Single spin asymmetries in the Drell-Yan process

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Abstract

We discuss single transverse spin asymmetries in the Drell-Yan process originating from so-called gluonic poles in twist-three hadronic matrix elements, as first considered by Qiu and Sterman. Even though time-reversal invariance is not broken, the effects of such poles cannot be distinguished from those of time-reversal odd distribution functions. We show the connection between gluonic poles and large distance gluon fields, in particular we focus on boundary conditions. We identify the possible single spin asymmetries in the Drell-Yan process.

1 Introduction

In the standard description of the Drell-Yan process (DY) in terms of distribution functions time-reversal symmetry implies the absence of single spin asymmetries at tree level, even including order $1/Q$ corrections \cite{1}. Additional time-reversal odd (T-odd) distribution functions (DFs) are present when the incoming hadrons cannot be treated as plane-wave states. This may occur due to some factorization breaking mechanism \cite{2}. We will show that, even apart from such mechanisms, the contributions of T-odd DFs may effectively arise due to the presence of so-called gluonic poles attributed to large distance gluon fields. The gluonic poles appearing in the twist-three hadronic matrix elements \cite{3, 4, 5} together with imaginary phases of hard subprocesses effectively give rise to the same single spin asymmetries as T-odd DFs, but without a violation of time-reversal invariance. This is the origin of the single spin asymmetry of Ref. \cite{6}. For a detailed account on these matters see \cite{7}.

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2 The DY process in terms of distribution functions

We employ methods originating from Refs. [8, 9, 10, 11, 12, 13, 14] in order to describe the soft (non-perturbative) parts of the scattering process in terms of correlation functions, which are (Fourier transforms of) hadronic matrix elements of non-local operators. We restrict ourselves to tree-level, but include $1/Q$ power corrections. The asymmetries under investigation are loosely referred to as 'twist-three' asymmetries, since they are suppressed by a factor of $1/Q$, where the photon momentum $q$ sets the scale $Q$, such that $Q^2 = q^2$. We do not take $Z$ bosons into account, since the asymmetries are likely to be negligible at or above the $Z$ threshold.

The Drell-Yan process consists of two soft parts (depicted in Fig. 1 for the leading order) and one of them is described (up to order $1/Q$) by the quark correlation functions $\Phi$ and $\Phi_A$, and the other soft part by the antiquark correlation functions, denoted by $\bar{\Phi}$ and $\bar{\Phi}_A$. The quark-quark correlation function,

$$\Phi_{ij}(P_1, S_1; p) = \int \frac{d^4z}{(2\pi)^4} e^{ip \cdot z} \langle P_1, S_1|\bar{\psi}_j(0)\psi_i(z)|P_1, S_1\rangle,$$

is a function of the momentum and spin vectors $P_1, S_1$ of the incoming hadron (spin-1/2), with $P_1 \cdot S_1 = 0$, and the quark momentum $p$. The hadron momentum $P_1$ is chosen to be predominantly along a light-like direction given by the vector $n_+$. Another light-like direction $n_-$ is chosen such that $n_+ \cdot n_- = 1$; both vectors are dimensionless. The second hadron is chosen to be predominantly in the $n_-$ direction, such that $P_1 \cdot P_2 = O(Q^2)$. We write $p^\pm = p \cdot n_\pm$ and approximate the parton momentum $p \approx x P_1 + p_T$ and the polarization vector $S_1 \approx \lambda_1 P_1/M_1 + S_{1T}$.

We will consider the case where one integrates over the transverse momentum $q_T$ of the photon. One then only encounters correlation functions integrated over all but the leading component, such that they are functions of the light-cone momentum fractions (e.g. $x$) only. So we consider the partly integrated quark correlation functions

$$\Phi_{ij}(x) \equiv \int \frac{d\lambda}{2\pi} e^{i\lambda x} \langle P_1, S_1|\bar{\psi}_j(0)\psi_i(\lambda n_-)|P_1, S_1\rangle,$$

$$\Phi_{Aij}(x, y) \equiv \int \frac{d\lambda d\eta}{2\pi 2\pi} e^{i\lambda x + i\eta (y-x)} \langle P_1, S_1|\bar{\psi}_j(0)g A_\tau^\alpha(\eta n_-)\psi_i(\lambda n_-)|P_1, S_1\rangle.$$
The anti-quark correlation function is defined as
\[ \Phi(\vec{x}) = \int \frac{d\lambda}{2\pi} e^{-i\lambda\vec{x}} \langle P_2, S_2 | \psi(\lambda n_+) \bar{\psi}(0) | P_2, S_2 \rangle \] (4)
and the function \( \Phi_A^\alpha(\vec{x}, \vec{y}) \) is defined analogously. Moreover, the inclusion of path-ordered exponentials, such as,
\[ \mathcal{L}(0, \lambda n) = \mathcal{P} \exp \left( -ig \int_0^{\lambda n} dz A_\mu(z) \right), \] (5)
which are needed in order to render the correlation functions gauge invariant, is implicit.

In the expression of the hadron tensor the fermion propagators appearing in the hard part of the subleading contributions (cf. Fig. 2) is approximated like (neglecting contributions that will appear suppressed by \( 1/Q^2 \))
\[ \frac{p_1 - q + m}{(p_1 - q)^2 - m^2 + i\epsilon} \approx -\frac{\not{p}_+}{Q\sqrt{2} \, x - y + i\epsilon}. \] (6)
Hence, a zero-momentum gluon \((x = y)\) is always accompanied by an on-shell quark propagator,

\[ \int dy \Phi^\alpha_A(x, y) \frac{x - y}{x - y + i\epsilon} \xrightarrow{\Phi^\alpha_A(x, y) = 0} \int dy \Phi^\alpha_A(x, y). \] (7)

For the correlation functions \( \Phi \) and \( \Phi^\alpha_A \) we need up to order \( 1/Q \) the following parametrizations in terms of distribution functions [14]:
\[ \Phi(x) = \frac{1}{2} \left[ f_1(x) P_1 + g_1(x) \lambda_1 \gamma_5 P_1 + h_1(x) \gamma_5 S_{1T} P_1 \right] \]
\[ + \frac{M_1}{2} \left[ e(x) 1 + g_T(x) \gamma_5 S_{1T} + h_L(x) \frac{\lambda_1}{2} \gamma_5 [\not{p}_+, \not{p}_-] \right], \] (8)
\[ \Phi_A^\alpha(x, y) = \frac{M_1}{2} \left[ G_A(x, y) i e_T^\alpha S_{1T} P_1 + \tilde{G}_A(x, y) S_{1T}^\alpha \gamma_5 P_1 \right. \]
\[ + \left. H_A(x, y) \lambda_1 \gamma_5 \gamma_T^\alpha P_1 + E_A(x, y) \gamma_T^\alpha P_1 \right], \] (9)
where \( \epsilon_{\mu\nu} = \epsilon^{\alpha\beta\mu\nu} n_+ \alpha n_- \).

The parametrization of \( \Phi(x) \) is consistent with requirements following from hermiticity, parity and time-reversal invariance,

\[
\Phi^\dagger(P_1, S_1; p) = \gamma_0 \Phi(P_1, S_1; p) \gamma_0 \quad \text{[Hermiticity]} \tag{10}
\]
\[
\Phi(P_1, S_1; p) = \gamma_0 \Phi(\bar{P}_1, -\bar{S}_1; \bar{p}) \gamma_0 \quad \text{[Parity]} \tag{11}
\]
\[
\Phi^*(P_1, S_1; p) = \gamma_5 C \Phi(\bar{P}_1, \bar{S}_1; \bar{p}) C^\dagger \gamma_5 \quad \text{[Time reversal]} \tag{12}
\]

where \( \bar{p} = (p^0, -\mathbf{p}), \) etc. For the one-argument functions in Eq. (8) it follows from hermiticity that they are real. Note that for the validity of Eq. (12) it is essential that the incoming hadron is a plane wave state. For \( \Phi_A^a \) hermiticity, parity and time-reversal invariance yield the following relations:

\[
[\Phi_A^a(P_1, S_1; p_1, p_2)]^\dagger = \gamma_0 \Phi_A^a(P_1, S_1; p_2, p_1) \gamma_0 \quad \text{[Hermiticity]} \tag{13}
\]
\[
\Phi_A^a(P_1, S_1; p_1, p_2) = \gamma_0 \Phi_{Aa}(\bar{P}_1, -\bar{S}_1; \bar{p}_1, \bar{p}_2) \gamma_0 \quad \text{[Parity]} \tag{14}
\]
\[
[\Phi_A^a(P_1, S_1; p_1, p_2)]^* = \gamma_5 C \Phi_{Aa}(\bar{P}_1, \bar{S}_1; \bar{p}_1, \bar{p}_2) C^\dagger \gamma_5 \quad \text{[Time reversal]} \tag{15}
\]

Hermiticity then gives for the two-argument functions in Eq. (9) the following constraints:

\[
G_A(x, y) = -G_A^*(y, x), \quad \tilde{G}_A(x, y) = \tilde{G}_A^*(y, x), \tag{16}
\]
\[
E_A(x, y) = -E_A^*(y, x), \quad H_A(x, y) = H_A^*(y, x). \tag{17}
\]

Hence, the real and imaginary parts of these two-argument functions have definite symmetry properties under the interchange of the two arguments. If we would impose time-reversal invariance all four functions must be real and \( \tilde{G}_A \) and \( \tilde{H}_A \) are then symmetric and \( G_A \) and \( E_A \) are antisymmetric under interchange of the two arguments, such that at \( x = y \) only \( \tilde{G}_A \) and \( \tilde{H}_A \) survive.

In the remainder of this section we do not impose time-reversal invariance and hence allow for imaginary parts of these functions. In addition, the following (T-odd) one-argument DFs then appear:

\[
\Phi(x)_{|_{T-\text{odd}}} = \frac{M_1}{2} \left[ f_T(x) \epsilon^{\mu\nu}_{T} S_{1T} \gamma_{T\nu} - e_L(x) \lambda_1 i \gamma_5 + h(x) \frac{i}{2} \left[ \not{n}_+, \not{n}_- \right] \right]. \tag{18}
\]

The two-argument functions and the one-argument functions are related by the classical e.o.m. \(((i\bar{\mathcal{D}} - m)\psi = 0)\), which hold inside hadronic matrix elements \([11]\). Using a similar parametrization for \( \Phi^a_D \) (defined like \( \Phi_A^a \), but with \( gA_T^a \) replaced by \( iD_T^a \)) as in Eq. (8), one has the following relations \([13, 14]\):

\[
\int dy \left[ \text{Re} \, G_D(x, y) + \text{Re} \, \tilde{G}_D(x, y) \right] = 2xg_T(x) - \frac{m}{M} h_1(x), \tag{19}
\]
\[
\int dy \left[ \text{Im} \, G_D(x, y) + \text{Im} \, \tilde{G}_D(x, y) \right] = 2ixf_T(x), \tag{20}
\]
\[
\int dy \left[ 2 \text{Re} \, H_D(x, y) \right] = xh_L(x) - \frac{m}{M} q_1(x), \tag{21}
\]
\[
\int dy \left[ 2 \text{Im} \, H_D(x, y) \right] = -ixe_L(x), \tag{22}
\]
\[
\int dy \left[ 2 \text{Re} \, E_D(x, y) \right] = xe(x) - \frac{m}{M} f_1(x), \tag{23}
\]
\[
\int dy \left[ 2 \text{Im} \, E_D(x, y) \right] = ixh(x). \tag{24}
\]
From this (and \( iD^\alpha = i\partial^\alpha + gA^\alpha \)) we see that the (T-odd) imaginary parts of the two-argument functions are related to the T-odd one-argument functions, as one expects. So if time-reversal invariance is imposed, the imaginary parts of the e.o.m. Eqs. (20), (23) and (24) become three trivial equalities. We like to point out that if one integrates Eqs. (19) and (20) over \( x \), weighted with some test-function \( \sigma(x) \), one arrives at the sum rules discussed in [13, 15].

3 Gluonic poles and time-reversal odd behavior

We are interested in the behavior of the quark-gluon correlation function \( \Phi_3^\alpha \) in case \( x = y \), when the gluon has zero-momentum. For this purpose, we define (\( \alpha \) is a transverse index)

\[
\Phi_{Fij}^\alpha(x, y) \equiv \int \frac{d\lambda \, d\eta}{2 \pi} e^{i\lambda x} e^{i\eta(y-x)} \langle P, S | \overline{\psi}_j(0) F^{+\alpha}(\eta n_-) \psi_i(\lambda n_-) | P, S \rangle
\]  

(25)

and \( F^{\rho\sigma}(z) = \frac{i}{2} [D^{\rho}(z), D^{\sigma}(z)] \). This matrix element has the same hermiticity, but the opposite time-reversal behavior as \( \Phi_3^\alpha \),

\[
[\Phi_{F}^\alpha(x, y)]^* = -\gamma_5 C \Phi_{F\alpha}(x, y) C^\dagger \gamma_5 \quad \text{[Time reversal]} \tag{26}
\]

and we will parametrize it identically with help of functions called \( G_F(x, y), \tilde{G}_F(x, y), H_F(x, y) \) and \( E_F(x, y) \), noting that time-reversal implies that \( G_F \) and \( E_F \) are symmetric and thus may survive at \( x = y \) (in contrast to \( G_A(x, x) \) and \( E_A(x, x) \)). In the gauge \( A^+ = 0 \) one has \( F^{+\alpha} = \partial^+ A_T^\alpha \) and one finds by partial integration

\[
(x-y) \Phi_3^\alpha(x, y) = -i \Phi_{F}^\alpha(x, y). \tag{27}
\]

If a specific Dirac projection of \( \Phi_{F}^\alpha(x, x) \) is nonvanishing, then the corresponding projection of \( \Phi_{A}^\alpha(x, x) \) has a pole, hence the name gluonic pole. An example is the function \( T(x, S_T) \equiv \pi \text{Tr} \left[ \Phi_{F}^\alpha(x, x) \epsilon T^\beta S_T^\beta \gamma_- \right] / P^+ = 2\pi i M S_T^2 G_F(x, x) \) discussed by Qiu and Sterman in Ref. [3].

In order to define Eq. (27) at the pole, one needs a prescription, which is related to the choice of boundary conditions on \( A_T^\alpha(\eta = \pm \infty) \) inside matrix elements. Possible inversions of \( F^{+\alpha} = \partial^+ A_T^\alpha \) are:

\[
A_T^\alpha(\eta n_-) = A_T^\alpha(\infty) - \int_{-\infty}^{\infty} dz^- \theta(z^- - \eta n_-) F^{+\alpha}(z^-)
\]

\[
= A_T^\alpha(-\infty) + \int_{-\infty}^{\infty} dz^- \theta(\eta n_- - z^-) F^{+\alpha}(z^-)
\]

\[
= \frac{A_T^\alpha(\infty) + A_T^\alpha(-\infty)}{2} - \frac{1}{2} \int_{-\infty}^{\infty} dz^- \epsilon(z^- - \eta n_-) F^{+\alpha}(z^-). \tag{28}
\]

One can use the representations for the \( \theta \) and \( \epsilon \) functions,

\[
\pm i\theta(\pm x) = \int \frac{dk}{2\pi} \frac{e^{ikx}}{k \mp i\epsilon}, \quad i\epsilon(x) = \int \frac{dk}{2\pi} P \frac{e^{ikx}}{k}, \tag{29}
\]

to obtain

\[
\Phi_{A}^\alpha(x, y) = \delta(x-y) \Phi_{A}^\alpha(\infty) + \frac{-i}{x-y+i\epsilon} \Phi_{F}^\alpha(x, y) \tag{30}
\]
\[
\delta(x - y) \Phi^\alpha_{A(\pm \infty)}(x) + \frac{-i}{x - y - i\epsilon} \Phi^\alpha_F(x, y)
\]
\[
= \delta(x - y) \Phi^\alpha_{A(\infty)}(x) + \frac{\Phi^\alpha_{A(-\infty)}(x)}{2} + P \frac{-i}{x - y} \Phi^\alpha_F(x, y),
\]
where
\[
\delta(x - y) \Phi^\alpha_{A(\pm \infty)}(x) \equiv \int \frac{d\lambda d\eta}{2\pi 2\pi} e^{i\lambda x} e^{i\eta(y - x)} \langle P, S | \bar{\psi}_j(0) g A^\alpha_T(\eta = \pm \infty) \psi_i(\lambda n_-) | P, S \rangle.
\]

So Eq. (31) shows the importance of boundary conditions in the inversion of Eq. (27), if matrix elements containing \( A^\alpha_T(\eta = \pm \infty) \) do not vanish. When such matrix elements vanish (implicitly assumed in [1]) the pole prescription does not matter. Also one obtains
\[
2\pi \Phi^\alpha_F(x, x) = \left[ \Phi^\alpha_{A(\infty)}(x) - \Phi^\alpha_{A(-\infty)}(x) \right],
\]
which shows the relation between the zero-momentum quark-gluon correlation function and the boundary conditions.

The behavior of \( \Phi^\alpha_{A(\pm \infty)}(x) \) under time-reversal is:
\[
\Phi^\alpha_{A(\pm \infty)}(x) = \gamma_5 C \Phi^\alpha_{A(\mp \infty)}(x) C^\dagger \gamma_5.
\]
This relation implies that time-reversal invariance only allows for symmetric or antisymmetric boundary conditions. Both situations (if nonvanishing) lead to a singularity in \( \Phi^\alpha_A(x, y) \) at the point \( x = y \), but only the antisymmetric case will be called a gluonic pole. The delta-function singularity in the case of nonvanishing symmetric boundary conditions will contribute to the functions \( \tilde{G}^\alpha_A(x, x) \) and \( H^\alpha_A(x, x) \) and hence, to T-even DFs. This would only affect the magnitude of double spin asymmetries. This case is also less interesting, because \( \Phi^\alpha_F(x, x) = 0 \).

We like to point out that so-called fermionic poles play a role in off-forward scattering, such as prompt photon production [14, 3, 4, 5]. Here a gluonic pole gives rise to an asymmetry proportional to \( T(x, S_T) g(\bar{x}) \) (see Fig. 3). Fermionic poles do not contribute in case of DY to this order.

Figure 3: A diagram yielding a single transverse spin asymmetry in prompt photon production
4 Effective T-odd distribution functions

To study the effect of gluonic poles we consider the (nonvanishing) antisymmetric boundary condition \( \Phi_{A(\infty)}(x) = -\Phi_{A(-\infty)}(x) \), which implies

\[
\pi \Phi_p^\alpha(x, x) = \Phi_{A(\infty)}(x),
\]

\[
\Phi_{A(\pm\infty)}(x) = -\gamma_5 C \Phi_{A(\pm\infty)\alpha}(x) C^\dagger \gamma_5. \tag{36}
\]

In the calculation of the cross-section one always encounters the pole of the matrix element (in this case in the principal value prescription) multiplied with the propagator in the hard subprocess (having a causal prescription),

\[
\Phi_{A}^{\text{eff}}(y, x) \equiv \frac{x - y}{x - y + i\epsilon} \Phi_{A}(y, x), \tag{37}
\]

\[
= \frac{-i}{x - y + i\epsilon} \Phi_{p}^\alpha(y, x), \tag{38}
\]

\[
= \Phi_{A}^\alpha(y, x) - \pi \delta(x - y) \Phi_{p}^\alpha(y, x). \tag{39}
\]

The time-reversal constraint applied to \( \Phi_{A}^\alpha(x, y) \) implies the analogue of Eq. (15), while \( \Phi_{p}^\alpha(x, y) \) has the opposite behavior under time-reversal compared to \( \Phi_{A}^\alpha(x, y) \). Thus for \( \Phi_{A}^{\text{eff}}(x, y) \) one does not have definite behavior under time-reversal symmetry. Specifically, the allowed T-even functions of \( \Phi_{p}^\alpha(x, x) \), \( G_{F}(x, x) \) and \( E_{F}(x, x) \), can be identified with T-odd functions in the effective correlation function \( \Phi_{A}^{\text{eff}} \). This implies that \( G_{A}^{\text{eff}}(x, y) \) and \( E_{A}^{\text{eff}}(x, y) \) will have an imaginary part and this gives rise to two "effective" time-reversal-odd DFs via the imaginary part of the e.o.m.

To say it again in a different way: by partial integration we find for instance

\[
G_{A}(x, y) = \frac{-i}{x - y} G_{F}(x, y). \tag{40}
\]

If one applies time-reversal invariance, \( G_{A}(x, y) \) will be a real function and \( G_{F}(x, y) \) imaginary. So one expects the pole prescription to be the principal value. But when convoluting the pole of the matrix element (with the principal value prescription) with the propagator in the hard subprocess (with a causal prescription), it is formally possible to shift the imaginary part from the pole of the latter to the pole of the former. This will effectively give rise to a causal prescription in Eq. (40), instead of a principal value (but without the additional boundary term required by time-reversal, cf. Eq. (30)). This implies that \( G_{A}(x, y) \) (and also \( E_{A}(x, y) \)) will effectively acquire an imaginary part.

For simplicity we neglect intrinsic tranverse momentum, thus we assume \( \Phi_{A(\infty)}(x) = \Phi_{D(\infty)}(x) \). So, by identification we have

\[
i\pi G_{F}(x, x) = \int dy \text{Im} G_{A}^{\text{eff}}(y, x), \tag{41}
\]

\[
i\pi E_{F}(x, x) = \int dy \text{Im} E_{A}^{\text{eff}}(y, x) \tag{42}
\]

and then it follows from the e.o.m. that

\[
x f_{T}^{\text{eff}}(x) = i\pi G_{F}(x, x) = \frac{1}{2MS_{T}^{2}} T(x, S_{T}), \tag{43}
\]

\[
x h^{\text{eff}}(x) = 2i\pi E_{F}(x, x) = \frac{-i\pi}{2MP^{+}} \text{Tr} [\Phi_{F}^{\alpha}(x, x) \gamma_{\tau a} \gamma_{\rho -}]. \tag{44}
\]

7
The function $e_L^{\text{eff}}$ receives no gluonic pole contribution, since time-reversal symmetry requires $H_F(x, x) = 0$.

Of course, the mechanism for generating finite projections of $\Phi_F^\rho(x, x)$ remains unknown. We just can conclude that if there is indeed a non-zero gluonic pole (in the case of non-zero antisymmetric boundary conditions), then at twist-three there are two non-zero “effective” T-odd DFs, namely $f_T$ and $h$. The first one generates the twist-three single spin asymmetry found by Hammon et al. [6], in their notation it is proportional to $T(x, x)$. The second one leads to a new asymmetry (see next section). Summarizing, we find for the T-even parametrization of $\Phi_{A(\infty)}^\alpha(x)$,

$$
\Phi_{A(\infty)}^\alpha(x) = -\frac{i x M}{2} \left[ f_T^{\text{eff}}(x) i \epsilon_T^\alpha S_T \beta P + \frac{1}{2} h^{\text{eff}}(x) \gamma_T^\alpha P \right].
$$

The antisymmetric nonvanishing boundary condition for $\Phi_{A(\pm \infty)}^\alpha(x)$ might arise from a linear A-field, giving a constant field strength (cf. e.g. [17, 18]). One might also think of an instanton background field. In both cases one should interpret infinity to mean ‘outside the proton radius’. Also, the constant field strength should be understood as an average value of the gluonic chromomagnetic field, which is non-zero due to a correlation with the direction of the proton spin. The large distance origin of the asymmetries arising from such a gluonic pole is apparent.

5 Single spin asymmetries in the Drell-Yan process

We will now discuss the single spin asymmetries in the Drell-Yan process in case one integrates over transverse photon momentum. So one uses the above parametrizations of the correlation functions in the expression for the integrated hadron tensor, which after contraction with the lepton tensor yields the cross-section.

$$
A_T = \frac{4 \sin(2\theta) \sin(\phi_{S_i}) |S_{1T}|}{1 + \cos^2 \theta} \times \sum_a e_a^2 \left[ M_1 x f_T^a(x) f_1^a(\bar{x}) + M_2 h_1^a(x) \bar{x} h_1^a(\bar{x}) \right] / \sum_a e_a^2 f_1^a(x) f_1^a(\bar{x}),
$$

Figure 4: Kinematics of the Drell-Yan process in the lepton center of mass frame

Under the assumption that $\Phi_{A(\infty)}^\alpha = \Phi_{D(\infty)}^\alpha$ we find the following single spin asymmetry (hadron-two unpolarized), given in the lepton center of mass frame:
where $\phi_{S_1}$ is the angle between $S_{1T}$ and the perpendicular part of the lepton momentum $l$, $\hat{\mu}_\perp \equiv \left( g^{\mu\nu} - \hat{t}^{\mu} \hat{t}^\nu + \hat{z}^{\mu} \hat{z}^\nu \right) l_\nu$. The first term in the asymmetry (proportional to $f_T$) is the one discussed in [3] (in their notation it is proportional to $T(x, x^T q(y)$), which will also be present in DIS ($f_1(\bar{x}) = \delta(1 - \bar{x})$). The second term is another, new single spin asymmetry arising in DY from a gluonic pole. It is not proportional to $T(x, S_T)$, but to a chiral-odd projection of $\Phi_\alpha^\rho$ in the point $x = y$, cf. Eq. (44).

6 Conclusions

We have demonstrated for the Drell-Yan process that the effects of so-called gluonic poles in twist-three hadronic matrix elements cannot be distinguished from those of T-odd distribution functions. Imaginary phases arising from hard subprocesses together with gluonic poles give rise to effective T-odd distribution functions. This leads to single spin asymmetries for the Drell-Yan process, such as the one found recently by Hammon et al. [6]. We have found a similar asymmetry arising from a gluonic pole, which involves chiral odd distribution functions. We have moreover shown that the presence of gluonic poles is in accordance with time-reversal invariance and requires large distance gluonic fields with antisymmetric boundary conditions.

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References

[1] R.D. Tangerman and P.J. Mulders, Phys. Rev. D 51 (1995) 3357 and hep-ph/9408305.
[2] M. Anselmino, M. Boglione and F. Murgia, Phys. Lett. B 362 (1995) 164.
[3] J. Qiu and G. Sterman, Phys. Rev. Lett. 67 (1991) 2264; Nucl. Phys. B 378 (1992) 52.
[4] V.M. Korotkiyan and O.V. Teryaev, Dubna preprint E2-94-200.
[5] A.V. Efremov, V.M. Korotkiyan and O.V. Teryaev, Phys. Lett. B 384 (1995) 577.
[6] N. Hammon, O. Teryaev and A. Schäfer, Phys. Lett. B 390 (1997) 409.
[7] D. Boer, P.J. Mulders and O.V. Teryaev, hep-ph/9710223.
[8] D.E. Soper, Phys. Rev. D 15 (1977) 1141; Phys. Rev. Lett. 43 (1979) 1847.
[9] J.P. Ralston and D.E. Soper, Nucl. Phys. B 152 (1979) 109.
[10] J.C. Collins and D.E. Soper, Nucl. Phys. B 194 (1982) 445.
[11] H.D. Politzer, Nucl. Phys. B 172 (1980) 349.
[12] R.K. Ellis, W. Furmanski and R. Petronzio, Nucl. Phys. B 212 (1983) 29.
[13] A.V. Efremov and O.V. Teryaev, *Sov. J. Nucl. Phys.* **39** (1984) 962.

[14] R.L. Jaffe and X. Ji, *Phys. Rev. Lett.* **67** (1991) 552; *Nucl. Phys. B* **375** (1992) 527.

[15] O.V. Teryaev, in *Proceedings of the 12th International Symposium on High-Energy Spin Physics*, C.W. de Jager *et al.* (eds), (World Scientific, 1997), p. 594.

[16] A.V. Efremov and O.V. Teryaev, *Phys. Lett. B* **150** (1985) 383.

[17] A. Schäfer, L. Mankiewicz, P. Gornicki and S. Güllenstern, *Phys. Rev. D* **47** (1993) R1.

[18] B. Ehrnsperger, A. Schäfer, W. Greiner, L. Mankiewicz, *Phys. Lett. B* **321** (1994) 121.