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Gluon tomography from deeply virtual Compton scattering at small x

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I. INTRODUCTION

The deeply virtual Compton scattering (DVCS) is one of the most important channels to study the partonic structure of a nucleon, in particular, to unveil the orbital angular momentum information for the quarks and gluons [1–4]. It has attracted tremendous interest from both the theory and experimental sides [5–10]. Experimentally, it is a simple high energy scattering process and is a major emphasis in the current and future lepton-nucleon collision facilities [9,10]. Among the observables in DVCS, it has been predicted that there exists a cos 2ϕ azimuthal correlation due to the so-called helicity-flip gluon generalized parton distributions (GPDs) [11–15]. In this paper, we investigate this physics in the small- x dipole formalism. We will show that the cos 2ϕ correlation can be accommodated in the dipole model through the so-called elliptic gluon distribution [16–19]. We also find the cos ϕ correlation, which has been missing in previous dipole-based studies.

In the small- x dipole factorization approach, the DVCS amplitude can be schematically calculated as [20–23]

$$A_{\text{DVCS}} \sim \int d^2 b_\perp e^{ib_\perp \cdot \Delta_\perp} \times \int dz d^2 r_\perp \Psi^*_\perp(z, r_\perp) \Psi^*_\perp(z, r_\perp) T(b_\perp, r_\perp),$$

where $\Psi$ and $\Psi^*$ are the wave functions for the incoming virtual photon and outgoing real photon, respectively. The physics behind this factorization can be understood as illustrated in Fig. 1, where the virtual photon fluctuates into a quark-antiquark pair to form a color dipole. The latter scatters on the nucleon target and merges into a real photon in the final state, whereas the nucleon recoils with momentum transfer $\Delta$. The wave functions depend on the momentum fraction of the photon carried by the quark $z$ and the dipole size $r_\perp$. For sufficiently hard scatterings, they are perturbatively calculable. In the DVCS amplitude, $T$ describes the elastic scattering of the dipole with the nucleon target. This is different from the inclusive deep inelastic scattering, which depends on the inelastic scattering amplitude. The elastic scattering amplitude can be written as

$$T = 1 - S,$$

where $S$ represents the dipole S-matrix (defined below). In the previous calculations of DVCS in the dipole formalism, the main focus was on the azimuthally symmetric cross section in which the photon helicity is conserved. In order to obtain the azimuthal cos 2ϕ correlation, we need to carry out the calculation on the helicity-flip amplitude. We perform our calculations in both coordinate space and momentum space and check their consistency.

An important aspect of our calculations is the comparison with the collinear factorization results. The key observation is the connection between the gluon GPDs at small $x$ and the dipole scattering amplitude. For the cos 2ϕ azimuthal correlation in the DVCS process, we show that the helicity-flip amplitude calculated from the elliptic gluon distribution reduces, in the collinear limit, to

FIG. 1. Deeply virtual Compton scattering $\gamma' p \rightarrow \gamma p$ in the small- x limit.
that from the helicity-flip gluon GPD in the collinear framework. Meanwhile, for the azimuthally symmetric cross section, the dipole formalism leads to divergence in the collinear limit. This can be interpreted as the \( \mathcal{O}(x_c) \) contribution to the quark GPD in the collinear framework, according to the relation between the quark GPD and the gluon GPD at small \( x \). These results establish a complete consistency between the dipole formalism and the collinear factorization framework.

The rest of the paper is organized as follows. In Sec. II, we establish the connection between the gluon GPDs at small \( x \) and the gluon Wigner distributions. In particular, the so-called elliptic gluon Wigner distribution will contribute to the helicity-flip gluon GPD. In Sec. III, we calculate the DVCS amplitude in the dipole framework in coordinate space and derive the cos \( 2\phi \) correlation. In Sec. IV, we perform the calculations in momentum space and demonstrate the consistency with the coordinate space derivations in Sec. III. The comparisons to the collinear factorization results will be made in Secs. III and IV and the Appendix. In Sec. V, we compute the contribution from the longitudinally polarized virtual photon and find the cos \( \phi \) correlation. Finally, we summarize our paper in Sec. VI.

II. DIPOLE S-MATRIX AND THE GLUON GPD

In this section, we introduce the basic ingredient to calculate the DVCS amplitude at small \( x \), namely, the dipole S-matrix. We shall clarify the relation between the gluon GPDs and the dipole S-matrix and show that the latter provides an efficient description of the DVCS amplitude which is free of collinear divergences.

In the dipole framework, the DVCS amplitude is represented by the diagram in Fig. 2 in coordinate space (left) and in momentum space (right). We work in a frame in which the virtual photon and the proton are collinear, with the proton moving fast in the positive \( z \)-direction. In coordinate space, we fix the transverse coordinates of the quark and antiquark to be \( x_{1\perp} = b_{\perp} + (1 - z)r_{\perp} \) and \( x_{2\perp} = b_{\perp} - zr_{\perp} \), respectively, with \( z \) defined as the longitudinal momentum fraction of the quark with respect to the incoming virtual photon. The “center of mass” of the \( \bar{q}q \) system coincides with the virtual photon coordinate \( x_{1\perp} + (1 - z)x_{2\perp} = b_{\perp} \). The size of the \( \bar{q}q \) system is \( r_{\perp} = x_{1\perp} - x_{2\perp} \). In this setup, the forward S-matrix for the \( \bar{q}q \) pair scattering off the target reads

\[
S_x(b_{\perp} + (1 - z)r_{\perp}, b_{\perp} - zr_{\perp}) = \left\langle \frac{1}{N_c} \text{Tr}[U(b_{\perp} + (1 - z)r_{\perp})U^\dagger(b_{\perp} - zr_{\perp})] \right\rangle_x,
\]

where \( x \) is the relevant momentum fraction of gluons in the target. In DVCS and in the small-\( x \) limit, it is related to the Bjorken variable \( x_B \) as \( x \approx \frac{x_B}{2} \), which is also the same as the skewness parameter \( \xi \) (defined below). \( U \) is the Wilson line

\[
U(x) = P \exp \left( -ig \int_{-\infty}^{\infty} dx^- A^+(x^-, x_\perp) \right),
\]

which represents the eikonal propagation of the quark. The brackets \( \langle \ldots \rangle \) denote the off-forward proton matrix element \( \langle p'|\lambda'|p\rangle \) with \( p' = p + \Delta \). In momentum space, we define

\[
F_x(q_{\perp}, \Delta_{\perp}, z) = \int \frac{d^2r_{\perp} d^2b_{\perp}}{(2\pi)^4} e^{ib_{\perp} \cdot \Delta_{\perp} + i\tilde{q}_{\perp} \cdot r_{\perp}} S_x(b_{\perp} + (1 - z)r_{\perp}, b_{\perp} - zr_{\perp})
\]

\[
= \int \frac{d^2r_{\perp} d^2b_{\perp}}{(2\pi)^4} e^{ib_{\perp} \cdot \Delta_{\perp} + i\tilde{q}_{\perp} \cdot r_{\perp}} e^{-i\delta_{\perp} \cdot r_{\perp}} S_x(b_{\perp} + \frac{r_{\perp}}{2}, b_{\perp} - \frac{r_{\perp}}{2})
\]

\[
= F_x(q_{\perp} \equiv \tilde{q}_{\perp} - \delta_{\perp}, \Delta_{\perp})
\]

where \( \delta_{\perp} \equiv \frac{1 - 2z}{2} \Delta_{\perp} \) and

\[
F_x(q_{\perp}, \Delta_{\perp}) = \int \frac{d^2r_{\perp} d^2b_{\perp}}{(2\pi)^4} e^{ib_{\perp} \cdot \Delta_{\perp} + i\tilde{q}_{\perp} \cdot q_{\perp}} \times S_x(b_{\perp} + \frac{r_{\perp}}{2}, b_{\perp} - \frac{r_{\perp}}{2}).
\]

In momentum space, we can also write

\[
F_x = \frac{1}{2\pi} \int d^2x_1 \cdot d^2x_2 \cdot e^{i(k_{1\perp} \cdot x_{1\perp} - k_{1\perp} \cdot x_{2\perp})} S_x(x_{1\perp}, x_{2\perp}) \quad \text{with} \quad k_{1\perp} \equiv \tilde{q}_{\perp} + z\Delta_{\perp} \quad \text{and} \quad k_{2\perp} \equiv \tilde{q}_{\perp} - (1 - z)\Delta_{\perp} \quad \text{conjugate to} \quad x_{1\perp} \quad \text{and} \quad x_{2\perp}, \quad \text{respectively.}
\]

The directions of the transverse momenta flow of exchanged gluons are labeled in Fig. 2. Following Ref. [16], we decompose \( F \) into the angular independent and “elliptic” parts.

![FIG. 2. Left diagram: DVCS amplitude in transverse coordinate space; Right diagram: DVCS amplitude in momentum space.](image-url)
Below, $F_e$ will be referred to as the elliptic gluon distribution. It is at most a few percent in magnitude compared to $F_0$ but has very different functional dependencies on $x$ and $q_{\perp}$ [17]. It can thus lead to distinct experimental signatures [16,18,19]. One of the main goals of this paper is to clarify the role of $F_e$ in DVCS.

Comments are in order regarding the phase factor $e^{-i\phi_{q_{\perp}}}$ in (5). In Ref. [23], the authors introduced a phase factor in the DVCS amplitude in the $b_{\perp}$-space,

$$\frac{d^2\sigma}{d^2b_{\perp}} = 2(1 - S(b_{\perp}, r_{\perp})) \frac{d^2\sigma}{d^2b_{\perp}} e^{-i(1-z)\Delta_{\perp} r_{\perp}},$$

and this prescription has been used in many subsequent works [24–28]. It is motivated by the explicit perturbative analysis in Ref. [21] that such a phase factor arises in nonforward amplitudes $\Delta_{\perp} \neq 0$. However, the result of Ref. [21] has been misinterpreted. To see the problem, note that (8) is not invariant under the combined transformation $z \rightarrow 1 - z$ and $r_{\perp} \rightarrow -r_{\perp}$. This transformation interchanges the quark and antiquark and has been emphasized in Ref. [21] as the exact symmetry of the dipole formalism. The phase factor discussed in Ref. [21] ensures that the effective transverse coordinates of the quark and antiquark are $b_{\perp} + (1 - z) r_{\perp}$ and $b_{\perp} - z r_{\perp}$, respectively, and this has been taken into account in (3). Equation (5) then shows that the correct phase factor should be $e^{-i\phi_{q_{\perp}} r_{\perp}} = e^{-i\frac{\Delta_{\perp}}{2} r_{\perp}}$, which is by itself invariant under the transformation $z \rightarrow 1 - z$ and $r_{\perp} \rightarrow -r_{\perp}$. As a nontrivial cross-check, in Sec. IV, we compute the DVCS amplitude in the momentum space and find the equivalent of this phase factor. We then show in Sec. V that this phase factor plays an important role in DVCS processes involving the longitudinally polarized virtual photon. We remark in passing that no phase factor is needed in the case of diffractive dijet production [16], though the process looks rather similar to DVCS.

A. Relation to GPD at small $x$

Let us point out the relation between $F_0$ and $F_e$ introduced above and the gluon GPDs which are defined as [12]

$$\frac{1}{P^+} \int \frac{d\zeta^-}{2\pi} e^{iP^+\zeta^-} \langle \langle p | F^{i+} F^{j+} | p \rangle \rangle = \delta^{ij} \frac{xH_g(x, \Delta_{\perp})}{2M^2} \left( \Delta_{\perp}^2 - \frac{\delta^{ij} \Delta_{\perp}^2}{2} \right) + \cdots,$$

where $M$ is the proton mass and $P = p^+ - P^+$ and we only kept the relevant terms for the nucleon-spin independent gluon GPDs at small $x$. Since we shall be interested later in matching between the dipole and GPD approaches in the limit $\Delta_{\perp} \ll Q$ ($Q$ is the photon virtuality), we assumed that $\Delta_{\perp}$ is small and approximated various spinor products $\bar{u}(p'') \bar{u}(ps)$ which appear in the usual parameterization of GPDs by their forward counterparts, e.g., $\bar{u}(p'') \gamma^+ u(PS) \approx 2P^+$. Our convention for the gluon GPDs is such that $H_g(x, \Delta_{\perp} \rightarrow 0) = G(x)$ (the unpolarized gluon PDF) in the forward limit. The helicity-flip$^1$ gluon GPD $E_{Tg}$ is also called the gluon transversity GPD, and the above normalization coincides with that of Ref. [12].$^2$ We suppress the dependence of GPDs on the skewness parameter $\xi = (p^+ - p'^+)/ (p^+ + p'^+)$, unless otherwise specified, it is understood that $H_g(x, \Delta_{\perp}) = H_g(x, \xi = x, \Delta_{\perp})$ and $E_{Tg}(x, \Delta_{\perp}) = E_{Tg}(x, \xi = x, \Delta_{\perp})$. This is because the imaginary part of the DVCS amplitude, which we assume to be dominant at small $x$, probes GPDs at $\xi = x$ to leading order. It is also known that, for the gluon GPDs at small $x$, this dependence has been found to be very mild, see for example the discussions in Ref. [6], which is consistent with the color-dipole formalism. The leading contribution of the S-matrix in the dipole formalism does not differentiate the dependence on $x$ and $\xi$.

At small $x$, the left-hand side of (9) can be approximately written as [16]

$$\frac{1}{P^+} \int \frac{d\zeta^-}{2\pi} e^{iP^+\zeta^-} \langle \langle p | F^{i+} F^{j+} | p \rangle \rangle$$

$$\approx \frac{2N_c}{\alpha_s} \int d^2q_{\perp} \left( \frac{\Delta_{\perp}^i}{2} \right) \left( \frac{\Delta_{\perp}^j}{2} + \frac{\Delta_{\perp}^i}{2} \right) F(q_{\perp}, \Delta_{\perp})$$

$$= \frac{2N_c}{\alpha_s} \int d^2q_{\perp} q_{\perp}^i q_{\perp}^j \left[ F_0(|q_{\perp}|, |\Delta_{\perp}|) + 2 \left( \frac{2(q_{\perp} \cdot \Delta_{\perp})^2}{q_{\perp}^2 \Delta_{\perp}^2} - 1 \right) F_e(|q_{\perp}|, |\Delta_{\perp}|) \right]$$

$$= \frac{2N_c}{\alpha_s} \left( \frac{\delta^{ij}}{2} \right) \int d^2q_{\perp} q_{\perp}^2 F_0$$

$$+ \frac{1}{\Delta_{\perp}^2} \left( \Delta_{\perp}^i \Delta_{\perp}^j - \frac{\delta^{ij} \Delta_{\perp}^2}{2} \right) \int d^2q_{\perp} q_{\perp}^2 F_e,$$

$^1$Note that “helicity flip” here refers to that of gluons (and the photon in DVCS) [12]. The nucleon helicity is not flipped. The GPDs associated with nucleon helicity flip are subleading at small $x$ and are not shown in (9). In other words, averaging over nucleon spin $s = \mp$ is implicit in (9). We do not consider polarized GPDs, either.

$^2$It differs from the normalization in Ref. [6] by a factor $-2x$. There exists another GPD called $H_T^g$ in Refs. [6,14], which was not noticed in Ref. [12]. Its contribution, if nonvanishing, can be trivially included in (9) and in all the formulas below by a simple shift $E_{Tg} \rightarrow E_{Tg} \equiv E_{Tg} + 2H_T^g$. 

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where we used the fact that \( \int d^2 q_\perp F(q_\perp, \Delta_\perp) = 0 \) for \( \Delta_\perp \neq 0 \). We thus obtain important relations between the gluon GPDs and the small-x dipole distributions as follows:

\[
xH_g(x, \Delta_\perp) = \frac{2N_c}{\alpha_s} \int d^2 q_\perp q_\perp^2 F_0, \tag{11}
\]

\[
xE_Tg(x, \Delta_\perp) = \frac{4N_c M^2}{\alpha_s \Delta_\perp} \int d^2 q_\perp q_\perp^2 F_c. \tag{12}
\]

These formulas will be used below to check the consistency with the collinear approach. The physical interpretation of the gluon GPDs and the above relations becomes manifest in the following computations of DVCS amplitudes.

### III. DVCS Amplitude and \( \cos 2\phi \) Azimuthal Angular Correlation

The differential cross section for DVCS can be written as

\[
d\sigma(ep \to e'p\gamma') = \frac{\alpha_{em}x_{Bj}}{4\pi Q^4} \frac{L_{\mu\nu} M_{\mu\nu}}{Q^4}, \tag{13}
\]

where \( L_{\mu\nu} \) is the lepton tensor and \( M_{\mu\nu} \) is the hadronic tensor. We use vectors \( l \) and \( l' \) for the initial and final state lepton momenta and \( p \) and \( p' = p + \Delta \) for the initial and final state proton momenta, respectively. The incoming virtual photon has momentum \( q = l - l' \) with virtuality \( q^2 = -Q^2 \) with vanishing transverse momentum. We use the standard variables \( x_{Bj} = Q^2/(2q \cdot p), \ y = q \cdot p / (l \cdot p) \), \( t = -\Delta_\perp^2 \) and \( W^2 = (q + p)^2 \approx Q^2/x_{Bj} \). In (13) and in the following, we only take into account the DVCS process and neglect the Bethe-Heitler (BH) contribution. In fixed-target experiments such as at COMPASS where \( Q \) is at most a few GeV or less at small \( x \), the cross section is dominated by the BH contribution. In collider experiments such as at HERA and the electron-ion collider (EIC), especially at large center-of-mass energies and small \( x \), there exist regions in kinematic variables where the cross section is dominated by the DVCS process [29,30]. However, we should keep in mind that, even in this latter situation, the azimuthally asymmetric part of the cross section may receive significant contributions from the interference with the BH amplitude. (As we commented above, the \( \cos 2\phi \) correlation from the elliptic Wigner distribution is a few percent effect.) In general, the BH amplitude can be highly suppressed when \( Q^2 \) is sufficiently large, since the target proton is unlikely to remain intact in the case of large momentum transfer. A detailed quantitative analysis of the impact of the BH amplitude is beyond the scope of this paper, and we leave it for future work.

The hadronic tensor can be decomposed as

\[
M_{\mu\nu} = M_{\mu\nu}^{TT} + M_{\mu\nu}^{TL} + M_{\mu\nu}^{L\perp L\perp}, \tag{14}
\]

where the subscripts \( T \) and \( L \) (transverse and longitudinal) denote the polarizations of the virtual photon in the amplitude and complex-conjugate amplitudes. (The outgoing real photon is always transversely polarized.) In this and the next sections, we will focus on \( M_{TT} \). The longitudinally polarized case will be treated in Sec. V. In the present frame, the lepton tensor can be decomposed into, for \( \mu, \nu \) transverse,

\[
L_{\mu\nu} = 2(l_\mu l'_\nu + l_\nu l'_\mu - g_{\mu\nu} l \cdot l')
\]

\[
= \frac{2Q^2}{y} \left[ \left( 1 - y + \frac{y^2}{2} \right) g_{\mu\nu} + (1 - y) \hat{h}_{\mu\nu} \right], \tag{15}
\]

where \( g_{\mu\nu} = -g_{\mu\nu} + (\hat{p}^\mu \hat{n}^\nu + \hat{p}^\nu \hat{n}^\mu)/\hat{p} \cdot \hat{n} \) and \( \hat{h}_{\mu\nu} = 2\epsilon_{\mu\nu} - g_{\mu\nu} \hat{n} \) and \( \hat{n} \) are two lightlike vectors: \( \hat{p}^2 = \hat{n}^2 = 0 \) and \( \hat{p} \cdot \hat{n} = 1 \). Here, \( l_\perp \) represents the transverse momentum of the lepton. It satisfies the relation \( l_\perp^2 = \frac{1}{2} Q^2 \). The hadronic tensor is calculated from the amplitude squared of \( \gamma' + p \to \gamma + p' \),

\[
M_{TT}^{\mu\nu} = W^2 g_{\perp\perp} A_T^{\mu\nu}(A_T^{\rho\sigma})^*, \tag{16}
\]

where \( \mu, \nu \) represent the (transverse) polarization indices for the incoming virtual photon, and \( \alpha, \beta \) represent the outgoing photon polarization indices. We have defined \( A^{\mu\nu} \) as the imaginary part of the amplitude. The real part is subleading at small \( x \) and can be retrieved through the dispersion relation, if necessary. It is convenient to decompose the tensor indices as

\[
A_T^{\mu\nu}(\Delta_\perp) = g_{\perp\perp}^{\mu\nu} A_0(\Delta_\perp) + h_{\perp\perp}^{\mu\nu} A_2(\Delta_\perp), \tag{17}
\]

where \( h_{\perp\perp}^{\mu\nu} = 2\Delta_\perp^{\mu\nu} - g_{\perp\perp}^{\mu\nu} \). The differential cross section then takes the form

\[
d\sigma_{TT} = \frac{\alpha_{em}}{\pi x_{Bj} Q^2} \left\{ \left( 1 - y + \frac{y^2}{2} \right) (A_0^2 + A_2^2) \right. \right.

\left. + (1 - y) 2A_0 A_2 \cos(2\phi_{\Delta\perp}) \right\}, \tag{18}
\]

where \( \phi_{\Delta\perp} \) is the azimuthal angle of the final state photon with respect to the lepton plane. The amplitudes \( A_{0,2} \) can be calculated from different projections of the tensor \( A_T^{\mu\nu} \). Alternatively, as noted in Refs. [11,12], they can also be obtained from the helicity-conserved and helicity-flip amplitudes as

\[
\frac{1}{2} \sum_\lambda A_T^{\mu\nu,\lambda}(\Delta_\perp) = A_0,
\]

\[
\frac{1}{2} \sum_\lambda A_T^{\mu\nu,\lambda'}(\Delta_\perp) = -A_2 \cos 2\phi_{\Delta\perp}, \tag{19}
\]

where \( \lambda \) and \( \lambda' \) represent the helicities of the incoming and outgoing photons. \( A_T^{\mu\nu,\lambda} \equiv \epsilon_\mu^{\rho\sigma} A_T^{\rho\sigma}(\epsilon_\nu^{\rho\sigma})^* \) can be conveniently expressed in coordinate space using the dipole S-matrix introduced in Sec. II.
\[ A_{\gamma}^{\perp}(\Delta_\perp) = 2 \int d^2 b_\perp e^{i b_\perp \cdot \Delta_\perp} N_c \sum_q \int d^2 r_\perp \]

\[ \times \int_0^1 \frac{dz}{4\pi} \Psi_f^\perp(z, r_\perp) \Psi_f^\perp(z, r_\perp) \]

\[ \times \left( 1 - S(b_\perp + (1 - z) r_\perp, b_\perp - z r_\perp) \right), \] (20)

where \( \Psi \) is the photon wave function. For the incoming virtual photon, it is given by

\[ \Psi_{\gamma q \perp}^\perp(z, r_\perp) = \frac{i e_q}{\pi} e_q K_1(e_q r_\perp) \]

\[ \times \left[ \frac{r_\perp e_q^{(i)}}{r_\perp} [\delta_{\alpha+\delta_{\beta+}z - \delta_{\alpha-} \delta_{\beta-} (1 - z)], \lambda = 1, \right. \]

\[ \left. \times \frac{r_\perp e_q^{(i)}}{r_\perp} [\delta_{\alpha-\delta_{\beta-}z - \delta_{\alpha+} \delta_{\beta+} (1 - z)], \lambda = 2, \right. \]

(21)

where \( \Psi \) is the photon wave function. For the incoming virtual photon, it is given by

\[ \Psi_{\gamma q \perp}^\perp(z, r_\perp) = \frac{i e_q}{\pi} e_q K_1(e_q r_\perp) \]

\[ \times \left[ \frac{r_\perp e_q^{(i)}}{r_\perp} [\delta_{\alpha+\delta_{\beta+}z - \delta_{\alpha-} \delta_{\beta-} (1 - z)], \lambda = 1, \right. \]

\[ \left. \times \frac{r_\perp e_q^{(i)}}{r_\perp} [\delta_{\alpha-\delta_{\beta-}z - \delta_{\alpha+} \delta_{\beta+} (1 - z)], \lambda = 2, \right. \]

(21)

\[ \mathcal{A}_0 = \frac{1}{2} \sum_q A_{\gamma}^{\perp}(\Delta_\perp) = -\sum_q e_q^2 N_c \int_0^1 d z [z^2 + (1 - z)^2] \int d^2 r_\perp \]

\[ \times \int d^2 q_\perp e^{-i q_\perp \cdot r_\perp} F(q_\perp, \Delta_\perp, z). \]

\[ = -\sum_q e_q^2 N_c \int_0^1 d z [z^2 + (1 - z)^2] \int d^2 r_\perp \]

\[ \times \int d^2 q_\perp e^{-i q_\perp \cdot r_\perp} [J_0(|q_\perp|, \Delta_\perp) + 2 \cos 2(\phi_{q_\perp} - \phi_{\Delta_\perp})] F_e(|q_\perp|, \Delta_\perp)]. \]

(24)

\[ \mathcal{A}_0(\Delta_\perp) \approx \sum_q e_q^2 N_c \int_0^1 d z [z^2 + (1 - z)^2] \]

\[ \times \int d^2 q_\perp \ln \left[ 1 + \frac{(q_\perp + \Delta_\perp)^2}{z(1 - z)Q^2} \right] F_0(|q_\perp|, \Delta_\perp)]. \]

(26)

A. Helicity-conserved amplitude

From (19) and (20), we immediately find

\[ \mathcal{A}_0(\Delta_\perp) \approx \sum_q e_q^2 N_c \int_0^1 d z [z^2 + (1 - z)^2] \]

\[ \times \int d^2 q_\perp \ln \left[ 1 + \frac{(q_\perp + \Delta_\perp)^2}{z(1 - z)Q^2} \right] F_0(|q_\perp|, \Delta_\perp)]. \]

(26)

Let us first consider the \( F_0 \) term in the last line. The \( r_\perp \)-integral looks divergent at first sight, since \( \int_0^\infty d r_\perp e_q K_1(e_q r_\perp) J_0(q_\perp r_\perp) \) is logarithmically divergent at \( r_\perp = 0 \). However, this divergence is not physical, and it can be removed easily. Using the fact that \( \int d^2 q_\perp F(q_\perp, \Delta_\perp) = 0 \), we obtain a convergent result:

\[ \int_0^\infty d r_\perp e_q K_1(e_q r_\perp) [J_0(q_\perp + \Delta_\perp) - 1] \]

\[ = -\frac{1}{2} \ln \left[ 1 + \frac{(q_\perp + \Delta_\perp)^2}{e_q} \right]. \]

(25)

The \( r_\perp \)-integral in the \( F_e \) term can also be done analytically in terms of the Appell function (see the formula 6.578-2 in Ref. [31]). We may, however, neglect this term as a higher order effect, \( J_2(\Delta_\perp) \sim \Delta_\perp^2 \), and obtain

\[ \Psi_{\gamma q \perp}^{\perp}\,(z, r_\perp) = \frac{e_q}{\pi} e_q \frac{1}{\Delta_\perp} K_0(e_q |r_\perp|) \delta_{\alpha \beta}, \]

(22)

where \( \alpha \) and \( \beta \) are the quark and antiquark helicities, \( e_q \) is the electric charge of the quark (in units of \( e \)), and \( e_q^2 = z(1 - z)Q^2 \). The quark mass has been neglected. For the outgoing real photon, we have

\[ \Psi_{\gamma q \perp}^{\perp}\,(z, r_\perp) = \frac{i}{\pi} \frac{e_q}{\Delta_\perp} K_1(e_q |r_\perp|) \delta_{\alpha \beta}, \]

(23)

\[ \Psi_{\gamma q \perp}^{\perp}\,(z, r_\perp) = \frac{i}{\pi} \frac{e_q}{\Delta_\perp} K_1(e_q |r_\perp|) \delta_{\alpha \beta}. \]

(23)

This is equivalent to the dimensional regularization with \( d = 2 - 2 e \) in the momentum space.
Expanding $J_0(q \cdot r_\perp) = 1 - \frac{1}{2} q^2 r^2_\perp + \cdots$ and keeping only the second term which is the leading twist contribution, we find

\[
-\frac{1}{2} \int_0^1 dz [z^2 + (1 - z)^2] \int \frac{d^2 r_\perp}{2\pi r_\perp} e_i K_1(e_i r_\perp) r^2_\perp \rightarrow -\frac{1}{Q^2} \left( \frac{Q^2 e^{-\gamma_E}}{\mu^2} \right)^{-\epsilon} \Gamma(2 - \epsilon) \Gamma(2 + \epsilon) \Gamma(-\epsilon) \frac{\Gamma(2 - 2\epsilon)}{\Gamma(2 - 2\epsilon)}
\]

\[
= -\frac{1}{Q^2} \left[ -\frac{1}{\epsilon} + \ln \frac{Q^2}{\mu^2} - 2 \right].
\]  

(27)
We further take the collinear limit \( Q^2 \gg q_\perp^2 \) and arrive at

\[
A_2(\Delta_\perp) = -\sum_q \frac{e_q^2 N_c}{Q^2} \int d^2 q_\perp q_\perp^2 F_\perp(q_\perp, \Delta_\perp)
= -\sum_q \frac{e_q^2 \alpha_s \Delta_\perp^2}{4M^2 Q^2} x E_{Tg}(x, \Delta_\perp),
\]

(35)

where Eq. (12) is used in the last step. This should be compared to the collinear factorization calculation by Ji-Hoodbhoy [12]. Their result reads, in the present normalization,

\[
A_2 = \frac{\sum_q e_q^2 N_c}{8\pi Q^2 M^2} \frac{\alpha_s}{\xi} \text{Im}
\times \left[ \int dx \left( \frac{1}{x - \xi + i\epsilon} + \frac{1}{x + \xi - i\epsilon} \right) E_{Tg}(x, \xi) \right].
\]

(36)

Noting that \( E_{Tg}(x, \xi) = -E_{Tg}(-x, \xi) \), we see that the above two are consistent with each other.

We thus see that the helicity-flip amplitude is proportional to the elliptic gluon distribution. Moreover, the collinear limit can be safely taken, as there is no divergence from the remaining \( z \)-integration. The resulting \( \cos 2\phi \) correlation should be measurable in the future experiments at the EIC. A similar observable in quasielastic scattering \( \gamma^*_p \rightarrow p'X \) has been proposed in Ref. [18]. Since these observables are associated with the correlation in the phase space Wigner distribution [16], such measurements will provide a unique perspective on the gluon tomography in nucleons at small \( x \).

**IV. MOMENTUM SPACE CALCULATION AND THE COLLINEAR LIMIT**

In this section, we repeat the calculation of the DVCS amplitude fully in momentum space and reproduce the results in the previous section. An advantage of the momentum space calculation is that it makes the connection to the collinear factorization approach more transparent. This is particularly important for the azimuthally symmetric part \( A_0 \) which, as we have already seen, contains divergence in the collinear limit. We show that this divergence can be interpreted as that of the quark GPD contribution to the DVCS amplitude. This is because the quark GPD can be calculated from the gluon GPD at small \( x \). When we substitute the quark GPD into the collinear formula for the DVCS amplitude, we are able to reproduce the result of the helicity-conserved DVCS amplitude in the previous section. This demonstrates the complete consistency of the dipole and collinear factorization approaches to DVCS.

In momentum space, the DVCS amplitude can be straightforwardly calculated from the right diagram in Fig. 2,

\[
A_{\mu\nu}^\perp = \sum_q \frac{e_q^2 N_c}{2\pi} \int dz d^2q_\perp d^2q_{\perp1} (-2) F_\perp(q_\perp, \Delta_\perp) \frac{2(z^2 + (1 - z)^2)q_{\perp1}^\mu k_{\perp1}^\nu - q_{\perp1}^\mu k_{\perp1}^\nu - q_{\perp1}^\mu k_{\perp1}^\nu + q_{\perp1} \cdot k_{\perp1} g_{\perp1}^\nu}{q_{\perp1}^2 (k_{\perp1}^2 + e_q^2)},
\]

(37)

where \( k_{\perp1} = q_{\perp1} + \frac{1}{z} \Delta_\perp - q_\perp \) with \( \bar{z} = 1 - z \) and \( F_\perp \) is defined as in (6). We have included a factor \(-2\) to adjust to the normalization \( A_T \sim -2S \) in (20). If we change variables as \( \bar{q}_\perp = q_\perp + \delta_\perp \), Eq. (37) takes the form

\[
A_{T}^\perp = \sum_q \frac{e_q^2 N_c}{2\pi} \int dz d^2\bar{q}_\perp d^2q_{\perp1} (-2) F_\perp(\bar{q}_\perp, \Delta_\perp, z) \times \frac{2(z^2 + (1 - z)^2)q_{\perp1}^\mu k_{\perp1}^\nu - q_{\perp1}^\mu k_{\perp1}^\nu - q_{\perp1}^\mu k_{\perp1}^\nu + q_{\perp1} \cdot k_{\perp1} g_{\perp1}^\nu}{q_{\perp1}^2 (k_{\perp1}^2 + e_q^2)},
\]

(38)

where \( k_{\perp1} = q_{\perp1} - \bar{q}_\perp \) and Eq. (5) is used. We thus see that this shift of loop momentum is related to the appearance of the phase factor \( e^{-i\theta_1 \cdot \tau_1} \) in coordinate space discussed in Sec. II. For the components (17), we obtain

\[
A_0 = \frac{q_{\perp1} \mu \nu A_{T}^\perp}{2}
= -\sum_q \frac{e_q^2 N_c}{\pi} \int d^2q_{\perp1} d^2q_{\perp1}
\times \frac{(z^2 + (1 - z)^2)q_{\perp1} \cdot k_{\perp1}}{q_{\perp1}^2 (k_{\perp1}^2 + e_q^2)} F_\perp(q_\perp, \Delta_\perp),
\]

(39)

and

\[
A_2 = \frac{h_{\perp1} \mu \nu A_{T}^\perp}{2}
= \sum_q \frac{2e_q^2 N_c}{\pi} \int d^2q_{\perp1} d^2q_{\perp1}
\times \frac{z(1 - z)[2q_{\perp1} \cdot \Delta_\perp k_{\perp1} \cdot \Delta_\perp - q_{\perp1} \cdot k_{\perp1}^2]}{q_{\perp1}^2 (k_{\perp1}^2 + e_q^2) \Delta_\perp^2}
\times F_\perp(q_{\perp1}, \Delta_\perp).
\]

(40)

It is interesting to notice that the \( A_2 \) depends on \( \cos(2\phi) \). For example, we can rewrite as
with respect to Feynman parametrization and performing the loop integral,

\[ \mathcal{A}_2 \text{ receives a contribution only from } \mathcal{A}_2, \text{ whereas } \mathcal{A}_0 \text{ receives a contribution from both terms. After applying the Feynman parametrization and performing the loop integral, } \mathcal{G}_2 \text{ can be written as} \]

\[ \mathcal{G}_2 = -\pi \int_0^1 \frac{da}{a q^2 + \epsilon_0^2}. \]  

Substituting the above result into \( \mathcal{A}_2 \), we obtain

\[ \mathcal{A}_2 = -2 \sum_q e_q^2 N_c \int dz \, d^2q \, q \Delta_\perp \frac{z(1-z)\alpha 2(\bar{q}_\perp \cdot \Delta_\perp)^2 - \bar{q}_\perp^2 \Delta_\perp^2}{\Delta_\perp^2} F_x(q_\perp, \Delta_\perp). \]  

By construction, Eq. (44) should be equivalent to (31), although it is difficult to see this analytically. We have checked this numerically for both the \( F_0 \) and \( F_e \) terms. In the DVCS limit \( \Delta_\perp \ll Q \), we can write

\[ \frac{2(\bar{q}_\perp \cdot \Delta_\perp)^2 - \bar{q}_\perp^2 \Delta_\perp^2}{\Delta_\perp^2} \approx q^2 \cos(2\phi_{q\Delta}), \]

and therefore

\[ \mathcal{A}_2 = -\sum_q e_q^2 N_c Q^2 \int d^2q \Delta_\perp F_e(q_\perp, \Delta_\perp) \]

\[ = -\frac{e_q^2 \alpha \Delta_\perp^3}{4Q^2M^2} E_{Tg}(x, \Delta_\perp), \]  

which is in agreement with (35).

We now return to \( \mathcal{A}_0 \) in (39) and take the DVCS limit

\[ \mathcal{A}_0 = -\sum_q \frac{e_q^2 N_c}{\pi} \int dz \, d^2q \, d^2k \]

\[ \times \frac{(z^2 + (1-z)^2)k_\perp \cdot (k_\perp + q_\perp)}{(k_\perp + q_\perp)^2(k_\perp^2 + e_0^2)} F_x(q_\perp, \Delta_\perp). \]  

In order to see the infrared behavior of the above integration more clearly, we examine the low transverse momentum region \( Q \gg k_\perp \sim q_\perp \) of the above integrand. We first notice that only the end points of the \( z \)-integral contribute. For example, if \( z \neq 1 \) or 0 so that \( e_q^2 \sim Q^2 \gg k_\perp^2 \), we immediately find that the above integral vanishes. Therefore, we have to separate out the dominant kinematic region of the above integration. To do that, we follow the trick of Ref. [33] and insert an identity: \( \int dx \delta(x - 1/(1 + \Lambda^2/\epsilon_0^2)) = 1 \) where \( \Lambda^2 = (1-z)k_\perp^2 + z(k_\perp + q_\perp)^2 \). In the region \( Q \gg k_\perp \sim q_\perp \), we can expand the \( \delta \)-function as

\[ \delta \left( x - \frac{1}{1 + \Delta \omega^2} \right) = \frac{1 - z}{x} \delta \left( 1 - z \right) \left( 1 - x \right) \frac{x^2 \Lambda^2}{z Q^2} \]

\[ = \frac{1 - z}{x} \left[ \delta(1 - z) \delta(1 - x) + \delta(1 - x) - \delta(1 - x) \ln \left( \frac{Q^2}{k_\perp^2} \right) \right]. \]

Let us show that only the first term contributes to \( \mathcal{A}_0 \) in the above expansion. For that purpose, we replace \( Q^2 \) and \( e_q^2 \) by applying the above \( \delta \)-function \( e_q^2 = \frac{1}{(1 - z) \Lambda^2} \), \( Q^2 = \frac{x}{(1 - z)(1 - x) \Lambda^2} \) and obtain

\[ \mathcal{A}_0 = -\sum_q \frac{e_q^2 N_c}{\pi Q^2} \int dx \, dz \, d^2q \, d^2k \]  

\[ \times \frac{k_\perp \cdot (k_\perp + q_\perp)}{(k_\perp + q_\perp)^2(1-x)k_\perp^2 + x \Lambda^2} \]

\[ \times \frac{\Lambda^2}{z(k_\perp + q_\perp)^2(1-x)} \frac{k_\perp \cdot (k_\perp + q_\perp)}{1-x} + \delta(1-x) \delta(1-z) \ln \left( \frac{Q^2}{k_\perp^2} \right) \]

\[ \times F_x(q_\perp, \Delta_\perp). \]  

First, we can easily check that the \( \delta(1-x) \delta(1-z) \) term vanishes. Second, the term proportional to \( \delta(1-x) \) also vanishes because the integrand can be simplified as

\[ \frac{k_\perp \cdot (k_\perp + q_\perp)}{z(1-z)(k_\perp + q_\perp)^2}, \]

and the azimuthal integration gives zero. Thus, the final result comes from the \( \delta(1-z) \) term
In the collinear limit, we can further simplify this as

$$\mathcal{A}_0 = \sum_q 4\pi e_q^2 N_c \int d^2 k_\perp \left[ \frac{1}{2\pi} \right] \int d^2 q_\perp F_x(q_\perp, \Delta_\perp)$$

In the above calculation, we picked up the leading contribution in the region of $z \sim 1$, which is similar to the current fragmentation contribution in semi-inclusive deep inelastic scattering (DIS) at small $x$ studied in Ref. [33]. For the $z \sim 0$ region, we can repeat the same procedure with $z \leftrightarrow 1 - z$. As a result, Eqs. (51) and (52) are doubled, and the divergent part of the latter agrees with (29). In the Appendix, we show that (52) can be interpreted as the quark GPD at small $x$.

**V. LONGITUDINALLY POLARIZED VIRTUAL PHOTON**

Finally, we study the contribution from the longitudinally $(L)$ polarized photon. The transition amplitude from the longitudinally polarized virtual photon to the transversely polarized real photon $\gamma^p \to \gamma p'$ is usually neglected in the dipole framework and actually vanishes unless one includes the phase factor $e^{-i\delta_{\perp r}}$ [21]. Here, we calculate its contribution to the DVCS cross section. The interference term between the transverse and longitudinal virtual photon amplitudes reads

$$\frac{L_{\mu_\mu} M_{\perp\perp}^{LL}}{W^4} = 2\text{Re} \sum_L L_{\mu_\mu} e^{L^{(i)*}} e^{L^{(i)}} \epsilon_{\nu} g_{\perp\mu\nu} A_T^{\perp\perp}(A_L^{\perp\perp})^{\nu} \epsilon_{\nu}^{L^{(i)}} \epsilon_{\nu}^{L^{(i)*}}.$$  (53)

Writing $e^{L^{(i)*}} A_L^{\perp\perp} = \frac{\epsilon_{\nu}^{L^{(i)}}}{(2\pi)^2} A_L$ and using

$$L_{\mu_\mu} e^{L^{(i)*}} e^{L^{(i)}} = \frac{2(2 - y)}{y} l_{\perp} \cdot e^{L^{(i)*}},$$  (54)

we obtain

$$\frac{L_{\mu_\mu} M_{\perp\perp}^{LL}}{W^4} = 4(2 - y) Q(A_0 + A_2) A_T \sum_L l_{\perp} \cdot e^{L^{(i)*}} \frac{\Delta_\perp \cdot e^{L^{(i)}}}{|\Delta_\perp|}$$

$$= \frac{4(2 - y)\sqrt{1 + y}}{y^2} Q^2(A_0 + A_2) A_L \cos \phi_{\Delta L}.$$  (55)

We immediately recognize the $\cos \phi_{\Delta L}$ angular distribution. $A_L$ can be evaluated as

$$A_L = - \sum_q \frac{2ie_q^2 N_c}{\pi|\Delta_\perp|} \int_0^1 dzz(1 - z)(1 - 2z)$$

$$\times \int d^2 r_\perp K_0(e_q r_\perp) \frac{r_\perp \cdot \Delta_\perp}{r_\perp}$$

$$\times \int d^2 q_\perp e^{-i(q_\perp + \delta_\perp \cdot r_\perp)} F_x(q_\perp, \Delta_\perp).$$  (56)

Naively, the $z$-integral vanishes because the integrand seems to be antisymmetric under $z \to 1 - z$. However, the phase $e^{-i\delta_\perp \cdot r_\perp}$ also depends on $z$, and this makes the integral finite. Performing angular integrations, we find

$$A_L = - \sum_q \frac{8\pi e_q^2 N_c}{\pi|\Delta_\perp|} \int_0^1 dzz(1 - z)(1 - 2z)$$

$$\times \int_0^\infty dr_\perp K_0(e_q r_\perp) \int_0^\infty dq_\perp q_\perp$$

$$\times (J_1(\delta_\perp r_\perp) J_0(q_\perp, \Delta_\perp) F_0(q_\perp, \Delta_\perp) - J_1(\delta_\perp r_\perp) J_2(q_\perp, \Delta_\perp) F_0(q_\perp, \Delta_\perp)),$$  (57)

where $\delta_\perp = \frac{1 - 2z}{\Delta_\perp}$ in the argument of the Bessel functions. Let us ignore the $J_2(\delta_\perp r_\perp)$ term and expand as $J_1(\delta_\perp r_\perp) \approx \frac{1}{\delta_\perp} \delta_\perp r_\perp$. We then get a nonzero result:

$$A_L \approx - \sum_q \frac{e_q^2 N_c |\Delta_\perp|}{Q} \int_0^1 dzz(1 - z)(1 - 2z) \int d^2 q_\perp$$

$$\times \frac{F_0(q_\perp, \Delta_\perp)}{e_q^2 + q_\perp^2} + F_x(q_\perp, \Delta_\perp)$$

$$\times \left( \frac{1}{e_q^2 + q_\perp^2} - \frac{1}{Q} \ln \left( 1 + \frac{q_\perp^2}{e_q^2} \right) \right).$$  (58)

If we do the collinear expansion, the $F_0$ term gives $xH_g(\Delta_\perp)$ via (11), but again the $z$-integral diverges at $z = 0, 1$. Similarly, the $F_x$ term gives $xE_Tg(\Delta_\perp)$ with a divergent coefficient. Regularizing this divergence as in (27), we find

$$A_L = - \sum_q \frac{e_q^2 N_c |\Delta_\perp|}{Q^3} \left( xH_g(\Delta_\perp) + \frac{\Delta_\perp^2}{4M^2} xE_Tg(\Delta_\perp) \right) \frac{1}{e} + \cdots.$$  (59)
The first term in (59) again comes from the quark GPD of which the contribution to the $\cos \phi$ part of the cross section is manifest in the collinear calculation (see the function called $F_{\text{eff}}$ in Ref. [15]). We are, however, unsure of the origin of the second term. Presumably, this arises from the twist-3 part of $F_{\text{eff}}$, but we have not been able to show this explicitly. In any case, this divergence is an artifact of the collinear expansion. At the level of (58), $A_L$ is finite and can be used in practical calculations.

For completeness, we also note the result for the longitudinal amplitude squared:

$$
\frac{L_{\mu \nu} \mathcal{M}_{L,\mu}^{\nu, L}}{W^2} = L_{\mu \nu} c_L^\mu c_L^\nu g_{\perp, \alpha \beta} A_{L, \alpha}^\mu (A_{L, \beta}^\rho)^* c_L^\rho c_L^\rho
$$

$$
= \frac{4(1 - y)}{y^2} Q^2 A_L^2.
$$

(60)

Adding all the components, we arrive at the complete DVCS cross section in the dipole framework,

$$
d\sigma(ep \rightarrow e'\gamma pp')
\frac{dx_B dQ^2 d^2 \Delta}{d^3 q^2}
= \frac{\alpha^3}{\pi x_B Q^2} \left\{ (1 - y + \frac{y^2}{2}) (A_0^2 + A_2^2)
+ 2(1 - y) A_0 A_2 \cos(2\phi_{\Delta z})
+ (2 - y) \sqrt{1 - y} (A_0 + A_2) A_L \cos \phi_{\Delta t} + (1 - y) A_L^2 \right\}.
$$

(61)

We remind the reader that, for practical applications, the Bether-Heitler contributions have to be included; see our comment at the beginning of Sec. III.

VI. CONCLUSION

In summary, we have studied the DVCS amplitudes at small $x$ in the dipole formalism. The final formula for the cross section (61) involves the $\cos \phi$ and $\cos 2\phi$ azimuthal angular correlations. While such correlations are known in the standard collinear approach to DVCS [6,7], it is nontrivial to retrieve them in the dipole framework. In order to obtain the $\cos \phi$ term, we have to include the (correct) phase factor $e^{i\frac{Q^2}{2\Delta^2} x \cdot \Delta}$ in the amplitude. As for the $\cos 2\phi$ term, it is essential to consider the elliptic gluon Wigner distribution [16–18] which represents the dominant angular dependence of the dipole S-matrix. In this regard, it is interesting to note that the elliptic gluon distribution has been recently proposed [19] as a possible underlying mechanism for the observed elliptic flow (cos $2\phi$ azimuthal correlation among the final state hadrons) in high energy $pp$ and $pA$ collisions (see also Ref. [18]). Experimental investigations of these novel phenomena will provide crucial information about the gluon tomography in the nucleon at small $x$.

We have also shown that, in the collinear limit, the dipole formalism reproduces the results obtained in the collinear factorization approach for both the angular symmetric and elliptic amplitudes. As $Q^2$ is lowered, the DVCS amplitudes are sensitive to the transverse momentum distribution in the target, and the dipole framework becomes more appropriate.

At last, we notice that the calculation on DVCS presented in this paper can be easily generalized to diffractive vector meson ($J/\psi$, $\rho$, and $\phi$) productions in DIS ($g^* + p \rightarrow V + p'$) (see, e.g., Refs. [10,23,35–37] and references therein), if we replace the transverse wave function of the final state real photon by the vector meson wave function. Similar conclusions can also be applied to this process.

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APPENDIX: COLLINEAR FACTORIZATION

RESULTS AND QUARK GPD AND PDF AT SMALL $x$

The DVCS amplitude is calculated in terms of the off-forward tensor $T_{\mu \nu}$,

$$
T_{\mu \nu} = i \int d^4 z e^{-i q \cdot z} \langle P' | j_\nu(z/2) j^\rho(-z/2) | P \rangle
\equiv g_{\perp, \mu \nu} T_0 + h_{\perp, \mu \nu} T_2.
$$

(A1)

The above two terms have been calculated in the literature. In the small-$x$ limit, they take the following forms [1,12],

$$
T_0 = - \sum_q e_q^2 \int d\alpha(x) H_q(x, \xi, \Delta^2_1),
$$

(A2)

$$
T_2 = \sum_q e_q^2 \frac{\alpha_q}{4\pi} \frac{\Delta^2}{4M^2} \int d\alpha(x) E_{\gamma_0}(x, \xi, \Delta^2_1),
$$

(A3)

where $H_q$ and $E_{\gamma_0}$ are the quark GPD and helicity-flip gluon GPD and $\alpha(x)$ is defined as

$$
\alpha(x) = \frac{1}{x - \xi + ic} + \frac{1}{x + \xi - ic}.
$$

(A4)

The other contribution in $T_2$ is suppressed at small $x$ and has been neglected in the above. We are particularly
interested in the imaginary part of the scattering amplitudes
\[
\text{Im}T_0 = \frac{\pi}{\xi} \sum_q e_q^2 [\xi H_q(\xi, \xi, \Delta_+^2) + \xi H_q(\xi, \xi, \Delta_-^2)]. \tag{A5}
\]
\[
\text{Im}T_2 = -\frac{\pi \alpha_s}{\xi} \frac{\Delta_+^2}{2\pi^2 M^2} \sum_q e_q^2 \xi ET_g(\xi, \xi, \Delta_+^2), \tag{A6}
\]
where we have taken into account the antiquark contribution, \(H_q(-x, \xi, \Delta_+^2) = -H_q(x, \xi, \Delta_+^2)\).

At small \(x\), the quark distribution comes from the gluon splitting. The forward quark distribution can be calculated as
\[
xq(x) = \frac{\alpha_s}{2\pi^2} \int_x^1 dx' (\zeta^2 + (1 - \zeta^2))x' G(x') \int \frac{dk_1^2}{k_1^2}, \tag{A7}
\]
where \(\zeta = x/x'\) and \(G(x')\) is the integrated forward gluon distribution. By applying the small-\(x\) approximation, the above can be simplified as
\[
xq(x) \approx xG(x) \frac{\alpha_s}{2\pi^2} \frac{1}{2} \frac{2}{3} \int \frac{dk_1^2}{k_1^2}, \tag{A8}
\]
where we assumed that \(x' G(x')\) is approximately constant at small \(x'\). For the quark GPD, the evolution equation depends on the skewness parameter \(\zeta\), which reads, for \(x > \xi\),
\[
xH_q(x, \xi, \Delta_+^2) = \frac{\alpha_s}{2\pi^2} \int_x^1 dx' \frac{\xi^2 (1 - \zeta^2) - \zeta^2 \xi^2}{(1 - \zeta^2 \xi^2)^2} \times x' H_q(x', \xi, \Delta_+^2) \int \frac{dk_1^2}{k_1^2}, \tag{A9}
\]
where \(H_q(x', \xi, \Delta_+^2)\) is the gluon GPD. The limit \(x \to \xi\) requires some care because of the singularity. If one naively sets \(\xi = x\) in the integrand and assumes that \(x' H_q(x', \xi)\) is a constant, the \(\zeta\)-integral gives \(\int_{\xi}^{\xi^2} \frac{dx'}{(1 + \zeta x')} \approx \frac{1}{2}\). However, this is incorrect. One has to first evaluate the \(\zeta\)-integral exactly and then take the limit \(x \to \xi\). This gives
\[
\lim_{x \to \xi} \int_{\xi}^{\xi^2} \frac{dx'}{(1 - \zeta^2 \xi^2)^2} = \frac{1}{1 + \xi} \approx 1. \tag{A10}
\]
We thus find
\[
\xi H_q(\xi, \xi) \approx \xi H_q(\xi, \xi) \frac{\alpha_s}{2\pi^2} \frac{1}{2} \int \frac{dk_1^2}{k_1^2}. \tag{A11}
\]
It is interesting to notice that here the prefactor is 1, instead of \(\frac{2}{3}\) for the forward quark distribution in Eq. (A8). Substituting the above results, we obtain the collinear factorization result for the DVCS amplitudes at small \(x\),
\[
\text{Im}T_0 = \frac{\alpha_s}{2\pi^2} \sum_q e_q^2 \xi H_g(\xi, \xi, \Delta_+^2) \int \frac{dk_1^2}{k_1^2}, \tag{A12}
\]
\[
\text{Im}T_2 = -\frac{\alpha_s}{2\pi^2 M^2} \sum_q e_q^2 \xi ET_g(\xi, \xi, \Delta_+^2), \tag{A13}
\]
where we have combined the quark and antiquark contributions together. To compare to our results in this paper, we note that the normalizations for the hadronic tensor are different,
\[
\text{Im}T^\mu = W^2 A^\mu = Q^2 \frac{Q^2}{x_B} A^\mu \approx Q^2 \frac{Q^2}{2\xi} A^\mu. \tag{A14}
\]
We thus find that (A12) agrees with (29) or (52) [the latter has to be multiplied by 2 as noted above (52)], and (A13) agrees with (35).


