Deeply virtual Compton scattering beyond next-to-leading order: the flavor singlet case

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Abstract
We study radiative corrections to deeply virtual Compton scattering in the kinematics of HERA collider experiments to next-to-leading and next-to-next-to-leading order. In the latter case the radiative corrections are evaluated in a special scheme that allows us to employ the predictive power of conformal symmetry. As observed before, the size of next-to-leading order corrections strongly depends on the gluonic input, as gluons start to contribute at this order. Beyond next-to-leading order we find, in contrast, that the corrections for an input scale of few GeV\(^2\) are small enough to justify the uses of perturbation theory. For \(\xi \gtrsim 5 \cdot 10^{-3}\) the modification of the scale dependence is also small. However, with decreasing \(\xi\) it becomes moderate or even large, in particular for the phase.

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1 Introduction

Deeply virtual Compton scattering (DVCS) is considered as the theoretically cleanest process to investigate generalized parton distributions (GPDs) \[1, 2, 3\]. These distributions are a hybrid of parton densities, form-factors, and distribution amplitudes and might be represented in terms of light–cone wave functions \[4, 5, 6\]. They are rather intricate functions depending on the longitudinal momentum fractions in the s– and t–channels, the momentum transfer squared, and the resolution scale. On the other hand they are phenomenologically very attractive, since they allow to combine information from different experiments in an optimal manner and since they encode non–perturbative information that cannot be extracted from either inclusive or elastic measurements alone. Their second moments provide e.g. the total angular momentum of partons in the nucleon \[2\] and the gravitational form factor of the nucleon. Moreover, a specific partial Fourier-transformation of GPDs is phenomenologically very interesting, as it provides functions which have a probabilistic interpretation. In the impact parameter space they can be viewed as parton densities in dependence of the longitudinal momentum fraction and the transverse distance from the proton center \[7, 8, 9\]. The knowledge of transverse parton distribution does not only add substantially to our understanding of hadron structure but is also relevant for the prediction of cross sections in dependence of the impact parameter. For proton–proton scattering this has been especially emphasized with respect to LHC physics in Ref. \[10\].

On one hand GPDs are thus a new window to study non–perturbative QCD and it has been already impressively demonstrated that they are experimentally accessible via the DVCS process \[11, 12, 13, 14, 15, 16\]. On the other hand for several theoretical and experimental reasons the extraction of GPDs from measurements remains quite challenging because typically one is sensitive only to convolutions containing GPDs and one has to disentangle different contributions. The analysis simplifies substantially at large energies, where both photon \[14, 15, 16\] and vector–meson leptoproduction, e.g., Refs. \[17, 18, 19\], have been measured by H1 and ZEUS. In HERA kinematics the photon–proton interaction starts to be flavor blind and so one mainly accesses flavor singlet GPDs. Moreover, spin flip effects are suppressed, too, and only one set of GPDs is relevant, namely, the proton helicity conserved and parton helicity averaged GPDs $H(x, \xi, t, Q^2)$. In this paper we concentrate on these singlet NNLO corrections, as the non-singlet case has already been studied in \[20\]. As usual, such analysis is only possible after adopting some parametrization for the dependence of GPDs on the s– and t–channel momentum fraction. We believe that realistic models can be most easily constructed by means of the partial wave decomposition of GPDs and amplitudes \[21, 22, 23\], where the dominant contributions arise from the leading Regge trajectories \[21, 24, 25\]. Support for this conjecture arises also from lattice calculations \[26, 27, 28, 29, 30\].

In this letter we study radiative corrections to DVCS at and beyond next–to–leading (NLO)
order. This investigation is partially motivated by the fact that within a certain class of GPD models perturbative corrections at this order were reported to be rather large \cite{31,32} and so one should worry about the justification of the perturbative QCD approach. Recently, the radiative corrections in the flavor non–singlet case have been studied and it has been concluded that the relative radiative corrections are moderate at NLO and become smaller at next–to–next–to–leading order (NNLO). Hence, these findings support the perturbative formalism in this sector. Considering the singlet case at hand, we recall that the leading order (LO) contribution is given by the quark handbag diagram and that at next–to–leading order the gluon distribution appears as a new entry. It is known from deeply inelastic scattering (DIS) that the gluonic contribution in the small \( x \)–region is much larger than that of the sea quarks. Hence, the size of the NLO corrections depends in particular on the gluonic GPD and the appearance of large NLO corrections does not necessarily mean that perturbation theory fails.

To clarify the situation, we employ here conformal symmetry to obtain the next–to–next–to–leading order corrections of the DVCS amplitude. In Sect. 2 we present, after a short introduction to the conformal approach, the analytic result for the DVCS amplitude in NNLO. In Sect. 3 relying on the pomeron pole as the dominant contribution at small momentum fraction, we numerically evaluate radiative corrections up to NNLO for the kinematics of HERA collider experiments. This analysis includes a comparison of the standard predictions in NLO with those of the conformal approach. We finally give our conclusions in Sect. 4.

2 The DVCS amplitude in NLO and NNLO

The DVCS amplitude is defined in terms of the hadronic tensor

\[
T_{\mu\nu}(q, P_1, P_2) = \frac{i}{e^2} \int d^4x e^{ix \cdot q} \langle P_2, S_2 | T j_\mu(x/2) j_\nu(-x/2) | P_1, S_1 \rangle,
\]

where \( q = (q_1 + q_2)/2 \) (\( \mu \) and \( q_2 \) refers to the outgoing real photon). To leading power in \( Q^2 = -q_1^2 \) (leading twist) and LO in the QCD coupling constant the hadronic tensor (1) is evaluated from the hand–bag diagram. In terms of the kinematical variables \( P = P_1 + P_2 \) and \( \Delta = P_2 - P_1 \), the result can be written as

\[
T_{\mu\nu}(q, P, \Delta) = -\left( \frac{\tilde{g}_{\mu\nu} - \frac{P_{\mu} q_\nu}{P \cdot q} - \frac{P_{\nu} q_\mu}{P \cdot q}}{P \cdot q} \right) \frac{q_\sigma V^\sigma}{P \cdot q} - \frac{i}{P \cdot q} \epsilon_{\mu\nu\rho\sigma} (P \cdot q)^2 + \cdots
\]

where the tilde–symbol denotes contraction \( \tilde{X}_{\mu\nu} \equiv \mathcal{P}_{\mu\rho} X^{\rho\sigma} \mathcal{P}_{\sigma\nu} \) with projectors

\[
\mathcal{P}^{\alpha\beta} = g^{\alpha\beta} - \frac{q_1^2 q_2^2}{q_1 \cdot q_2},
\]
to ensure current conservation [33]. The ellipsis indicate terms that are finally power suppressed in the DVCS amplitude or are determined by the gluon transversity GPD, which is suppressed by \( \alpha_s/\pi \) and is not considered here. Note that to leading twist accuracy the parenthesis in (2) can be replaced by \( \tilde{g}_{\mu\nu} \). In the parity even sector the vector

\[
V^\sigma = \mathcal{U}(P_2, S_2) \left( \mathcal{H} \gamma^\sigma + \mathcal{E} \frac{i\sigma^{\alpha\rho} \Delta_\rho}{2M} \right) U(P_1, S_1) + \cdots ,
\]

is decomposed into the target helicity conserving Compton form factor (CFF) \( \mathcal{H} \) and the helicity flip one \( \mathcal{E} \). Analogously, the axial-vector

\[
A^\sigma = \mathcal{U}(P_2, S_2) \left( \tilde{\mathcal{H}} \gamma^\sigma \gamma_5 + \tilde{\mathcal{E}} \frac{\Delta^\sigma \gamma_5}{2M} \right) U(P_1, S_1) + \cdots ,
\]

is parametrized in terms of \( \tilde{\mathcal{H}} \) and \( \tilde{\mathcal{E}} \), where again higher twist contributions are neglected. The normalization of the spinors is \( \mathcal{U}(p, S) \gamma^\sigma U(p, S) = 2p^\sigma \). We also introduce the scaling variables

\[
\xi = \frac{Q^2}{P \cdot q}, \quad \eta = -\frac{\Delta \cdot q}{P \cdot q},
\]

where \( Q^2 = -q^2 \). In DVCS kinematics and twist–two accuracy we have \( \xi = \eta \), while \( Q^2 = 2Q^2 \).

Before we proceed, let us decompose the CFFs, denoted by the set \( \mathcal{F} = \{ \mathcal{H}, \mathcal{E}, \tilde{\mathcal{H}}, \tilde{\mathcal{E}} \} \), in flavor non–singlet (NS) and singlet (S) ones:

\[
\mathcal{F} = Q^2_{NS}\mathcal{F} + Q^2_S\mathcal{F}, \quad S\mathcal{F} = \Sigma\mathcal{F} + G\mathcal{F},
\]

where the singlet piece contains the quark flavor singlet \( \Sigma\mathcal{F} \) and gluon \( G\mathcal{F} \) CFFs. The charge factors \( Q^2_i \) with \( i = \{ \text{NS}, S \} \) are given as linear combination of squared quark charges, e.g., the singlet one is given by the average of the squared charges for \( n_f \) active quarks:

\[
Q^2_S = \frac{1}{n_f} \sum_{i=u,d,...} Q^2_i.
\]

In the momentum fraction representation the Compton form factors are represented as convolution of the coefficient function with the corresponding GPD. In the singlet sector in which quark \( (\Sigma\mathcal{O}) \) and gluon \( (G\mathcal{O}) \) operators mix under renormalization, we might introduce the vector notation:

\[
S\mathcal{F}(\xi, \Delta^2, Q^2) = \int_{-1}^{1} \frac{dx}{x} C(x/\xi, Q^2/\mu^2, \alpha_s(\mu)|\xi) \mathcal{F}(x, \eta = \xi, \Delta^2, \mu^2).
\]

Here the column vector

\[
\mathcal{F} = \left( \begin{array}{c} \Sigma\mathcal{F} \\ G\mathcal{F} \end{array} \right), \quad F = \{ H, E, \tilde{H}, \tilde{E} \}
\]
contains the GPDs, and the row one, defined as $C = (\Sigma^C, (1/\xi)^G C)$, consists of the hard scattering part that to LO accuracy reads

$$\frac{1}{\xi} C(x/\xi, Q^2/\mu^2, \alpha_s(\mu)|\xi) = \left( \frac{1}{\xi - x - i\epsilon} \right) + O(\alpha_s). \quad (11)$$

We remark that the $\xi$ dependence in $\Sigma^C$ and $G^C$ enters only via the ratio $x/\xi$. Note also that the $u$–channel contribution in the quark entry (11) has been reabsorbed into the symmetrized quark singlet distribution

$$\Sigma^F(x, \eta, \Delta^2, \mu^2) = \sum_{q=u,d,\cdots} \left[ qF(x, \eta, \Delta^2, \mu^2) \mp qF(-x, \eta, \Delta^2, \mu^2) \right]. \quad (12)$$

Here the second term in the square brackets with $-(+)$–sign for $H, E$ ($\tilde{H}, \tilde{E}$)-type GPDs is for $x > \eta$ related to the $s$–channel exchange of an anti–quark. The gluon GPDs have definite symmetry property under the exchange of $x \rightarrow -x$: $G^H$ and $G^E$ are even, while $\tilde{G}H$ and $\tilde{G}E$ are odd.

The convolution formula (9) has already at LO the disadvantage that it contains a singularity at the cross–over point between the central region ($-\eta \leq x \leq \eta$) and the outer region ($\eta \leq x \leq 1$), i.e., for $x = \xi = \eta$. Its treatment is defined by the $i\epsilon$ prescription, coming from the Feynman propagator. The GPD is considered smooth at this point, but will generally not be holomorphic [34]. The fact that both regions are dual to each other, up to a so–called $D$-term contribution [35], makes the numerical treatment even more complicated. This motivated our development of a more suitable formalism in [21]. The factorization scale $\mu$ in the GPDs is ambiguous and at LO order this induces the main uncertainty. Beyond LO this factorization scale dependence will be cancelled in the considered order of perturbation theory. The NLO corrections to the coefficient functions [36] and to the evolution kernels [37, 38, 39] were predicted from conformal constraints, where the rotation to the standard $\overline{MS}$ scheme has been taken into account. Note that the conformal symmetry in $\overline{MS}$ scheme is broken and that the predicted results coincide with the diagrammatic evaluation [40, 41, 42]. To this order and in this scheme a numerical code has been made accessible that includes evolution, see, e.g., [32]. As already mentioned above, it was found that the perturbative corrections to NLO can be quite large.

At present it seems hardly possible to study perturbative corrections beyond NLO accuracy in the standard scheme, since the diagrammatical evaluation would require enormous effort. Fortunately, we can employ conformal symmetry to relate the perturbative corrections at NNLO to those for DIS [43, 44, 45], where the NNLO corrections in the vector case has been completed by the substantial effort of Vogt, Moch and Vermaseren [46]. From these calculations we get the normalization of the Wilson coefficients and anomalous dimensions. The conformal predictions arise from the application of the conformal operator product expansion and are valid as long as
the twist–two operators behave covariantly under conformal transformation \([47, 48, 49]\). This is certainly true at tree level and it also can be ensured for vanishing \(\beta\)-function in any order of perturbation theory within a special renormalization scheme \([48]\). To make contact with the conformal OPE, we expand the hard–scattering amplitude in terms of Gegenbauer polynomials with indices \(3/2\) and \(5/2\) for quarks and gluons, respectively, and introduce the conformal GPD moments, which formally leads to

\[
S\mathcal{F}(\xi, \Delta^2, Q^2) = 2 \sum_{j=0}^{\infty} \xi^{-j-1} C_j(Q^2/\mu^2, \alpha_s(\mu)) \mathcal{F}_j(\xi, \Delta^2, \mu^2).
\]  

(13)

The expansion coefficients \(C_j\) can be calculated by the projection:

\[
C_j(Q^2/\mu^2, \alpha_s(\mu)) = \frac{2^{j+1} \Gamma(j+5/2)}{\Gamma(3/2) \Gamma(j+4)} \times \frac{1}{2} \int_{-1}^{1} dx \left. C(x, Q^2/\mu^2, \alpha_s(\mu)|\xi = 1 \right) \left( \begin{array}{cc} (j+3)[1-x^2]C^{3/2}_j & 0 \\ 0 & 3[1-x^2]^2 C^{5/2}_{j-1} \end{array} \right) \right)(x).
\]  

(14)

Note that we have here rescaled the integration variable with respect to \(\xi\) and that the integral runs only over the rescaled central region. The conformal moments of the singlet GPDs are defined as

\[
\mathcal{F}_j(\eta, \Delta^2, \mu^2) = \frac{\Gamma(3/2) \Gamma(j+1)}{2 \Gamma(j+3/2)} \frac{1}{2} \int_{-1}^{1} dx \eta^{j-1} \left. C(x, Q^2/\mu^2, \alpha_s(\mu)|\xi = 1 \right) \left( \begin{array}{cc} \eta C^{3/2}_j & 0 \\ 0 & (3/2) C^{5/2}_{j-1} \end{array} \right) \right)(x) \mathcal{F}(x, \eta, \Delta^2, \mu^2).
\]  

(15)

Here \(j\) is an odd (even) non-negative integer for the (axial–vector) case.

In the forward kinematics \((\Delta \to 0)\), our conventions are such that the helicity conserved GPDs coincide with the flavor singlet quark distribution and with \(x\) times the gluon distribution. Hence, for the moments we have agreement with the common Mellin–moments of parton densities, e.g., for the helicity averaged GPD:

\[
\lim_{\eta \to 0} H(x, \eta) = \left( \frac{\sum \Sigma}{xG} \right) (x), \quad q_j \equiv \lim_{\eta \to 0} H_j(\eta) = \int_{0}^{1} dx \left( \sum \Sigma \right) (x).
\]  

(16)

Unfortunately, the series \([13]\) does not converge for DVCS kinematics, in particular not in the outer region, and one has to resum the OPE \([21, 20]\) or, equivalently, one can use a dispersion relation \([50, 51]\). The result for \(S\mathcal{H}\) in terms of a Mellin–Barnes integral reads

\[
S\mathcal{H}(\xi, \Delta^2, Q^2) = \frac{1}{2\pi} \int_{-\infty}^{c+\infty} dj \xi^{-j-1} \left[ i + \tan \left( \frac{\pi j}{2} \right) \right] C_j(Q^2/\mu^2, \alpha_s(\mu)) H_j(\xi, \Delta^2, \mu^2).
\]  

(17)

In the following we write the perturbative expansion as

\[
C_j(Q^2/\mu^2, Q^2/\mu_r^2, \alpha_s(\mu_r)) = \frac{2^{j+1} \Gamma(j+5/2)}{\Gamma(3/2) \Gamma(j+3)} \left[ C_j^{(0)} + \frac{\alpha_s(\mu_r)}{2\pi} C_j^{(1)}(Q^2/\mu^2) + \frac{\alpha_s^2(\mu_r)}{(2\pi)^2} C_j^{(2)}(Q^2/\mu^2, Q^2/\mu_r^2) + O(\alpha_s^3) \right].
\]  

(18)
where corresponding to our conventions, the LO Wilson coefficients are normalized as
\[ C_j^{(0)} = (1, 0), \] (19)
and here we choose to distinguish the renormalization \((\mu_r)\) and factorization \((\mu)\) scales.

Let us first give here the DVCS NLO corrections in the \(\overline{\text{MS}}\) scheme. We restrict ourselves to the analysis of the kinematically dominant contribution, i.e., \(S^H\), and so we provide here only the results for the vector case. The conformal moments \([36, 40, 41, 42]\) can be obtained from Refs. \(36\) \(40\) \(41\) \(42\). Using the representation of Ref. \(31\), the integrals which are needed are evaluated in a straightforward manner for integer conformal spin, see Appendix C of Ref. \(52\) for the quark entries. The analytic continuation to complex \(j\) leads to
\[ \Sigma C_j^{(1)}(Q/\mu^2) = C_F \left[ 2S_1^2(1+j) - \frac{9}{2} + \frac{5}{2} - 4S_1(j+1) + \frac{1}{(j+1)^2(j+2)^2} \right] + \frac{\Sigma \gamma_j^{(0)}}{2} \ln \frac{\mu^2}{Q^2}, \] (20)
\[ G C_j^{(1)}(Q/\mu^2) = -2n_f T_F \left[ \frac{4 + 3j + j^2}{(1+j)(2+j)(3+j)} \right] + \frac{\Sigma G_j^{(0)}}{2} \ln \frac{\mu^2}{Q^2}, \] (21)
where \(C_F = 4/3\) and \(T_F = 1/2\). The entries of the anomalous dimension matrix read at LO:
\[ \Sigma \gamma_j^{(0)} = -C_F \left( 3 + \frac{2}{(j+1)(j+2)} - 4S_1(j+1) \right), \] (22)
\[ \Sigma G_j^{(0)} = -4n_f T_F \frac{4 + 3j + j^2}{(j+1)(j+2)(j+3)}, \] (23)
\[ G \gamma_j^{(0)} = -2C_F \frac{4 + 3j + j^2}{j(j+1)(j+2)}, \] (24)
\[ G \gamma_j^{(0)} = -C_A \left( -\frac{4}{(j+1)(j+2)} + \frac{12}{j(j+3)} - 4S_1(j+1) \right) + \beta_0, \] (25)
where \(\beta_0 = 2n_f/3 - 11C_A/3\), \(C_A = 3\). In the \(\overline{\text{MS}}\) scheme also the complete anomalous dimension matrix is known to two-loop accuracy \([53, 54]\). However, the conformal moments will mix with each other and the solution of the evolution equation has so far not been given in terms of a Mellin–Barnes integral.

The advantage of the conformal symmetry is that it predicts the Wilson coefficients. However, the symmetry is only valid in a special conformal scheme \((\overline{\text{CS}})\). In such a scheme the structure
\footnote{The treatment of the terms proportional to \(\beta\), that break conformal symmetry, is of course ambiguous. We employ here the so-called \(\overline{\text{CS}}\) scheme in which the running of the coupling is implemented in the form of the conformal operator product expansion (COPE) that is valid for a hypothetical fixed point. In particular, conformal moments are multiplicatively renormalizable to NLO. For details see Refs. \(52\) \(20\).}
of the Wilson coefficients up to NNLO is

\[ C_j^{(1)}(Q^2/\mu^2) = c_j^{(1)} + \frac{s_j^{(1)}(Q^2/\mu^2)}{2} - c_j^{(0)} \gamma_j^{(0)}, \]

\[ C_j^{(2)}(Q^2/\mu^2, Q^2/\mu_0^2) = c_j^{(2)} + \frac{s_j^{(1)}(Q^2/\mu^2)}{2} \left[ c_j^{(0)} \gamma_j^{(1)} + c_j^{(1)} \gamma_j^{(0)} \right] + \frac{s_j^{(2)}(Q^2/\mu^2)}{8} c_j^{(0)} \left( \gamma_j^{(0)} \right)^2 \]

\[ + \frac{\beta_0}{2} \left[ C_j^{(1)}(Q^2/\mu^2) \ln \frac{Q^2}{\mu_\gamma^2} + \frac{1}{4} c_j^{(0)} \gamma_j^{(0)} \ln^2 \frac{Q^2}{\mu^2} \right], \]

where the so-called shift coefficients \( s_j^{(i)}(Q^2/\mu^2) \) can be expressed in terms of harmonic sums \( S_p(n) = \sum_{k=1}^{n} 1/k^p \) as

\[ s_j^{(1)}(Q^2/\mu^2) = S_1(j + 3/2) - S_1(j + 2) + 2 \ln(2) - \ln \frac{Q^2}{\mu^2}, \]

\[ s_j^{(2)}(Q^2/\mu^2) = \left( s_j^{(1)}(Q^2/\mu^2) \right)^2 - S_2(j + 3/2) + S_2(j + 2), \]

and \( c_j^{(i)} = (\xi_j^{(i)}, \gamma_j^{(i)}) \) are the Wilson coefficients known from DIS. We have at LO \( c_j^{(0)} = (1, 0) \), at NLO

\[ \xi c_j^{(1)} = C_F \left[ S_1^2(1 + j) + \frac{3}{2} S_1(j + 2) - \frac{9}{2} + \frac{5 - 2 S_1(j)}{2(j + 1)(j + 2)} - S_2(j + 1) \right], \]

\[ \gamma c_j^{(1)} = -2n_F T_F \frac{(4 + 3j + j^2) S_1(j) + 2 + 3j + j^2}{(1 + j)(2 + j)(3 + j)}, \]

and at NNLO they are given by the Mellin moments of the DIS partonic structure functions \[ \text{[43, 44]. To simplify their evaluation, we take for } c_j^{(2)} \text{ a fit, given in [45], rather than the exact expression.} \]

The evolution of the singlet (integer) conformal moments in this \( \CS \) scheme is governed by

\[ \mu \frac{d}{d\mu} F_j(\xi, \Delta^2, \mu^2) = - \left[ \frac{\alpha_s(\mu)}{2\pi} \gamma_j^{(0)} + \frac{\alpha_s^2(\mu)}{(2\pi)^2} \gamma_j^{(1)} + \frac{\alpha_s^3(\mu)}{(2\pi)^3} \gamma_j^{(2)} + \mathcal{O}(\alpha_s^4) \right] F_j(\xi, \Delta^2, \mu^2) \]

\[ - \frac{\beta_0}{2} \frac{\alpha_s^3(\mu)}{(2\pi)^3} \sum_{k=0}^{j-2} \left[ \Delta_{jk}^{\CS} + \mathcal{O}(\alpha_s) \right] F_k(\xi, \Delta^2, \mu^2), \]

where the mixing matrix \( \Delta_{jk}^{\CS} \) is not completely known. In the vector case the anomalous dimensions are known to NNLO \[ \text{[46]. In absence of the mixing term, the solution of the renormalization group equation } F_j(\xi, \Delta^2, \mu^2) = \mathcal{E}_j(\mu, \mu_0) F_j(\xi, \Delta^2, \mu_0^2) \text{ can be given using the path–ordered exponential evolution operator} \]

\[ \mathcal{E}_j(\mu, \mu_0) = \mathcal{P} \exp \left\{ - \int_{\mu_0}^{\mu} \frac{d\mu'}{\mu'} \gamma_j(\alpha_s(\mu')) \right\}. \]
In the numerical analysis we will only resum the leading logarithms and expand the non–leading ones

$$E_j(\mu, \mu_0) = \sum_{a,b=\pm} \left[ \delta_{ab} a^b P_j + \frac{\alpha_s(\mu)}{2\pi} a^b A_j^{(1)}(\mu, \mu_0) + \frac{\alpha_s^2(\mu)}{(2\pi)^2} a^b A_j^{(2)}(\mu, \mu_0) + O(\alpha_s^3) \right] \left[ \frac{\alpha_s(\mu)}{\alpha_s(\mu_0)} \right]^{b_j/\gamma_0}.$$  

(34)

Here the projectors on the \{+, −\} modes are

$$\pm P_j = \frac{\pm 1}{\lambda_j + \lambda_j} \left( \gamma_j^{(0)} - \epsilon \lambda_j 1 \right),$$  

(35)

where the eigenvalues of the LO anomalous dimension matrix are

$$\pm \lambda_j = \frac{1}{2} \left( \Sigma \gamma_j^{(0)} + GG_{\gamma_j^{(0)}} \right) \mp \frac{1}{2} \left( \Sigma \gamma_j^{(0)} - GG_{\gamma_j^{(0)}} \right) \sqrt{1 + \frac{4 \Sigma_{g_j^{(0)}} GG_{\gamma_j^{(0)}}}{(\Sigma \gamma_j^{(0)} - GG_{\gamma_j^{(0)}})^2}}.$$  

(36)

A straightforward calculation leads to the matrix valued coefficients

$$a^b A_j^{(1)} = a^b R_j(\mu, \mu_0|1) a^b P_j \left[ \frac{1}{2\beta_0} \gamma_j^{(0)} - \gamma_j^{(1)} \right] b^P_j$$  

(37)

$$a^b A_j^{(2)} = \sum_{c=\pm} 1 \beta_0 + \gamma_j - b\lambda_j a^b R_j(\mu, \mu_0|2) - a^c R_j(\mu, \mu_0|1) \left[ \frac{\alpha_s(\mu_0)}{\alpha_s(\mu)} \right]^{-\beta_0 + \gamma_j - b\lambda_j} a^b P_j \left[ \frac{1}{2\beta_0} \gamma_j^{(0)} - \gamma_j^{(1)} \right]$$

$$\times \left[ \frac{\beta_0}{2\beta_0} \gamma_j^{(0)} - \gamma_j^{(1)} \right] b^P_j - a^b R_j(\mu, \mu_0|2) a^c P_j \left[ \frac{\beta_0}{2\beta_0} \gamma_j^{(0)} - \gamma_j^{(1)} \right]$$

$$\times \left[ \frac{\beta_0}{2\beta_0} \gamma_j^{(0)} + \gamma_j^{(2)} \right] b^P_j,$$  

(38)

where the \mu dependence is accumulated in the following functions:

$$a^b R_j(\mu, \mu_0|n) = \frac{1}{n\beta_0 + a\lambda_j - b\lambda_j} \left[ 1 - \left( \frac{\alpha_s(\mu_0)}{\alpha_s(\mu)} \right)^{n\beta_0 + a\lambda_j - b\lambda_j} \right].$$  

(39)

The expansion coefficients of the \beta function are defined as

$$\frac{\beta}{g} = \frac{\alpha_s(\mu)}{4\pi} \beta_0 + \frac{\alpha_s^2(\mu)}{(4\pi)^2} \beta_1 + \frac{\alpha_s^3(\mu)}{(4\pi)^3} \beta_2 + O(\alpha_s^4),$$  

(40)

$$\beta_0 = \frac{2}{3} n_f - 11, \quad \beta_1 = \frac{38}{3} n_f - 102, \quad \beta_2 = -\frac{325}{54} n_f^2 + \frac{5033}{18} n_f - \frac{2857}{2}.$$  

The expansion of the evolution operator (34) will then be consistently combined with the Wilson–coefficients (18), see for instance Ref. 52.
3 Numerical evaluation of radiative corrections

Here we study the perturbative corrections in the small–ξ region. We adopt a simple ansatz for the conformal moments that is inspired by the dominance of the pomeron and by the assumption that $O(\xi^2)$ terms in the conformal polynomials are insignificant:

$$H_j(\xi, \Delta^2, Q^2) = \left( \frac{N_{\text{sea}}{F(\Delta^2) B(1 + j - \alpha_{\text{sea}}(\Delta^2), 8)/B(2 - \alpha_{\text{sea}}(0), 8)}}{N_{G} G F(\Delta^2) B(1 + j - \alpha_G(\Delta^2), 6)/B(2 - \alpha_G(0), 6)} \right) + \ldots, \quad (41)$$

where $B(x, n) = \frac{\Gamma(x) \Gamma(n)}{\Gamma(x + n)}$, and the ellipsis denotes the neglected $O(\xi^2)$ terms, as well as valence components whose contributions are also small for small $\xi$. Here the normalization is $\text{sea} F(\Delta^2 = 0) = G F(\Delta^2 = 0) = 1$ and for $\alpha_i(\Delta^2)$ we use the “effective” pomeron trajectory $\alpha(\Delta^2) = \alpha(0) + \alpha' \Delta^2$. We remind that in deeply inelastic scattering the structure function $F_2 \sim (1/x_{Bj})^{\lambda(Q^2)}$ grows with increasing $Q^2$. Here the exponent is related to the intercept of the Regge trajectory $\lambda = \alpha(0) - 1$. The values for $\alpha(0)$ will be specified below. Although also the slope $\alpha'$ is scale dependent \[24\], we choose here the standard value of the soft pomeron $\alpha' = 0.25$.

In the forward case the moments (41) arise from the parton densities

$$\Sigma = \frac{N_{\text{sea}}}{B(2 - \alpha_{\text{sea}}(0), 8)} x^{-\alpha_{\text{sea}}(0)} (1 - x)^7 + u_v(x) + d_v(x), \quad G = \frac{N_{G}}{B(2 - \alpha_G(0), 6)} x^{-\alpha_G(0)} (1 - x)^5, \quad (42)$$

for which we have adopted a generic realistic parametrization. Note that the valence component $u_v + d_v$ is not taken into account in Eq. (41), since it is a non–leading contribution for small $x$. The normalization factors are related by the momentum sum rule

$$\int_0^1 dx x [\Sigma(x) + G(x)] = 1 \Rightarrow N_G + N_{\text{sea}} + \int_0^1 dx x [u_v + d_v](x) = 1. \quad (43)$$

In the asymptotic limit $Q \rightarrow \infty$, the evolution equation tells us that $N_G = 4C_F/(4C_F + n_f)$ i.e. that more than 50% of the longitudinal proton momentum is carried by gluons. At experimentally accessible large scales the gluons already carry about 40% of the momentum. For the momentum of the valence quarks we choose the generic value $1/3$ and so $N_{\text{sea}} = 2/3 - N_G$.

Usually, the GPDs are taken from some non–perturbative (model) calculation or ansatz and plugged into the CFFs at a given input scale. Since the scale and the perturbative scheme are usually not specified, the matching of perturbative and non–perturbative frameworks has its own uncertainties. Let us first study the changes of the CFF (17) in a given scheme and input scale that appear when one includes the next order. The changes to the modulus and phase are appropriately measured by the $K$-factors:

$$K_\lambda^P = \frac{\ln \left| \mathcal{H}^{NP, \text{LO}} \right|}{\ln \left| \mathcal{H}^{NP-1, \text{LO}} \right|}, \quad K_{\arg}^P = \frac{\arg \left( \mathcal{H}^{NP, \text{LO}} \right)}{\arg \left( \mathcal{H}^{NP-1, \text{LO}} \right)}. \quad (44)$$
Here $\mathcal{SH}^{N^P\text{LO}}$ denotes the $N^P\text{LO}$ approximation, e.g., $P = 0$ for LO. One should bear in mind that $K$-factors actually measure the necessary reparameterization of the GPD to fit the given experimental data. We take the ansatz (14) with $\text{sea} F(\Delta^2) = G F(\Delta^2)$ and so the factorized $\Delta^2$ dependence essentially drops out in the $K$–factors. The ratio of gluon GPD to quark one is controlled by the factor $N_G/N_{\text{sea}}$ and, more importantly, by the differences of intercepts $\alpha_G(0) - \alpha_{\text{sea}}(0)$. To study the influence of this ratio, we distinguish two cases:

H) “hard” gluon: $N_G = 0.4, \quad \alpha_G(0) = \alpha_{\text{sea}}(0) + 0.1$ \hspace{1cm} (45)

S) “soft” gluon: $N_G = 0.3, \quad \alpha_G(0) = \alpha_{\text{sea}}(0).$ \hspace{1cm} (46)

We will use these parameters and $\alpha_{\text{sea}}(0) = 1.1$ at the input scale $Q^2 = 2.5\text{ GeV}^2$. Moreover, we set $\mu = Q$ and independently of the considered approximation we choose $\alpha_s(\mu_r^2 = 2.5\text{ GeV}^2) = 0.1\pi$. 

Figure 1: The relative radiative corrections, defined in Eq. (44), are plotted versus $\xi$ for the logarithm of the modulus [(a) and (b)] and phase [(c) and (d)] of $\mathcal{S}\mathcal{H}$, see Eqs. (17) and (41), for $\Delta^2 = 0$ [(a) and (c)] and $\Delta^2 = -0.5\text{ GeV}^2$ [(b) and (d)]: NNLO (solid) as well as in NLO for the ${\overline{\text{MS}}}$ (dashed) and MS (dotted) scheme. Thick (thin) lines refer to the “hard” (“soft”) gluon parameterization and we always set $\mu = Q$ and $\alpha_s(Q^2 = 2.5\text{ GeV}^2)/\pi = 0.1$. 

$P=2$ (NNLO), $CS$

$P=1$ (NLO), $CS$

$P=1$ (NLO), $MS$
and set the number of active flavors to three.

In Fig. 1 we depict for the typical kinematics of HERA collider experiments, i.e., $10^{-5} \lesssim \xi \lesssim 5 \cdot 10^{-2}$, the resulting $K$ factors for the logarithm of the modulus [(a) for $\Delta^2 = 0$ and (b) for $\Delta^2 = -0.5\, \text{GeV}^2$] and phase [(c) for $\Delta^2 = 0$ and (d) for $\Delta^2 = -0.5\, \text{GeV}^2$]. Here the thick and thin lines correspond to the “hard” and “soft” gluon parameterizations, respectively. We observe an almost flat $\xi$ dependence of the $K_\lambda$ factors in panels (a) and (b). This is not surprising, since the essential contribution arises from the pomeron pole and the CFF behaves as:

$$S^H \sim \left(\frac{1}{\xi}\right)^{\alpha(\Delta^2)} \left[i + \tan\left(\frac{\pi}{2}(\alpha(\Delta^2) - 1)\right)\right] \Rightarrow \ln|S^H| \approx \alpha(\Delta^2) \ln(1/\xi) + \text{const.} \quad (47)$$

For small $\xi$ this leads to the flatness we observe. The size of perturbative NLO corrections, see dashed (CS scheme) and dotted (MS scheme) lines, essentially depends on the ratio of gluon to quark GPDs. Since the gluons are a new entry, formally counted as NLO contribution, this finding is obvious and goes along with the observation that the perturbative corrections strongly vary within the used parameterization of parton densities in the Radyushkin GPD ansatz [55]. In the “hard” gluon parameterization the logarithm of the modulus reduces about 7–11% [5–8%] in the MS [CS] scheme, corresponding to the reduction of the modulus itself in the range of 40–70% [30–55%], where the drastic upper values correspond to $\xi = 10^{-5}$.

The relative radiative corrections to the phase grow in the small $\xi$ region with decreasing $\xi$ and can be of the order of up to 24% [13%] in the MS [CS] scheme. These effects are related to the signs for NLO Wilson coefficients, see Eqs. (20), (21), (30), and (31). In the “soft” gluon parameterization the NLO corrections are quite moderate for the modulus [(a) and (b)] and negligible for the phase [(c) and (d)]. From all four panels it can be realized that compared to MS scheme in the CS one the NLO corrections are typically reduced by 30–50%. This reduction has been also observed in the flavor non–singlet case [21]. The NNLO corrections (solid), compared to the NLO (dashed) ones, are drastically reduced. For the “soft” gluon parameterization they are practically negligible while for the “hard” gluon input they are reduced to the 1–2% level, except for the phase with $\Delta^2 = -0.5\, \text{GeV}^2$, where 5% are reached at $\xi = 10^{-5}$.

Let us finally address the modification of the scale dependence due to the higher order corrections. We only consider here the CS scheme and, analogously as in Eq. (44), we quantify the relative changes due to the evolution by the ratios

$$\dot{K}_\lambda^P = \frac{d \ln \left| S^H_{NP} \right|}{d \ln Q^2} / \frac{d \ln \left| S^H_{NP-1} \right|}{d \ln Q^2}, \quad \dot{K}_{\text{arg}} = \frac{d \arg \left( S^H_{NP} \right)}{d \ln Q^2} / \frac{d \arg \left( S^H_{NP-1} \right)}{d \ln Q^2}. \quad (48)$$

For the (exact) evolution of $\alpha_s(Q)$ we take the same scale setting and initial condition as above. However, the conformal moments (11) are evolved in the CS scheme, starting at the input scale.
Figure 2: The relative change of scale dependence, cf. Eq. (48), in the CS scheme at NLO (dashed, dotted) and NNLO (solid, dash–dotted) versus $\xi$ is depicted for the logarithm of the modulus (a) and phase (b) of the CFF (17) with $\Delta^2 = 0$ (dashed, solid) and $\Delta^2 = -0.5 \text{GeV}^2$ (dotted, dash–dotted) and $Q^2 = 4 \text{GeV}^2$. We set $\mu = Q$, $\alpha_s(\mu^2 = 2.5 \text{GeV}^2)/\pi = 0.1$ and took the input (41) at the scale $Q^2_0 = 1 \text{GeV}^2$. Thick and thin lines correspond again to “hard” and “soft” gluonic input.

$Q^2_0 = 1 \text{GeV}^2$, to $Q^2 = 4 \text{GeV}^2$. The non–leading logs in the solution of the evolution equation (32) are expanded with respect to $\alpha_s$ and are consistently combined with the Wilson–coefficients (18) in the considered order. The unknown NNLO mixing term $\Delta_{jk}^{\text{CS}}$ in Eq. (32) is neglected. This mixing can be suppressed at the input scale by an appropriate initial condition and so we expect only a minor numerical effect; see also Ref. [56]. The dashed and dotted lines in Fig. 2 show that in NLO the scale dependence changes can be rather large even of about 100% or more. In general the relative radiative corrections to NNLO are getting smaller. For instance, the NNLO corrections in panel (a) are almost negligible for the “soft” gluonic input with $\Delta^2 = 0$ (thin solid), they increase, however, for $\Delta^2 = -0.5 \text{GeV}^2$ and are becoming large for the “hard” gluonic input, e.g., about $-35\%$ at $\xi = 10^{-5}$ (thick dash–dotted). Note that these large corrections at very small $\xi$ are essentially caused by those of the anomalous dimensions in the vicinity of $j = 0$, corresponding to the large corrections of the gluon splitting kernels at small $x$, reported in [46]. The same sources also cause the huge NNLO corrections to the phase in panel (b). We remark that the modulus of $^8H$ is dominated by its imaginary part for which radiative corrections are milder than for the real part. The real part and so also the phase at very small $\xi$ are rather strongly affected by the NNLO corrections to anomalous dimensions. On the other hand for $5 \cdot 10^{-4} \lesssim \xi$ and $5 \cdot 10^{-3} \lesssim \xi$ the radiative NNLO corrections to the logarithm of the modulus (a) and phase (b), respectively, are rather mild (solid and dash–dotted lines). Restricted to these kinematics our findings support the convergence of the perturbative series.
4 Summary

In this letter we have studied NLO and NNLO corrections to deeply virtual Compton scattering in the small $\xi$ region. We confirmed that large radiative corrections at NLO can appear, reported before, and clarified their source which is entirely tied to the gluonic sector. In particular, if the gluon distribution starts to have a steeper increase at small $\xi$ than the quark ones, the NLO corrections will be dominated by the negative NLO gluon contribution and so the modulus of $S^gH$ will drastically reduce. On the other hand, if the gluon contribution is relatively small, already the NLO corrections are moderate. In any case the NNLO corrections are becoming moderate or even small at a given input scale, even at a few GeV$^2$. This fact supports the perturbative framework of DVCS.

The situation with respect to the scale dependence is not so conclusive. Going from LO to NLO we observe in general a big enhancement that arises from the large corrections to the anomalous dimensions, cf. [46]. To NNLO they will be reduced and the relative changes for the logarithm of the modulus are getting reasonable but grows to be large with decreasing $\xi$. Note that in this region the NNLO gluonic evolution effects are comparable in size with the NLO ones [46]. Also the NLO radiative corrections to the scale dependence of the phase of $S^gH$ are rather large, in particular for $\Delta^2 = 0$, at the scale of 4 GeV$^2$. To NNLO accuracy the convergency improves for $5 \cdot 10^{-3} \lesssim \xi$. Unfortunately, at smaller values of $\xi$ the convergency is lost. These large corrections due to evolution at small $\xi$ are certainly related to those found in DIS [46].

If one is interested to access GPDs from the DVCS cross section measurement at small $\xi$, only the modulus of $S^gH$ is essential. In that case perturbation theory seems to work in the sense that NNLO corrections of the Wilson–coefficients are negligible. They are, however, important for the scale violating effects for $\xi \lesssim 5 \cdot 10^{-4}$ (at relatively low $Q^2 \sim 4$ GeV$^2$). We also conclude that the photon and vector–meson leptonproduction data taken by the H1 and ZEUS collaborations should be perturbatively analyzed at NLO [57]. To our best knowledge a common perturbative analysis has not been done so far. To achieve this in a simple and numerical stable manner, the Mellin–Barnes integral representation seems to be preferred.

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