# Unified Picture for Single Transverse-Spin Asymmetries in Hard-Scattering Processes

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Using lepton-pair production in hadron-hadron collisions as an example, we explore the relation between two well-known mechanisms for single-transverse-spin asymmetries in hard processes: twist-three quark-gluon correlations when the pair's transverse momentum is large, $q_\perp \gg \Lambda_{\rm QCD}$, and time-reversal-odd and transverse-momentum-dependent parton distributions when $q_\perp$ is much less than the pair's mass. We find that, although the two mechanisms each have their own domain of validity, they describe the same physics in the kinematic region where they overlap. This unifies the two mechanisms and imposes an important constraint on phenomenological studies.

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*I. Introduction.*—When two objects collide, one with spin and another without, the result most likely depends on the direction of the spin. In high-energy scattering of two hadrons (protons, for instance), when one of them has a nonzero transverse polarization, a strong spin dependence has been observed since three decades ago [1]. The physical quantity of interest is the so-called single-transverse-spin asymmetry (SSA) $A_N \equiv (\sigma(S_\perp) - \sigma(-S_\perp))/(\sigma(S_\perp) + \sigma(-S_\perp))$, defined as the ratio of the difference and the sum of the cross sections when the spin $S_\perp$ is flipped. SSAs in hadronic physics have attracted much interest in recent years in both experiment and theory, particularly after publication of data [2] by the HERMES Collaboration at Deutsches Elektronen-Synchrotron and the polarized Relativistic Heavy-Ion Collider at Brookhaven National Laboratory coming into operation. Although it was realized some time ago [3] that perturbative quantum chromodynamics (QCD) can be used to study the effects of transverse spin, the size of the observed asymmetries came as a surprise and has posed a challenge for researchers in this field [4].

Two mechanisms have been proposed in QCD to explain the observed large size of SSAs. One follows the collinear (CO) QCD factorization approach, attributing the SSAs to spin-dependent twist-three quark-gluon correlation functions and to a quantum interference between different partonic scattering amplitudes [Efremov-Teryaev-Qiu-Sterman (ETQS) mechanism] [5,6]. The other explicitly connects the SSAs to spin dependence of the transverse motions of quarks and gluons in a polarized proton and expresses the SSAs in terms of time-reversal-odd ($T$-odd) and transverse-momentum-dependent (TMD) parton distributions [7]. Much progress has been made in understanding the TMD parton distributions and their gauge properties [8–13] and in establishing corresponding QCD factorization formulas [8,9,14,15]. Both mechanisms have been used in phenomenological studies of the available data [6,16]. Recently, a relation between the two types of parton distribution functions was derived [13,17], indicating a certain dynamical connection between the two mechanisms. However, a clear relationship between the two apparently different physical mechanisms, particularly at the level of physical observables, was not established so far. This will be achieved in this Letter.

In this Letter, we explore the connection between the two mechanisms by studying the SSA in Drell-Yan hadronic lepton-pair production. We consider the scattering of a transversely polarized proton of spin $S_\perp$ and momentum $P$ on an unpolarized hadron (another proton, for example) of momentum $P'$, producing a virtual photon that subsequently decays into a pair of leptons with invariant mass $Q$, transverse momentum $q_\perp$, and a positive rapidity $y$ (in the forward direction of the polarized proton). We calculate the *spin-dependent* differential cross section $d\Delta\sigma(S_\perp)/dQ^2dyd^2q_\perp$, with $\Delta\sigma(S_\perp) = [\sigma(S_\perp) - \sigma(-S_\perp)]/2$, for both mechanisms: the collinear QCD factorization formalism in Sec. II and the TMD QCD factorization formalism in Sec. III, respectively. The collinear QCD factorization formalism for the SSAs works when both $q_\perp, Q \gg \Lambda_{\rm QCD}$, the strong interaction scale, and the TMD QCD factorization formalism is valid when $q_\perp \ll Q$. Therefore, we expect that, in the common kinematic region $\Lambda_{\rm QCD} \ll q_\perp \ll Q$, the two mechanisms coincide and, hence, describe the same physics. In Sec. IV, we expand the spin-dependent cross section calculated in the collinear QCD factorization formalism at small $q_\perp/Q$ and find that the leading-order result in $q_\perp/Q$ is indeed the same as that obtained within the TMD QCD factorization formalism. This conclusion imposes a rigorous constraint on phenomenological studies of the SSA data: If the data can be described in both ways, the two mechanisms must produce identical results.

*II. Collinear QCD factorization.*—When $q_\perp, Q \gg \Lambda_{\rm QCD}$, the parton momenta entering the hard scattering can be approximated to be parallel to either of the incoming hadron's light-cone momenta $P^+$ or $P'^-$ [the light-cone components of a space-time vector $v^\mu$ are defined as $v^\pm = (v^0 \pm v^3)/\sqrt{2}, v_\perp = (v^1, v^2)$]. The lepton-pair yield is proportional to the square of the scattering amplitude. The latter is equal to a sum of all possible partonic scattering amplitudes, including, for example, the antiquark-quark annihilation in Fig. 1(a) and the same annihilation in the presence of a polarized "color electromagnetic" field (the gluon) in Fig. 1(b). The spin-dependent cross section in the collinear approximation results mainly from a nonvanishing quantum interference between the scattering amplitudes in Figs. 1(a) and 1(b) [5,6]. At large $y$, it can be calculated in terms of a twist-three quark-gluon correlation, the shaded oval in the lower part of the diagram in Fig. 1(c). The QCD expression for the correlation is [6]

$$T_F(x_1, x_2) = \int \frac{d\xi^- d\eta^-}{4\pi} e^{i(k_{q1}^+ \eta^- + k_g^+ \xi^-)} \epsilon_\perp^{\beta\alpha} S_{\perp\beta} \langle PS|\bar{\psi}(0) \times \mathcal{L}(0, \xi^-)\gamma^+ g F_\alpha^+(\xi^-) \mathcal{L}(\xi^-, \eta^-)\psi(\eta^-)|PS\rangle, \quad (1)$$

where $x_1 = k_{q1}^+/P^+$ and $x_2 = k_{q2}^+/P^+$ are the fractions of the polarized proton's light-cone momentum carried by the quark in Fig. 1, while $x_g = k_g^+/P^+ = x_2 - x_1$ is the fractional momentum carried by the gluon. $\epsilon_\perp^{\beta\alpha}$ is the two-dimensional Levi-Civita tensor with $\epsilon_\perp^{12} = 1$. In Eq. (1), $\psi$ and $F$ are quark field and gluon field strength, respectively, and $\mathcal{L}$ is a gauge link operator that makes the correlation gauge invariant [6]. The spin dependence of this correlation can be seen as follows: When a transversely polarized proton is traveling at nearly the speed of light, its internal color electric and magnetic fields have preferred orientations in the transverse plane. By parity invariance, the color electric field must be orthogonal to the spin of the proton. If averaged over the proton wave function, the field vanishes because the proton is color-neutral (also because of time-reversal symmetry). However, if one multiplies the color electric field with the quark color current, the average may be nonzero. This average defines a quark-gluon correlation function, and its spin dependence characterizes a property of a polarized proton.

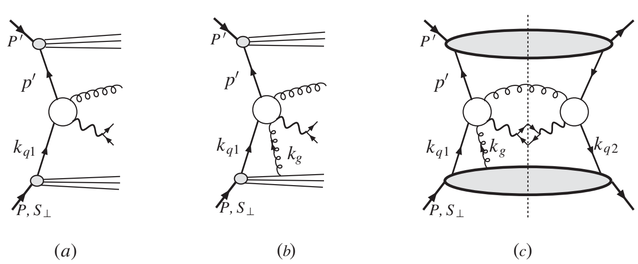


FIG. 1. Scattering amplitudes for Drell-Yan dilepton production via (a) $q + \bar{q} \to \gamma^* + g$ and (b) $q + \bar{q} + g \to \gamma^* + g$. (c) A typical diagram, from the interference of the amplitudes in (a) and (b), that gives a contribution to the SSA.

A nonvanishing SSA requires the presence of a nontrivial strong interaction phase, because it is proportional to $\vec{S} \cdot \vec{p} \times \vec{q}$, which is odd under a naive time-reversal transformation. For the ETQS mechanism, the phase arises from the interference between the two scattering amplitudes in Figs. 1(a) and 1(b). Already at the lowest nontrivial order in the strong coupling constant, the amplitude with the extra gluon in Fig. 1(b) may have an imaginary part. It is generated by the poles (or on-shell conditions) of parton propagators in the diagram, when the integration over the gluon's momentum $x_g$ is performed. If such a pole occurs at $x_g = 0$, it is called a "soft" (gluon) pole [6]; otherwise, it is referred to as a "hard" pole [18,19].

Summing over all contributions by the leading-order diagrams of the type shown in Fig. 1(c) and their complex conjugates, we derive, within the collinear QCD factorization, the spin-dependent cross section [19]

$$\frac{d^3\Delta\sigma_{\rm CO}^{q\bar{q}\to\gamma^* g}(S_\perp)}{dQ^2 dy d^2q_\perp} = \int \frac{dx}{x}\frac{dx'}{x'} \bar{q}(x')\delta(\hat{s} + \hat{t} + \hat{u} - Q^2) \times C\left\{\frac{1}{2N_c}\left[\left(x\frac{\partial}{\partial x}T_F(x, x)\right)\frac{D_{q\bar{q}}^s}{-\hat{u}} + T_F(x, x)\frac{N_{q\bar{q}}^s}{-\hat{u}}\right] + T_F(x, x - \bar{x}_g) \times \frac{N_{q\bar{q}}^h}{-\hat{u}}\left[\frac{1}{2N_c} + C_F\frac{\hat{s}}{\hat{s} + \hat{u}}\right]\right\}, \quad (2)$$

where the sum over all quark flavors, weighted with their electric charge squared, is implicit. The factor $C$ is $\sigma_0\epsilon^{\alpha\beta}S_{\perp\alpha}q_{\perp\beta}\alpha_s/2\pi^2$, with $\sigma_0 = 4\pi\alpha_{\rm em}^2/3N_c sQ^2$, $N_c = 3$, and $s = (P + P')^2$. $x'$ is the momentum fraction of the unpolarized quark, and $x$ is the total parton momentum fraction from the polarized proton when $x_g$ is fixed using the pole condition. The partonic Mandelstam variables are defined as $\hat{s} = (xP + x'P')^2$, $\hat{t} = (xP - q)^2$, and $\hat{u} = (x'P' - q)^2$. In Eq. (2), the terms proportional to $T_F(x, x)$ and $T_F(x, x - \bar{x}_g)$, with $\bar{x}_g = -x\hat{t}/(Q^2 - \hat{t})$, correspond to the soft and hard pole contributions, respectively. The coefficients associated with them are $D_{q\bar{q}}^s = \hat{u}/\hat{t} + \hat{t}/\hat{u} + 2Q^2\hat{s}/(\hat{u}\hat{t})$, $N_{q\bar{q}}^s = [Q^2(\hat{u}^2 - \hat{t}^2) + 2Q^2\hat{s}(Q^2 - 2\hat{t}) - (\hat{u}^2 + \hat{t}^2)\hat{t}]/\hat{t}^2\hat{u}$, and $N_{q\bar{q}}^h = [(Q^2 - \hat{t})^3 + Q^2\hat{s}^2]/(-\hat{t}^2\hat{u})$. In the real-photon limit $Q^2 = 0$, the above result reduces to the spin-dependent direct-photon cross section.

*III. TMD QCD factorization.*—When $q_\perp \ll Q$, the observed $q_\perp$ could be sensitive to the transverse momenta of the scattering partons, while $Q$ sets the scale of the hard collision. The Drell-Yan cross section in leading order $q_\perp/Q$ can then be factorized in terms of TMD quark distributions. When an antiquark of the unpolarized proton scatters off a quark with a nonvanishing transverse momentum $k_\perp$ from the polarized proton, the lepton-pair yield is proportional to a spin-dependent TMD quark distribution $q(x, k_\perp, S_\perp)$, and the spin-dependent cross section is then proportional to the difference $q(x, k_\perp, S_\perp) - q(x, k_\perp, -S_\perp)$. This difference, after factoring out the explicit spin dependence, is defined to be a new spin-dependent TMD quark distribution [19] $q_T(x, k_\perp)$, which was originally proposed by Sivers [7] and is often referred to as the "Sivers function." It describes an asymmetric dependence of the quark density on the relative orientation of the quark transverse momentum $\vec{k}_\perp$ and the direction of the proton spin, and it may generate a nonvanishing SSA. $q_T(x, k_\perp)$ would vanish, however, without a gauge link connecting the quark field to infinity [10–13] along a certain direction $v^\mu$, which is usually chosen to be off the light cone to avoid light-cone singularities. The gauge link, which is required by gauge invariance, describes the eikonal phase accumulated through the quark propagation in the background field produced by gluons moving collinear to the proton's momentum. Thus, in the TMD framework, the strong interaction phase required for SSAs is already included in $q_T(x, k_\perp)$. Because of Lorentz symmetry, the $v$ dependence of the TMD distributions is of the form $\zeta^2 = (2v \cdot P)^2/v^2$ [8].

In terms of the $q_T(x, k_\perp)$, the spin-dependent Drell-Yan cross section was shown to have the following TMD factorization [14]:

$$\frac{d^3\Delta\sigma_{\rm TMD}(S_\perp)}{dQ^2dyd^2q_\perp} = \sigma_0\frac{\epsilon^{\alpha\beta}S_{\perp\alpha}q_{\perp\beta}}{M_P}\int d^2\vec{k}_{1\perp}d^2\vec{k}_{2\perp}d^2\vec{\lambda}_\perp\frac{\vec{k}_{1\perp}\cdot\vec{q}_\perp}{\vec{q}_\perp^{\,2}}\delta^{(2)}(\vec{k}_{1\perp}+\vec{k}_{2\perp}+\vec{\lambda}_\perp-\vec{q}_\perp)q_T(z_1,k_{1\perp})\bar{q}(z_2,k_{2\perp}) \times (S(\lambda_\perp))^{-1}H, \tag{3}$$

where $z_1 = Q/\sqrt{s}e^y$ and $z_2 = Q/\sqrt{s}e^{-y}$ are the momentum fractions of the colliding hadrons associated with the observed lepton pair. $H$ is a hard factor and is entirely perturbative. $\bar{q}$ is the TMD antiquark distribution of the unpolarized proton. $M_P$ is a hadron mass, used to normalize the $q_T$ and $\bar{q}$ TMD distributions to the same mass dimension. The soft factor $S$ is a vacuum matrix element of Wilson lines and captures the effects of soft gluon radiation with a total transverse momentum $\lambda_\perp$ [8].

In order to compare the two mechanisms for the SSAs, we need to calculate the explicit $q_\perp$ dependence of the TMD factorization in perturbative QCD in the region $\Lambda_{\rm QCD} \ll q_\perp \ll Q$. At leading order in $q_\perp/Q$, the $q_\perp$ dependence in Eq. (3) is approximately generated by a sum of various contributions, each of which is represented by one of the reduced diagrams shown in Fig. 2 [9]. This factorization property can be pictured from the diagram in Fig. 1(c) by considering the real gluon emitted along the direction of either $P$ or $P'$. For calculating each contribution in Fig. 2, we let one of the transverse momenta $\vec{k}_{i\perp}$ and $\vec{\lambda}_\perp$ in the $\delta$ function in Eq. (3) be of the order of $\vec{q}_\perp$ and the others small. When $\vec{\lambda}_\perp$ is large, for example, in Fig. 2(c), we neglect the $\vec{k}_{i\perp}$ in the $\delta$ function. The integrations over these momenta yield the ordinary antiquark distribution $\int d^2\vec{k}_\perp\bar{q}(x,k_\perp) = \bar{q}(x)$ and the diagonal part of the twist-three quark-gluon correlation $\int d^2\vec{k}_\perp(\vec{k}_\perp^2/M_P)q_T(x,k_\perp) = T_F(x,x)$, a relation first derived in Ref. [13] (where the $q_T$ was referred to as $f_{1T}^\perp$). When one of the $\vec{k}_{i\perp}$ is taken to be of order $\vec{q}_\perp$ and $\vec{\lambda}_\perp$ is neglected in the $\delta$ function, we need the normalization $\int d^2\vec{\lambda}_\perp S(\lambda_\perp) = 1$.

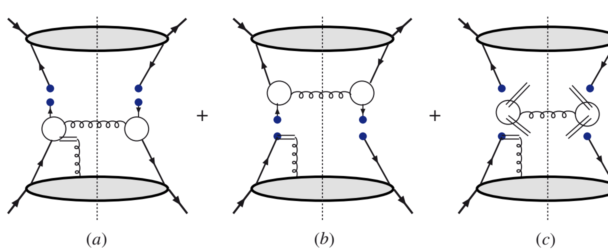


FIG. 2 (color online). Decomposition of the generic diagram in Fig. 1 into different regions in the transverse-momentum-dependent factorization approach. The transverse momentum of the Drell-Yan pair may come from (a) the Sivers function, (b) the antiquark TMD distribution, and (c) the soft factor.

What remains is to calculate the large-$k_\perp$ behavior of the TMD distributions and the soft factor. When $k_\perp$ becomes large, the leading contribution comes from perturbative one-gluon exchanges. The soft factor at large $k_\perp$ has been calculated to order $\alpha_s$ in Ref. [14]. To the same order, the unpolarized antiquark TMD distribution can be calculated from the top part of the diagram in Fig. 2(b),

$$\bar{q}(z_2,k_\perp) = \bar{q}(z_2)\delta^{(2)}(\vec{k}_\perp) + \frac{\alpha_s}{2\pi^2}\frac{1}{\vec{k}_\perp^2}C_F\int\frac{dx'}{x'}\bar{q}(x') \times\left[\frac{1+\xi_2^2}{(1-\xi_2)_+}+\delta(\xi_2-1)\left(\ln\frac{z_2^2\zeta_2^2}{\vec{k}_\perp^2}-1\right)\right], \tag{4}$$

where $\xi_2 = z_2/x'$ and where an additional contribution involving the gluon distribution has been neglected [19]. The $T$-odd TMD distribution $q_T(x, k_\perp)$ at large $k_\perp$ can be calculated in terms of the twist-three quark-gluon correlation from the bottom part of the diagram in Fig. 2(a):

$$q_T(z_1,k_\perp) = \frac{\alpha_s}{4\pi^2}\frac{2M_P}{(\vec{k}_\perp^2)^2}\int\frac{dx}{x}\left\{A + C_FT_F(x,x)\delta(\xi_1-1) \times\left(\ln\frac{z_1^2\zeta_1^2}{\vec{k}_\perp^2}-1\right)\right\}, \tag{5}$$

where

$$A = \frac{1}{2N_c}\left[\left(x\frac{\partial}{\partial x}T_F(x,x)\right)(1+\xi_1^2) + T_F(x,x-\hat{x}_g)\frac{1+\xi_1}{(1-\xi_1)_+} + T_F(x,x)\frac{(1-\xi_1)^2(2\xi_1+1)-2}{(1-\xi_1)_+}\right] + C_FT_F(x,x-\hat{x}_g)\frac{1+\xi_1}{(1-\xi_1)_+}, \tag{6}$$

with $\hat{x}_g = (1-\xi_1)x$ and $\xi_1 = z_1/x$. This relation is new and more general than the moment relation found in Ref. [13].

Plugging the above TMD distributions at large $k_\perp$ into Eq. (3), we obtain, in the TMD QCD factorization, the spin-dependent cross section for the $q + \bar{q}$ channel

$$\frac{d^3\Delta\sigma_{\rm TMD}^{q\bar{q}\to\gamma^* g}(S_\perp)}{dQ^2 dy d^2 q_\perp} = \frac{C}{(q_\perp^2)^2}\int \frac{dx}{x}\frac{dx'}{x'}\bar{q}(x')\{\delta(\xi_2 - 1)A + \delta(\xi_1 - 1)B\}, \tag{7}$$

where $A$ and $C$ have been given above and where $B = C_F T_F(x,x)\{[(1+\xi_2^2)/(1-\xi_2)_+] + 2\delta(\xi_2 - 1)\times \ln(Q^2/\vec{q}_\perp^2)\}$, with $\xi_2 = z_2/x'$. As expected, the physical cross section in Eq. (7) is independent of $v$ and $\zeta^2$.

*IV. Summary.*—In order to compare the spin-dependent cross sections in Eqs. (2) and (7), which we have calculated for the two different mechanisms, in their common region $\Lambda_{\rm QCD} \ll q_\perp \ll Q$, we need to derive an asymptotic form of the spin-dependent cross section in Eq. (2) when $q_\perp \ll Q$.

At leading order in $q_\perp/Q$, the partonic Mandelstam variables in Eq. (2) become $\hat{s} = \vec{q}_\perp^2/(1-\xi_1)(1-\xi_2)$, $\hat{t} = -\vec{q}_\perp^2/(1-\xi_2)$, and $\hat{u} = -\vec{q}_\perp^2/(1-\xi_1)$. Using the phase space $\delta$ function $\delta(\hat{s}+\hat{t}+\hat{u}-Q^2) \approx [\delta(\xi_2 - 1)/(1-\xi_1)_+ + \delta(\xi_1 - 1)/(1-\xi_2)_+ + \delta(\xi_1 - 1)\delta(\xi_1 - 1)\times \ln(Q^2/q_\perp^2)]$, we find that the leading-order (in $q_\perp/Q$) expansion of the cross section in Eq. (2) indeed reproduces Eq. (7):

$$\left.\frac{d^3\Delta\sigma_{\rm CO}^{q\bar{q}\to\gamma^* g}(S_\perp)}{dQ^2 dy d^2 q_\perp}\right|_{q_\perp\ll Q} = \frac{d^3\Delta\sigma_{\rm TMD}^{q\bar{q}\to\gamma^* g}(S_\perp)}{dQ^2 dy d^2 q_\perp}. \tag{8}$$

The same conclusion applies to the contribution to the SSA generated by gluon-quark scattering [19].

In summary, we have studied the single-transverse-spin asymmetry in Drell-Yan lepton-pair production at both large and small transverse momenta $q_\perp$ of the lepton pair. At large $q_\perp$, the spin-dependent cross section is calculated in the collinear QCD factorization formalism and expressed in terms of a twist-three quark-gluon correlation function. At small $q_\perp$, the cross section is given by a factorization formula in terms of TMD parton distributions. We have demonstrated that, in the intermediate region $\Lambda_{\rm QCD} \ll q_\perp \ll Q$, both approaches give the same answer. This explicitly establishes a connection between the two mechanisms in a physical process. This connection unifies the physical pictures for the underlying dynamics of single-transverse-spin asymmetries. Therefore, phenomenological studies using the two mechanisms in the kinematic region where they are both valid are constrained to yield identical results for them. Our results also provide a scheme for the analysis of SSAs over the whole kinematic regime of transverse momentum: the TMD QCD factorization formalism for low $q_\perp$, and the collinear QCD factorization formalism for high $q_\perp$, while both formalisms produce the same result where they overlap. An interesting extension to semi-inclusive deep inelastic scattering will be presented elsewhere.

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