Finite-\(t\) and target mass corrections in off-forward hard reactions

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Abstract

We describe a systematic approach \cite{1} to the calculation of kinematic corrections \(\propto t/Q^2, m^2/Q^2\) in hard exclusive processes which involve momentum transfer from the initial to the final hadron state. As an example, the complete expression is derived for the time-ordered product of two electromagnetic currents that includes all kinematic corrections due to the quark distribution to twist-four accuracy. The results are applicable e.g. to the studies of deeply-virtual Compton scattering.

1 Introduction

There is hope that hard exclusive scattering processes in Bjorken kinematics can provide one with a three-dimensional picture of the proton in longitudinal and transverse plane \cite{2}, encoded in generalized parton distributions (GPDs) \cite{3,4}. One of the most important reactions in this context is Compton scattering with one real and one highly-virtual photon (DVCS) which has received a lot of attention. The QCD description of DVCS is based on the operator product expansion (OPE) of the time-ordered product of two electromagnetic currents. In this language the GPDs appear as leading-twist operator matrix elements. In order to probe the transverse proton structure one needs to measure the dependence of the amplitude on the momentum transfer to the target \(t = (P' - P)^2\) in a broad range. Since the available photon virtualities \(Q^2\) are limited to a few GeV\(^2\) range, corrections of the type \(\propto t/Q^2\) (which are formally higher-twist effects), are significant and have to be taken into account.

Such corrections are usually dubbed “kinematic” since they only involve ratios of kinematic variables and at first sight have nothing to do with nonperturbative effects (e.g. one may
consider a theoretical limit $\Lambda_{\text{QCD}}^2 \ll t \ll Q^2$. The separation of kinematic corrections $\propto t/Q^2$ from generic twist-four corrections $\mathcal{O}(\Lambda_{\text{QCD}}^2/Q^2)$ proves, however, to be surprisingly difficult. The problem is well known and its importance for phenomenology has been acknowledged by many authors [4, 5, 6, 7, 8, 9, 10, 11, 12, 13].

The challenge is that, unlike target mass corrections in inclusive reactions [15], which are determined solely by the contributions of leading twist operators, the $\sim t/Q^2$ corrections to off-forward processes (and for spin-1/2 targets also $\sim m^2/Q^2$ corrections) also receive contributions from higher-twist-four operators that can be reduced to total derivatives of the twist-two ones. Indeed, let $\mathcal{O}_{\mu_1...\mu_n}$ be a multiplicatively renormalizable (conformal) local twist-two operator, symmetrized and traceless over all indices. The operators

$$\mathcal{O}_1 = \partial^2 \mathcal{O}_{\mu_1...\mu_n}, \quad \mathcal{O}_2 = \partial^{\mu_1} \mathcal{O}_{\mu_1...\mu_n}$$

are, on the one hand, twist-four, and on the other hand their matrix elements are related to the leading twist matrix elements times the momentum transfer squared (up to, possibly, target mass corrections). Thus, both operators contribute to the $\propto t/Q^2$, $\propto m^2/Q^2$ accuracy and must be taken into account.

Moreover, all these contributions are intertwined by electromagnetic gauge and Lorentz invariance. Implementation of the electromagnetic gauge invariance beyond the leading twist accuracy has been at the center of many discussions, starting from Ref. [14]. By contrast, importance of the translation invariance condition has never been emphasized, to the best of our knowledge. In particular the distinction between the kinematic corrections of Nachtmann’s type, i.e. due to contributions of leading-twist [6, 7, 10, 11, 12, 13], and of higher-twist operators in Eq. (1) is not invariant under translations along the line connecting the electromagnetic currents in the $T$-product. Hence this distinction has no physical meaning; the existing estimates of kinematic effects, e.g. in DVCS, by the contributions of leading twist operators alone can be misleading.

On a more technical level, the problem arises because $\mathcal{O}_2$ has rather peculiar properties: the divergence of a conformal operator vanishes in the free theory (the Ferrara-Grillo-Parisi-Gatto theorem [16]). A related feature is that using QCD equations of motion (EOM) $\mathcal{O}_2$ can be expressed in terms of quark-antiquark-gluon operators. The simplest example of such a relation is known for many years [17, 18, 19]:

$$\partial^\mu O_{\mu\nu} = 2\bar{q}i\gamma_\nu G_{\gamma\mu} q,$$  \hspace{1cm} (2)

where $O_{\mu\nu} = (1/2)[\bar{q}\gamma_\mu \gamma_\nu q + (\mu \leftrightarrow \nu)]$ is the quark part of the energy-momentum tensor. The operator on the r.h.s. of Eq. (2) involves the gluon field strength and, naively, its hadronic matrix elements are of the order of $\Lambda_{\text{QCD}}^2$, which is in fact not the case. More complicated examples can be found in [20, 21].

The general structure of such relations is, schematically

$$(\partial \mathcal{O})_N = \sum_k a_k^{(N)} G_{Nk},$$  \hspace{1cm} (3)

where $G_{Nk}$ are twist-four quark-antiquark-gluon operators and $a_k^{(N)}$ are the numerical coefficients. The subscript $N$ stands for the number of derivatives in $\mathcal{O}_N$ and the summation goes
over all contributing operators which may include total derivatives (so that in practice \(k\) is a certain multi-index). The same operators, \(G_{N,k}\), also appear in the OPE for the product of currents of interest at the twist-four level:

\[
T\{j(x)j(0)\}^{tw=4} = \sum_{N,k} c_{N,k}(x) G_{N,k}.
\] (4)

A separation of “kinematic” and “dynamical” contributions to the OPE implies that one attempts to reassemble this expansion in such a way that the contribution of a particular combination appearing in (3) is separated from the remaining twist-four contributions. The “kinematic” power correction would correspond to taking into account this term only, and discarding contributions of “genuine” quark-gluon operators.

The guiding principle is that the separation of kinematic and dynamical effects is only physically meaningful (e.g. they are separately gauge- and Lorentz-invariant) if they have autonomous scale dependence. Different twist-four operators of the same dimension mix with each other and satisfy a certain renormalization group (RG) equation which can be solved, at least in principle. Let \(G_{N,k}\) be the set of multiplicatively renormalizable twist-four operators so that

\[
G_{N,k} = \sum_{k'} \psi^{(N)}_{k,k'} G_{N,k'}.
\] (5)

Eq. (3) tells us that one of the solutions of the RG equation is known without calculation. Indeed, it provides one with an explicit expression for a twist-four operator with the anomalous dimension equal to the anomalous dimension of the leading twist operator. (For simplicity we ignore the contributions of \(\partial^2 \mathcal{O}_N\) in this discussion; they do not pose a problem and can be taken into account using conventional methods.)

Let us assume that this special solution corresponds to \(k = 0\), i.e. \(G_{N,k=0} \equiv (\partial \mathcal{O})_N\) and \(\psi^{(N)}_{k=0,k'} = \alpha_{k'}\). Inverting the matrix of coefficients, \(\psi^{(N)}_{k,k'}\), and separating the term with \(k = 0\) we can write the expansion of an arbitrary twist-four operator in terms of the multiplicatively renormalizable ones

\[
G_{N,k} = \phi^{(N)}_{k,0} (\partial \mathcal{O})_N + \sum_{k' \neq 0} \phi^{(N)}_{k,k'} G_{N,k'}.
\] (6)

Inserting this expansion into Eq. (4) one obtains

\[
T\{j(x)j(0)\}^{tw=4} = \sum_{N,k} c_{N,k}(x) \phi^{(N)}_{k,0} (\partial \mathcal{O})_N + \ldots,
\] (7)

where the ellipses stand for the “genuine” twist-four quark-gluon operators (e.g. with different anomalous dimensions). This is the solution we want to have, but the problem with it is that finding the coefficients \(\phi^{(N)}_{k,0}\) in general requires knowledge of the full matrix \(\psi^{(N)}_{k,k'}\), in other words the explicit solution of the twist-four RG equations, which is not available.

Our starting observation is that twist-four operators in QCD come in two big groups: the so-called quasipartonic [22], that only involve “plus” components of the fields, and non-quasipartonic which also include “minus” light-cone projections. Quasipartonic operators are not relevant for the present discussion since they have an autonomous evolution (to one-loop
accuracy). As a consequence, \((\partial O)_N\) does not appear in the expansion of quasipartonic operators in multiplicatively renormalizable ones, Eq. (6): the corresponding coefficients \(\phi_{k,0}^{(N)}\) vanish. Hence the kinematic power correction \(\sim (\partial O)_N\) originates entirely from contributions of non-quasipartonic operators.

Renormalization of twist-four non-quasipartonic operators was studied recently in [23, 24]. The main result is that in a suitable operator basis the corresponding RG equations can be written in terms of several \(SL(2)\)-invariant kernels. Using \(SL(2)\)-invariance we are able to prove that the anomalous dimension matrix for non-quasipartonic operators is hermitian with respect to a certain scalar product. This implies that different eigenvectors are mutually orthogonal, i.e.

\[
\sum_k \mu_k^{(N)} \psi_{l,k}^{(N)} \psi_{m,k}^{(N)} \sim \delta_{l,m},
\]

where \(\mu_k^{(N)}\) is the corresponding (nontrivial) measure. From this orthogonality relation and the expression (3) for the relevant eigenvector one obtains, for the non-quasipartonic operators \(\phi_{k,0}^{(N)} = a_k^{(N)} ||a^{(N)}||^{-2}\), where \(||a^{(N)}||^2 = \sum_k \mu_k^{(N)} (a_k^{(N)})^2\). Inserting this expression into (7) one ends up with the desired separation of kinematic effects.

The actual derivation is done using the two-component spinor formalism in intermediate steps and requires some specific techniques of the \(SL(2)\) representation theory. This talk is based on the results presented in Ref. [1]; details of the derivation will be given in a forthcoming paper.

2 T-product of two electromagnetic currents

We have been able to find the contributions related to the leading-twist operator (11) in the \(T\)-product of two electromagnetic currents \(T_{\mu\nu} = i T\{j_{\mu}^{\text{em}}(x)j_{\nu}^{\text{em}}(0)\}\) to twist-four accuracy. The result can be brought to the form

\[
T_{\mu\nu} = -\frac{1}{\pi^2 x^4} \left\{ x^\alpha \left[ S_{\mu\nu} \gamma^\beta + i \epsilon_{\mu\nu\alpha\beta} A_\beta \right] + x^2 \left[ (x_\mu \partial_\nu + x_\nu \partial_\mu) X + (x_\mu \partial_\nu - x_\nu \partial_\mu) Y \right] \right\},
\]

where \(\partial_\mu = \partial/\partial x^\mu\), \(S_{\mu\nu} = g_{\mu\alpha} g_{\nu\beta} + g_{\nu\alpha} g_{\mu\beta} - g_{\mu\nu} g_{\alpha\beta}\) and a totally antisymmetric tensor is defined such that \(\epsilon_{0123} = 1\). The expansion of invariant functions \(V_\beta\) and \(A_\beta\) starts from twist two, whereas \(X\) and \(Y\) are already twist-four. In order to write the result we first need to introduce some notations.

We define nonlocal (light-ray) vector \(O_\nu\) and axial-vector \(O_A\) operators of the leading-twist-two as the generating functions for local twist-two operators

\[
O(z_1 x, z_2 x) = \left[ \bar{q}(z_1 x) f (\gamma_5) Q^2 q(z_2 x) \right]_{l.t.},
\]

where \(x_\mu\) is an arbitrary four-vector (not necessarily light-like), \(z_1\) and \(z_2\) are real numbers and \(Q\) is the matrix of quark electromagnetic charges. Here and below the Wilson line between the
quark fields is implied. The leading-twist projector \([\ldots]_{l.t.}\) stands for the subtraction of traces of the local operators so that by definition
\[
[q(z_1 x) \not\! q^2 q(z_2 x)]_{l.t.} = \sum_N \frac{1}{N!} x_\mu x_\mu \ldots x_\mu N \left\{ q(0) \gamma_\mu [z_1 \tilde{D}_\mu + z_2 \tilde{D}_\mu] \ldots [z_1 \tilde{D}_\mu N + z_2 \tilde{D}_\mu N] Q^2 q(0) - \text{traces}\right\}.
\]
The leading-twist light-ray operators satisfy the Laplace equation \(\partial_\mu^2 O(z_1 x, z_2 x) = 0\). The explicit form of the projector \([\ldots]_{l.t.}\) is irrelevant for what follows. Useful representations can be found e.g. in [9, 25].

Thanks to crossing symmetry the vector and axial-vector operators always appear to be antisymmetrized and symmetrized over the quark and antiquark positions, respectively, so we define the corresponding combinations:
\[
O_V^{(\pm)}(z_1, z_2) = [q(z_1 x) \not\! q^2 q(z_2 x)]_{l.t.} - (z_1 \leftrightarrow z_2),
\]
\[
O_A^{(\pm)}(z_1, z_2) = [q(z_1 x) \gamma_5 Q^2 q(z_2 x)]_{l.t.} + (z_1 \leftrightarrow z_2).
\]
The leading-twist expressions are well known and can be written as (cf. [25])
\[
\Psi^{l.t. = 2}_\mu = \frac{1}{2} \partial_\mu \int_0^1 du O_V^{(-)}(u, 0), \quad A^{l.t. = 2}_\mu = \frac{1}{2} \partial_\mu \int_0^1 du O_A^{(+)}(u, 0).
\]
Note that the separation of the leading-twist terms \([\ldots]_{l.t.}\) from the nonlocal operators produces a series of kinematic power corrections to the amplitudes, which are similar to Nachtmann target mass corrections in deep-inelastic lepton-nucleon scattering [15]. Such corrections are discussed in detail in [8, 9, 7, 10, 11, 12, 13].

For the twist-three terms we obtain
\[
\Psi^{l.t. = 3}_\mu = \left[ i \mathcal{P}_\nu, \int_0^1 du \left\{ i e_{\mu \alpha \beta \gamma} x_\alpha \partial_\beta \tilde{O}_V^{(+)}(u) + \left( S_{\mu \nu \beta \gamma} x_\alpha \partial_\beta + \ln u \partial_\mu x_2 \partial_\nu \right) \tilde{O}_V^{(-)}(u) \right\} \right],
\]
\[
A^{l.t. = 3}_\mu = \left[ i \mathcal{P}_\nu, \int_0^1 du \left\{ i e_{\mu \alpha \beta \gamma} x_\alpha \partial_\beta \tilde{O}_A^{(-)}(u) + \left( S_{\mu \nu \beta \gamma} x_\alpha \partial_\beta + \ln u \partial_\mu x_2 \partial_\nu \right) \tilde{O}_A^{(+)}(u) \right\} \right].
\]
Here \(\mathcal{P}_\nu\) is the momentum operator \([i \mathcal{P}_\nu, q(y)] = \frac{\partial}{\partial y^\nu} q(y)\), and we used the notation
\[
\tilde{O}_a^{(\pm)}(z) = \frac{1}{4} \int_0^2 dw O_a^{(\pm)}(z, w).
\]
One can easily verify that \(x_\mu \Psi^{l.t. = 3}_\mu = \partial_\mu \Psi^{l.t. = 3}_\mu = 0\) and similarly \(x_\mu A^{l.t. = 3}_\mu = \partial_\mu A^{l.t. = 3}_\mu = 0\). Note that the terms in \(\ln u\) in Eqs. (15) are themselves twist-four and can be omitted if the calculation is done to twist-three accuracy. The resulting simplified expression is in agreement with Refs. [6, 7]. These terms must be included, however, in order to ensure correct separation of twist-three and twist-four contributions.

The flavor-nonsinglet twist-four contributions to Eq. (10) present our main result. In this case we prefer to write the answer in terms of integrals over the position of the local conformal
operators, cf. Eq. (17). This form is usually referred to as the conformal OPE [26]. For example, a light-ray operator can be written as the conformal expansion

\[ O(z_1x, z_2x) = \sum_N \chi_N z_{12}^N \int_0^1 du (u\bar{u})^{N+1} [O_N(z_{12}^N x)]_{l.t.}, \]  

(17)

where

\[ \chi_N = 2(2N + 3)/(N + 1)! \]

and we use the shorthand notation \( \bar{u} = 1 - u \), \( z_{12} = z_1 - z_2 \), \( z_{12}^N = \bar{u}z_1 + u z_2 \). The conformal operator \( O_N \) is defined as

\[ O_N(y) = (\partial_{z_1} + \partial_{z_2})^N C_N^{3/2} \left( \frac{\partial_{z_1} - \partial_{z_2}}{\partial_{z_1} + \partial_{z_2}} \right) O(z_1x + y, z_2x + y) \Big|_{z_1 = 0}, \]  

(18)

where \( C_N^{3/2}(x) \) is the Gegenbauer polynomial.

The leading-twist contribution to the OPE of two electromagnetic currents can be written in the same form, for comparison:

\[ \mathbb{V}_{\mu}^{\mu} = \partial_\mu \sum_{N, o d d} \frac{\chi_N}{N + 2} \int_0^1 du u^N \bar{u}^{N+2} [O_N^V(ux)]_{l.t.}. \]  

(19)

Here \( O_N^V(ux) \) is the conformal operator (18) at the space-time position \( ux \).

We obtain

\[ \mathbb{V}_{\mu}^{\mu} = \frac{1}{2} \sum_{N, o d d} \frac{\chi_N}{(N + 2)^2} \int_0^1 du \left\{ (u\bar{u})^{N+1} x_\mu [\tilde{O}_N^V(ux)]_{l.t.} \right\}, \]

\[ + \frac{N}{2} u^{N-1} \bar{u}^{N+2} \left[ u + \frac{1}{N + 2} \right] x^2 \partial_\mu [\tilde{O}_N^V(ux)]_{l.t.} \right\}, \]

\[ \tilde{A}_{\mu}^{\mu} = \frac{1}{4} \sum_{N, e v e n} \frac{\chi_N N}{(N + 2)^2} \int_0^1 du u^{N-1} \bar{u}^{N+2} \left[ 1 - \frac{N + 1}{N + 2} \bar{u} \right] x^2 \partial_\mu [\tilde{O}_N^A(ux)]_{l.t.} \right\}, \]

\[ \tilde{X}_{\mu}^{\mu} = \frac{1}{4} \sum_{N, o d d} \frac{\chi_N}{(N + 2)^2} \int_0^1 du u^{N-1} \bar{u}^{N+1} \left[ 1 - 2 \frac{N + 1}{N + 2} \bar{u} \right] [\tilde{O}_N^V(ux)]_{l.t.} \right\}, \]

\[ \tilde{Y}_{\mu}^{\mu} = - \frac{1}{4} \sum_{N, o d d} \frac{\chi_N}{(N + 2)^2} \int_0^1 du u^{N-1} \bar{u}^{N+1} \left[ 1 - 2 \frac{N + 1}{N + 2} \bar{u} + 2 \frac{N + 1}{N + 3} \bar{u}^2 \right] [\tilde{O}_N^V(ux)]_{l.t.}. \]  

(20)

Here \( \tilde{O}_N \) is defined as the divergence of the leading-twist conformal operator, cf. \( O_2 \) in Eq. (1):

\[ \tilde{O}_N(y) = \frac{1}{N + 1} \partial_{\mu} [iP_\mu, O_N(y)] = [iP_\mu, O_{\mu_1 \cdots \mu_N}(y)] x^{\mu_1} \cdots x^{\mu_N}. \]  

(21)

Note that the operator \( O_1 \) in Eq. (1), which corresponds to \([iP_\mu, iP_\mu, O_N]\) in our present notation, does not contribute to the answer for our special choice of the correlation function.
A short calculation yields $F$ where twist GPD and we used the notations $P$ Hadronic matrix elements of the twist-4 operator $3$ typical matrix elements advantageous in certain applications.

Conservation of the electromagnetic current implies that $\partial^{\mu} T_{\mu\nu}(x) = 0$ and $\partial^{\nu} T_{\mu\nu}(x) = i[P^{\nu}, T_{\mu\nu}(x)]$. We have checked that these identities are satisfied up to twist-5 terms.

For completeness we give the relation for the operator $[iP_\mu, \partial^{\mu} O(z_1, z_2)]$ entering the twist-three functions $V_{\mu}^{t-3}$, $A_{\mu}^{t-3}$ in terms of $\hat{O}_N$:

$$[iP_\mu, \partial^{\mu} O(z_1, z_2)] = \frac{1}{2} S^+ \int_0^1 u du [iP_\mu [iP^\mu, O(uz_1, uz_2)]] + \sum_N z_N (N+1)^2 z_1 z_2 \int_0^1 dv v^{N-1} \int_0^1 du (u\bar{u})^{N+1} \hat{O}_N (vz_1 vz_2),$$

where $S^+ = z_1^2 \partial_{z_1} + z_2^2 \partial_{z_2} + 2z_1 + 2z_2$. It is also possible to rewrite, v.v., all contributions of local operators $\hat{O}_N$ in terms of the nonlocal light-ray operator $[iP_\mu, \partial^{\mu} O(z_1, z_2)]$, which can be advantageous in certain applications.

### 3 Typical matrix elements

Hadronic matrix elements of the twist-4 operator $\hat{O}_N$ are of course related to those of the leading twist, $O_N$. For illustration, we present the corresponding explicit expressions for the two proton states with momenta $p' \neq p$, which are relevant e.g. for virtual Compton scattering. The leading-twist matrix elements can be parametrized as (cf. [34])

$$\langle p' | O_N(n) | p \rangle = \bar{u}(p')\bar{u}(p) \sum_{k=even}^N F_{N,k}(t) \Delta^k P^N - k + \frac{1}{m} \bar{u}(p') u(p) \sum_{k=even}^{N+1} H_{N,k}(t) \Delta^k P^N - 1 - k,$$

where $F_{N,k}(t)$ and $H_{N,k}(t)$ are generalized form factors corresponding to moments of the leading-twist GPD and we used the notations $P = (p + p')/2$, $\Delta = p' - p$, $p'^2 = m^2$; $t = \Delta^2$; $u(p)$ is the nucleon spinor. By analogy, we define

$$\langle p' | \hat{O}_N(n) | p \rangle = \bar{u}(p')\bar{u}(p) \sum_{k=even}^N \hat{F}_{N,k}(t) \Delta^k P^N - k + \frac{1}{m} \bar{u}(p') u(p) \sum_{k=even}^{N+1} \hat{H}_{N,k}(t) \Delta^k P^N - 1 - k.$$

A short calculation yields

$$\hat{F}_{N,k}(t) = t F_{N,k}(t) \frac{k(2N + 3 - k)}{2(N + 1)^2} - \left( m^2 - \frac{t}{4} \right) F_{N,k-2} \frac{(N - k + 2)(N - k + 1)}{2(N + 1)^2}$$

$$\hat{H}_{N,k}(t) = t H_{N,k}(t) \frac{k(2N + 3 - k)}{2(N + 1)^2} - \left( m^2 - \frac{t}{4} \right) H_{N,k-2} \frac{(N - k + 3)(N - k + 2)}{2(N + 1)^2} - m^2 \frac{(N - k + 2)}{(N + 1)^2} F_{N,k-2}(t).$$

Note that the twist-4 matrix elements involve both finite-$t$ and target (nucleon) mass corrections. Concrete applications will be considered elsewhere.
4 Conclusions

To summarize, we have given a complete expression for the time-ordered product of two electromagnetic currents that resums all kinematic corrections related to quark GPDs to twist-four accuracy. The results can be applied to various two-photon processes, e.g. to the studies of deeply-virtual Compton scattering and $\gamma^* \rightarrow (\pi, \eta, \ldots) + \gamma$ transition form factors. The twist-four terms calculated in this work give rise to both a $\propto t/Q^2$ correction and the target mass correction $\propto m^2/Q^2$ for DVCS, whereas for the transition form factors these two effects are indistinguishable as there is only one scale. The main remaining question is whether QCD factorization itself is valid in such reactions to twist-four accuracy, at least for kinematic contributions. Clarification of this issue goes beyond the tasks of this study.

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