Azimuthal and single spin asymmetry in deep-inelastic lepton-nucleon scattering

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The collinear expansion technique is generalized to the factorization of unintegrated parton distributions and other higher twist parton correlations from the corresponding collinear hard parts that involve multiple parton final state interaction. Such a generalized factorization provides a consistent approach to the calculation of inclusive and semi-inclusive cross sections of deep-inelastic lepton-nucleon scattering. As an example, the azimuthal asymmetry is calculated to the order of $1/Q$ in semi-inclusive deeply inelastic lepton-nucleon scattering with transversely polarized target. A non-vanishing single-spin asymmetry in the “triggered inclusive process” is predicted to be $1/Q$ suppressed with a part of the coefficient related to a moment of the Sivers function.

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Many interesting phenomena have been observed in semi-inclusive deep-inelastic lepton-nucleon scattering (SIDIS), in particular the azimuthal asymmetries in the momentum distribution of the final hadrons and their spin dependence. Since most of the studies involve hadrons with $p_{\perp} \sim 1$ GeV/c, the intrinsic parton transverse momenta and multiple parton scattering become critical in the perturbative QCD (pQCD) approach.

The effects of intrinsic parton transverse momenta and multiple parton scattering in inclusive and the $p_{\perp}$-integrated semi-inclusive processes are normally higher-twist and have been studied extensively in the past. An elegant and practical framework in terms of collinear expansion has been developed and applied to different reactions. Within the framework, a factorized form for the cross section of inclusive and semi-inclusive deep-inelastic scattering is obtained as a convolution of the calculable hard parts with the universal parton distributions and correlation functions (hereafter referred generally as parton correlations) that can be measured in different reactions. The gauge invariant parton correlations contain contributions from the initial and final state interactions with soft gluons. They have many remarkable properties in particular when extended to the polarized case. In this note, we extend the generalized factorization to the transverse momentum dependence of SIDIS. Such a framework will provide a consistent and systematic approach to the pQCD study of SIDIS beyond the leading twist.

FIG. 1: Feynman diagrams for the three cases we considered. The gluon momenta in (c) are: $k_3 = k - k_1$ and $k_4 = k - k_2$.

For simplification, we do not consider the fragmentation and thus start with the singly polarized semi-inclusive process $e^- p^1 \rightarrow e^- q X$ with one photon exchange. The cross section can be written as,

$$d\sigma = \frac{e^4 e^2}{2s Q^4} L^{\mu\nu}(l,l') W^{(si)}_{\mu\nu}(q,p,S,k') \frac{d^3 l'}{(2\pi)^3 2E_{l'}} \frac{d^3 k'}{(2\pi)^3 2E_{k'}}$$

where $l,l',k,k',p$ are the 4-momenta of the leptons, quarks and nucleon respectively (see Fig.1); $S$ is the spin of the nucleon. We neglect the masses and use the light-cone coordinates. Unit vectors are $\hat{n}$ and $\hat{X}$. The leptonic and hadronic tensors are given by,

$$L^{\mu\nu}(l,l') = 4[(l'^\nu l'^\mu + l'^\mu l'^\nu) - (l \cdot l') g^{\mu\nu}],$$

$$W^{(si)}_{\mu\nu}(q,p,S,k') = \frac{1}{2\pi} \sum X \langle p,S,J_\mu(0)|k'|X \rangle \langle k'|J_\nu(0)|p,S \rangle (2\pi)^4 \delta^4(p + q - k' - p_X),$$

where the superscript $(si)$ denotes that it is for SIDIS.

We consider final state interaction in pQCD up to two gluon exchanges thus have contributions from the three diagrams shown in Fig.1(a-c). Hence, $W^{(si)}_{\mu\nu} = \sum_{j=0,1,2} W^{(si),j}_{\mu\nu}$, where $j$ denotes the number of gluons.

The lowest order in pQCD is $W^{(0,si)}_{\mu\nu}$ from Fig. 1(a),

$$W^{(0,si)}_{\mu\nu} = \frac{1}{2\pi} \int \frac{d^4 k}{(2\pi)^4} Tr[\tilde{H}^{(0,si)}_{\mu\nu}(k,k',q)\phi^{(0)}(k,p,S)],$$

$$\tilde{H}^{(0,si)}_{\mu\nu}(k,k',q) = \gamma_\mu(k + \hat{q})\gamma_\nu(2\pi)^4 \delta^4(k' - k - q).$$

This is to compare with $W^{(0)}_{\mu\nu}$ for inclusive reaction, where an integration over $d^3 k'$ is carried out so that,

$$W^{(0)}_{\mu\nu}(q,p,S) = \frac{1}{2\pi} \int \frac{d^4 k}{(2\pi)^4} Tr[\tilde{H}^{(0)}_{\mu\nu}(k,k,q)\phi^{(0)}(k,p,S)],$$

$$\tilde{H}^{(0)}_{\mu\nu}(k,q) = \gamma_\mu(k + \hat{q})\gamma_\nu(2\pi)^4 \delta^4((k + q)^2),$$

where $\delta_+$ means that only the positive solution is taken. Defining $K(k',k,q) = 2E_{k'}(2\pi)^3 \delta^3(\vec{k}' - \vec{k} - \vec{q})$, one has,
\[ W_{\mu\nu}^{(0,si)}(q,p,S,k') = \frac{1}{2\pi} \int \frac{d^4k}{(2\pi)^4} K(k',k,q) \text{Tr} [\hat{H}_{\mu\nu}^{(0)}(k,q)\hat{\phi}^{(0)}(k,p,S)]. \]  

The difference between \( W_{\mu\nu}^{(0,si)}(q,p,S,k') \) for SIDIS and \( W_{\mu\nu}^{(0)}(q,p,S) \) for inclusive DIS is the factor \( K(k',k,q) \) in the integrand. Similarly, corresponding to Figs. 1(b) and (c), we have contributions from one and two gluon exchanges,

\begin{align*}
W_{\mu\nu}^{(1,si)}(q,p,S,k') &= \frac{1}{2\pi} \int \frac{d^4k_1}{(2\pi)^4} \frac{d^4k_2}{(2\pi)^4} \sum_{c=L,R} K(k',k_1,q) \text{Tr} [\hat{H}_{\mu\nu}^{(1,c)}(k_1,k_2,q)\hat{\phi}_\rho^{(1)}(k_1,k_2,p)]; \\
W_{\mu\nu}^{(2,si)}(q,p,S,k') &= \frac{1}{2\pi} \int \frac{d^4k_1}{(2\pi)^4} \frac{d^4k_2}{(2\pi)^4} \frac{d^4k_3}{(2\pi)^4} \sum_{c=L,R,M} K(k',k_1,q) \text{Tr} [\hat{H}_{\mu\nu}^{(2,c,\sigma)}(k_1,k_2,k_3,q)\hat{\phi}_\rho^{(2)}(k_1,k_2,k_3,p,S)];
\end{align*}

where \( c \) denotes the different cuts of the diagrams and \( k_L = k_R = k_2 = k_3 = k \). The hard parts are given by,

\[ \hat{H}_{\mu\nu}^{(1,L)}(k_1,k_2,q) = \gamma_\mu(\bar{k}_1 + \bar{q})\gamma_\nu \frac{k_2 + \bar{q}}{(k_2 + q)^2} - i\epsilon \gamma_\nu(2\pi)\delta_+(k_1 + q)^2, \]

\[ \hat{H}_{\mu\nu}^{(2,L)}(k_1,k_2,q) = \gamma_\mu(\bar{k}_1 + \bar{q})\gamma_\nu \frac{k_2 + \bar{q}}{(k_2 + q)^2 + i\epsilon} \gamma_\nu \gamma_\sigma \frac{k_3 + \bar{q}}{(k_3 + q)^2 + i\epsilon} \gamma_\rho(2\pi)\delta_+(k_1 + q)^2, \]

and similarly for \( \hat{H}_{\mu\nu}^{(1,R)} \), \( \hat{H}_{\mu\nu}^{(2,R)} \) and \( \hat{H}_{\mu\nu}^{(2,M)} \). The \( \hat{\phi} \)'s are defined as,

\[ \hat{\phi}^{(0)}(k,p,S) = \int d^4ze^{ik\cdot x}(p,S)|\psi(0)(\psi(z)|p,S), \]

\[ \hat{\phi}^{(1)}(k_1,k_2,p,S) = \int d^4yd^4ze^{ik_1\cdot y+k_2\cdot (z-y)}(p,S)|\psi(0)(gA_p(y)\psi(z)|p,S), \]

\[ \hat{\phi}^{(2)}_\rho(k_1,k_2,k_3,p,S) = \int d^4yd^4yd^4ze^{i(k_1\cdot y+k_2\cdot (z-y)+k_3\cdot (z'-y'))}(p,S)|\psi(0)(gA_p(y)gA_{p'}(y')\psi(z)|p,S). \]

We note that the \( \hat{\phi} \)'s defined above are not gauge invariant. To use the above results in terms of gauge invariant parton correlations, we apply a collinear expansion to the hard parts as in the inclusive processes [11, 13, 17],

\[ \hat{H}_{\mu\nu}^{(0)}(k,q) = \frac{\partial \hat{H}_{\mu\nu}^{(0)}(xp)}{\partial k_\rho} \omega_\rho' k_\rho' + \frac{\partial^2 \hat{H}_{\mu\nu}^{(0)}(xp)}{2 \partial k_\rho \partial k_\sigma} \omega_\rho' \omega_\sigma' k_\rho' k_\sigma' + ... \]

where \( x = k^+/p^+ \), \( \hat{H}_{\mu\nu}^{(0)}(xp) \equiv \hat{H}_{\mu\nu}^{(0)}(k,q)|_{k=xp} \), \( \partial \hat{H}_{\mu\nu}^{(0)}(xp) / \partial k_\rho \equiv \partial \hat{H}_{\mu\nu}^{(0)}(k,q) / \partial k_\rho|_{k=xp} \), and so on; \( \omega_\rho' \equiv g_\rho'\gamma_\sigma - \bar{n}_\rho n^\sigma \) is a projection operator so that \( \omega_\rho' k_\rho' = (k-x)p_\rho \). Next, we decompose the gluon field \( A_p \) into \( A_p(y) = \omega_\rho' A_{p'}(y) + p_\rho n \cdot A(y)/n \cdot p \). Using some generalized Ward identities one can relate the derivatives of the hard parts to the hard parts of a higher order. After rearranging all the terms by adding the contributions with the same hard part together we can obtain \( \tilde{W}_{\mu\nu}^{(0,si)}(q,p,S,k') = \sum_{j=0,1,2} \tilde{W}_{\mu\nu}^{(j,si)}(q,p,S,k') \), and,

\[ \tilde{W}_{\mu\nu}^{(0,si)}(q,p,S,k') = \frac{1}{2\pi} \int \frac{d^4k_1}{(2\pi)^4} \frac{d^4k_2}{(2\pi)^4} \sum_{c=L,R} K(k',k_1,q) \text{Tr} [\hat{H}_{\mu\nu}^{(1,L)}(k_1,k_2,q)\hat{\phi}_\rho^{(1)}(k_1,k_2,p,S)]; \]

\[ \tilde{W}_{\mu\nu}^{(1,si)}(q,p,S,k') = \frac{1}{2\pi} \int \frac{d^4k_1}{(2\pi)^4} \frac{d^4k_2}{(2\pi)^4} \frac{d^4k_3}{(2\pi)^4} \sum_{c=L,R,M} K(k',k_1,q) \text{Tr} [\hat{H}_{\mu\nu}^{(2,c,\sigma)}(k_1,k_2,k_3,q)\hat{\phi}_\rho^{(2)}(k_1,k_2,k_3,p,S)]; \]

The gauge invariant matrix elements \( \hat{\phi} \)'s have contributions from all the three diagrams in Fig. 1,

\[ \hat{\phi}^{(0)}(k_1,k_2,p,S) = \int d^4yd^4ze^{ik_1\cdot y+k_2\cdot (z-y)}(p,S)|\psi(0)(\psi(z)|p,S), \]

\[ \hat{\phi}^{(1)}_\rho(k_1,k_2,p,S) = \int d^4yd^4zd^4ze^{ik_1\cdot y+k_2\cdot (z-y)}(p,S)|\psi(0)(\psi(z)|p,S), \]

\[ \hat{\phi}^{(2)}_\rho(k_1,k_2,k_3,p,S) = \int d^4yd^4yd^4zd^4ze^{i(k_1\cdot y+k_2\cdot (z-y)+k_3\cdot (z'-y'))}(p,S)|\psi(0)(\psi(z)|p,S), \]

\[ \mathcal{L}(-\infty,z) = 1 + ig \int_{-\infty}^{z} dy^- A^+(z^+,y^-,z_L) + (ig)^2 \int_{-\infty}^{z} dy^- A^+(z^+,y^-,z_L) \int_{-\infty}^{z} dy^- A^+(z^+,y^-,z_L), \]

where \( D_\rho(y) = -i\partial_\rho + gA_\rho(y) \) is the covariant derivative. Including higher order contributions in \( g \),

\[ \mathcal{L}(-\infty,z) = Pe^{ig} \int_{-\infty}^{z} dy^- n^\rho A(z^+,y^-,z_L), \]

is the path integral representation of the gauge link with transverse displacement \( z_L \). Note that the leading twist matrix elements \( \hat{\phi}^{(0)} \) is related to the unintegrated quark distributions in a nucleon.

Eq.(11) can be considered as a generalized factorization of transverse momentum dependent SIDIS as a result of the collinear expansion. Consequently, there are two distinct properties in the factorized form: (A) All the hard parts in the Eq. (11) are only functions of the longitudinal parton momenta. This corresponds to hard scattering of partons with only longitudinal momentum. All the information of transverse momentum is contained only in the matrix elements \( \hat{\phi} \)'s. (B) The operators \( \omega_\rho' \) and \( \omega_\sigma' \) in \( \tilde{W}_{\mu\nu}^{(1,si)}, \tilde{W}_{\mu\nu}^{(2,si)} \) project away the longitudinal components of the
gauged fields that go into the gauge links. We want to point out that the final result for the cross section is not simply a convolution of the transverse momentum dependent lepton-quark scattering cross section with the unintegrated (or transverse momentum dependent) quark distributions. Such a naive convolution can result in double-counting of the effects from the transverse momentum.

The properties (A) and (B) mentioned above lead to a great simplification of \( \hat{W}^{(1,s)}_{\mu \nu} \). It can be shown that,

\[
H^{(1,L)}_{\mu \nu}(x_1 p, x_2 p)\omega_{\rho}^{\prime} = \frac{\pi}{2} \delta(x_1 - x_B)\omega_{\rho}^{\prime} \gamma^\mu \gamma^\rho \gamma^\nu \gamma^\nu,
\]

\[
H^{(1,R)}_{\mu \nu}(x_1 p, x_2 p)\omega_{\rho}^{\prime} = \frac{\pi}{2} \delta(x_2 - x_B)\omega_{\rho}^{\prime} \gamma^\mu \gamma^\rho \gamma^\nu \gamma^\nu,
\]

which are independent of \( x_2 \) or \( x_1 \) respectively. Hence, we can carry out the integration over \( d^4k_2 \) or \( d^4k_1 \) in \( \Phi^{(0)}_{\rho}(k_1, k_2, p, S) \) and obtain,

\[
\hat{W}^{(1,s)}_{\mu \nu} = \frac{1}{\pi} \text{Re} \int \frac{dk}{(2\pi)^3} K(k, q) \text{Tr} \left[ H^{(1,R)}_{\mu \nu}(x)p \Phi^{(0)}_{\rho}(k, p, S) \right],
\]

where \( H^{(1,R)}_{\mu \nu}(x) \) is simply a convolution of the \( \rho \) dependence of the formalism, we next calculate the azimuthal asymmetries to the order of \( \frac{1}{Q} \), viz.

\[
A(y)\left[ f_1 + \frac{B(y)}{M} f_{1T} \sin(\phi - \phi_s) \right] - B(y)\frac{M}{\tilde{k}^2 L} \left( x_B f_L - \frac{1}{Q} \right) \cos \phi + \frac{E^2}{2M} (\varphi_{sL2} - 3\varphi_{sL2}) \sin \phi_s + \frac{E^2}{2M} (\varphi_{sL1} + \varphi_{sL1}) \sin(2\phi - \phi_s) \right],
\]

where \( B(y) = 2(2-y)\sqrt{1-y} \), \( \phi \) and \( \phi_s \) are the azimuthal angle of \( \bar{k}_L \) and \( S_L \), respectively, relative to the lepton plane; \( \bar{k}_L = k^2 - q^- \) and \( k^+ = x_B p^+ \).

From Eq. (20), one notes that the leading contribution to the azimuthal angle dependence comes only from the leading twist quark distribution \( \Phi_0(x) \) while the higher order \( (1/Q) \) receives contribution from both the leading and the next leading twist parton matrix elements. In view of the fact that the momentum of the current is small compared to the lepton momentum, the leading twist parton matrix elements are much smaller than the next leading twist parton matrix elements. One can see more clearly from the cross section integrated over \( \phi \),

\[
d\sigma = \frac{\alpha_s^2 \pi^2}{Q^4} \frac{dx d\eta dQ^2 dz}{(2\pi)^3} \int \frac{dy}{(2\pi)^4} K(k, q) \delta(x - x_B) \cdot \left\{ A(y) f_1 + \frac{B(y)}{2M} f_{1T} \sin(\phi - \phi_s) \right\}.
\]

where \( A(y) = 1 + (1-y)^2 \), and \( \alpha_s = e^2/4\pi \). The Lorentz structure of \( \Phi_0(x) \) is given by [10],

\[
\Phi^{(0)}_{\alpha} = p_\alpha f_1 + \omega_\alpha^{\prime} k_\alpha f_+ + \varepsilon_\alpha^{\prime} \delta^\rho \gamma_k S^\rho f_{1T}/M,
\]

where \( M \) is the nucleon mass. A complete examination of the Lorentz structures of \( \varphi_{\rho\rho}^{(1)} \) and \( \varphi_{\rho\rho}^{(1)} \) shows that, to the order of \( 1/Q \), the contributing terms are,

\[
\varphi_{\rho\rho}^{(1)} = k_\rho p_\alpha \varphi_{\rho\rho}^{(1)} + M p_\alpha sL sL S^\rho \varphi_{\rho\rho}^{(1)} + p\rho sL sL S^\rho \varphi_{\rho\rho}^{(1)},
\]

(17)

where \( \varepsilon_{\rho\rho} = \varepsilon_{\rho\rho\gamma} \delta n^\gamma n^\delta \), and \( f_1, f_2, f_{1T}, \varphi_{\rho\rho}^{(1)}, \varphi_{\rho\rho}^{(1)} \) are all Lorentz scalars and functions of \( k \cdot p \) and \( k^2 \). Note that by integrating \( \Phi_0(x) \) over \( k^2 \), we obtain the \( k \cdot p \)-dependent quark distributions \( \Phi_0(x, k^2) \) in particular at HERMES [5], are not very high, whereas the energies of the current polarized experiments discussed in connection with single-spin asymmetries. Equation of motion can relate the \( \varphi \)'s and \( \varphi \)'s to \( f_1 \)'s as,

\[
x_f = -\varphi_{\rho\rho}^{(1)} + \varphi_{\rho\rho}^{(1)}\left( \frac{y - 1}{y - 1} \right),
\]

(18)

\[
\frac{x_f}{M} f_{1T} = -\left( \varphi_{\rho\rho}^{(1)} + \varphi_{\rho\rho}^{(1)} \right)\left( \frac{y - 1}{y - 1} \right).
\]

(19)

Using the above expansion of the parton matrix elements in Eq. (16) and (17), one has the SIDIS cross section to the order of \( 1/Q \),

\[
d\sigma = \frac{2\alpha_s^2 \pi^2}{Q^4} \frac{dx d\eta dQ^2 dz}{(2\pi)^3} \int \frac{dy}{(2\pi)^4} K(k, q) \delta(x - x_B) \cdot \left\{ A(y) f_1 + \frac{B(y)}{M} f_{1T} \sin(\phi - \phi_s) \right\}.
\]

(20)

where \( \varepsilon_{\rho\rho} = \varepsilon_{\rho\rho\gamma} \delta n^\gamma n^\delta \), and \( f_1, f_2, f_{1T}, \varphi_{\rho\rho}^{(1)}, \varphi_{\rho\rho}^{(1)} \) are all Lorentz scalars and functions of \( k \cdot p \) and \( k^2 \). Note that by integrating \( \Phi_0(x) \) over \( k^2 \), we obtain the \( k \cdot p \)-dependent quark distributions \( \Phi_0(x, k^2) \) in particular at HERMES [5], are not very high, whereas the energies of the current polarized experiments discussed in connection with single-spin asymmetries. Equation of motion can relate the \( \varphi \)'s and \( \varphi \)'s to \( f_1 \)'s as,
tensor. This implies that the integration of $k^{-}f_{1T}$, $\varphi_{\perp s}^{(1)}$ and $\varphi_{\perp s}^{(1)}$ over $k^{-}$ should vanish for unrestricted range of $k^{-}$, leading to zero single-spin asymmetry for inclusive DIS. However, if we integrate over $k^{-}$ only for a restricted range in which the outgoing partons are timelike, i.e. for $0 < k^{-} < \infty$, the result might be non-zero. In practice, this is equivalent to identifying a large momentum final hadron in the photon direction to guarantee that the outgoing parton is timelike. We call such events as “triggered inclusive process” and denotes it by $e^{-}p \to e^{-} + h_{\text{trig}} + X$. The averaged single-spin asymmetry could be finite for such triggered inclusive DIS.

At the hadron level, such a single-spin asymmetry corresponds to a spin-dependent symmetric term in the hadronic tensor $W_{\mu\nu}^{(trig)}(q, p, S)$, i.e.,

$$\frac{1}{2M} W_{\mu\nu}^{(S,trig)}(q, p, S) = \frac{1}{2p_q} \left[ \epsilon_{\mu\nu\gamma}(q_{\nu} + 2x_{B}p_{\rho}) + \epsilon_{\mu\nu\gamma}(q_{\mu} + 2x_{B}p_{\rho}) \right] S^{\gamma} G_{s}(x_{B}, Q^{2}),$$

(22)

and the single-spin asymmetry $A_{N}^{(trig)}$ is given by,

$$A_{N}^{(trig)} = \frac{d\sigma^{+}}{d\sigma^{-}} - \frac{d\sigma^{-}}{d\sigma^{+}} = B(y) \frac{x_{B} M G_{s}}{Q} \frac{Q}{F_{1}},$$

(23)

where $F_{1}$ is the normal spin averaged structure function.

In terms of the parton correlations discussed above, $G_{s}$ is given by,

$$x_{B}G_{s} = -\int \frac{d^{4}k}{(2\pi)^{4}} \delta(x - x_{B}) \frac{k_{2}^{2}}{2M^{2}} \left( \varphi_{\perp s}^{(1)} - 3\varphi_{\perp s}^{(2)} \right).$$

(24)

The integration limit for $k^{-}$ is $-q^{-} < k^{-} < \infty$. Experimental measurements of the above averaged single-spin asymmetry in the triggered DIS would lead to useful information on the spin structure of nucleon.

If we neglect the final-state soft gluon interactions, i.e. set $g = 0$, the covariant derivatives in Eq. (20) become normal derivatives. We then have $\varphi_{\perp}^{(1)} = -k_{\rho}g^{(0)}$, and $\varphi_{\perp}^{(1)} = f_{1} = x f_{\perp}$, $\varphi_{\perp s}^{(1)} = \varphi_{\perp s}^{(1)} = 0$, $\varphi_{\perp s}^{(2)} = 0$. Hence, only $A(y) f_{1} + 2B(y)(k_{\perp} / Q) f_{1} \cos \phi$ are left in Eq. (20) which is the result in Ref. [19].

The difference between this and the full result given by Eq.(20) comes from the final-state soft gluon interactions.

Data on $\langle \cos \phi \rangle$ obtained in unpolarized experiments [1], [2], [4] suggest the existence of such QCD contributions. It is also obvious that fragmentation can contribute to the azimuthal asymmetries discussed above. A complete formalism for SIDIS should include fragmentation. Such an extension is underway.

In summary, we derived the factorized form of the cross section for semi-inclusive deep-inelastic lepton-nucleon scattering as a convolution of the hard parts with the gauge invariant unintegrated parton distributions and higher twist parton correlations. As a consequence of the collinear expansions, the hard parts depend only on the longitudinal components of the parton momenta. The

gauge invariant parton distributions receive contributions from initial and final state interactions and provide the only dependence on the initial parton transverse momentum. Results for azimuthal angle dependence to the order of $1/Q$ in reactions with transversely polarized targets are given. A single-spin asymmetry for the triggered inclusive process $e^{-}p \to e^{-} + h_{\text{trig}} + X$ is predicted to the order of $M/Q$.

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[19] We note that, in [11], the authors started with an expression where the hard part contains transverse components which lead to double-counting of the transverse contributions. But they have either the projection operator $\omega_{\rho}^{\rho}$ or the $\partial_{\rho}$-term in $\varphi^{(1)}$. This is why they came also to this result at $g = 0$. 