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Abstract

We study the transverse momentum dependent factorization for single spin asymmetries in Drell-Yan and semi-inclusive deep inelastic scattering processes at one-loop order. The next-to-leading order hard factors are calculated in the Ji-Ma-Yuan factorization scheme. We further derive the QCD resummation formalisms for these observables following the Collins-Soper-Sterman method. The results are expressed in terms of the collinear correlation functions from initial and/or final state hadrons coupled with the Sudakov form factor containing all order soft-gluon resummation effects. The scheme independent coefficients are calculated up to one-loop order.

Single transverse spin asymmetry in hadronic reactions have attracted great attentions from both experiment and theory sides in recent years. It promises strong connection to the three-dimensional partonic tomography of the nucleon, and provides unique opportunities to study QCD dynamics, such as the factorization and universality of parton distributions. In particular, the short distance partonic interactions resulting into different Wilson lines in the relevant transverse momentum dependent (TMD) parton distributions between the Drell-Yan lepton pair production in *pp* collisions and the semi-inclusive hadron production in deep inelastic scattering (SIDIS) predicts opposite signs for the Sivers-type single transverse spin asymmetries in these two processes [1, 2]. This nontrivial universality prediction has stimulated strong interests around the world to measure, especially, the SSA in the Drell-Yan (DY) process, since that in the SIDIS has been observed in various experiments [3].

However, the numerical predictions for the SSAs in these processes are all based on a leading order naive TMD factorization [4]. Although a general TMD factorization has been argued for these processes [5–9], we need to know the next-to-leading-order (NLO) perburbative QCD corrections to have more reliable predictions, which have not yet been calculated. Another important issue is the energy dependence. Current DIS experiments cover the Q^2 range about 2-5GeV^2, whereas the planned DY process will be measured at relative larger Q^2 about 20-25GeV^2. Here, Q represents the large momentum scale, i.e., the virtuality of the photon in these processes. To accurately describe the Q^2 evolution of the transverse momentum dependent observables, the QCD resummation effects have to be taken into account [6].

In this paper, we will build a theoretical framework to address these important questions. We carry out, at the first time, the complete NLO perturbative correction to the single transverse spin dependent cross sections in DY and SIDIS processes in the TMD factorization. One of the important implications of the explicit one-loop calculations is to help to construct the *correct* resummation formalism for the SSA observables. Earlier attempts to formulate these effects in SSAs have been made in various forms [10, 11]. A resummation formula close to ours was used in Ref. [10], and a significant suppression effect were found when Q^2 is very large. In the following, based on the one-loop calculation results, we will derive the complete soft gluon resummation formalism following the Collins-Soper-Sterman (CSS) method [6], which can be easily implemented in the phenomenological studies.

We take the DY process as an example to demonstrate our procedure and present the main results,

$$A(P_A, S_\perp) + B(P_B) \to \gamma^*(q) + X \to \ell^+ + \ell^- + X, \tag{1}$$

where P_A and P_B represent the momenta of hadrons A and B, and $S_\perp$ for the transverse polarization vector of A, respectively. The single transverse spin dependent differential cross section can be expressed as

$$\frac{d\Delta\sigma(S_\perp)}{dydQ^2d^2q_\perp} = \sigma_0 \epsilon^{\alpha\beta} S_\perp^\alpha W_{UT}^\beta(Q; q_\perp) \ , \tag{2}$$

where $q_\perp$ and y are transverse momentum and rapidity of the lepton pair, $\sigma_0 = 4\pi\alpha_{em}^2/3N_c sQ^2$ with $s = (P_A + P_B)^2$, and $\epsilon^{\alpha\beta}$ is defined as $\epsilon^{\alpha\beta\mu\nu} P_{A\mu} P_{B\nu}/P_A \cdot P_B$. At low transverse momentum ($q_\perp \ll Q$) the structure function W_{UT} can be formulated in terms of the TMD factorization where the quark Sivers function is involved [1, 2], whereas at large transverse momentum ($q_\perp \gg \Lambda_{\rm QCD}$) it can be calculated in the collinear factorization approach in terms of the twist-three quark-gluon-quark correlation functions [12–14]. It has been shown that the TMD and collinear twist-three approaches give the consistent results in the intermediate transverse momentum region: $\Lambda_{\rm QCD} \ll q_\perp \ll Q$ [14, 15]. This consistency allows us to separate W_{UT} into two terms [6],

$$W^\alpha_{UT}(Q;q_\perp) = \int \frac{d^2b}{(2\pi)^2} e^{i\vec{q}_\perp\cdot\vec{b}} \widetilde{W}^\alpha_{UT}(Q;b) + Y^\alpha_{UT}(Q;q_\perp) \ ,$$

where the first term dominates at $q_\perp \ll Q$ region, and the second term at $q_\perp \sim Q$. The latter is obtained by subtracting the low $q_\perp$ expansion from the full perturbative calculation.

According to the TMD factorization, we have [7, 11],

$$\widetilde{W}^\alpha_{UT}(Q;b) = \tilde{f}^{(\perp\alpha)}_{1T}(z_1,b,\zeta_1)\bar{q}(z_2,b,\zeta_2) \times H_{UT}(Q)\left(S(b,\rho)\right)^{-1} \ , \tag{3}$$

where $z_{1,2} = Q/\sqrt{s}e^{\pm y}$ and the sums over flavor weighting with the charge squared of the quarks are implicit, $f^\perp_{1T}$ and $\bar{q}$ are the TMD quark Sivers function of A and antiquark distribution of B, H and S are hard and soft factors, respectively. In this paper, we follow the Ji-Ma-Yuan factorization scheme [7], where two off-light-front Wilson lines (along off-light-front vectors v_1 and v_2) are introduced to regulate the light-cone singularities for the TMD quark distributions. We further define $\zeta_1^2 = 4(v_1\cdot P_A)^2/v_1^2$ and $\zeta_2^2 = 4(v_2\cdot P_B)^2/v_2^2$, and the rapidity cut-off parameter ρ: $z_1^2\zeta_1^2 = z_2^2\zeta_2^2 = \rho Q^2$. The Sivers function in the impact parameter $b_\perp$-space is defined as $\tilde{f}^{(\perp\alpha)}_{1T}(x,b_\perp) = \int d^2k_\perp e^{-i\vec{k}_\perp\cdot\vec{b}_\perp} k^\alpha_\perp f^{\perp({\rm DY})}_{1T}(x,k_\perp)/M_P$, where $f^\perp_{1T}$ follows the definition of Ref. [16]. Our results can be translated to other factorization schemes where different regularizations for the light-cone singularities are used [9].

We investigate the above factorization formula Eq. (3) in the perturbative region of $1/b_\perp \gg \Lambda_{\rm QCD}$. The explicit calculations at one-loop order will verify the TMD factorization, from which we obtain the NLO correction to the hard factor. In the calculation of the spin-average (or double spin asymmetry) TMD factorization, it is convenient to choose a quark (or gluon) target, where every factor in the factorization formula can be calculated perturbatively [7]. However, the SSA vanishes with on-shell quark. To get nonzero effect, we have to go beyond the simple quark target picture.

Our calculations are based on the collinear correlation functions from the incoming hadrons. In this framework, the SSA is naturally a twist-three effect, and involves the twist-three quark-gluon-quark correlation function from the polarized nucleon [12, 13]. In particular, the transverse spin-dependent Qiu-Sterman matrix element $T_F(x_1,x_2)$ [13] is responsible for the SSA in the process of (1), and is related to the quark Sivers function: $T_F(x,x) = \int d^2k_\perp |k_\perp^2| f^{\perp({\rm DY})}_{1T}(x,k_\perp)/M_P$ [14, 17]. There have been great theoretical developments in this framework in recent years, see, e.g., Refs. [14, 17–20]. We will utilize these techniques to compute the structure function $\widetilde{W}^\alpha_{UT}(Q,b)$. First, let us write down a general form,

$$\widetilde{W}^\alpha_{UT}(Q,b) = \left(\frac{-ib^\alpha_\perp}{2}\right)\int \frac{dx_1dx_2dx'}{x_1x_2x'} T_F(x_1,x_2)\bar{q}(x') \times\mathcal{H}(x_1,x_2,x';Q,b) \ , \tag{4}$$

where T_F follows the definition of Ref. [17] and $\bar{q}(x')$ is the integrated anti-quark distribution. Other twist-three quark-gluon-quark correlation functions contributing to the SSA can be included as well. For simplicity, we focus on T_F contributions in this paper, including both

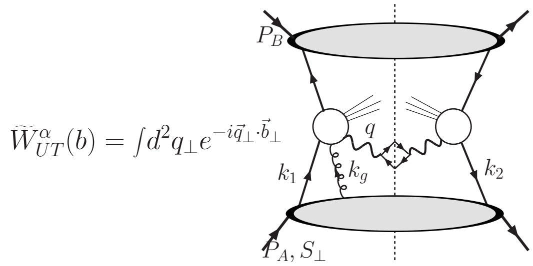

FIG. 1. *Generic Feynman diagram contribution to the impact parameter space structure function* $\widetilde{W}_{UT}(Q,b)$.

soft-gluon-pole and hard-gluon-pole contributions. The soft-fermion-pole contributions can be included accordingly.

The above formula is similar to that of the $q_\perp$-weighted asymmetry in the same process calculated in Ref. [17]. There are two major differences: (1) instead of weighting with $q_\perp^\alpha$, here we weight with $e^{-i\vec{q}_\perp\cdot\vec{b}_\perp}$; (2) we will restrict our calculations of the differential cross section in the TMD domain, i.e., $q_\perp \ll Q$ ($1/b_\perp \ll Q$). We calculate Eq.(4) in a covariant gauge, where a generic diagram is illustrated in Fig. 1. Because of the twist-three effect, an additional gluon attachment has to be taken into account for hard partonic scattering amplitude. A collinear expansion will be performed to calculate the hard part $\mathcal{H}(Q,b)$. In particular, the amplitude is expanded in terms of $k_{g\perp}^\alpha = k_{2\perp}^\alpha - k_{1\perp}^\alpha$ where $k_{i\perp} \ll 1/b_\perp$. Combining this expansion with matrix element from the polarized nucleon, it will lead to the spin dependent cross section expressed in terms of $T_F(x_1,x_2)$.

The leading order diagrams are shown in Fig. 2(a) and (b). From the kinematics, we find that $q_\perp$ is related to the transverse momenta of the two quark lines as: $q_\perp = k_{2\perp}$ for Fig. 2(a) and $q_\perp = k_{1\perp}$ for Fig. 2(b). Therefore, the contributions from these two diagrams will be

$$\begin{aligned}
\text{Fig.2} &= \int d^2q_\perp e^{-i\vec{q}_\perp\cdot\vec{b}_\perp}\left(\frac{ig}{-(k_2^+-k_1^+)-i\epsilon}\right) \\
&\quad\times\left[\delta(q_\perp-k_{2\perp})-\delta(q_\perp-k_{1\perp})\right] \\
&= \left(\frac{ig}{-(k_2^+-k_1^+)-i\epsilon}\right)\left[e^{-i\vec{k}_{2\perp}\cdot\vec{b}_\perp}-e^{-i\vec{k}_{1\perp}\cdot\vec{b}_\perp}\right] \\
&= \frac{ig}{-(k_2^+-k_1^+)-i\epsilon}\left(-ib_\perp^\alpha\right)k_{g\perp}^\alpha\ ,
\end{aligned} \tag{5}$$

where the last equation comes from the collinear expansion of the exponential factor. The initial state interactions represented by the propagator leads to a pole contribution. After taking the pole, we will obtain the leading order contribution,

$$\widetilde{W}_{UT}^{\alpha(0)}(Q,b) = \left(\frac{-ib_\perp^\alpha}{2}\right)T_F(z_1,z_1)\bar{q}(z_2)\ . \tag{6}$$

This also normalizes the leading order hard factor as $H_{UT}^{(0)} = 1$ in Eq. (3), because $f_{1T}^{\perp\alpha(\mathrm{DY})}(z_1,b_\perp) = T_F(z_1,z_1)(-ib_\perp^\alpha/2)$ at this order at small $b_\perp$.

Order α_s corrections come from real and virtual gluon radiation contributions. Similar to that calculated in Ref. [17], the virtual diagrams contain soft and collinear divergences.

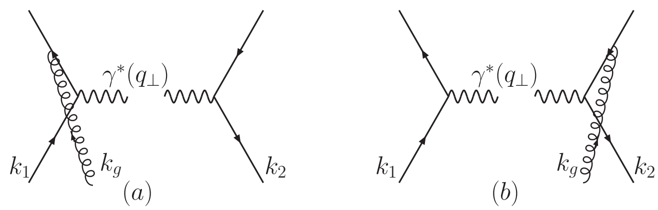

FIG. 2. *Leading order Born diagram calculation of* $\widetilde{W}_{UT}(Q, b)$.

The soft divergence will be cancelled out by the real gluon radiation contributions, and the collinear divergences will be absorbed into the parton distributions in both TMD and collinear factorizations in Eqs. (3) and (4). The real diagrams are similar to those calculated in Refs. [14, 17]. In order to obtain their contributions to $\widetilde{W}_{UT}^{\alpha}(b_\perp)$, we need to first take the low transverse momentum limit, and then Fourier transform into the impact parameter $b_\perp$-space. After summing up both real and virtual diagrams contributions, we indeed find that the soft divergences cancel out between each other, and the total contribution only contains the collinear divergences,

$$\frac{\alpha_s}{2\pi}\left(\frac{-ib_\perp^\alpha}{2}\right)\left\{\left[-\frac{1}{\epsilon}+\ln\frac{4}{b^2\mu^2}e^{-2\gamma_E}\right]\left(\mathcal{P}_{q/q}\otimes\bar{q}(z_2')\right.\right.$$
$$\left.+\mathcal{P}_{qg\to qg}^T\otimes T_F(z_1',z_1'')\right)+C_F(1-\xi_2)\delta(1-\xi_1)$$
$$+\left(-\frac{1}{2N_c}\right)(1-\xi_1)\delta(1-\xi_2)+\delta(1-\xi_1)\delta(1-\xi_2)$$
$$\left.\times C_F\left[-\ln^2\left(\frac{Q^2b^2}{4}e^{2\gamma_E-\frac{3}{2}}\right)-\frac{23}{4}+\pi^2\right]\right\}\ , \quad (7)$$

where $\xi_i = z_i/z_i'$, $\mathcal{P}_{qq}$ is the quark splitting function, and $\mathcal{P}_{qg\to qg}^T$ the splitting function for the Qiu-Sterman matrix element. Our calculations provide an important cross check for this splitting kernel [17, 20]. We will present the detailed comparison with previous calculations in a separate publication. After subtracting the collinear divergences from the splitting of T_F and $\bar{q}$, we demonstrate the factorization form in Eq. (4) up to one-loop order.

The above result can also be casted into the TMD factorization formula Eq. (3), where we have to subtract the TMD quark Sivers function, antiquark distribution and the soft factor. The latter two have been calculated before [7]. The quark Sivers function can be calculated similarly, and we find that,

$$\tilde{f}_{1T}^{\alpha}(z_1,b_\perp)=\frac{\alpha_s}{2\pi}\left(\frac{-ib_\perp^\alpha}{2}\right)\left\{\left[-\frac{1}{\epsilon}+\ln\frac{4}{b^2\mu^2}e^{-2\gamma_E}\right]\right.$$
$$\times\mathcal{P}_{qg\to qg}^T\otimes T_F(z_1',,z_1'')+\delta(1-\xi_1)C_F$$
$$\left[-\frac{3}{2}\ln\frac{4}{b^2\mu^2}e^{-2\gamma_E}-\frac{1}{2}\ln^2\left(\frac{z_1^2\zeta_1^2b^2}{4}e^{2\gamma_E-1}\right)-\frac{3+\pi^2}{2}\right]$$
$$\left.+\left(-\frac{1}{2N_c}\right)(1-\xi_1)\right\}\ . \quad (8)$$

After subtracting these factors out, we find that the hard factor H_{UT} is free of infrared divergence, and it is the same as that for the spin average one calculated in Ref. [7],

$$
\begin{aligned}
H_{UT}^{(1)\mathrm{DY}} &= H_{UU}^{(1)}|_{\mathrm{DY}} \\
&= \frac{\alpha_s}{2\pi} C_F \left[\ln \frac{Q^2}{\mu^2}(1+\ln\rho^2) - \ln\rho^2 + \ln^2\rho + 2\pi^2 - 4 \right],
\end{aligned}
$$

which is very interesting and suggests that the hard factors are spin-independent. Following the same procedure, we calculate the hard factor for the single transverse spin dependent cross section in SIDIS. Again, it is the same as the spin-average case, $H_{UT}^{(1)}|_{\mathrm{SIDIS}} = H_{UU}^{(1)}|_{\mathrm{SIDIS}}$ [7]. Clearly, the hard factors for DY and SIDIS differ by a factor of π^2. The hard factors in other TMD factorization scheme can be calculated similarly.

The factorization formula Eq. (3) contains large logarithms in terms of $\ln^2(Q^2b^2)$ [11] which is also shown in the total result of Eq. (7). Follow the CSS method [6], one could resum these large logarithms. Here it is important to realize that Sivers function in b-space $\tilde{f}_{1T}^{(\perp\alpha)}(z_1, b, \zeta_1)$ obeys the same energy evolution equation as the spin-averaged quark distribution [11]. Solving these evolution equations, we will have

$$
\begin{aligned}
\widetilde{W}_{UT}^{\alpha}(Q;b) &= e^{-\mathcal{S}_{UT}(Q^2,b)} \widetilde{W}_{UT}^{\alpha}(C_1/b, b) \\
&= (-ib_\perp^\alpha/2)\, e^{-\mathcal{S}_{UT}(Q^2,b)} \Sigma_{i,j} \\
&\quad \times \Delta C_{qi}^T \otimes f_{i/A}^{(3)}(z_1', z_1'') C_{\bar{q}j} \otimes f_{j/B}(z_2') , \qquad (9)
\end{aligned}
$$

where $f_{j/B}$ represents the integrated parton distribution from hadron B, and $f_{i/A}^{(3)}$ the twist-three function from hadron A, of which $T_F(z_1, z_2)$ is the most relevant one as $T_F(z_1, z_1)$ appears in the leading order contribution in Eq. (6). The last step of the above equation comes from further applying the collinear factorization formula Eq. (4) at lower energy scale C_1/b. From the above one-loop calculations, we find that the perturbative Sudakov factor $\mathcal{S}_{UT}$ have the same form as that for the spin-average case,

$$
\begin{aligned}
\mathcal{S}_{UT}(Q^2, b) = \int_{C_1^2/b^2}^{C_2^2Q^2} \frac{d\mu^2}{\mu^2} &\left[\ln\left(\frac{C_2^2 Q^2}{\mu^2}\right) A_{UT}(C_1; g(\mu)) \right. \\
&\left. + B_{UT}(C_1, C_2; g(\mu)) \right] , \qquad (10)
\end{aligned}
$$

up to one-loop order, where C_1 and C_2 are constants in the order of 1.

The A, B and C functions can be calculated in perturbation theory: $A = \sum_{n=1} A^{(n)}(\alpha_s/\pi)^n$. From the explicit one-loop calculations, we obtain the following results for these coefficients,

$$
\begin{aligned}
A_{UT}^{(1)} = C_F,\ B_{UT}^{(1)} = -3/2 C_F,\ \Delta C_{qq}^{T(0)} = \delta(1-x) , \\
\Delta C_{qq}^{T(1)} = -\frac{1}{4N_c}(1-x) + \frac{C_F}{2}\delta(x-1)\left[\frac{\pi^2}{2} - 4\right] , \qquad (11)
\end{aligned}
$$

where C_{qq} follows the spin-average case [6], and we have chosen the canonical values for $C_1 = 2e^{-\gamma_E}$ and $C_2 = 1$ to simplify the above expressions. For SIDIS, A, B remain the same, whereas ΔC^T have opposite sign and there is no π^2 term in $\Delta C_{qq}^{T(1)}$. The above coefficients can also be calculated by comparing the fixed order calculations of the differential cross section depending on transverse momentum [14] to the expansion of the resummation formula Eq. (9). We have checked that this gives the consistent results.

Eq.(9) is our main result for the QCD resummation for the single spin asymmetry in the DY process. The structure, in particular, the pre-factor $(ib_\perp^\alpha)$ is the unique feature for the single transverse spin azimuthal angel dependent cross section. This pre-factor comes from the explicit one-loop calculation above. It also guarantees the proper behaviors for the spin asymmetry at small and large transverse momentum. At very small $q_\perp$, the asymmetry has to vanish, so that the differential cross section and W_{UT}^α will be proportional to $q_\perp^\alpha$. On the other hand, at large $q_\perp$, the asymmetry is power suppressed by $1/q_\perp$ from the perturbative calculation, which leads to $W_{UT}^\alpha \propto q_\perp^\alpha/q_\perp^4$. With the correct coefficients extracted from our one-loop calculation, these behaviors are satisfied in Eq. (9). We would like to emphasize that the TMD and collinear factorizations for $\widetilde{W}_{UT}^\alpha(b)$ in Eqs. (3) and (4) are crucial to obtain the final resummation formula of Eq. (9). Without these factorization results, we can not apply the CSS resummation.

Another important feature is that Eq. (9) depends on the integrated parton distributions, which are universal. The opposite sign for the SSAs in DY and SIDIS is reflected by the opposite C coefficients in this formula. Moreover, our resummation formula is scheme-independent, although the TMD factorization Eq. (3) depends on the scheme of how to regulate the light-cone singularities in the TMD distributions. This can be clearly seen from the disappearance of ρ and ζ_i^2 in Eq. (9) with the above coefficients.

In summary, we have derived the CSS resummation formalism for the single spin asymmetries in DY and SIDIS processes. The relevant coefficients are calculated up to one-loop order. These results shall be further studied to understand the energy dependence of the SSAs in these processes, and provide more accurate predictions for the DY process which is actively pursued by several experiments. Our results shall shed light on all other $k_\perp$-odd observables, and should be applied to the azimuthal angular dependent observables in DY, SIDIS, and e^+e^- annihilation processes. We performed our calculations in the framework of the collinear correlation functions of hadrons. This allows us to compute the cross sections in a consistent and rigorous way. We noticed that recently, a different framework has been developed where a gluonic degree of freedom is included in the leading order quark target [21]. It will be interesting to apply this method and compare with our results.

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