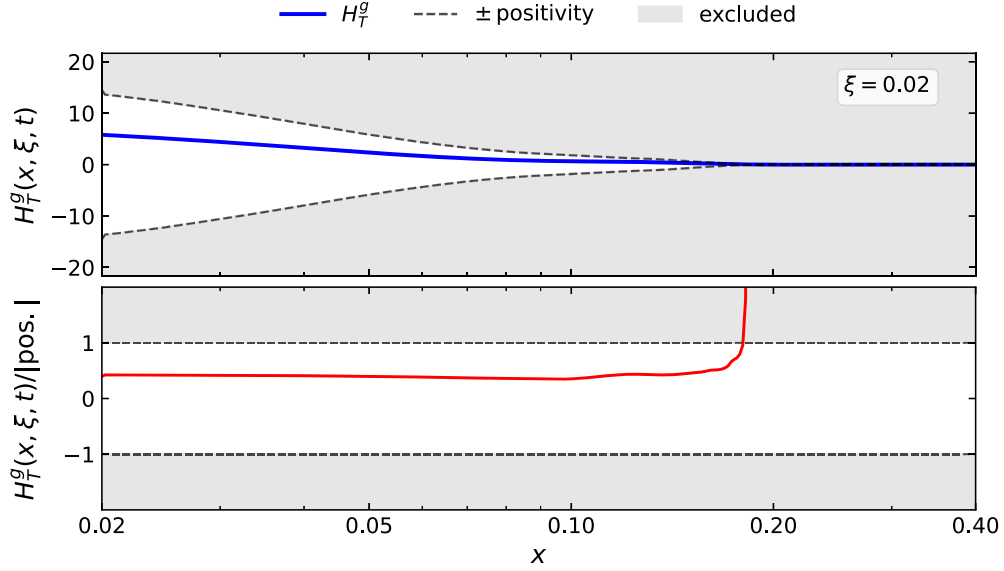


FIG. 5. Our model for the GPD H_T^g compared against the positivity bound, see Eq. (37), for $\xi = 0.02$ and $t = -10 \text{ GeV}^2$.

for densities:

$$f_i^A(x, \mathbf{b}_T) = \int \frac{d^2 \Delta_T}{(2\pi)^2} e^{-i\mathbf{b}_T \cdot \Delta_T} H^i(x, 0, -\Delta_T^2). \quad (42)$$

The impact parameter, $\mathbf{b}_T = (b_x, b_y)$, is defined in a coordinate system whose origin is set at the center of momentum of all proton constituents [12].

As a first step of our fit procedure, we determine the free parameters of Eq. (31) from the helium-4 elastic form factor data. We model the first two diffractive minima ($n=2$). The fit, based on the data from Refs. [39–43], is shown in Fig. 6, together with the corresponding spatial charge distribution in helium-4, obtained by Fourier transforming the fitted form factor to impact-parameter space. No constraints are imposed on the fit parameters, and we obtain $p_0 = (0.088 \pm 0.016)/\text{GeV}^2$, $p_1 = 39.5 \pm 6.4$, $p_{2,1} = (0.3833 \pm 0.0024) \text{ GeV}^2$, $p_{3,1} = 3.49 \pm 0.16$, $p_{4,1} = 0.985 \pm 0.027$, $p_{2,2} = (1.993 \pm 0.018) \text{ GeV}^2$, $p_{3,2} = 2.38 \pm 0.86$, and $p_{4,2} = 0.988 \pm 0.099$. Due to the isospin symmetry, we use the same parameters for valence up and down quarks.

The last remaining parameter of the model, namely the slope p , is fitted to the CLAS data [21], for the beam-spin asymmetry, defined as

$$A_{LU}(\phi) = \frac{\sigma_{-}^{++} + \sigma_{-}^{+-} - \sigma_{-}^{-+} - \sigma_{-}^{--}}{\sigma_{-}^{++} + \sigma_{-}^{+-} + \sigma_{-}^{-+} + \sigma_{-}^{--}}, \quad (43)$$

where the differential cross section, $\sigma_{\chi}^{s\lambda}$, for a beam of charge sign χ and helicity s , producing a real photon with polarization λ was defined in Eq. (2).

We perform three distinct fits, calculating asymmetries using LO/LT, NLO/LT, and NLO/HT coefficient functions. Each calculation relies on specific types of GPDs, yields distinct amplitudes, and leads to varying levels of agreement with the experimental data, which we summarize in Table I. We note that in the fit utilizing LO/LT coefficient functions, the χ^2 function does not reach a minimum. Instead, it saturates for large values of p . Therefore, we are only able to provide a lower limit for this parameter. This indicates that when using LO/LT coefficient functions, the CLAS data prefers to have the sea contribution totally suppressed.

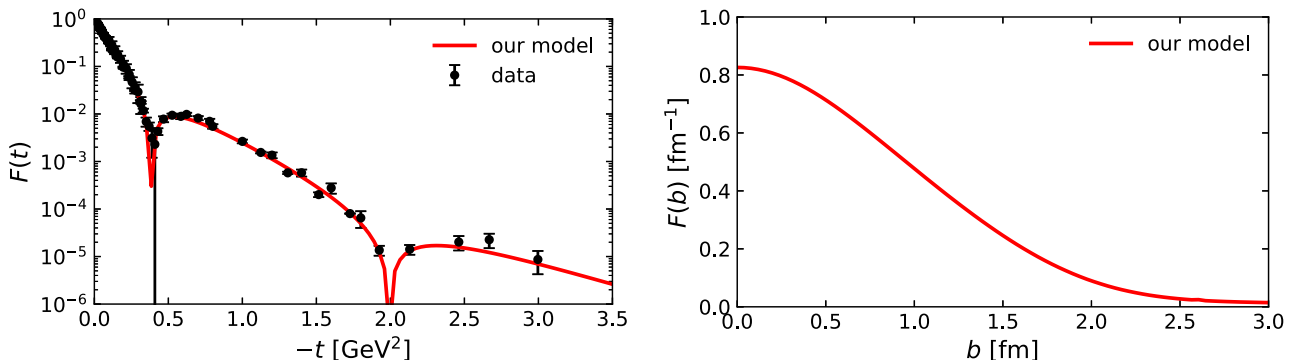


FIG. 6. (left) Fit to helium-4 elastic form factor data. (right) Fourier transform of fitted helium-4 elastic form factor, where $\int_0^\infty db F(b) = 1$.

TABLE I. Comparison of our A_{LU} asymmetry estimates with CLAS data [21]. The three rows represent fits using different coefficient functions, corresponding to varying levels of precision in the α_s and kinematic twist expansions. These fits utilize the specific GPD types and amplitudes listed in the table. The agreement between constrained models and the experimental data is quantified by the global χ^2/n value, where n is the number of points measured.

Precision in		GPDs involved	Non-vanishing amplitudes	Fitted value of p [GeV ⁻²]	Global χ^2/n
α_s expansion	Kinematic twist expansion				
LO	LT	H^q	$\mathcal{A}^{++}$	>70	1.31
NLO	LT	H^q, H^g, H_T^g	$\mathcal{A}^{++}, \mathcal{A}^{+-}$	$23.4^{+7.4}_{-3.9}$	1.28
NLO	HT	H^q, H^g, H_T^g	$\mathcal{A}^{++}, \mathcal{A}^{+-}, \mathcal{A}^{0+}$	$22.0^{+5.5}_{-2.1}$	1.00

For our nominal fit utilizing the NLO/HT coefficient functions we obtain $p = 22.0^{+5.5}_{-2.1} \text{ GeV}^{-2}$. Such a value implies a steeper dependence on t than in the valence sector and, consequently, a broader spatial distribution of sea quarks and gluons, as expected. More robust constraints on the tomography of sea quarks and gluons should become possible with upcoming JLab data and, in the longer term, with data from the Electron Ion Collider.

Figure 7 compares the A_{LU} asymmetries measured by CLAS [21] with our predictions. For each (x_B, t, Q^2) kinematic bin, three curves are shown, obtained from the same underlying GPD parametrization but using coefficient functions evaluated at LO/LT, NLO/LT, and NLO/HT accuracy.

To illustrate the main properties of the fitted model, we show its dependence on the relevant kinematic variables in the following figures. The dependence of our model on the variable x is shown in Fig. 8 for several values of ξ . For small skewness, the support is similar to that of an ordinary PDF, with the distributions vanishing for $x \gtrsim 0.25$, while for larger ξ it becomes broader, reflecting the support properties of the underlying double-distribution representation. In Fig. 9, we illustrate the dependence on the variable t at fixed x and $\xi = 0$, including the corresponding hadron tomography images. The figure clearly demonstrates the steep fitted slope of the sea-quark and gluon contributions.

Figure 10 shows the unpolarized cross sections corresponding to the kinematic bins used in the CLAS measurement. The figure also includes the contribution from the BH process, which has been scaled down by a factor of 20 for visual clarity. These plots illustrate the impact of NLO and HT corrections on both the shape and magnitude of the pure DVCS cross section. Consequently, constraining the unpolarized pure DVCS cross section would provide a stringent test of these perturbative frameworks, which is a task that remains challenging due to the BH dominance.

The three helicity amplitudes contributing to the DVCS process, shown as functions of x_B , t , and Q^2 in Fig. 11, clearly exhibit the dominance of the imaginary part of $\mathcal{A}^{++}$. The figure also confirms the expected pattern of higher-twist (HT) corrections, which become suppressed at large Q^2 and small t . It furthermore shows that the gluon-transversity contribution to $\mathcal{A}^{+-}$ remains small compared with the HT contribution.

Figure 12 presents three-dimensional (3D) tomographic images of the ^4He nucleus, shown separately for valence quarks, sea quarks, and gluons. We see that the valence quarks, on average, carry higher momenta (as expected from

PDF analyses) and are embedded within a broader spatial distribution of softer sea quarks and gluons.²

Finally, Fig. 13 presents our predictions for the A_{LU} asymmetries and unpolarized cross sections at typical JLab12 kinematic settings. We select two values of Q^2 , namely 1.2 GeV^2 and 3 GeV^2 , while keeping y fixed, demonstrating the suppression of HT effects in the latter case.

V. CONCLUSIONS

Our phenomenological study of DVCS on the helium-4 nucleus has highlighted the importance of including both higher-order corrections in α_s and kinematical twist-3 and twist-4 contributions to the QCD amplitude. This improved level of accuracy allowed for a successful description of the beam-spin asymmetry data and for constraining the gluon, valence-quark, and sea-quark GPDs of the helium-4 nucleus. It has also demonstrated that a first analysis of the impact-parameter dependence of GPDs is already possible, thereby opening the way to a first quark and gluon tomography of a light nucleus. More generally, our results indicate that such an extended treatment is likely to be fruitful, and perhaps even necessary, also in phenomenological analyses of Compton processes on the proton and neutron. At the same time, the role of genuine twist-3 and twist-4 effects remains an open theoretical issue that must be clarified before fully definitive conclusions can be drawn from experimental data at relatively low values of the hard scale Q^2 .

The experimental data used in our analysis originate from an early Jefferson Lab measurement with a 6 GeV electron beam. More precise data at higher beam energy and larger Q^2 are expected from an ongoing experiment [44], for which the present study already provides phenomenological predictions. A similar analysis for the higher-energy regime of future electron-ion colliders [45,46] lies beyond the scope of the present work. In parallel, several complementary exclusive reactions may provide important additional information on the tomography of helium-4. The crossed process, namely timelike Compton scattering in exclusive lepton-pair photo-production [47], is known [48] to be a valuable complement to DVCS, in particular because gluon contributions enter the NLO amplitudes of the two processes in quite different ways

²A complementary representation is included in the companion paper [20].

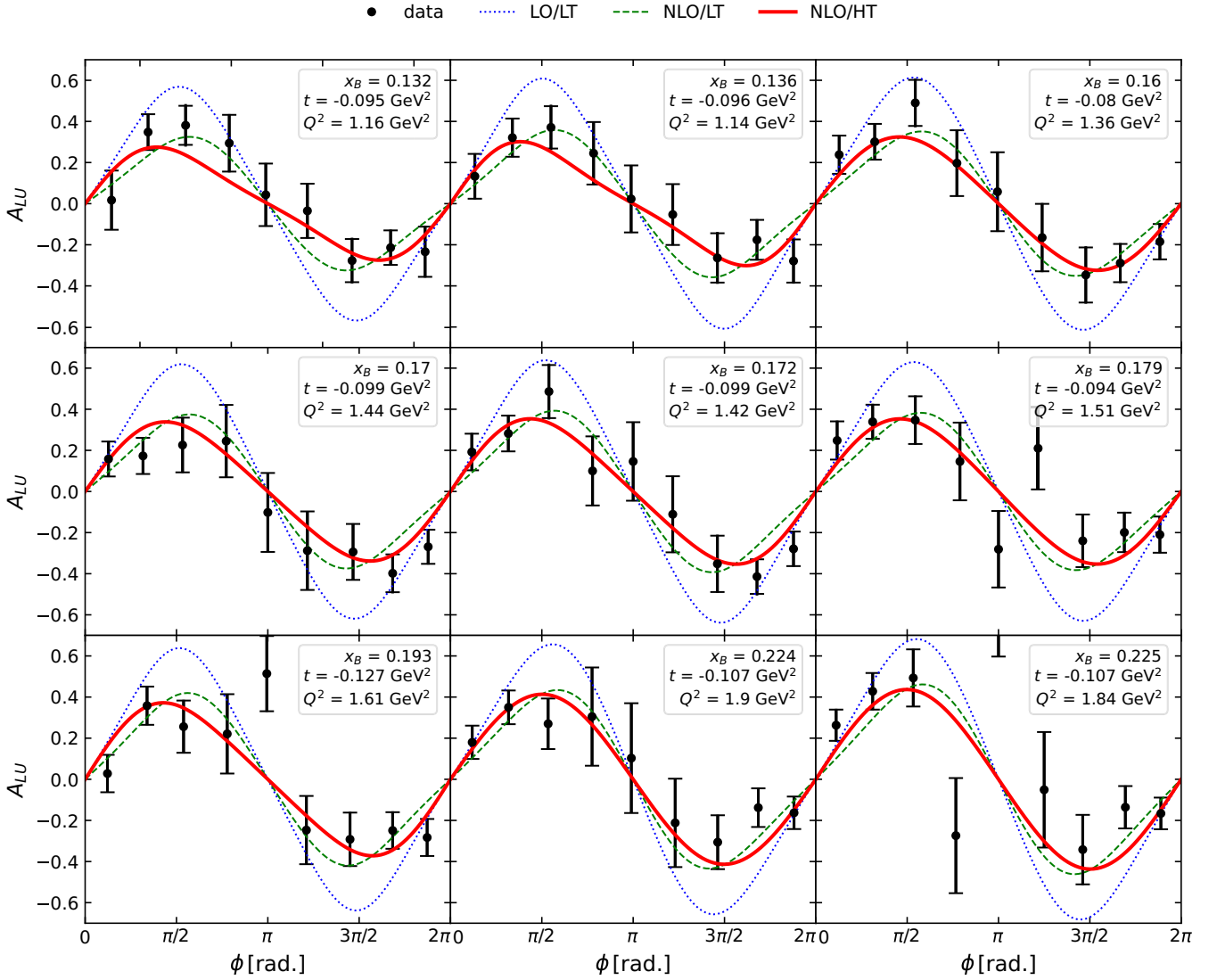


FIG. 7. Experimental data for the A_{LU} asymmetry [21] compared with our model. See the legend at the top of the figure for more information.

[18,49,50]. The more challenging process of double DVCS [51], i.e., the electroproduction of a lepton pair, for which higher-twist calculations are already available [11], may also offer new insight into the tomography of helium-4. Because of the charge-conjugation symmetry properties of DVCS and TCS, these reactions probe only the charge-even part of quark GPDs. Access to the charge-odd sector requires considering other exclusive channels, such as diphoton production at large invariant mass [52–54] or photon-plus-vector-meson production [55,56]. The theoretical and experimental study of such processes on helium-4 will be necessary to complete the quark tomographic picture of this nucleus.

ACKNOWLEDGMENTS

We acknowledge useful discussions and correspondence with R. Dupré, V. Guzey, and C. Mezrag. This research was funded in whole or in part by the National Science Centre, Poland (Grant IMPRESS-U No. 2024/06/Y/ST2/00155 and

Grant SONATA BIS-15 No. 2025/58/E/ST2/00045). The research of V.M.-F. was funded in part by l'Agence Nationale de la Recherche (ANR), Project No. ANR-23-CE31-0019.

DATA AVAILABILITY

There are no publicly available research data or software supporting this manuscript. Requests for further information or data should be sent to the authors.

APPENDIX A: KINEMATICS

In this appendix, we describe the kinematics of the process (1). The core of our calculation is performed in the “target rest frame” (TRF) that is related to the usual Trento frame [26] by a 180° rotation around the y axis. This TRF frame coincides with TRF-I of the DDVCS studies in Refs. [51,57] and it is depicted in Fig. 2, allowing us to stay consistent with previous literature on exclusive processes.

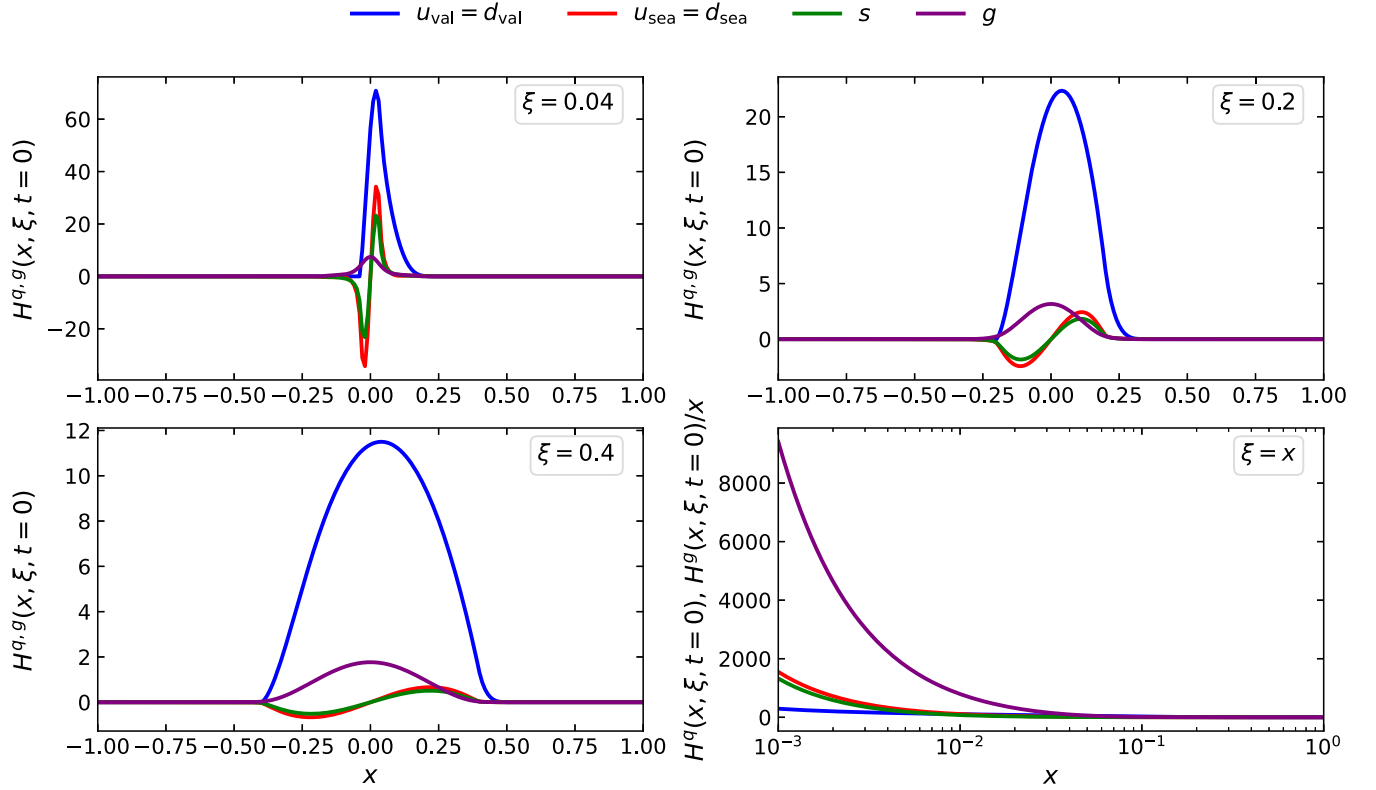


FIG. 8. Dependence of our GPD models for helium-4 on the variable x for $\xi = \{0.04, 0.2, 0.4, x\}$ at $|t| = 0.2 \text{ GeV}^2$ and $\mu^2 = 2 \text{ GeV}^2$. In the plot for $\xi = x$ the gluon GPD is divided by x .

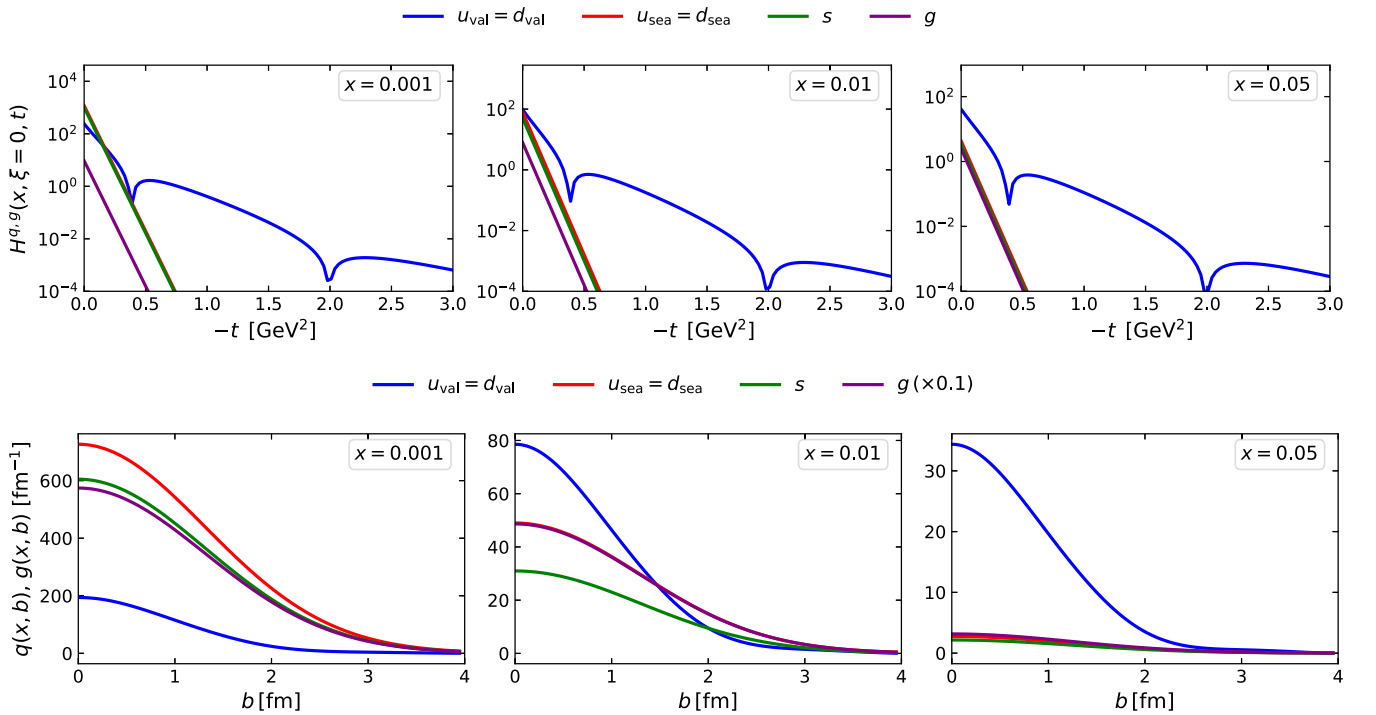


FIG. 9. Dependence of our GPD models for helium-4 on the variable t for $\xi = 0$, $\mu^2 = 2 \text{ GeV}^2$ and $x = \{0.001, 0.01, 0.05\}$, and corresponding tomographic pictures, where $\int_0^\infty db q(x, b) = q(x)$ and similar for gluons. In the lower plots, the gluon curves are scaled by a factor of 0.1.

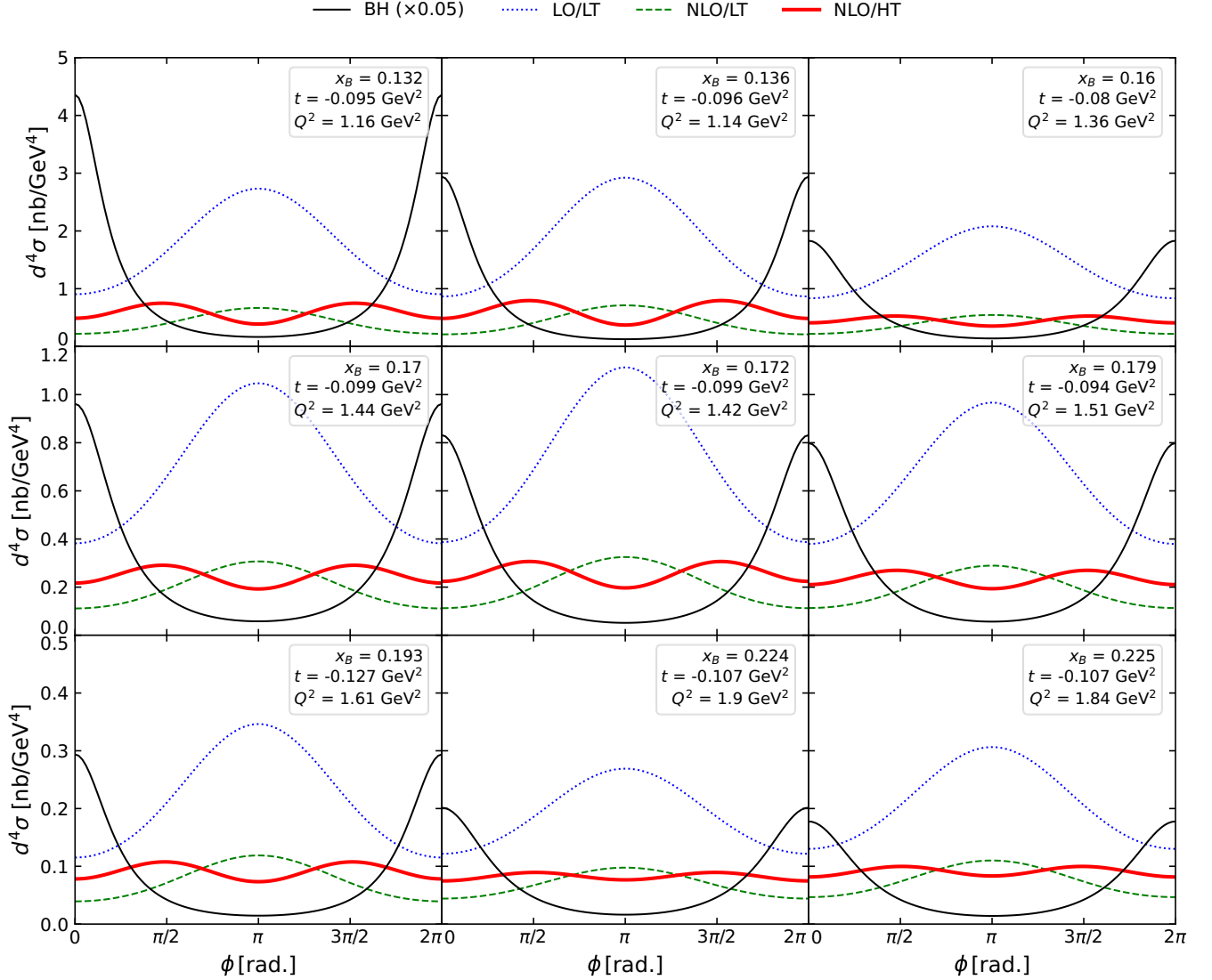


FIG. 10. Differential cross section, $d^4\sigma/(dx_B dt dQ^2 d\phi)$, for the BH and DVCS subprocesses. For better visibility, the BH curves have been scaled down by a factor of 20.

1. Momenta parametrization

In TRF, the target is at rest so that the initial-state momentum is given by

$$p^\mu = (M, 0, 0, 0), \quad (\text{A1})$$

while the z axis is opposite to the incoming photon momentum,

$$q^\mu = (q^0, 0, 0, -|q^3|). \quad (\text{A2})$$

Accounting for its spacelike virtuality, $q^2 = -Q^2 < 0$, as well as the nuclear Bjorken variable, $x_A = Q^2/(2pq)_{\text{TRF}} = Q^2/(2Mq^0)$, we find

$$q^\mu = \frac{Q}{\omega}(1, 0, 0, -\sqrt{1 + \omega^2}), \quad \omega = \frac{2x_A M}{Q}. \quad (\text{A3})$$

The final-state nucleus has a momentum

$$p'^\mu = (p'^0, |\vec{p}'|s_N c_\varphi, |\vec{p}'|s_N s_\varphi, |\vec{p}'|c_N), \quad (\text{A4})$$

where $s_N = \sin \theta_N$, $c_N = \cos \theta_N$ for θ_N the polar angle with respect to TRF, and equivalently $s_\varphi = \sin \varphi$, $c_\varphi = \cos \varphi$ for its azimuth φ . With the definition of the Mandelstam $t = (p' - p)^2$ and the above parametrization of the initial- and final-state momenta, we find

$$p'^0 = M - \frac{t}{2M}, \quad |\vec{p}'| = \sqrt{-t \left(1 - \frac{t}{4M^2}\right)}. \quad (\text{A5})$$

Momentum conservation fixes the angle θ_N . Indeed, the momentum of the the real outgoing photon is related to the momentum of incoming photon and the momentum transfer ($\Delta = p' - p$) by

$$q' = q - \Delta \Rightarrow q'^2 = 0 = -Q^2 + t - 2q\Delta, \quad (\text{A6})$$

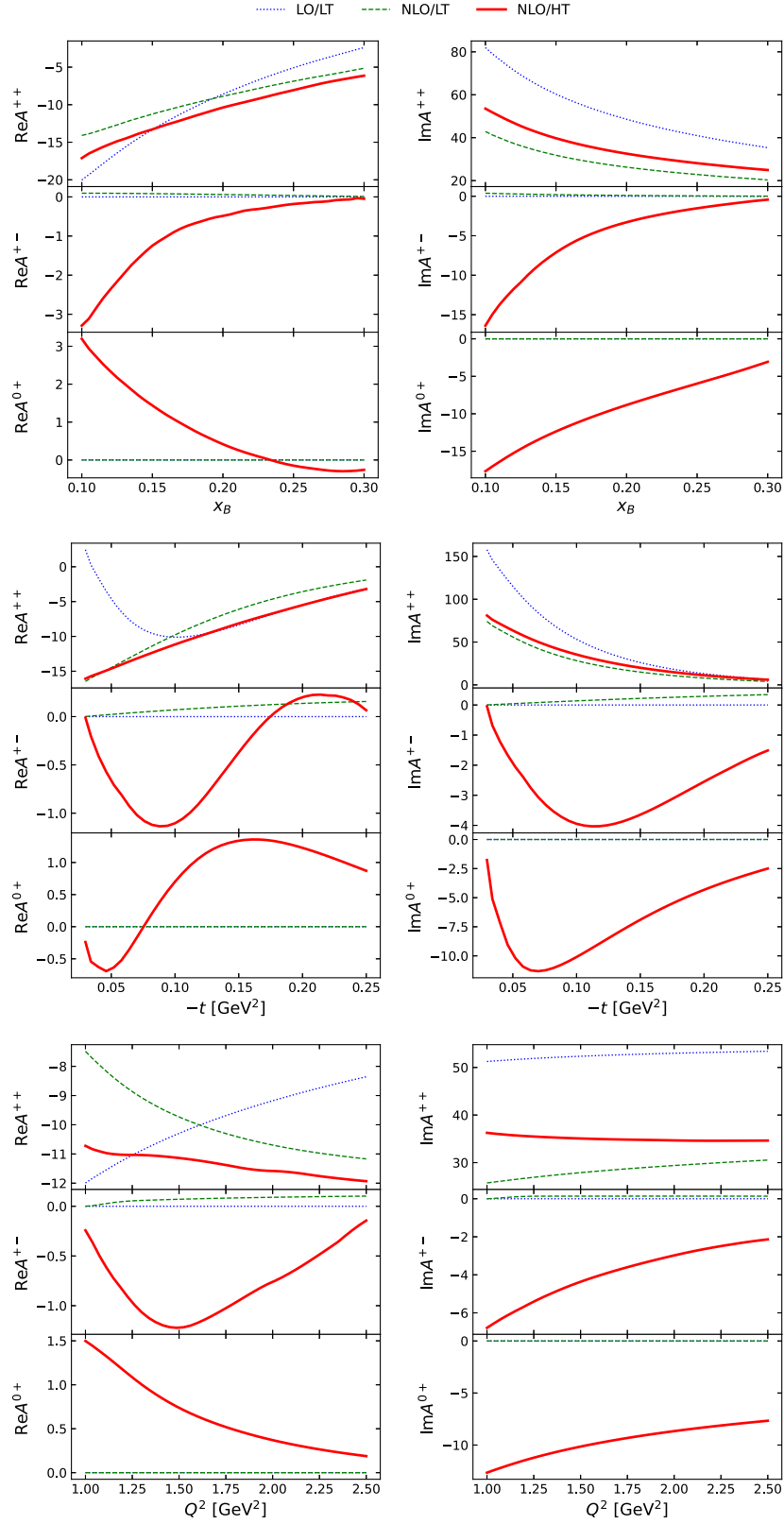


FIG. 11. Amplitudes as a function of x_B (top row) for $t = -0.1 \text{ GeV}^2$ and $Q^2 = 1.5 \text{ GeV}^2$; $-t$ (middle row) for $x_B = 0.18$ and $Q^2 = 1.5 \text{ GeV}^2$; and Q^2 (bottom row) for $x_B = 0.18$ and $t = -0.1 \text{ GeV}^2$.

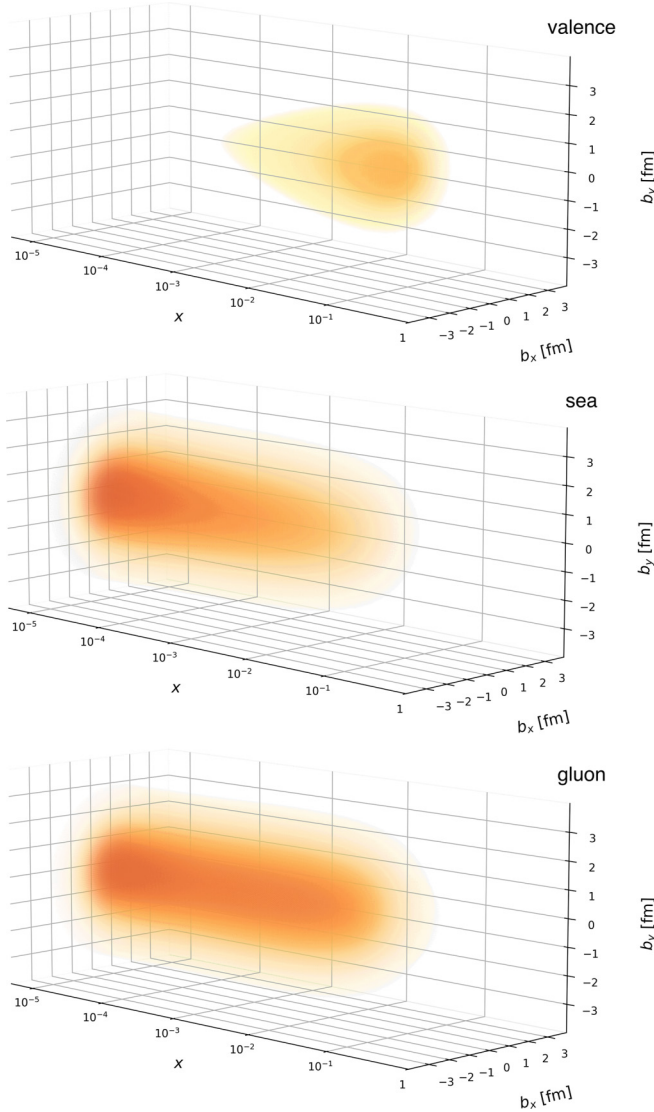


FIG. 12. Tomographic pictures of ^4He nuclei, $xq(x, b_x, b_y)$ and $xg(x, b_x, b_y)$. The valence distribution indicates contributions coming from both up and down quarks, while the sea distribution includes contributions from up, down and strange quarks and antiquarks.

where

$$\begin{aligned}\Delta^\mu &= (p' - p)^\mu = \left(-\frac{t}{2M}, \vec{p}'\right) \Rightarrow q\Delta \\ &= \frac{Q}{\omega} \left(-\frac{t}{2M} + \sqrt{1 + \omega^2 |\vec{p}'|^2} c_N\right).\end{aligned}\quad (\text{A7})$$

Introducing this product into Eq. (A6), we can solve c_N to be

$$c_N = -\frac{\omega^2(Q^2 - t) - 2x_A t}{4Mx_A \sqrt{1 + \omega^2 |\vec{p}'|^2}}. \quad (\text{A8})$$

Thanks to the null virtuality of the outgoing photon, $q'^\mu = (q^0, \vec{q}')$ such that $(q^0)^2 = |\vec{q}'|^2$. Again, due to momentum conservation

$$t = (p' - p)(q - q') = p'\Delta - pq + pq'. \quad (\text{A9})$$

With the above parametrization,

$$\left. \begin{aligned} p'\Delta &= \underbrace{\vec{p}'\Delta}_0 - p\Delta = -M\Delta^0 = \frac{t}{2}, \\ pq &= p^0 q^0 = \frac{MQ}{\omega}, \\ pq' &= p^0 q'^0 = Mq'^0, \end{aligned} \right\} \Rightarrow q'^0 = \frac{t}{2M} + \frac{Q}{\omega}, \quad (\text{A10})$$

which leads us to

$$q'^\mu = \left(\frac{t}{2M} + \frac{Q}{\omega}, -\vec{p}' + \vec{q}\right). \quad (\text{A11})$$

In TRF and for a massless lepton beam,

$$k^\mu = E(1, \sin \theta_e, 0, \cos \theta_e), \quad E = \frac{Q}{y\omega}, \quad (\text{A12})$$

where the inelasticity variable has been defined as $y = pq/(pk)_{\text{TRF}} = Q/(\omega E)$.

Finally, the momentum of the scattered electron reads

$$k' = -\Delta + k - q' \Rightarrow k'^\mu = \left(\frac{(1-y)Q}{y\omega}, \vec{k} - \vec{q}\right). \quad (\text{A13})$$

Note that only the momenta p' , q' have a dependence on the azimuthal angle ϕ whose relation to its Trento counterpart ϕ is

$$\phi = \begin{cases} \pi - \phi, & \text{if } \phi \in [0, \pi], \\ 3\pi - \phi, & \text{if } \phi \in (\pi, 2\pi). \end{cases} \quad (\text{A14})$$

2. Longitudinal vs transverse plane

The collinear factorization that gives rise to GPDs and CFFs requires a separation of dominant longitudinal kinematics (characterized by lightlike vectors) and sub-leading transverse components that are orthogonal to the former. In such a framework, any four-vector can be written as

$$v^\mu = v^+ n'^\mu + v^- n^\mu + v_\perp^\mu, \quad (\text{A15})$$

where $n^2 = n'^2 = 0$, $v_\perp n = v_\perp n' = 0$, and $nn' \neq 0$. Following the procedure of Refs. [11,27] and to obtain a leading-twist limit for the Compton tensor that is compatible with charge conservation ($q_\nu T_{\text{LT}}^{\mu\nu} = T_{\text{LT}}^{\mu\nu} q'_\mu = 0$), we choose purely longitudinal photon momenta

$$n^\mu = \alpha q^\mu + \beta q'^\mu, \quad n'^\mu = \alpha' q^\mu + \beta' q'^\mu, \quad (\text{A16})$$

where α , β , α' , β' are real parameters constrained by $n^2 = n'^2 = 0$, together with the normalization

$$nn' = -\Delta q' = \frac{t + Q^2}{2} \neq 0. \quad (\text{A17})$$

The solution after rescaling $Q^2 n \rightarrow n$ and $n'/Q^2 \rightarrow n'$ is

$$n^\mu = Q^2 q'^\mu, \quad (\text{A18})$$

$$n'^\mu = -\frac{1}{Q^2} q^\mu + \frac{1}{Q^2 + t} q'^\mu. \quad (\text{A19})$$

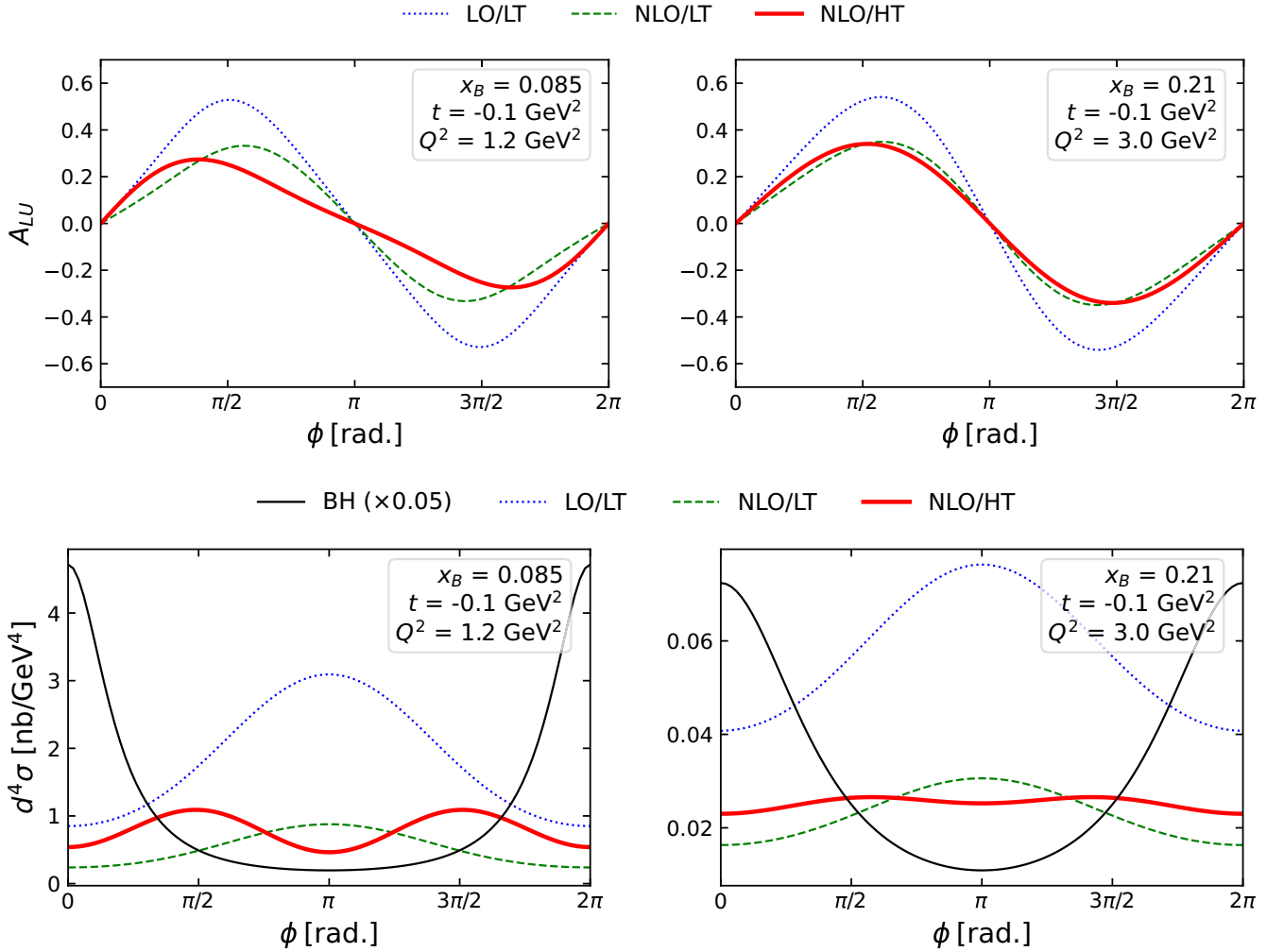


FIG. 13. Predictions for JLab12: differential cross sections (upper row) and A_{LU} asymmetries (lower row) for $Q^2 = 1.2 \text{ GeV}^2$ (left) and $Q^2 = 3 \text{ GeV}^2$ (right), at $y = 0.75$, $t = -0.1 \text{ GeV}^2$, and an electron beam energy of 10.06 GeV .

Alternatively,

$$q^\mu = \frac{1}{Q^2 + t} n^\mu - Q^2 n'^\mu, \quad (\text{A20})$$

$$q'^\mu = \frac{1}{Q^2} n^\mu. \quad (\text{A21})$$

An immediate consequence is the vanishing of the transverse component of the momentum transfer, this is $\Delta_\perp = (q - q')_\perp = 0$, while $p_\perp = p'_\perp = \bar{p}_\perp$ is the only momentum with transverse components. Taking into account that $\Delta \bar{p} = 0$, the definition of the skewness $\xi = -\Delta n / (2\bar{p}n)$ and the square $\bar{p}^2 = M^2 - t/4$, we find

$$\bar{p}^\mu = \frac{(p + p')^\mu}{2} = \frac{-t}{2\xi Q^2(t + Q^2)} n^\mu + \frac{Q^2}{2\xi} n'^\mu + \bar{p}_\perp^\mu, \quad (\text{A22})$$

where

$$\bar{p}_\perp^2 = M^2 - \frac{t}{4} \left(1 - \frac{1}{\xi^2}\right). \quad (\text{A23})$$

Taking into account that $-\bar{p}_\perp^2 = |\bar{p}_\perp|^2 \geq 0$, one can get the value of t for which its modulus is minimal, namely t_0 , for a fixed value of the skewness:

$$t_0 = -\frac{4M^2\xi^2}{1 - \xi^2} \quad (|t| \geq |t_0|). \quad (\text{A24})$$

This way, the vector $\bar{p}_\perp^\mu$ together with its dual $\tilde{\bar{p}}_\perp^\mu = \epsilon_\perp^{\mu\nu} \bar{p}_\nu$ span the transverse plane, allowing to define transverse-polarization vectors as [11]

$$\varepsilon'^\mu(\lambda) = -\frac{1}{\sqrt{2}|\bar{p}_\perp|} (\bar{p}_\perp^\mu - i\lambda \tilde{\bar{p}}_\perp^\mu), \quad \lambda = \pm. \quad (\text{A25})$$

The “prime” notation indicates that this vector corresponds to the outgoing photon with momentum q' such that $q' \varepsilon'(\pm) = 0$. For the incoming photon, the same polarization vectors apply, this is $\varepsilon'(\pm) = \varepsilon(\pm)$ with $q\varepsilon(\pm) = 0$. In terms of q and q' we can introduce the following longitudinal-polarization vector for the incoming photon [11]

$$\varepsilon^\mu(0) = \frac{1}{Q} \left[\frac{2Q^2}{Q^2} q'^\mu - q^\mu \right]. \quad (\text{A26})$$

Here, we denoted the hard scale of the process as

$$\mathbb{Q}^2 = Q^2 + t. \quad (\text{A27})$$

Finally, the skewness, which accounts for the fractional longitudinal kick to the target, takes the form

$$\xi = \frac{-\Delta n}{2\bar{p}n} = \frac{Q^2 + t}{2Q^2/x_A - Q^2 + t} \quad (\text{A28})$$

$$= \frac{x_A}{2 - x_A} + O\left(\frac{|t|}{Q^2}\right) \quad (\text{A29})$$

$$\stackrel{\text{Eq. (3)}}{\approx} \frac{x_B}{2A - x_B} + O\left(\frac{|t|}{Q^2}\right). \quad (\text{A30})$$

APPENDIX B: BH AMPLITUDE

1. First Bethe-Heitler amplitude, middle diagram in Fig. 1

For a configuration of beam and outgoing-photon polarizations given by s and λ , respectively, the amplitude of BH reads

$$i\mathcal{M}_{\text{BH}}^{s\lambda} = \frac{-ie^3}{(k - \Delta)^2 t} \varepsilon'_\rho(-\lambda) J_\alpha \underbrace{\bar{u}(k', s) \gamma^\rho (\not{k} - \not{\Delta}) \gamma^\alpha u(k, s)}_{\mathcal{T}^{\rho\alpha}}, \quad (\text{B1})$$

where e is the proton electric charge, J is the nonelementary hadron current

$$J^\mu = \langle p' | j^\mu(0) | p \rangle = 2\bar{p}^\mu F(t),$$

$$j^\mu(0) = \sum_f e_f \bar{q}_f(0) \gamma^\mu q_f(0) \quad (\text{B2})$$

with $F(t)$ the electromagnetic form factor (33), and $\varepsilon'(\lambda)$ is the polarization vector of the real outgoing-photon with momentum q' and polarization $\lambda = \pm$, vid. Eq. (A25). Note that we employed the relation $(\varepsilon'(\lambda))^* = \varepsilon'(-\lambda)$.

In order to use the KS techniques, we need to write down the quark propagator as a superposition of lightlike vectors. Considering a massless lepton beam

$$\not{k} - \not{\Delta} = \not{k}' + \not{q}' = \sum_{h=\pm} [u(k', h) \bar{u}(k', h) + u(q', h) \bar{u}(q', h)]. \quad (\text{B3})$$

Then,

$$\mathcal{T}^{\rho\alpha} = \sum_{h=\pm} \sum_{i=1}^2 \bar{u}(k', s) \gamma^\rho u(R_i, h) \bar{u}(R_i, h) \gamma^\alpha u(k, s), \quad (\text{B4})$$

where $R_i \in \{k', q'\}$ and contracting with ε' and J ,

$$\begin{aligned} \varepsilon'_\rho(-\lambda) J_\alpha \mathcal{T}^{\rho\alpha} &= \sum_{i=1}^2 \varepsilon'_\rho(-\lambda) J_\alpha \bar{u}(k', s) \gamma^\rho u(R_i, s) \bar{u}(R_i, s) \gamma^\alpha \\ &\quad \times u(k, s) \\ &= -\frac{F_1 \sqrt{2}}{|\bar{p}_\perp|} \sum_{i=1}^2 (p_{\perp, \rho} + i\lambda \tilde{p}_{\perp, \rho}) \bar{p}_\alpha \bar{u}(k', s) \gamma^\rho \\ &\quad \times u(R_i, s) \bar{u}(R_i, s) \gamma^\alpha u(k, s). \end{aligned} \quad (\text{B5})$$

As ε' and J are given by means of $p_\perp$, $\tilde{p}_\perp$, and $\bar{p}$, we need to write these momenta in terms of lightlike vectors in order to apply the KS techniques. To this purpose,

$$\begin{aligned} p^\mu &= p^+ n'^\mu + p^- n^\mu + p_\perp^\mu \\ &= \underbrace{(p^+ - \alpha) n'^\mu}_{r_1^\mu} + \underbrace{(\alpha n'^\mu + p^- n^\mu + p_\perp^\mu)}_{r_2^\mu}, \quad r_i^2 = 0. \end{aligned} \quad (\text{B6})$$

With $p^2 = M^2$ and $r_1^2 = r_2^2 = 0$, the parameter α takes the form

$$\alpha = p^+ - \frac{M^2}{2p^-(nn')}. \quad (\text{B7})$$

Consequently, the transverse component of p reads

$$p_\perp^\mu = r_2^\mu - p^- Q^2 q'^\mu - \alpha n'^\mu \Rightarrow \tilde{p}_\perp = \sum_{j=1}^3 B_j \mathbb{L}_j, \quad (\text{B8})$$

where $B_j \in \{1, -p^- Q^2, -\alpha\}$ and $L_j \in \{r_2, q', n'\}$.

In a similar fashion, for the momentum of the final-state target

$$\begin{aligned} p'^\mu &= p'^+ n'^\mu + p'^- n^\mu + p'_\perp^\mu \\ &= \underbrace{(p'^+ - \alpha') n'^\mu}_{r_1'^\mu} + \underbrace{(\alpha' n'^\mu + p'^- n^\mu + p'_\perp^\mu)}_{r_2'^\mu}, \\ r_i'^2 &= 0 \text{ and } \alpha' = p'^+ - \frac{M^2}{2p'^-(nn')}, \end{aligned} \quad (\text{B9})$$

so that

$$2\bar{p} = p + p' = r_1 + r_2 + r_1' + r_2'. \quad (\text{B10})$$

Let us recover from Ref. [51] the function $g(s, \ell, a, k)$ and the current $j_\mu(s, k', k)$ given for any lightlike vectors ℓ, a, k, k' and helicity $s = \pm$,

$$\begin{aligned} g(s, \ell, a, k) &= \bar{u}(\ell, s) \not{\ell} u(k, s) \\ &= \delta_{s+} s(\ell, a) t(a, k) + \delta_{s-} t(\ell, a) s(a, k), \end{aligned} \quad (\text{B11})$$

$$\begin{aligned} j_\mu(s, k', k) &= \bar{u}(k', s) \gamma_\mu u(k, s) \\ &= \frac{1}{N_{k'k}} \{ k'_\mu (k\kappa_0) + k_\mu (k'\kappa_0) - \kappa_{0,\mu} (k'k) \\ &\quad + is \epsilon_{\mu\alpha\beta\gamma} k'^\alpha k^\beta \kappa_0^\gamma \}, \end{aligned} \quad (\text{B12})$$

where $N_{k'k} = \sqrt{(k'\kappa_0)(k\kappa_0)}$ with the auxiliary vector $\kappa_0^\mu = (1, 1, 0, 0)$. The KS functions $s(k, k')$, $t(k, k')$ represent the bilinears

$$s(k, k') = \bar{u}(k, +) u(k', -) = -s(k', k), \quad (\text{B13})$$

$$t(k, k') = \bar{u}(k, -) u(k', +) = [s(k', k)]^*, \quad (\text{B14})$$

that with the aid of κ_0 , and as long as $k\kappa_0 \neq 0$ and $k'\kappa_0 \neq 0$, acquire the simple form

$$s(k, k') = (k^2 + ik^3) \sqrt{\frac{k'^0 - k'^1}{k^0 - k^1}} - (k \leftrightarrow k'). \quad (\text{B15})$$

Making use of these expressions in Eq. (B5), we can fully remove the Dirac spinors

$$\begin{aligned} \varepsilon'_\rho(-\lambda)J_\alpha \mathcal{T}^{\rho\alpha} = & -\frac{F(t)\sqrt{2}}{2|\bar{p}_\perp|} \sum_{i=1}^2 \sum_{\ell=1}^4 \\ & \times \left[\sum_{j=1}^3 B_j g(s, k', L_j, R_i) g(s, R_i, \tilde{r}_\ell, k) \right. \\ & \left. - i\lambda p_v \epsilon_\perp^{\nu\rho} j_\rho(s, k', R_i) g(s, R_i, \tilde{r}_\ell, k) \right], \end{aligned} \quad (\text{B16})$$

and write the amplitude *à la* KS:

$$\begin{aligned} i\mathcal{M}_{\text{BH}}^{s\lambda} = & \frac{ie^3 F(t)}{(k' + q')^2 t} \left(\frac{\sqrt{2}}{2|\bar{p}_\perp|} \right) \sum_{i=1}^2 \sum_{\ell=1}^4 g(s, R_i, \tilde{r}_\ell, k) \\ & \times \left[\sum_{j=1}^3 B_j g(s, k', L_j, R_i) - i\lambda p_v \epsilon_\perp^{\nu\rho} j_\rho(s, k', R_i) \right], \end{aligned} \quad (\text{B17})$$

where $B_j \in \{1, -p^- Q^2, -\alpha\}$, $L_j \in \{r_2, q', n'\}$, $R_i \in \{k', q'\}$ and $\tilde{r}_\ell \in \{r_1, r_2, r'_1, r'_2\}$.

2. Second Bethe-Heitler amplitude, right diagram in Fig. 1

As for the case before, for a configuration of beam and outgoing-photon polarizations given by s and λ , respectively,

the amplitude of the BHX subprocess reads

$$i\mathcal{M}_{\text{BHX}}^{s\lambda} = \frac{-ie^3}{(k - q')^2 t} J_\alpha \varepsilon'_\rho(-\lambda) \underbrace{\bar{u}(k', s) \gamma^\alpha (\not{k} - \not{q}') \gamma^\rho u(k, s)}_{\mathcal{T}_X^{\alpha\rho}} \quad (\text{B18})$$

Following the same steps as for BH, we find

$$\begin{aligned} J_\alpha \varepsilon'_\rho(-\lambda) \mathcal{T}_X^{\alpha\rho} = & -\frac{F(t)\sqrt{2}}{2|\bar{p}_\perp|} \sum_{i=1}^2 \sum_{\ell=1}^4 \sigma(\tilde{R}_i) g(s, k', \tilde{r}_\ell, \tilde{R}_i) \\ & \times \left[\sum_{j=1}^3 B_j g(s, \tilde{R}_i, L_j, k) \right. \\ & \left. - i\lambda p_v \epsilon_\perp^{\nu\rho} j_\rho(s, \tilde{R}_i, k) \right]. \end{aligned} \quad (\text{B19})$$

So that the amplitude *à la* KS takes the form

$$\begin{aligned} i\mathcal{M}_{\text{BHX}}^{s\lambda} = & \frac{ie^3 F(t)}{(k - q')^2 t} \left(\frac{\sqrt{2}}{2|\bar{p}_\perp|} \right) \sum_{i=1}^2 \sum_{\ell=1}^4 \sigma(\tilde{R}_i) g(s, k', \tilde{r}_\ell, \tilde{R}_i) \\ & \times \left[\sum_{j=1}^3 B_j g(s, \tilde{R}_i, L_j, k) - i\lambda p_v \epsilon_\perp^{\nu\rho} j_\rho(s, \tilde{R}_i, k) \right], \end{aligned} \quad (\text{B20})$$

where, again $B_j \in \{1, -p^- Q^2, -\alpha\}$, $L_j \in \{r_2, q', n'\}$, and $\tilde{r}_\ell \in \{r_1, r_2, r'_1, r'_2\}$, but $\tilde{R}_i \in \{k, q'\}$ with the signature function $\sigma(\tilde{R}_i) \in \{+1, -1\}$, respectively.

APPENDIX C: COEFFICIENT FUNCTIONS

In this appendix, we collect the expressions for the helicity amplitudes employed in the main text to parametrize the Compton tensor and compute cross sections and asymmetries. These formulas were published in earlier works [11,19,38]:

$$\begin{aligned} \mathcal{A}^{++}(\xi, t) = & \int_{-1}^1 \frac{dx}{\xi} \left\{ \left[C_0(x/\xi) + \frac{\alpha_s C_F}{4\pi} C_{\text{NLO}}^q(x/\xi) + \frac{t}{Q^2} (C_{0,\text{HT}}^{++}(x/\xi) - \xi \partial_\xi C_{1,\text{HT}}^{++}(x/\xi)) + \frac{2\xi^2 \bar{p}_\perp^2}{Q^2} \xi^2 \partial_\xi^2 C_{1,\text{HT}}^{++}(x/\xi) \right] \right. \\ & \left. \times H^\Sigma(x, \xi, t) + \frac{\alpha_s T_F}{4\pi} \frac{1}{\xi} C_{\text{NLO}}^g(x/\xi) H^g(x, \xi, t) \right\}, \end{aligned} \quad (\text{C1})$$

$$\mathcal{A}^{+-}(\xi, t) = \int_{-1}^1 \frac{dx}{\xi} \left\{ \frac{2\xi^2 \bar{p}_\perp^2}{Q^2} \xi^2 \partial_\xi^2 C_{0,\text{HT}}^{+-}(x/\xi) H^\Sigma(x, \xi, t) + \frac{\alpha_s T_F}{4\pi} \frac{2\xi^2 \bar{p}_\perp^2}{M^2} \frac{1}{\xi} C_{\text{NLO}}^{+-}(x/\xi) H_T^g(x, \xi, t) \right\}, \quad (\text{C2})$$

$$\mathcal{A}^{0+}(\xi, t) = -\frac{Q\sqrt{2}\xi|\bar{p}_\perp|}{Q^2} \int_{-1}^1 \frac{dx}{\xi} (1 - \xi \partial_\xi) C_{0,\text{HT}}^{0+}(x/\xi) H^\Sigma(x, \xi, t), \quad (\text{C3})$$

where the derivatives with respect to the skewness ($\partial_\xi, \partial_\xi^2$) act on everything to the right, including the GPDs. Note that the higher twists proportional to $\bar{p}_\perp$ contain Nachman-like mass corrections and vanish as $t \rightarrow t_0$: $2\xi^2 \bar{p}_\perp^2 = (1 - \xi^2)(t - t_0)/2$. The leading-twist coefficient functions (μ_F stands for the factorization scale) read

$$C_0(u) = \frac{-1}{u - 1 + i0} - (u \rightarrow -u), \quad (\text{C4})$$

$$C_{\text{NLO}}^{q/g}(u) = C_1^{q/g}(u) + \ln\left(\frac{Q^2}{\mu_F^2}\right) C_{\text{coll}}^{q/g}(u) \mp (u \rightarrow -u) \quad (- \text{ for } q, + \text{ for } g), \quad (\text{C5})$$

$$C_1^q(u) = \frac{1}{u + 1 - i0} \left[9 - 3 \frac{u+1}{u-1} \ln\left(\frac{u+1}{2} - i0\right) - \ln^2\left(\frac{u+1}{2} - i0\right) \right], \quad (\text{C6})$$

$$C_{\text{coll}}^q(u) = \frac{1}{u+1-i0} \left[-3 - 2 \ln \left(\frac{u+1}{2} - i0 \right) \right], \quad (\text{C7})$$

$$C_1^g(u) = \left(\sum_f e_f^2 \right) \frac{1}{(u+1-i0)(u-1+i0)} \left[2 \frac{u+3}{u-1} \ln \left(\frac{u+1}{2} - i0 \right) - \frac{u+1}{u-1} \ln^2 \left(\frac{u+1}{2} - i0 \right) \right], \quad (\text{C8})$$

$$C_{\text{coll}}^g(u) = \left(\sum_f e_f^2 \right) \frac{2}{(u+1-i0)(u-1+i0)} \left[-\frac{u+1}{u-1} \ln \left(\frac{u+1}{2} - i0 \right) \right], \quad (\text{C9})$$

$$C_{\text{NLO}}^{+-}(u) = \left(\sum_f e_f^2 \right) \frac{4}{(u-1+i0)(u+1-i0)}, \quad (\text{C10})$$

while the higher-twist kernels are given by

$$C_{0,\text{HT}}^{++}(u) = \tilde{\mathbb{P}}_{(\text{iii})}(u) - \frac{\mathcal{L}(u)}{2} - (u \rightarrow -u), \quad (\text{C11})$$

$$C_{1,\text{HT}}^{++}(u) = \frac{\tilde{\mathbb{P}}_{(\text{iii})}(u) - \mathcal{L}(u)}{2} - (u \rightarrow -u), \quad (\text{C12})$$

$$C_{0,\text{HT}}^{+-}(u) = u \tilde{\mathbb{P}}_{(\text{iii})}(u) - (u \rightarrow -u), \quad (\text{C13})$$

$$C_{0,\text{HT}}^{0+}(u) = \tilde{\mathbb{P}}_{(\text{iii})}(u) - (u \rightarrow -u), \quad (\text{C14})$$

$$\tilde{\mathbb{P}}_{(\text{iii})}(u) = \frac{-2}{u+1} \ln \left(\frac{u-1+i0}{-2+i0} \right), \quad (\text{C15})$$

$$\mathcal{L}(u) = \frac{4}{u-1} \left[\text{Li}_2 \left(\frac{u+1}{2-i0} \right) - \text{Li}_2(1) \right]. \quad (\text{C16})$$

-
- [1] E. R. Berger, F. Cano, M. Diehl, and B. Pire, Generalized parton distributions in the deuteron, *Phys. Rev. Lett.* **87**, 142302 (2001).
- [2] A. Kirchner and D. Mueller, Deeply virtual Compton scattering off nuclei, *Eur. Phys. J. C* **32**, 347 (2003).
- [3] V. Guzey and M. Strikman, Deeply virtual Compton scattering on spinless nuclear targets in the impulse approximation, *Phys. Rev. C* **68**, 015204 (2003).
- [4] F. Cano and B. Pire, Deep electroproduction of photons and mesons on the deuteron, *Eur. Phys. J. A* **19**, 423 (2004).
- [5] S. Scopetta, Generalized parton distributions of ^3He , *Phys. Rev. C* **70**, 015205 (2004).
- [6] A. V. Belitsky, D. Mueller, and A. Kirchner, Theory of deeply virtual Compton scattering on the nucleon, *Nucl. Phys. B* **629**, 323 (2002).
- [7] X. Ji, Gauge-invariant decomposition of nucleon spin, *Phys. Rev. Lett.* **78**, 610 (1997).
- [8] A. V. Belitsky and A. V. Radyushkin, Unraveling hadron structure with generalized parton distributions, *Phys. Rep.* **418**, 1 (2005).
- [9] M. Diehl, Generalized parton distributions, *Phys. Rep.* **388**, 41 (2003).
- [10] V. M. Braun, Y. Ji, and A. N. Manashov, Kinematic power corrections to deeply virtual Compton scattering to twist-six accuracy, *Phys. Rev. D* **111**, 076011 (2025).
- [11] V. Martinez-Fernandez, B. Pire, P. Sznajder, and J. Wagner, Timelike Compton scattering on a spin-0 target with kinematic twist-4 precision, *Phys. Rev. D* **111**, 074034 (2025).
- [12] M. Burkardt, Impact parameter space interpretation for generalized parton distributions, *Int. J. Mod. Phys. A* **18**, 173 (2003).
- [13] J. P. Ralston and B. Pire, Femtophotography of protons to nuclei with deeply virtual Compton scattering, *Phys. Rev. D* **66**, 111501 (2002).
- [14] M. Diehl, Generalized parton distributions in impact parameter space, *Eur. Phys. J. C* **25**, 223 (2002); Erratum: **31**, 277 (2003).
- [15] C. Lorcé, H. Moutarde, and A. P. Trawiński, Revisiting the mechanical properties of the nucleon, *Eur. Phys. J. C* **79**, 89 (2019).
- [16] V. Martínez-Fernández and C. Mezrag, Dispersion relations of deeply virtual Compton scattering: investigating twist-4 kinematic power corrections, *Phys. Rev. D* **113**, 094004 (2026).
- [17] V. Martínez-Fernández, D. Binosi, C. Mezrag, and Z.-Q. Yao, Constraining the energy-momentum tensor through the DVCS dispersion relation beyond leading power, *Phys. Rev. D* **113**, 094003 (2026).
- [18] B. Pire, L. Szymanowski, and J. Wagner, Next-to-leading order corrections to timelike, spacelike, and double deeply virtual Compton scattering, *Phys. Rev. D* **83**, 034009 (2011).
- [19] A. V. Belitsky and D. Mueller, Off-forward gluonometry, *Phys. Lett. B* **486**, 369 (2000).
- [20] V. Martínez-Fernández, B. Pire, P. Sznajder, and J. Wagner, companion paper, Quark and gluon tomography of the helium-4 nucleus, *Phys. Rev. Lett.* **137**, 111905 (2026).
- [21] M. Hattawy, N. A. Baltzell, R. Dupré, K. Hafidi, S. Stepanyan, S. Bültmann, R. De Vita, A. El Alaoui, L. El Fassi, H. Egiyan, *et al.* (CLAS Collaboration), First exclusive measurement of deeply virtual Compton scattering off ^4He : Toward the 3D tomography of nuclei, *Phys. Rev. Lett.* **119**, 202004 (2017).
- [22] R. Dupré, M. Hattawy, N. A. Baltzell, S. Bültmann, R. De Vita, A. El Alaoui, L. El Fassi, H. Egiyan, F. X. Girod, M. Guidal,

- et al.* (CLAS Collaboration), Measurement of deeply virtual Compton scattering off ^4He with the CEBAF large acceptance spectrometer at Jefferson lab, *Phys. Rev. C* **104**, 025203 (2021).
- [23] S. Fucini, S. Scopetta, and M. Viviani, Coherent deeply virtual Compton scattering off ^4He , *Phys. Rev. C* **98**, 015203 (2018).
- [24] R. Kleiss, The cross section for $e^+e^- \rightarrow e^+e^-e^+e^-$, *Nucl. Phys. B* **241**, 61 (1984).
- [25] R. Kleiss and W. Stirling, Spinor techniques for calculating, *Nucl. Phys. B* **262**, 235 (1985).
- [26] A. Bacchetta, U. D'Alesio, M. Diehl, and C. A. Miller, Single-spin asymmetries: The Trento conventions, *Phys. Rev. D* **70**, 117504 (2004).
- [27] V. M. Braun, A. N. Manashov, and B. Pirnay, Finite- t and target mass corrections to deeply virtual Compton scattering on a scalar target, *Phys. Rev. D* **86**, 014003 (2012).
- [28] X.-D. Ji and J. Osborne, One-loop corrections and all order factorization in deeply virtual Compton scattering, *Phys. Rev. D* **58**, 094018 (1998).
- [29] V. M. Braun, A. N. Manashov, D. Müller, and B. M. Pirnay, Deeply virtual Compton scattering to the twist-four accuracy: Impact of finite- t and target mass corrections, *Phys. Rev. D* **89**, 074022 (2014).
- [30] S. V. Goloskokov and P. Kroll, The longitudinal cross section of vector meson electroproduction, *Eur. Phys. J. C* **50**, 829 (2007).
- [31] A. V. Radyushkin, Double distributions and evolution equations, *Phys. Rev. D* **59**, 014030 (1998).
- [32] R. A. Khalek, R. Gauld, T. Giani, E. R. Nocera, T. R. Rabemananjara, and J. Rojo, nNNPDF3.0: Evidence for a modified partonic structure in heavy nuclei, *Eur. Phys. J. C* **82**, 507 (2022).
- [33] M. V. Polyakov and C. Weiss, Skewed and double distributions in the pion and the nucleon, *Phys. Rev. D* **60**, 114017 (1999).
- [34] V. Bertone, H. Dutrieux, C. Mezrag, J. M. Morgado, and H. Moutarde, Revisiting evolution equations for generalised parton distributions, *Eur. Phys. J. C* **82**, 888 (2022).
- [35] V. Bertone, APFEL++: A new PDF evolution library in C++, [arXiv:1708.00911](https://arxiv.org/abs/1708.00911).
- [36] B. Pire, J. Soffer, and O. Teryaev, Positivity constraints for off-forward parton distributions, *Eur. Phys. J. C* **8**, 103 (1999).
- [37] M. Kirch, P. V. Pobylitsa, and K. Goeke, Inequalities for nucleon generalized parton distributions with helicity flip, *Phys. Rev. D* **72**, 054019 (2005).
- [38] V. M. Braun, Y. Ji, and A. N. Manashov, Next-to-leading-power kinematic corrections to DVCS: A scalar target, *J. High Energy Phys.* **01** (2023) 078.
- [39] J. P. Repellin, P. Lehmann, J. Lefrançois, and D. B. Isabelle, Elastic electron scattering on helium 4, *Phys. Lett.* **16**, 169 (1965).
- [40] R. F. Frosch, J. S. McCarthy, R. E. Rand, and M. R. Yearian, Structure of the ^4He nucleus from elastic electron scattering, *Phys. Rev.* **160**, 874 (1967).
- [41] R. G. Arnold *et al.*, Elastic electron scattering from ^3He and ^4He at high momentum transfer, *Phys. Rev. Lett.* **40**, 1429 (1978).
- [42] C. R. Ottermann, G. Kobschall, K. Maurer, K. Rohrich, C. Schmitt, and V. H. Walther, Elastic electron scattering from ^3He and ^4He , *Nucl. Phys. A* **436**, 688 (1985).
- [43] A. Camsonne *et al.* (Jefferson Lab Hall A Collaboration), JLab measurement of the ^4He charge form factor at large momentum transfers, *Phys. Rev. Lett.* **112**, 132503 (2014).
- [44] W. Armstrong *et al.*, Partonic structure of light nuclei, [arXiv:1708.00888](https://arxiv.org/abs/1708.00888).
- [45] R. Abdul Khalek *et al.*, Science requirements and detector concepts for the electron-ion collider, *Nucl. Phys. A* **1026**, 122447 (2022).
- [46] D. P. Anderle *et al.*, Electron-ion collider in China, *Front. Phys.* **16**, 64701 (2021).
- [47] E. R. Berger, M. Diehl, and B. Pire, Timelike Compton scattering: Exclusive photoproduction of lepton pairs, *Eur. Phys. J. C* **23**, 675 (2002).
- [48] O. Grocholski, H. Moutarde, B. Pire, P. Sznajder, and J. Wagner, Data-driven study of timelike Compton scattering, *Eur. Phys. J. C* **80**, 171 (2020).
- [49] D. Mueller, B. Pire, L. Szymanowski, and J. Wagner, Timelike and spacelike hard exclusive reactions, *Phys. Rev. D* **86**, 031502 (2012).
- [50] H. Moutarde, B. Pire, F. Sabatie, L. Szymanowski, and J. Wagner, Timelike and spacelike deeply virtual Compton scattering at next-to-leading order, *Phys. Rev. D* **87**, 054029 (2013).
- [51] K. Deja, V. Martinez-Fernandez, B. Pire, P. Sznajder, and J. Wagner, Phenomenology of double deeply virtual Compton scattering in the era of new experiments, *Phys. Rev. D* **107**, 094035 (2023).
- [52] A. Pedrak, B. Pire, L. Szymanowski, and J. Wagner, Hard photoproduction of a diphoton with a large invariant mass, *Phys. Rev. D* **96**, 074008 (2017); **100**, 039901(E) (2019).
- [53] A. Pedrak, B. Pire, L. Szymanowski, and J. Wagner, Electroproduction of a large invariant mass photon pair, *Phys. Rev. D* **101**, 114027 (2020).
- [54] O. Grocholski, B. Pire, P. Sznajder, L. Szymanowski, and J. Wagner, Phenomenology of diphoton photoproduction at next-to-leading order, *Phys. Rev. D* **105**, 094025 (2022).
- [55] R. Boussarie, B. Pire, L. Szymanowski, and S. Wallon, Exclusive photoproduction of a $\gamma\rho$ pair with a large invariant mass, *J. High Energy Phys.* **02** (2017) 054; Erratum: **10** (2018) 029.
- [56] G. Duplanić, S. Nabeebaccus, K. Passek-Kumerički, B. Pire, L. Szymanowski, and S. Wallon, Probing chiral-even and chiral-odd leading twist quark generalized parton distributions through the exclusive photoproduction of a $\gamma\rho$ pair, *Phys. Rev. D* **107**, 094023 (2023).
- [57] A. V. Belitsky and D. Mueller, Exclusive electroproduction of lepton pairs as a probe of nucleon structure, *Phys. Rev. Lett.* **90**, 022001 (2003).