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Transverse spin asymmetries for W-production in proton–proton collisions

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ABSTRACT

We study parity-even and parity-odd polarization observables for the process $pp \to l^{\pm}X$, where the lepton comes from the decay of a W-boson. By using the collinear twist-3 factorization approach, we consider the case when one proton is transversely polarized, while the other is either unpolarized or longitudinally polarized. These observables give access to two particular quark–gluon–quark correlation functions, which have a direct relation to transverse momentum dependent parton distributions. We present numerical estimates for RHIC kinematics. Measuring, for instance, the parity-even transverse single spin correlation would provide a crucial test of our current understanding of single spin asymmetries in the framework of QCD.

1. Introduction

It has long been recognized that production of W-bosons in hadronic collisions can provide new insights into the partonic structure of hadrons, with polarization observables being of particular interest. In this context the parity-odd longitudinal single spin asymmetry (SSA) in proton–proton scattering plays a very important role, both for leptonic as well as hadronic final states (see [1–14] and references therein). A major aim of looking into this observable is to get new and complementary information on the quark helicity distributions inside the proton.

In the meantime, also a few theoretical studies for W-production with transversely polarized protons are available [15–18]. These papers mainly focus on a particular parity-even transverse single spin effect in $pp \to W^{\pm}X$ (with a subsequent decay of the $W^{\pm}$ into a lepton pair) that is related to the transverse momentum dependent Sivers function $f_{1T}^{\perp}$ [19] in the polarized proton. Such an observable could, in principle, be measured at the Relativistic Heavy Ion Collider (RHIC) in Brookhaven. In order to have clean access to transverse momentum dependent parton distributions (TMDs) like the Sivers function, one has to reconstruct the W-boson in the experiment. However, what one measures is $pp \to l^{\pm}X$, and the detectors at RHIC do not allow to fully determine the momentum of the W.

The kinematics for inclusive production of a single lepton in proton–proton collisions coincides with the one for inclusive production of a jet or a hadron, for which mostly collinear factorization is used in the literature. In this Letter, we compute transverse spin observables for $pp \to l^{\pm}X$ in the collinear twist-3 formalism at the level of Born diagrams. The machinery of collinear twist-3 factorization was pioneered already in the early 1980s [20,21], and in the meantime frequently applied to transverse spin effects in hard semi-inclusive reactions (see [22–25] and references therein).

If one of the protons in $pp \to l^{\pm}X$ is transversely polarized, and the other is either unpolarized or longitudinally polarized, one can identify two parity-even and two parity-odd spin observables. We will discuss below that, in the collinear twist-3 approach, these four observables contain two specific twist-3 quark–gluon–quark correlation functions. One is the so-called ETQS (Efremov–Teryaev–Qiu–Sterman) matrix element [20,22,23], which is related to a particular moment of the transverse momentum dependent Sivers function as shown in [26,27]. The second is related to the TMD g_{1T} [25,28], where we use the TMD-notation of Refs. [29–32].

In addition to the analytical results, we provide numerical estimates for typical RHIC kinematics ($\sqrt{s} = 500$ GeV). All the observables are peaked around $l_T \approx M_W/2$, with l_T representing the transverse momentum of the lepton and M_W the W-boson mass. In each case we predict clearly measurable effects. For the parity-even transverse SSA A_{TU}^{e} our numerical results are very close to those obtained in Ref. [17] on the basis of factorization in terms of transverse momentum dependent parton correlators.

Before presenting our results we emphasize that measuring A_{TU}^{e} would provide a crucial test of our present understanding of transverse SSAs in QCD. In particular, this means that such a measurement would test the same physics—the gluon exchange between the remnants of the hadrons and the active partons—which underlies the famous process-dependence of the Sivers function and of related time-reversal odd parton distributions [33]. In other words, experimental results for A_{TU}^{e} in $pp \to l^{\pm}X$, even if analyzed in terms of collinear parton correlators, would check a crucial

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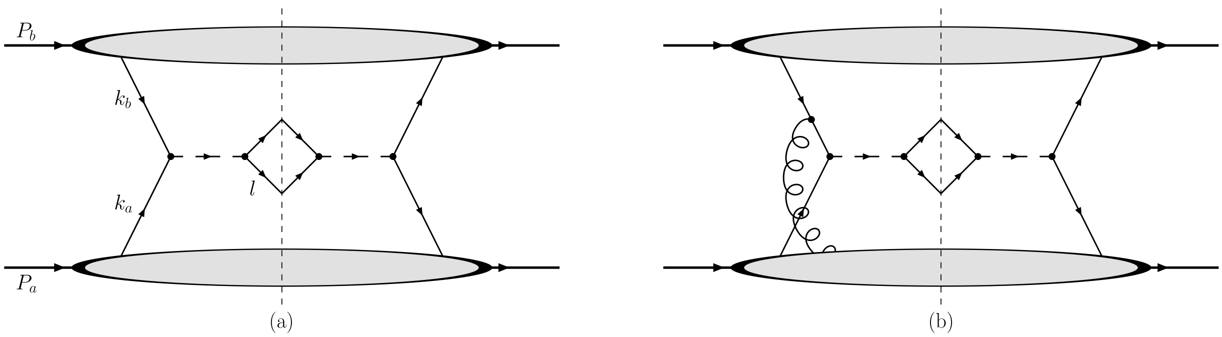

Fig. 1. Diagram (a): parton model representation for $pp \to l^\pm X$, where the lepton is produced in the decay of a W-boson. The final state (anti-)neutrino goes unobserved. Diagram (b): contribution from quark–gluon–quark correlation. This diagram, together with its Hermitian conjugate which is not displayed, needs to be taken into account when computing twist-3 observables.

ingredient of TMD-factorization [34–36]. Such a check, in essence, can be considered to be as fundamental as measuring the sign of the Sivers asymmetry in the Drell–Yan process.

2. Analytical results

We start by fixing the kinematical variables for the process $pp \to l^\pm X$, and assign 4-momenta to the particles according to

$$p(P_a) + p(P_b) \to l^\pm(l) + X. \tag{1}$$

By means of these momenta we specify a coordinate system through $\hat{e}_z = \hat{P}_a = -\hat{P}_b$, $\hat{e}_x = \hat{l}_T$ (with $\vec{l}_T$ representing the transverse momentum of the jet), and $\hat{e}_y = \hat{e}_z \times \hat{e}_x$. Mandelstam variables are defined by

$$s = (P_a + P_b)^2, \qquad t = (P_a - l)^2, \qquad u = (P_b - l)^2, \tag{2}$$

while on the partonic level one has

$$\hat{s} = (k_a + k_b)^2 = x_a x_b s, \qquad \hat{t} = (k_a - l)^2 = x_a t,$$
$$\hat{u} = (k_b - l)^2 = x_b u, \tag{3}$$

where k_a and k_b denote the momentum of the active quark/antiquark in the protons; see also Fig. 1(a). The momentum fraction x_a characterizes the (large) plus-momentum of the quark/antiquark in the proton moving along $\hat{e}_z$ through $k_a^+ = x_a P_a^+$.[1] Likewise, one has $k_b^- = x_b P_b^-$. The relation $\hat{s} + \hat{t} + \hat{u} = 0$ implies

$$x_a = -\frac{x_b u}{x_b s + t} = \frac{x_b \sqrt{s}\, l_T e^{\eta}}{x_b s - \sqrt{s}\, l_T e^{-\eta}}. \tag{4}$$

In the second step in (4) we express x_a, for a given $\sqrt{s}$, through $l_T = |\vec{l}_T|$ and the pseudo-rapidity $\eta = -\ln\tan(\vartheta/2)$ of the lepton, since transverse momenta and (pseudo-)rapidities are commonly used to describe the kinematics of a final state particle in proton–proton collisions.

Next, we turn to the polarization observables for $pp \to l^\pm X$, which we compute in the collinear factorization framework. As already mentioned, we focus on the situation when one proton is transversely polarized, while the other is either unpolarized or longitudinally polarized. One finds the following expression for the cross section[2]:

$$\begin{aligned}
l^0 \frac{d^3\sigma}{d^3 l}
&= \frac{\alpha_{em}^2}{12 s \sin^4\vartheta_W} \sum_{a,b} |V_{ab}|^2 \int_{x_b^{\min}}^{1} \frac{dx_b}{x_a x_b} \frac{1}{x_b s + t} \Big\{ H^{ab} f_1^a(x_a) f_1^b(x_b) \\
&\quad + 2\pi M \varepsilon_T^{ij} l_T^i S_{aT}^j \tilde{H}^{ab} \bigg[\bigg(T_F^a(x_a, x_a) - x_a \frac{d}{dx_a} T_F^a(x_a, x_a) \bigg) \\
&\quad + K(\hat{s}) T_F^a(x_a, x_a) \bigg] f_1^b(x_b) \\
&\quad + 2M \vec{l}_T \cdot \vec{S}_{aT} \tilde{H}^{ab} \bigg[\bigg(\tilde{g}^a(x_a) - x_a \frac{d}{dx_a} \tilde{g}^a(x_a) \bigg) \\
&\quad + K(\hat{s}) \tilde{g}^a(x_a) + 2x_a g_T^a(x_a) \bigg] f_1^b(x_b) \\
&\quad - 2\pi M \lambda_b \varepsilon_T^{ij} l_T^i S_{aT}^j \tilde{H}^{ab} \bigg[\bigg(T_F^a(x_a, x_a) - x_a \frac{d}{dx_a} T_F^a(x_a, x_a) \bigg) \\
&\quad + K(\hat{s}) T_F^a(x_a, x_a) \bigg] g_1^b(x_b) \\
&\quad - 2M \lambda_b \vec{l}_T \cdot \vec{S}_{aT} \tilde{H}^{ab} \bigg[\bigg(\tilde{g}^a(x_a) - x_a \frac{d}{dx_a} \tilde{g}^a(x_a) \bigg) \\
&\quad + K(\hat{s}) \tilde{g}^a(x_a) + 2x_a g_T^a(x_a) \bigg] g_1^b(x_b) + \cdots \Big\},
\end{aligned} \tag{5}$$

with $K(\hat{s}) = \dfrac{2M_W^2(\hat{s} - M_W^2 - \Gamma_W^2)}{(\hat{s} - M_W^2)^2 + M_W^2 \Gamma_W^2}$.

In Eq. (5), ϑ_W is the weak mixing angle, V_{ab} is a CKM matrix element, M is the proton mass, M_W is the W-mass and Γ_W its decay width. We also use $\varepsilon_T^{ij} \equiv \varepsilon^{-+ij}$ with $\varepsilon^{0123} = 1$. The transverse spin vector of the proton moving along $\hat{e}_z$ is denoted by $\vec{S}_{aT}$, whereas λ_b represents the helicity of the second proton. The lower limit of the x_b-integration is given by $x_b^{\min} = -t/(s+u)$. One can project out the four spin-dependent components of the cross section in (5), in order, through

$$\sigma_{TU}^e = \frac{1}{4}\big(\big[\sigma(\uparrow_y, +) - \sigma(\downarrow_y, +)\big] + \big[\sigma(\uparrow_y, -) - \sigma(\downarrow_y, -)\big] \big), \tag{6}$$

$$\sigma_{TU}^o = \frac{1}{4}\big(\big[\sigma(\uparrow_x, +) - \sigma(\downarrow_x, +)\big] + \big[\sigma(\uparrow_x, -) - \sigma(\downarrow_x, -)\big] \big), \tag{7}$$

$$\sigma_{TL}^o = \frac{1}{4}\big(\big[\sigma(\uparrow_y, +) - \sigma(\downarrow_y, +)\big] - \big[\sigma(\uparrow_y, -) - \sigma(\downarrow_y, -)\big] \big), \tag{8}$$

1 For a generic 4-vector v, we define light-cone coordinates according to $v^\pm = (v^0 \pm v^3)/\sqrt{2}$ and $\vec{v}_T = (v^1, v^2)$.

2 Polarization degrees are suppressed in the cross section formula (5).

$$\sigma^e_{TL} = \frac{1}{4}\big(\big[\sigma(\uparrow_x,+) - \sigma(\downarrow_x,+)\big] - \big[\sigma(\uparrow_x,-) - \sigma(\downarrow_x,-)\big]\big). \tag{9}$$

In these formulas, '$\uparrow_{x/y}$' ('$\downarrow_{x/y}$') denotes transverse polarization along $\hat{e}_{x/y}$ ($-\hat{e}_{x/y}$) for the proton moving in the $\hat{e}_z$-direction, whereas '+' and '−' represent the helicities of the second proton.

The dots in Eq. (5) indicate longitudinal single spin and double spin observables, as well as four possible correlations for double transverse polarization. In collinear factorization, the latter are at least twist-4 effects in the Standard Model. Note that double transverse polarization observables for W-production were also discussed in connection with potential physics beyond the Standard Model (see [37,38] and references therein).

We computed the (twist-2) unpolarized cross section in the first line of (5) on the basis of diagram (a) in Fig. 1 by applying the collinear approximation to the momenta k_a and k_b of the active partons. The result contains the ordinary unpolarized quark distribution f_1^a for a quark flavor a. The hard scattering coefficients H^{ab} and $\tilde{H}^{ab}$ in Eq. (5), expressed through the partonic Mandelstam variables in (3), read

$$H^{ab} = \frac{\hat{u}^2}{(\hat{s} - M_W^2)^2 + M_W^2 \Gamma_W^2},$$
$$\tilde{H}^{ab} = \frac{1}{\hat{u}} H^{ab}, \quad \text{for } ab = d\bar{u}, s\bar{u}, \bar{d}u, \bar{s}u. \tag{10}$$

In Eq. (10), one has to replace $\hat{u}$ by $\hat{t}$ for $ab = \bar{u}d, \bar{u}s, u\bar{d}, u\bar{s}$.

The four cross sections in (6)–(9) represent twist-3 observables. Calculational details for such observables in collinear factorization can be found in various papers; see, e.g., Refs. [23–25,39–41]. We merely mention that one has to expand the hard scattering contributions around vanishing transverse parton momenta. While for twist-2 effects only the leading term of that expansion matters, in the case of twist-3 the second term is also relevant. In addition, the contribution from quark–gluon–quark correlations, as displayed in diagram (b) in Fig. 1, needs to be taken into consideration. The sum of all the terms can be written in a color gauge invariant form, which provides a consistency check of the calculation.

The quark–gluon–quark correlator showing up in σ^e_{TU} and σ^o_{TL} is the aforementioned ETQS matrix element $T_F^a(x,x)$ [20,22,23]. The peculiar feature of this object is the vanishing gluon momentum—that's why it is also called "soft gluon pole matrix element". If the gluon momentum becomes soft one can hit the pole of a quark propagator in the hard part of the process, providing an imaginary part (nontrivial phase) which, quite generally, can lead to single spin effects [20,22,23]. Note also that in our lowest order calculation no so-called soft fermion pole contribution (see [42] and references therein) emerges. For σ^o_{TU} and σ^e_{TL} another quark–gluon–quark matrix element—denoted as $\tilde{g}^a$; see, in particular, Refs. [25,28,43]—appears, together with the familiar twist-3 quark–quark correlator g_T^a (and, in the case of σ^o_{TL}, together with the quark helicity distribution g_1^a).

We use the common definitions for f_1, g_1, and g_T. The quark–gluon–quark correlators T_F and $\tilde{g}$ are specified according to[3]

$$-i\varepsilon_T^{ij} S_T^j T_F(x,x)$$
$$= \frac{1}{2M}\int \frac{d\xi^- d\zeta^-}{(2\pi)^2} e^{ixP^+\xi^-} \langle P, S_T | \bar{\psi}(0)\gamma^+ igF^{+i}(\zeta^-)\psi(\xi^-)$$
$$\times |P, S_T\rangle, \tag{11}$$

$$S_T^i \tilde{g}(x)$$
$$= \frac{1}{2M}\int \frac{d\xi^-}{2\pi} e^{ixP^+\xi^-} \langle P, S_T | \bar{\psi}(0)\gamma_5\gamma^+$$
$$\times \left(iD_T^i - ig\int_0^\infty d\zeta^- F^{+i}(\zeta^-)\right)\psi(\xi^-)|P, S_T\rangle, \tag{12}$$

with $F^{\mu\nu}$ representing the gluon field strength tensor, and $D^\mu = \partial^\mu - igA^\mu$ the covariant derivative. Eqs. (11) and (12) hold in the light-cone gauge $A^+ = 0$, while in a general gauge Wilson lines need to be inserted between the field operators.

It is important that T_F and $\tilde{g}$ are related to moments of TMDs. To be explicit, one has [25–28]

$$\pi T_F(x,x) = -\int d^2k_T \frac{\vec{k}_T^2}{2M^2} f_{1T}^\perp(x,\vec{k}_T^2)\Big|_{DIS}, \tag{13}$$

$$\tilde{g}(x) = \int d^2k_T \frac{\vec{k}_T^2}{2M^2} g_{1T}(x,\vec{k}_T^2), \tag{14}$$

where we use the conventions of Refs. [29–32] for the TMDs $f_{1T}^\perp$ and g_{1T}. In Eq. (13) we take into account that the Sivers function $f_{1T}^\perp$ depends on the process in which it is probed [33,44]. In order to make numerical estimates we will exploit the relations in (13), (14).

Finally, note that, due to the pure vector–axialvector coupling of the W-boson, no chiral-odd parton correlator shows up in any of the four spin correlations in (5), which makes those observables rather clean. The situation is different if one considers single lepton production from the decay of a virtual photon or of a Z-boson.

3. Numerical results

Now we move on to discuss numerical results for the polarization observables by limiting ourselves to the transverse single spin effects. This means, we consider the two spin asymmetries A^e_{TU} and A^o_{TU},

$$A^e_{TU} = \frac{\sigma^e_{TU}}{\sigma_{UU}}, \qquad A^o_{TU} = \frac{\sigma^o_{TU}}{\sigma_{UU}}, \tag{15}$$

with σ^e_{TU} and σ^o_{TU} from Eqs. (6) and (7), respectively, and σ_{UU} denoting the unpolarized cross section. Note that the definition of A^e_{TU} corresponds to the one of the transverse SSA A_N, which has been extensively studied in one-particle inclusive production for hadron–hadron collisions; see [45–47] for recent experimental results from RHIC.

To compute σ_{UU} we use the unpolarized parton densities from the CTEQ6-parameterization [48]. For the ETQS matrix element we use the relation (13) between T_F and the Sivers function, and take $f_{1T}^\perp$ from the recent fit provided in Ref. [49] on the basis of data from semi-inclusive DIS. (For experimental studies of the Sivers effect we refer to [50,51], while extractions of the Sivers function from data can be found in [49,52–56].) In the case of A^o_{TU} one needs input for g_T and $\tilde{g}$. For g_T we resort to the frequently used Wandzura–Wilczek approximation [57] (see [58] for a recent study of the quality of this approximation)

$$g_T(x) \approx \int_x^1 \frac{dy}{y} g_1(y), \tag{16}$$

whereas for $\tilde{g}$ we use (14) and a Wandzura–Wilczek-type approximation for the particular k_T-moment of g_{1T} in (14) [59], leading

[3] Note that in the literature different conventions for T_F exist.

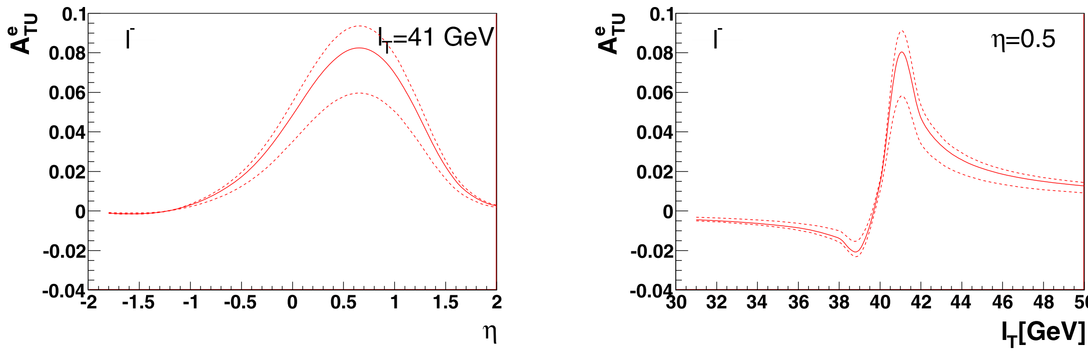

Fig. 2. A^e_{TU} for $pp \to l^- X$ as a function of η (left) and l_T (right) for $\sqrt{s} = 500$ GeV. The solid line represents the central value, while the error is indicated by the dashed lines (see also text).

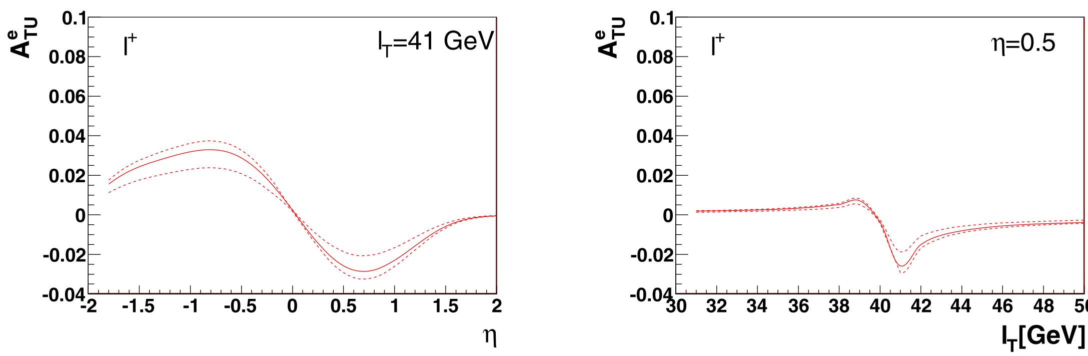

Fig. 3. A^e_{TU} for $pp \to l^+ X$ as a function of η (left) and l_T (right) for $\sqrt{s} = 500$ GeV. The solid line represents the central value, while the error is indicated by the dashed lines (see also text).

to

$$\tilde{g}(x) \approx x \int_x^1 \frac{dy}{y} g_1(y). \tag{17}$$

We mention that (17) and a corresponding relation between chiral-odd parton distributions were used in [60,61] in order to estimate certain spin asymmetries in semi-inclusive DIS. The comparison to data discussed in [61] looks promising, though more experimental information is needed for a thorough test of approximate relations like the one in (17). Measuring the SSA A^o_{TU} could provide such a test. The helicity distributions g_1^a in (16) and (17) are taken from the DSSV-parameterization [62]. The transverse momentum of the lepton l_T serves as the scale for the parton distributions.

The numerical estimates are for typical RHIC kinematics, i.e., $\sqrt{s} = 500$ GeV. We present the asymmetries either as function of η for fixed l_T or vice versa.

We start by discussing the parity-even asymmetry A^e_{TU}. As shown in the right plot in Figs. 2 and 3, this observable is peaked around $l_T \approx M_W/2$—a feature that does not depend on the value of η. To be more precise, the peak is at $l_T = 41$ GeV, i.e., slightly above $M_W/2$. The peak in the polarized cross section σ^e_{TU} gets enhanced in the asymmetry, because the unpolarized cross section drops rather fast when going beyond $l_T = M_W/2$. (As a side-remark we point out that the asymmetry in the peak region is completely dominated by the term in the 4th line in (5) containing the factor $K(\hat{s})$.) Nevertheless, in this kinematical region we expect A^e_{TU} to be measurable. As discussed in the introduction, in this context it is important to recall that information on the sign of the asymmetry is already sufficient for a crucial test of our current understanding of transverse SSAs.

In particular in the peak region, the asymmetry is larger for l^--production (W^--production) than for l^+-production, which is partly due to the rather large Sivers function for d-quarks obtained in the fit of Ref. [49]. The l^--asymmetry and l^+-asymmetry come with opposite sign because the Sivers functions for u-quarks and d-quarks have an opposite sign. Note also that both asymmetries change sign as function of l_T. Therefore, whether the sign of the asymmetry can be measured unambiguously may critically depend on the l_T-resolution in the experiment.

The plots also show the error of our estimate, where we merely took into account the uncertainty due to the Sivers function. In order to estimate the error we followed the method outlined in the Appendix of Ref. [49].

As the η-dependence of A^e_{TU} in left plot in Fig. 2 shows, the l^--asymmetry is maximal in the positive η-range, when a large-x parton from the polarized proton participates in the hard scattering. Obviously, by integrating over a suitable η-range one may optimize between magnitude of the asymmetry on the one hand and the size of the statistical error bars on the other. Moreover, it is worthwhile to mention that the contributions from the anti-quark Sivers functions are not negligible in the backward region. (Here we refer to a corresponding discussion on the Sivers asymmetry in the Drell–Yan process for proton–proton collisions in [63], where the strong sensitivity to the Sivers function for antiquarks was already pointed out.)

It is also interesting that for both l^+-production and l^--production the overall magnitude of A^e_{TU} is very similar to the predictions presented in Ref. [17], where TMD-factorization was used.

The fact that, in an experiment, a single lepton can originate from various background processes poses a challenge for the extraction of the desired signal. In the region of the large lepton transverse momenta, required in the present case, essentially leptons from Z-boson decay are relevant. (See Refs. [12–14] for a detailed discussion in the case of the parity-odd longitudinal

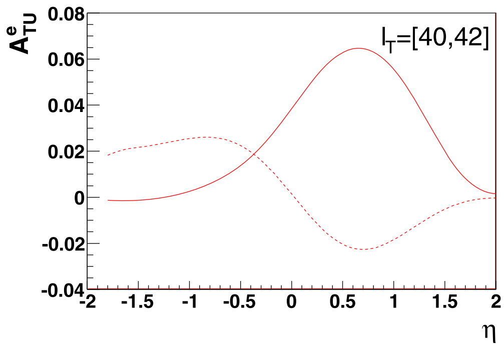

Fig. 4. A^e_{TU} for $pp \to l^\pm X$ as a function of η and integrated over the indicated l_T-range for $\sqrt{s} = 500$ GeV. The solid line is for l^--production, and the dashed line is for l^+-production.

SSA.) Since the number of produced Z-bosons at RHIC is smaller than the one for W-bosons one may expect that this background will not strongly modify the asymmetry. On the other hand, it is known that leptons originating from Z-boson decays cannot be neglected [12]. In the case of the transverse SSA, however, one may exploit that the asymmetry has this very pronounced l_T behavior. Potential contributions from Z-boson decay should be negligible around $l_T = 41$ GeV. Fig. 4 shows how the asymmetry as a function of η looks like for a specific l_T bin. A complete calculation of the asymmetry originating from Z-decays not only includes the Sivers effect but also, as mentioned earlier, a second contribution involving chiral-odd correlators.

Let us now turn to the parity-odd transverse SSA A^o_{TU}, which is displayed in Fig. 5. Again, this asymmetry has a pronounced peak at $l_T = 41$ GeV, and it is largest for l^+-production (up to about 8%). As outlined above, our prediction for A^o_{TU} is based on the Wandzura–Wilczek-type approximation leading to (17), which probably represents the most uncertain part of our calculation. At present, it is difficult to assign a quantitative error to this observable. However, from the study given in [61], in which the chiral-odd counterpart of (17) was used, we consider ±50% as a conservative error estimate. Like in the case of the parity-even SSA, also A^o_{TU} is almost entirely determined by the $K(\hat{s})$-term in the 6th line in (5). This implies that, due to the relation (14), it gives rather clean access to the TMD g_{1T}, which so far is experimentally unconstrained. Therefore, in any case, a measurement of A^o_{TU} would provide very interesting new information.

4. Summary

We have studied transverse spin asymmetries for the process $pp \to l^\pm X$, where the lepton is produced in the decay of a W-boson. If one of the protons is transversely polarized, and the other is either unpolarized or longitudinally polarized, there exist two parity-even and two parity-odd spin asymmetries. We computed these asymmetries in collinear twist-3 factorization at the level of Born diagrams. Moreover, for the two transverse single spin asymmetries A^e_{TU} and A^o_{TU}—defined through Eqs. (15) and (6), (7)—we made numerical estimates for typical kinematics at RHIC ($\sqrt{s} = 500$ GeV). In the following we summarize our main results:

- The analytical results for all four spin-dependent cross sections are given by two particular quark–gluon–quark correlators, which have a direct relation to transverse momentum dependent parton distributions: the Sivers function $f_{1T}^\perp$ and the TMD g_{1T}; see Eqs. (13), (14). Measuring these observables could therefore provide new information on the structure of the proton that goes beyond the collinear parton model.
- The parity-even SSA A^e_{TU} is largest for l^--production (up to about 8%), and it is peaked for transverse momenta l_T of the lepton slightly above $M_W/2$. (Actually, all the asymmetries studied in this Letter are significant only in a relatively narrow region around $l_T \approx M_W/2$.) Measuring the sign of this asymmetry can, in essence, provide an as crucial test as measuring the sign of the Sivers asymmetry in Drell–Yan would do: it can test our present understanding of the underlying dynamics of transverse SSAs and at the same time check an important ingredient of TMD-factorization, namely the influence of the Wilson-line which is generated by the interaction between the active partons and the remnants of the protons. (For related work we refer to [17,33,44,52,64–66].)
- To the best of our knowledge the parity-odd SSA A^o_{TU} was never before explored in the literature. We find A^o_{TU} to be largest for l^+-production (also up to about 8%, like A^e_{TU} for l^--production). This observable is directly related to (a moment of) the TMD g_{1T}, for which at this time no experimental information exists.

In general, we believe that W-physics for polarized proton–proton collisions is very promising not only in the case of longitudinally polarized protons, but has also a considerable discovery potential for transverse polarization.

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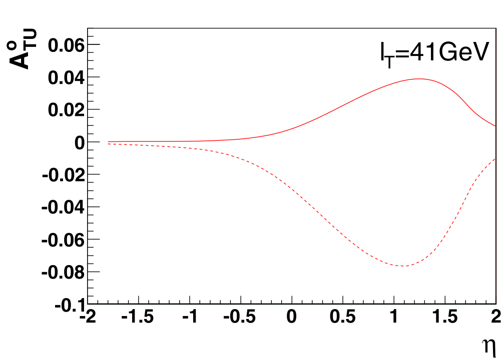

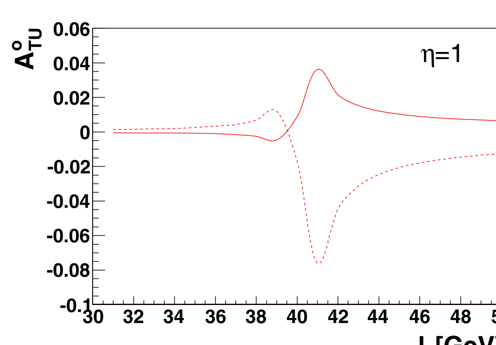

Fig. 5. A^o_{TU} for $pp \to l^\pm X$ as a function of η (left) and l_T (right) for $\sqrt{s} = 500$ GeV. The solid line is for l^--production, and the dashed line is for l^+-production.

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