

Sivers effect in the Drell-Yan process at BNL RHICJ. C. Collins,^{1,3} A. V. Efremov,² K. Goetze,³ M. Grosse Perdekamp,^{4,5} S. Menzel,³ B. Meredith,⁴
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On the basis of a fit to the Sivers effect in deep-inelastic scattering, we make predictions for single-spin asymmetries in the Drell-Yan process at BNL RHIC.

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I. INTRODUCTION

The Sivers effect [1] was originally suggested to explain the large single spin asymmetries (SSA) observed in $p^\dagger p \rightarrow \pi X$ (and $\bar{p}^\dagger p \rightarrow \pi X$) at FNAL [2] and recently at higher energies in the RHIC experiment [3]. The effect considers a nontrivial correlation between (the transverse component of) the nucleon spin \mathbf{S}_T and intrinsic transverse parton momenta \mathbf{p}_T in the nucleon. It is proportional to the “(naively or artificially) T-odd” structure $(\mathbf{S}_T \times \mathbf{p}_T) \cdot \mathbf{p}_N$ and quantified in terms of the Sivers function $f_{1T}^\perp(x, \mathbf{p}_T^2)$ [4] whose precise definition in QCD was worked out only recently [5–7].

A particularly interesting feature of the Sivers function (and other “T-odd” distributions) concerns the universality property. On the basis of time-reversal arguments it was predicted [6] that f_{1T}^\perp in semi-inclusive deeply inelastic scattering (SIDIS) and the Drell-Yan process (DY) have opposite sign (our definition of the Sivers function follows the Trento Conventions [8]),

$$f_{1T}^\perp(x, \mathbf{p}_T^2)_{\text{SIDIS}} = -f_{1T}^\perp(x, \mathbf{p}_T^2)_{\text{DY}}. \quad (1)$$

The experimental check of Eq. (1) would provide a thorough test of our understanding of the Sivers effect within QCD. In particular, the experimental verification of (1) is a crucial prerequisite for testing the factorization approach to the description of processes containing p_T -dependent correlators [9–11].

Recent data on SSA from SIDIS [12–14], and, in particular, those from transversely polarized targets [15–17], provide first measurements of the Sivers effect in SIDIS. On the basis of this information it was shown that the Sivers effect leads to sizeable SSA in $p^\dagger \pi^- \rightarrow l^+ l^- X$, which could be studied at COMPASS [18], and in $p^\dagger \bar{p} \rightarrow l^+ l^- X$ or $p \bar{p}^\dagger \rightarrow l^+ l^- X$ in the planned PAX experiment at GSI [19,20], making the experimental check of Eq. (1) feasible and promising [21]. Both experiments, which could be performed in the medium or long term, have the advantage of being dominated by annihilation of valence quarks (from p) and valence antiquarks (from \bar{p} or π^-), which yields sizeable counting rates. Moreover, the pro-

cesses are not very sensitive to the Sivers antiquark distributions, which are not constrained by the present SIDIS data, see [21–25].

On a shorter term the Sivers effect in DY can be studied at RHIC in $p^\dagger p \rightarrow l^+ l^- X$. Similar and other spin physics prospects at RHIC are discussed in [26–32]. Other earlier predictions for SSAs in DY were given in [33,34].

In pp -collisions inevitably antiquark distributions are involved, and the counting rates are smaller. In this work we shall demonstrate that the Sivers effect SSA in DY can nevertheless be measured at RHIC with an accuracy sufficient to unambiguously test Eq. (1). In particular, by focusing on certain kinematic regions the effect of the unknown Sivers antiquark distribution function can be minimized. And, by focusing on the opposite kinematic regions one can gain first information on the Sivers antiquark distribution itself. For our estimates we use the Sivers function extracted from HERMES data [16] in [22], see also [35,36].

It remains to be noted that the theoretical understanding of SSA in $p^\dagger p \rightarrow \pi X$, which originally motivated the introduction of the Sivers effect, is more involved and less lucid compared to SIDIS or DY, as here also other mechanisms such as the Collins effect [37] and/or dynamical twist-3 effects [38,39] could generate SSA. Phenomenological studies indicate, however, that in a picture based on p_T -dependent correlators the data [2,3] can be explained in terms of the Sivers effect alone [40–42] with the other effects playing a less important role [43,44]. For recent discussions of hadron-hadron collisions with more complicated final states (like, e.g., $p^\dagger p \rightarrow \text{jet}_1 \text{jet}_2 X$) we refer to Refs. [45].

II. SIVERS EFFECT IN SIDIS

The longitudinal SSA in SIDIS observed first [12–14] unfortunately cannot be unambiguously interpreted in terms of a unique (Sivers [1], Collins [37] or twist-3 [46]) effect [47–51], see Ref. [52] for a recent review. All that is clear at the present stage is that these SSAs are dominated by subleading twist effects [53]. The situ-

ation changed, however, with the transverse SSA in SIDIS observed at HERMES and COMPASS [15–17], where the Siverson and Collins effect can be distinguished due to the different azimuthal angle distribution of the produced hadrons [4].

In particular, in the SSA due to the Siverson effect the produced hadrons exhibit an azimuthal distribution $\propto \sin(\phi - \phi_S)$ in the plane transverse to the beam axis. Here ϕ and ϕ_S are the azimuthal angles of, respectively, the produced hadron and of the target polarization vector, with respect to the lepton scattering plane.

The Siverson SSA measured at HERMES [16] is defined as sum over SIDIS events i according to

$$A_{UT}^{\sin(\phi - \phi_S)} = \frac{\sum_i \sin(\phi_i - \phi_{S,i}) \{N^{\uparrow}(\phi_i, \phi_{S,i}) - N^{\downarrow}(\phi_i, \phi_{S,i} + \pi)\}}{\frac{1}{2} \sum_i \{N^{\uparrow}(\phi_i, \phi_{S,i}) + N^{\downarrow}(\phi_i, \phi_{S,i} + \pi)\}}. \quad (2)$$

$N^{\uparrow(\downarrow)}(\phi_i, \phi_{S,i})$ are the event counts for the respective target polarization (corrected for depolarization effects). In order to describe the Siverson SSA defined in Eq. 2 in Ref. [22] two major simplifications are made. First, soft factors [9–11] are neglected. Second, for the distribution of transverse momenta in $D_1^q(z, K_T^2)$ and $f_{1T}^{\perp q}(x, p_T^2)$ the Gaussian model is assumed.

The Gaussian model certainly oversimplifies the description of “unintegrated” distribution or fragmentation functions which are an involved issue in QCD [54]. On a longer term, an approach to the p_T -dependence of the Siverson SSA along the lines of the formalism in Ref. [55] would be desirable. The Gaussian model, however, provides a good effective description of SIDIS and DY data within a certain range of low transverse momenta, and is sufficient for the purpose of the present work. The free parameters, namely, the Gaussian widths, are consistently constrained in [22] by the HERMES data.

As regards the flavor dependence of the Siverson functions, there are no strong constraints from the present SIDIS data [16,17]. In fact, in Ref. [23] where this has been attempted, all fitted distributions but $f_{1T}^{\perp u}$ were found consistent with zero. In this situation it is appealing to invoke additional theoretical constraints. In particular, here we use predictions from the QCD limit of a large number of colors N_c .

In this limit the nucleon appears as N_c quarks bound by a mean field [56], which exhibits certain spin-flavor symmetries [57]. By exploring these symmetry properties it was proven in a model independent way that in the large- N_c limit [58]

$$f_{1T}^{\perp u}(x, \mathbf{p}_T^2) = -f_{1T}^{\perp d}(x, \mathbf{p}_T^2) \quad \text{modulo } 1/N_c \text{ corrections}, \quad (3)$$

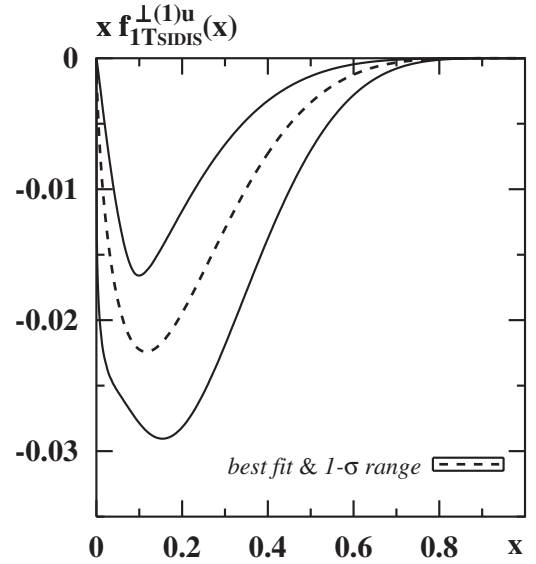


FIG. 1. The u -quark Siverson function $x f_{1T}^{\perp(1)u}(x)$ vs x at a scale of about 2.5 GeV^2 , as obtained from a fit to the HERMES data [16]. Shown are the best fit and its $1\text{-}\sigma$ uncertainty.

for not too small and not too large x satisfying $xN_c = \mathcal{O}(N_c^0)$. Analog relations hold for the Siverson antiquark distributions.¹ Imposing the large- N_c relation (3) as an additional constraint, and neglecting effects of antiquarks and heavier flavors, it is shown [22] that the HERMES data [16] can be described by the following 2-parameter ansatz and best fit

$$f_{1T}^{\perp(1)u}(x) = -f_{1T}^{\perp(1)d}(x) \stackrel{\text{ansatz}}{=} Ax^b(1-x)^5 \stackrel{\text{fit}}{=} -0.17x^{0.66}(1-x)^5, \quad (4)$$

with a χ^2 per degree of freedom of about 0.3. The fit and its $1\text{-}\sigma$ uncertainty due to the statistical error of the data [16] are shown in Fig. 1. Several comments are in order concerning the ansatz and fit result (4).

The fit (4) refers to a scale of 2.5 GeV^2 roughly set by the $\langle Q^2 \rangle$ in the HERMES experiment [16]. For the extraction we used the parametrizations for $f_1^q(x)$ and $D_1^q(x)$ from [62,63] at the corresponding scale.

Within the Gaussian model, of course, one does not need to work with the “transverse moment” of the Siverson function defined as, and in the Gaussian ansatz given by

¹For historical correctness we mention that previously (3) was discussed in the framework of (simple versions of) chiral models [59]. But the way in which (3) was obtained there was shown to be incorrect [60]. Recently, in Ref. [61] (a more sophisticated version of) a chiral model with vector mesons obeying a hidden local flavor symmetry was discussed, in which the Siverson function would obey (3).

$$f_{1T}^{\perp(1)a}(x) \equiv \int d^2\mathbf{p}_T \frac{\mathbf{p}_T^2}{2M_N^2} f_{1T}^{\perp a}(x, \mathbf{p}_T^2) \stackrel{\text{Gauss}}{=} \frac{\langle p_T^2 \rangle_{\text{Siv}}}{2M_N^2} f_{1T}^{\perp a}(x), \quad (5)$$

where $\langle p_T^2 \rangle_{\text{Siv}}$ denotes the Gaussian width of the Siverson function. However, doing so one benefits from the fact that the fit for the moment $f_{1T}^{\perp(1)a}(x)$ (in contrast to $f_{1T}^{\perp a}(x, \mathbf{p}_T^2)$) is nearly insensitive [22] to the value of $\langle p_T^2 \rangle_{\text{Siv}}$ which is poorly constrained by the data and the positivity bound [64].

The shape of the Siverson function at large x is not constrained by the data [16]. Our ansatz of $(1-x)^5$ dependence has some theoretical justification [21], although there are also arguments [24] for a $(1-x)^4$ behavior.

We remark that the result (4) is in good agreement with the fit of Ref. [21] to the *preliminary* HERMES data [15] on the Siverson SSA weighted with a power of the transverse momentum $P_{h\perp}$ of the produced hadron. With such a weighting the SSA can be interpreted (neglecting soft factors) unambiguously in terms of $f_{1T}^{\perp(1)a}(x)$ independently of any model for transverse momentum dependence [4]. Keeping in mind the *preliminary* status of the data [15], this agreement indicates the consistency of the Gaussian ansatz within the accuracy of the data.

The large- N_c motivated ansatz (4) is confirmed by results from the COMPASS experiment [17] where a solid polarized ${}^6\text{LiD}$ target [65,66] was used. Neglecting nuclear binding effects and exploring isospin symmetry, in deuterium $f_{1T}^{u/D} \approx f_{1T}^{u/p} + f_{1T}^{u/n} \approx f_{1T}^u + f_{1T}^d$ and analogously for d , \bar{q} , etc. Thus, the deuterium target is sensitive to the flavor combination which is suppressed in the large- N_c limit, see (3), and for which our ansatz (4) yields zero. This is in agreement with the COMPASS data [17] showing a Siverson effect from deuterium target compatible with zero within error bars.

The reason why the large N_c picture of the Siverson function works at the present stage of art, is due to the fact that the current precision of the first data [16,17] is comparable to the theoretical accuracy of the large- N_c relation (3). Thus, $1/N_c$ corrections (and antiquarks effects) cannot be resolved within the error bars of the data [16,17]. In future, when the precision of the data will increase, it will certainly be necessary to refine the ansatz (4) to include $1/N_c$ corrections and antiquark effects.

The fit (4) described above is to the final HERMES data in Ref. [16]. We observe that the fit is compatible within the $1-\sigma$ range with more precise, but *preliminary*, data that has recently become available [67]. Incorporating into our fit these data, which are not yet corrected for smearing and acceptance effects, would tend to increase somewhat the Siverson function, cf. [24,25,35,36].

The sign of $f_{1T}^{\perp(1)u}_{\text{SIDIS}}(x)$ in Eq. (4) is in agreement with the physical picture of the Siverson functions discussed in Ref. [68].

We remark that, in principle, one could include into the fit in addition to SIDIS data also the data on SSA in the hadronic processes $p^\uparrow p \rightarrow \pi X$ or $p^\uparrow \bar{p} \rightarrow \pi X$ [2,3]. This possibility was not explored in [22] as these SSA could also be due to other mechanisms [37–39]. Furthermore, as shown in [45] it is not clear that factorization holds in these processes in terms of p_T -dependent correlators, for reasons that do not apply to SIDIS and to the Drell-Yan process. One also has to keep in mind that it is *a priori* not clear whether the leading twist factorization approach is accurate at the $\langle Q^2 \rangle = 2.5 \text{ GeV}^2$ of the HERMES experiment as assumed in [22]. Only careful analyses (which will include soft factors) of future data from experiments performed at different Q^2 will reveal to what extent this assumption is justified.

III. SIVERS EFFECT IN SSA IN DY AT RHIC

The process $p^\uparrow p \rightarrow l^+ l^- X$ is characterized by the variables $s = (p_1 + p_2)^2$, the invariant mass of the lepton pair $Q^2 = (k_1 + k_2)^2$, and the rapidity

$$y = \frac{1}{2} \ln \frac{p_2 \cdot (k_1 + k_2)}{p_1 \cdot (k_1 + k_2)}, \quad (6)$$

where $p_{1/2}$ (and $k_{1/2}$) indicate the momenta of the incoming proton (and the outgoing lepton) pair.

Let us consider the azimuthal SSA defined as a sum over the events i according to

$$A_{\text{UT}}^{\sin(\phi - \phi_S)} = \frac{\sum_i \sin(\phi_i - \phi_{S,i}) \{N^\uparrow(\phi_i, \phi_{S,i}) - N^\downarrow(\phi_i, \phi_{S,i} + \pi)\}}{\frac{1}{2} \sum_i \{N^\uparrow(\phi_i, \phi_{S,i}) + N^\downarrow(\phi_i, \phi_{S,i} + \pi)\}}, \quad (7)$$

where \uparrow, \downarrow denote the transverse polarizations of the proton, the polarized proton moves in the positive z -direction, and $(\phi - \phi_S)$ is the azimuthal angle between the virtual photon and the polarization vector. Neglecting again soft factors and assuming the Gaussian model for the distribution of transverse momenta, to leading order the SSA is given by

$$A_{\text{UT}}^{\sin(\phi - \phi_S)} = 2 \frac{a_{\text{Gauss}}^{\text{DY}} \sum_a e_a^2 x_1 f_{1T}^{\perp(1)a}(x_1) x_2 f_1^{\bar{a}}(x_2)}{\sum_a e_a^2 x_1 f_1^a(x_1) x_2 f_1^{\bar{a}}(x_2)}, \quad (8)$$

where $a = u, \bar{u}, d, \bar{d}$, etc. and the parton momenta $x_{1,2}$ are given by $x_{1,2} = (Q^2/s)^{1/2} e^{\pm y}$. The dependence on the Gaussian model is contained in the factor

$$a_{\text{Gauss}}^{\text{DY}} = \frac{\sqrt{\pi}}{2} \frac{M_N}{\sqrt{\langle p_T^2 \rangle_{\text{Siv}} + \langle p_T^2 \rangle_{\text{unp}}}}. \quad (9)$$

If one introduced in the numerator of (7) an additional weight q_T/M_N , where q_T denotes the modulus of the

transverse momentum of the lepton pair with respect to the collision axis, the resulting SSA would be independent of a particular model for transverse momentum distribution and given by the expression on the right-hand-side of (8) but without the factor $a_{\text{Gauss}}^{\text{DY}}$. It was argued that such a ‘‘transverse momentum weighted’’ SSA might be protected against Sudakov dilution effects [69]. These effects need not be negligible, considering the fact that at RHIC we deal with SSA at considerably higher scales than in the HERMES experiment: $(4 \text{ GeV})^2 \leq Q^2 \leq (20 \text{ GeV})^2$ typically at RHIC [26] vs $\langle Q^2 \rangle = 2.5 \text{ GeV}^2$ at HERMES.

The DY process in pp -collisions is sensitive to $f_{1\text{T DY}}^{\perp(1)q}(x_1)f_1^{\bar{q}}(x_2)$ and to $f_{1\text{T DY}}^{\perp(1)\bar{q}}(x_1)f_1^q(x_2)$ on an equal footing. Thus, even though the effects of the Sivers antiquark distributions seem small and presently not observable in SIDIS, one cannot expect this to be the case in DY at RHIC.

In order to have a rough idea which regions of y (for a given Q^2) could be more and which regions less sensitive to Sivers antiquark distribution functions, let us introduce two *models*:

$$f_{1\text{T DY}}^{\perp(1)\bar{q}}(x) = \epsilon(x)f_{1\text{T DY}}^{\perp(1)q}(x), \quad (10)$$

with

$$\epsilon(x) = \begin{cases} \pm 0.25 = \text{const} & \text{model I} \\ \frac{(f_1^u + f_1^d)(x)}{(f_1^u - f_1^d)(x)} & \text{model II,} \end{cases} \quad (11)$$

where for $f_{1\text{T DY}}^{\perp(1)q}(x)$ we will use the result from Eq. (4) taking into account the change of sign in Eq. (1).

Model I is that the Sivers antiquark distribution is $\pm 25\%$ of the corresponding Sivers quark distribution at all x . Model II is that the ratio of the Sivers antiquark over quark distribution is fixed by the ratio of the unpolarized distribution functions. The particular ansatz for model II ensures compatibility with the large- N_c limit, and it preserves the following sum rule [70]—see also [21]—

$$\sum_{a=g,u,d,\dots} \int dx f_{1\text{T}}^{\perp(1)a}(x) = 0. \quad (12)$$

(Note that in the large- N_c limit the gluon Sivers distribution is suppressed with respect to the quark one [21].)

Equations (10) and (11) represent rough models for $f_{1\text{T DY}}^{\perp(1)\bar{q}}(x)$ which are, however, consistent with the HERMES data [16], and compatible with the presently known theoretical constraints, namely, the large- N_c relations (3), positivity constraints [64] and the Burkardt-sum rule (12). This makes them sufficient for our purposes.

In order to present estimates for RHIC we strictly speaking should use the fitted Sivers function (and $f_1^a(x)$) evolved to the relevant scale. Instead, let us assume that

$$\frac{\sum_a e_a^2 f_{1\text{T DY}}^{\perp(1)a} f_1^{\bar{a}}}{\sum_a e_a^2 f_1^a f_1^{\bar{a}}} \Big|_{\text{HERMES } Q^2} \approx \frac{\sum_a e_a^2 f_{1\text{T DY}}^{\perp(1)a} f_1^{\bar{a}}}{\sum_a e_a^2 f_1^a f_1^{\bar{a}}} \Big|_{\text{RHIC } Q^2}. \quad (13)$$

It is difficult to quantify exactly the error introduced in this way. However, we believe it to be smaller than other uncertainties in our analysis.

When dealing with unintegrated distribution functions at higher energies one must take into account a broadening of the average transverse momentum. Considering that this effect due to radiating off soft gluons does not depend on polarization, we roughly estimate

$$\begin{aligned} \langle p_T^2 \rangle_{\text{Siv}}^{\text{RHIC}} &\approx 2 \langle p_T^2 \rangle_{\text{Siv}}^{\text{HERMES}}, \\ \langle p_T^2 \rangle_{\text{unp}}^{\text{RHIC}} &\approx 2 \langle p_T^2 \rangle_{\text{unp}}^{\text{HERMES}} \approx 0.66 \text{ GeV}^2. \end{aligned} \quad (14)$$

Similar values for $\langle p_T^2 \rangle_{\text{unp}}$ at high energies were obtained in [42]. Next, let us assume that the transverse moment $f_{1\text{T DY}}^{\perp(1)a}$ is little affected by that. The combined effect of our assumptions is that the SSA decreases at higher energies – in qualitative agreement with [69], though not as fast as discussed there.

We recall, that from the HERMES data [16] and the positivity bound [64] the Gaussian width of the Sivers function is only poorly constrained. Under the assumption (14) one may expect it at RHIC to be in the range $\langle p_T^2 \rangle_{\text{Siv}}^{\text{RHIC}} \approx (0.20 \dots 0.64) \text{ GeV}^2$ [22]. However, the only place where this value numerically matters is the Gauss factor in Eq. (9). Varying $\langle p_T^2 \rangle_{\text{Siv}}^{\text{RHIC}}$ in the above range yields $a_{\text{Gauss}}^{\text{DY}} = 0.81 \cdot (1 \pm 10\%)$, i.e. it alters our estimates for RHIC only within $\pm 10\%$.

On the basis of the above assumptions and taking into account the change of sign for the Sivers distribution function in Eq. (1) we obtain the results shown in Fig. 2. We observe that the Drell-Yan SSA is noticeably sensitive to the Sivers \bar{q} -distribution functions, as modeled in Eq. (10). The effect is more pronounced at larger dilepton masses Q . However, we note that at $Q = 4 \text{ GeV}$ in the region of positive rapidities y the effect is moderate, and *does not* alter the sign of the asymmetry.

The region of negative rapidities y is strongly sensitive to the Sivers \bar{q} -distribution. For positive $\epsilon(x)$ in Eqs. (10) and (11) the SSA is positive, for negative $\epsilon(x)$ it is negative. As already mentioned, the effect of Sivers- \bar{q} is more pronounced at larger Q^2 . Thus, by focusing on these regions of y , Q^2 one could gain the first information on the magnitude and sign of the Sivers- \bar{q} distributions.

In order to see explicitly in which kinematical region our predictions are constrained by data, we show in Fig. 2 the x -range covered in the HERMES experiment which constrained the fit of the Sivers function. In this context it is worthwhile remarking that our estimates for RHIC based on SIDIS data are complementary to those made in Ref. [31]. There information on the Sivers function was used from the data on SSA in $p^{\uparrow}p \rightarrow \pi X$ [2]. Assuming factorization for this process and considering that other mechanisms cannot explain the observed SSA at large x_F [43,44], one finds that the data [2] constrain the Sivers function at large $x > 0.4$ above the region explored at

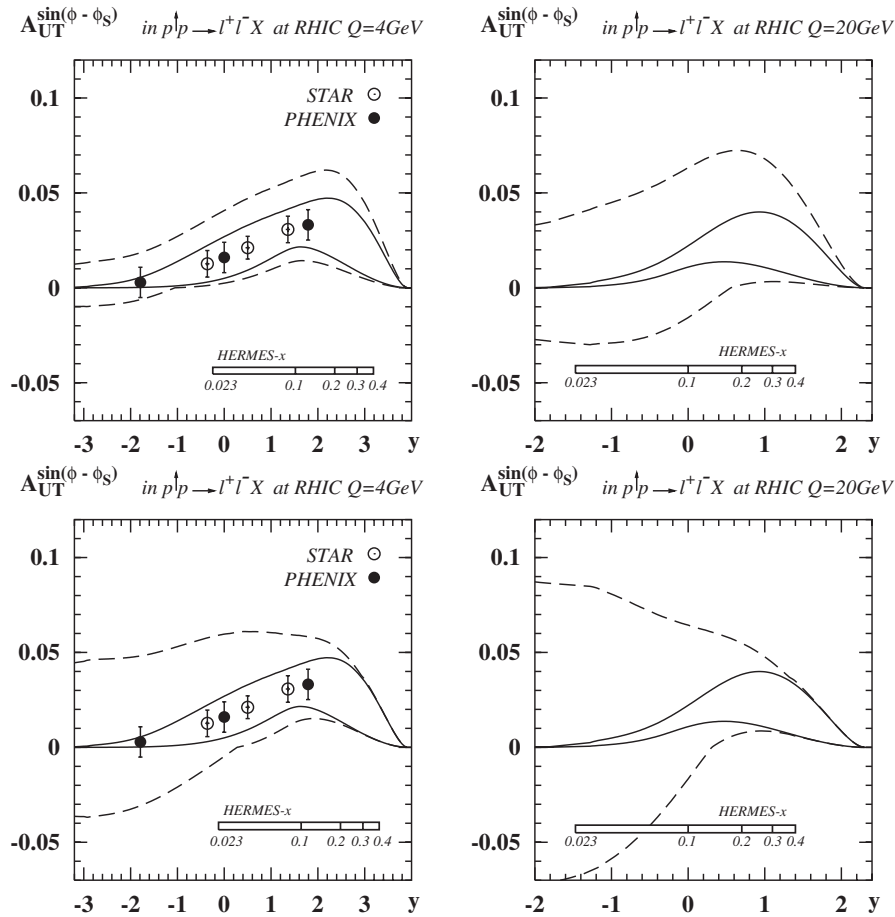


FIG. 2. The azimuthal SSA $A_{\text{UT}}^{\sin(\phi - \phi_S)}$ in Drell-Yan lepton pair production, $p^{\uparrow}p \rightarrow l^+l^-X$, as function of y for the kinematics of the RHIC experiment with $\sqrt{s} = 200$ GeV. The left plots show $Q^2 = (4 \text{ GeV})^2$, and the right plots $Q^2 = (20 \text{ GeV})^2$. The upper plots show the effects of varying the Sivers antiquark distributions within the range of model I—see Eqs. (10) and (11). The lower plots show instead the results of model II. In all cases, the central estimates are based on our fit (4) for the Sivers quark distribution functions from the HERMES data [16]. The inner error band (solid lines) shows the $1\text{-}\sigma$ uncertainty of the fit. The outer error bands (dashed lines) show the error from varying the Sivers antiquark distribution functions within the ranges specified in Eqs. (10) and (11), model I for the upper plots, and model II for the lower plots. The x -region explored in the HERMES kinematics is indicated. For $Q = 4$ GeV we show the estimated statistical error for STAR and PHENIX. For $Q = 20$ GeV the statistical error at STAR and PHENIX is comparable to the asymmetry. See text and Table I for a discussion of the planned RHIC II.

HERMES. Thus, in the region of large positive y , where our estimates are not constrained by SIDIS data, see Fig. 2, the estimates of Ref. [31] could be more reliable.

In Table I we show the statistical errors δA for the Sivers SSA in DY estimated with PYTHIA considering the ac-

ceptance of the STAR and PHENIX detectors. Detector acceptance is conveniently specified in terms of pseudorapidity $\eta = \ln(\tan \frac{\theta}{2})$, which is directly related to the scattering angle θ of the lepton pair with respect to the beam-axis, and thus to the geometry of the detector. In the

TABLE I. Statistical errors δA for the Sivers SSA in Drell-Yan for the PHENIX and STAR detectors at RHIC: Errors are shown for dilepton masses of $Q = 4$ GeV and 20 GeV assuming an integrated luminosity of $\int L dt = 125 \text{ pb}^{-1}$ and a beam polarization of $P = 0.7$. Error estimates have been carried out using the event generator PYTHIA. Projected errors are also shown for a possible future dedicated experiment for transverse spin with large acceptance at RHIC II (luminosity upgrade); see text for details.

	STAR	PHENIX		RHIC II					
	y	4 GeV	20 GeV	y	4 GeV	20 GeV	y	4 GeV	20 GeV
δA	-0.5	0.007	0.09	-1.8	0.008	0.2	± 2.5	0.003	0.03
	0.5	0.006	0.06	0.0	0.017	0.13	± 1.5	0.001	0.01
	1.5	0.007	0.11	1.8	0.008	0.2	± 0.5	0.001	0.01
$\int L dt$		125 pb^{-1}			125 pb^{-1}			$10 \times 125 \text{ pb}^{-1}$	

following acceptance cuts imposed on leptons will be given in pseudorapidity. However, asymmetries and their errors are analyzed in bins of photon rapidity y . We assume an integrated luminosity $\int L dt = 125 \text{ pb}^{-1}$ and a beam polarization of $P = 0.7$. We use these parameters as an upper estimate for the statistics these experiments could acquire with transverse beam polarization before RHIC detector and luminosity upgrades will become available in 2012 (RHIC II). The statistical errors can be easily scaled to different parameters for integrated luminosity and beam polarization.

For STAR, which covers the range $-1 < \eta < 2$ for e^+e^- , estimates for δA are presented for bins centered at $y = -0.5, 0.5$ and 1.5 . For PHENIX, which covers $|\eta| < 0.35$ (for e^+e^-) and $1.2 < |\eta| < 2.4$ (for $\mu^+\mu^-$) we have chosen bins centered at $y=0$ and $|y|=1.8$. The bins in the dilepton mass Q are chosen, respectively, at 4 GeV and 20 GeV.

For $Q = 4$ GeV the Siverson SSA can be measured at STAR and PHENIX. For illustrative purposes the estimated δA for STAR and PHENIX are shown in Fig. 2. The δA at $Q = 20$ GeV is of the order of magnitude of the asymmetry itself.

The region of higher Q could, however, be addressed taking advantage of the higher luminosity available at RHIC II. We consider a new large acceptance experiment dedicated to Drell-Yan physics with transverse spin. The new detector could be located at one of the RHIC interaction regions which are not equipped with spin rotator magnets and therefore always provide proton-proton collisions with transverse beam polarization.

We assume that at RHIC II the luminosities will be higher by a factor 2.5 through electron cooling and that an additional factor 4 in luminosity will result for the new experiment from special focusing magnets close to the interaction region. We expect that for a multiyear run with the new RHIC II detector an integrated luminosity of $10 \times 125 \text{ pb}^{-1}$ becomes easily accessible. The proposed detector covers an acceptance of $|\eta| < 3.0$. In Table I we present estimates for δA for bins centered at $|y| = 2.5, 1.5, 0.5$. Clearly, RHIC II could access the Siverson effect also at large dilepton masses, where the effects of Siverson antiquarks are expected to be more pronounced.

A ‘‘symmetric’’ (with respect to η) detector will allow to gain a factor 2 in statistics for SSA compared to double spin asymmetries, as either beam polarization can be used. This additional factor is not reflected in the estimates presented in Table I.

IV. q_T -DEPENDENCE OF SSA

In the previous Sections we have discussed SSA integrated over transverse momenta, namely, over q_T of the lepton pair in the DY process (or, $P_{h\perp}$ of the produced hadron in SIDIS). However, SSA as functions of q_T are equally interesting observables.

In the case of the DY Siverson SSA (8) we were able to minimize the model dependence to a certain extent (within the Gaussian model), e.g., by showing that the estimated SSA varied only within $\pm 10\%$ when the Gaussian width of the Siverson function was allowed to vary over a wide range $\langle p_T^2 \rangle_{\text{Siv}}^{\text{RHIC}} \approx (0.20 \dots 0.64) \text{ GeV}^2$. When discussing the DY Siverson SSA as a function of q_T our results are model-dependent. Nevertheless certain features are of a general character and it is worthwhile addressing them here. In what follows we assume for illustrative purposes that $\langle p_T^2 \rangle_{\text{Siv}}^{\text{RHIC}} \approx 0.3 \text{ GeV}^2$.

It is worthwhile stressing that unpolarized DY data on the q_T dependence can be well described by means of the Gaussian model up to $q_T \lesssim (2-3) \text{ GeV}$ (depending to some extent on the invariant mass Q and the kinematics of the respective experiment, see the detailed study in Ref. [42]). (For a description of the q_T dependence of unpolarized DY in terms of the QCD-based formalism of [55] we refer to the work [71] and references therein.) It is by no means clear whether the transverse parton momentum dependence of the Siverson function can be described equally satisfactory by the Gauss ansatz. This, in fact, is among what we are going to learn from RHIC and other experiments.

Let us first mention an unrealistic property of the Gauss model. Equation (8) was obtained upon integrating the relevant transverse momentum dependent cross sections over q_T from 0 to ∞ . In QCD the corresponding diagrams diverge, and physically it is clear that the large q_T must be cutoff at a scale $\sim \mathcal{O}(Q)$ [54].

Experimentally no artificial large- q_T cutoff is needed, since correct kinematics imposes a kinematic limit. What is more interesting in our context is to impose a low- q_T cutoff. This makes sense, and could be even preferable, because it could allow to increase the asymmetry, for the Siverson SSA has a kinematical zero, i.e. it is $\text{SSA} \propto q_T$.

Let us define the q_T -dependent DY Siverson SSA as

$$A_{\text{UT}}^{\sin(\phi - \phi_S)}(q_T) = \frac{\text{num}(q_T)}{\text{den}(q_T)} \quad (15)$$

with the numerator and denominator given by

$$\begin{aligned} \text{num}(q_T) &= \frac{\sigma_{\text{UT}}^{\text{Siv}} M_N}{(\langle p_T^2 \rangle_{\text{Siv}}^{\text{RHIC}} + \langle p_T^2 \rangle_{\text{unp}}^{\text{RHIC}})^2} q_T \\ &\quad \times \exp\left(-\frac{q_T^2}{\langle p_T^2 \rangle_{\text{Siv}}^{\text{RHIC}} + \langle p_T^2 \rangle_{\text{unp}}^{\text{RHIC}}}\right), \\ \text{den}(q_T) &= \frac{\sigma_{\text{UU}}}{2\langle p_T^2 \rangle_{\text{unp}}^{\text{RHIC}}} \exp\left(-\frac{q_T^2}{2\langle p_T^2 \rangle_{\text{unp}}^{\text{RHIC}}}\right), \end{aligned} \quad (16)$$

where

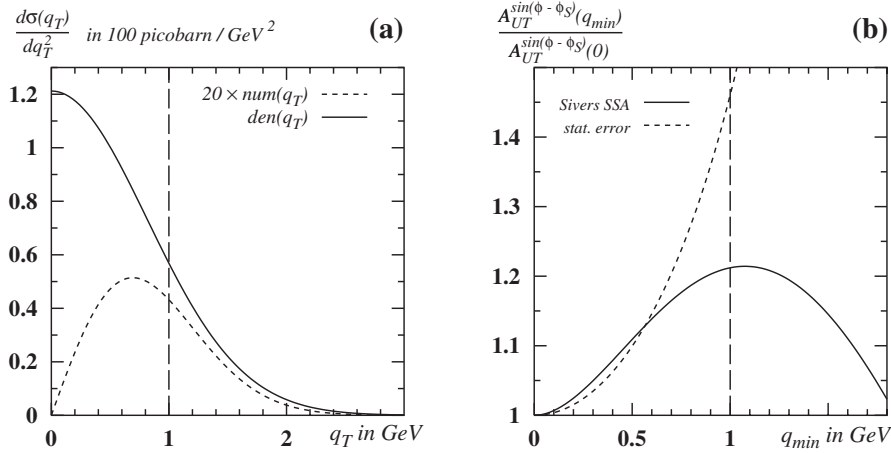


FIG. 3. a) The numerator and denominator, see Eqs. (15)–(17), for the azimuthal SSA $A_{UT}^{\sin(\phi-\phi_s)}$ in the DY process as functions of the dilepton transverse momentum q_T for $1.2 < y < 2.4$ at RHIC with $\sqrt{s} = 200$ GeV and $Q^2 = (4 \text{ GeV})^2$. The absolute numbers for the cross sections are somewhat altered by the assumption (13). b) The azimuthal SSA $A_{UT}^{\sin(\phi-\phi_s)}$ and its estimated statistical uncertainty as functions of the low- q_T cut q_{\min} , i.e. the ratio of the numerator and denominator in Fig. 3(a) integrated, respectively, over $q_T \in [q_{\min}, \infty]$. Both the SSA and its uncertainty are normalized with respect to their values for $q_{\min} = 0$.

$$\sigma_{UT}^{\text{Siv}} = \iint_{\text{cuts}} dQ^2 dy \frac{4\pi\alpha^2}{9Q^4} \sum_a e_a^2 x_1 f_{1TDY}^{\perp(1)a}(x_1) x_2 f_1^{\bar{a}}(x_2),$$

$$\sigma_{UU} = \iint_{\text{cuts}} dQ^2 dy \frac{4\pi\alpha^2}{9Q^4} \sum_a e_a^2 x_1 f_1^a(x_1) x_2 f_1^{\bar{a}}(x_2).$$
(17)

Integrating den (or num) over q_T^2 from 0 to ∞ yields the total unpolarized DY cross section σ_{UU} (or $a_{\text{Gauss}}^{\text{DY}} \sigma_{UT}^{\text{Siv}}$).

Considering a polarization-independent $K_{\text{factor}} \approx 1.5$ and under the above discussed assumptions we obtain for the (quark dominated) kinematics in the forward region $1.2 < y < 2.4$ for $Q = (4\text{--}5)$ GeV the results shown in Fig. 3(a). (The uncertainty of $\text{num}(q_T)$ due to the $1\text{-}\sigma$ region of the HERMES fit and the effect of antiquarks is: $K_{\text{factor}} \sigma_{UT}^{\text{Siv}} = (6.0 \pm 3.5) \cdot 10^{-12}$ barn. For better visibility Fig. 3(a) shows only the central value.)

Clearly, by including the low- q_T region the SSA is diminished. One could increase the SSA by applying a low- q_T cut. The effect of such a cut is illustrated in Fig. 3(b). E.g., choosing $q_{\min} = 1 \text{ GeV}^2$ increases the SSA by about 20% which is close to the optimum for our choice of parameters.² Of course, there is a price to pay for. The statistical error $\delta A_{UT}^{\sin(\phi-\phi_s)} \propto 1/\sqrt{N}$ grows by $\sim 50\%$ due to loss of statistics, i.e. fewer counts N . In principle, one can find a favored q_{\min} , which would allow to achieve the maximal relative accuracy in the experiment. This is a good illustration of the old “golden rule” of spin physics: the best analyzing power is in the kinematical region where

²Had we chosen $\langle p_T^2 \rangle_{\text{Siv}}^{\text{RHIC}} = 0.6 \text{ GeV}^2$ the gain would be much more spectacular, namely, up to a factor of 2.5. The optimum would then, however, appear for $q_{\min} = 3 \text{ GeV}^2$, i.e. beyond the presumed range of applicability of the Gauss model. Also the loss of statistics would then be considerable.

(spin asymmetry)² \times (number of counts) reaches a maximal value.

The above numbers are strongly model-dependent. The quantitative results change already by choosing different parameters within the Gaussian model, cf. footnote ². However, as an important qualitative conclusion of this exercise, we point out that in principle the introduction of an appropriately optimized low- q_T cut may be a helpful device to improve the experimental signal for the Siverts SSA in DY.

Since in this context the crucial feature is the presence of a kinematical zero, the same procedure may turn out helpful in studies of other SSA or other phenomena related to parton transverse momenta in DY, SIDIS and other processes.

V. SIVERS SSA IN DY AT PAX AND COMPASS

Previously transverse momentum weighted Siverts SSA were estimated for the PAX and the COMPASS (hadron mode) experiments [21] on the basis of a Siverts function fit to the *preliminary* HERMES data on the $P_{h\perp}$ -weighted Siverts SSA [15]. Here, we solidify the estimates and conclusions of Ref. [21] on the basis of the Siverts function fit (4) to the *final* HERMES data [16].

The fixed target mode at PAX would make available $s = 45 \text{ GeV}^2$, which would allow to access lepton pairs in $p^1 \bar{p} \rightarrow l^+ l^- X$ in the region $Q^2 = 2.5 \text{ GeV}^2$ below the J/ψ -resonance and well above the region of dileptons from $\phi(1020)$ decays. This Q^2 corresponds to the average scale of the HERMES experiment [16], thus at PAX it is unnecessary to consider p_T -broadening effects as we have modeled in Eq. (14).

On the basis of the fit (4) we obtain for the PAX experiment the estimates plotted in Fig. 4. Nearly the entire

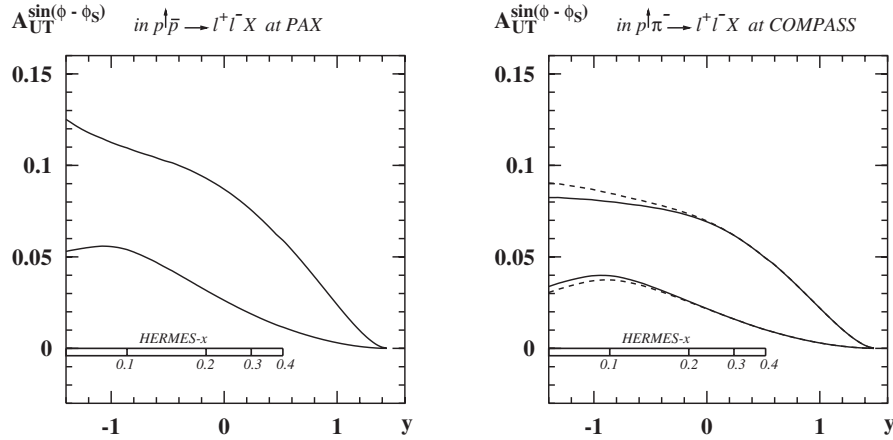


FIG. 4. The azimuthal SSA $A_{\text{UT}}^{\sin(\phi - \phi_S)}$ as function of y in Drell-Yan lepton pair production in $p^\dagger p \rightarrow l^+ l^- X$ for the kinematics of the PAX fixed target experiment with $s = 45 \text{ GeV}^2$ and $Q^2 = (2.5 \text{ GeV})^2$, and in Drell-Yan lepton pair production in $p^\dagger p \rightarrow l^+ l^- X$ for the kinematics of the COMPASS experiment with $s = 400 \text{ GeV}^2$ and $Q^2 = 20 \text{ GeV}^2$, respectively. The estimates are based on the fit for the Siverson q -distribution functions, see Eq. (4), obtained from the HERMES data [16]. The inner error band (solid lines) shows the $1\text{-}\sigma$ uncertainty of the fit. The outer error band (dashed lines) arises from assuming that the Siverson \bar{q} -distribution functions are proportional to the unpolarized antiquarks see Eqs. (10) and (11). For the PAX experiment the uncertainty due to Siverson antiquarks is not visible on this scale. The x -region explored in the HERMES kinematics is shown.

range of y for this kinematics is constrained by results from HERMES data. The SSA is sizeable, (5%–10%) in the central region, and could be measured, see [19]. The uncertainty due to the unknown Siverson antiquark distribution functions is completely negligible, as expected [21].

In the COMPASS experiment, using the possibility of a hadron beam, one could study SSA in $p^\dagger \pi^- \rightarrow l^+ l^- X$ at $s = 400 \text{ GeV}^2$ and, e.g., $Q^2 = 20 \text{ GeV}^2$ well above the resonance region of J/ψ and other charmonia. In this kinematics p_T -broadening effects can be expected to be comparable to RHIC, and we use estimates analogous to (14).

Our estimate for the COMPASS experiment is shown in Fig. 4. Again we observe that for this kinematics one probes the Siverson function in the x -region, where it is well constrained by the HERMES data [16]. Also at COMPASS, the Siverson SSA is sizeable, about (4%–8%) in the central region, i.e. sufficiently large to allow a test of the prediction (1). In order to see this, let us make the following crude estimate. The differential cross section for $\pi^- p \rightarrow \mu^+ \mu^- X$ is about $d\sigma/dQ \approx 50 \text{ pb/GeV}$ in the region $y > 0$ in a kinematics comparable to that considered here [72]. Assuming an integrated luminosity of 10^{39} cm^{-2} , which is realistic [18], one could measure with an ideal detector and a target polarization of $P = 100\%$ the SSA with an accuracy of $\delta A = 1/(P\sqrt{N}) \sim 0.5\%$ in the forward region $y > 0$ in a bin of size $\Delta Q \sim 1 \text{ GeV}$. Of course, this accuracy cannot be achieved with a realistic detector and polarization — but this is not necessary for a first test of the change of sign in Eq. (1).

The impact of Siverson antiquark distributions is weak and negligible as noted in Ref. [21]. In Fig. 4 we demonstrate this for the model II in Eq. (11). For the estimates in Fig. 4

we used the parametrizations [62,73] (using [74,75] confirms the results).

Thus, though at the present stage one cannot provide dedicated estimates of δA for these experiments as we did for RHIC in Sec. III, we conclude that also PAX and COMPASS could test the prediction (1) for quarks. However, both experiments are insensitive to Siverson antiquark distributions [21], which underlines the unique feature of RHIC with respect to this point.

We stress that quantitative analyses and extractions of the Siverson function from COMPASS, PAX and RHIC data will require a good theoretical control on NLO-QCD-corrections, soft factors and evolution effects. For discussions of some of these issues in the context of (collinear) double spin asymmetries, in particular, for the PAX kinematics see [76,77].

VI. CONCLUSIONS

We conclude that PHENIX and STAR can confirm the change of sign of the Siverson q -distribution function in SIDIS and the Drell-Yan process in Eq. (1) (as can PAX with an antiproton and COMPASS with a pion beam). Both PHENIX and STAR can provide first information on the Siverson \bar{q} -distribution functions which are not constrained by the current SIDIS data from HERMES or COMPASS, unlike PAX and COMPASS.

There are more recent (preliminary!) HERMES data [67]. Our fit is compatible with the new data, but the tendency is to increase the Siverson function. This tendency would result in even more optimistic estimates for DY at RHIC (and COMPASS and PAX).

We confirm Refs. [24,25] with respect to the capability of RHIC to access the Siverson q -distribution. In addition we

point out the possibility to learn about Siverson \bar{q} -distribution from RHIC.

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- [1] D. W. Siverson, Phys. Rev. D **41**, 83 (1990); **43**, 261 (1991).
 [2] D. L. Adams *et al.*, Phys. Lett. B **261**, 201 (1991); **264**, 462 (1991); Z. Phys. C **56**, 181 (1992).
 [3] J. Adams *et al.* (STAR Collaboration), Phys. Rev. Lett. **92**, 171801 (2004).
 [4] D. Boer and P. J. Mulders, Phys. Rev. D **57**, 5780 (1998).
 [5] S. J. Brodsky, D. S. Hwang, and I. Schmidt, Phys. Lett. B **530**, 99 (2002); Nucl. Phys. B **642**, 344 (2002).
 [6] J. C. Collins, Phys. Lett. B **536**, 43 (2002).
 [7] A. V. Belitsky, X. Ji, and F. Yuan, Nucl. Phys. B **656**, 165 (2003); X. D. Ji and F. Yuan, Phys. Lett. B **543**, 66 (2002); D. Boer, P. J. Mulders, and F. Pijlman, Nucl. Phys. B **667**, 201 (2003).
 [8] A. Bacchetta, U. D'Alesio, M. Diehl, and C. A. Miller, Phys. Rev. D **70**, 117504 (2004).
 [9] J. C. Collins and D. E. Soper, Nucl. Phys. B **193**, 381 (1981); **213**, 545(E) (1983).
 [10] X. D. Ji, J. P. Ma, and F. Yuan, Phys. Rev. D **71**, 034005 (2005); Phys. Lett. B **597**, 299 (2004).
 [11] J. C. Collins and A. Metz, Phys. Rev. Lett. **93**, 252001 (2004).
 [12] H. Avakian (HERMES Collaboration), Nucl. Phys. B, Proc. Suppl. **79**, 523 (1999).
 [13] A. Airapetian *et al.* (HERMES Collaboration), Phys. Rev. Lett. **84**, 4047 (2000); Phys. Rev. D **64**, 097101 (2001); Phys. Lett. B **562**, 182 (2003).
 [14] H. Avakian *et al.* (CLAS Collaboration), Phys. Rev. D **69**, 112004 (2004); E. Avetisyan, A. Rostomyan, and A. Ivanilov (HERMES Collaboration), hep-ex/0408002.
 [15] N. C. Makins (HERMES Collaboration), Transversity Workshop, Athens, Greece, 2003; R. Seidl (HERMES Collaboration), Proc. of DIS'2004, Štrbské Pleso, Slovakia, 2004; I. M. Gregor (HERMES Collaboration), Acta Phys. Pol. B **36**, 209 (2005).
 [16] A. Airapetian *et al.* (HERMES Collaboration), Phys. Rev. Lett. **94**, 012002 (2005).
 [17] V. Y. Alexakhin *et al.* (COMPASS Collaboration), Phys. Rev. Lett. **94**, 202002 (2005); P. Pagano (COMPASS Collaboration), hep-ex/0501035.
 [18] COMPASS Collaboration, CERN Report No. CERN/SPLC 96-14, SPSC/P 297.
 [19] P. Lenisa and F. Rathmann *et al.* (PAX Collaboration), hep-ex/0505054; F. Rathmann *et al.*, Phys. Rev. Lett. **94**, 014801 (2005); hep-ex/0412063.
 [20] M. Anselmino, V. Barone, A. Drago, and N. N. Nikolaev, Phys. Lett. B **594**, 97 (2004); A. V. Efremov, K. Goeke, and P. Schweitzer, Eur. Phys. J. C **35**, 207 (2004); hep-ph/0412427.
 [21] A. V. Efremov, K. Goeke, S. Menzel, A. Metz, and P. Schweitzer, Phys. Lett. B **612**, 233 (2005).
 [22] J. C. Collins, A. V. Efremov, K. Goeke, S. Menzel, A. Metz, and P. Schweitzer, Phys. Rev. D **73**, 014021 (2006).
 [23] M. Anselmino, M. Boglione, U. D'Alesio, A. Kotzinian, F. Murgia, and A. Prokudin, Phys. Rev. D **71**, 074006 (2005).
 [24] M. Anselmino, M. Boglione, U. D'Alesio, A. Kotzinian, F. Murgia, and A. Prokudin, Phys. Rev. D **72**, 094007 (2005).
 [25] W. Vogelsang and F. Yuan, Phys. Rev. D **72**, 054028 (2005).
 [26] G. Bunce, N. Saito, J. Soffer, and W. Vogelsang, Annu. Rev. Nucl. Part. Sci. **50**, 525 (2000).
 [27] D. Boer, Phys. Rev. D **60**, 014012 (1999).
 [28] B. Dressler, K. Goeke, M. V. Polyakov, P. Schweitzer, M. Strikman, and C. Weiss, Eur. Phys. J. C **18**, 719 (2001).
 [29] D. Boer and W. Vogelsang, Phys. Rev. D **69**, 094025 (2004).
 [30] D. Boer and J. W. Qiu, Phys. Rev. D **65**, 034008 (2002).
 [31] M. Anselmino, U. D'Alesio, and F. Murgia, Phys. Rev. D **67**, 074010 (2003).
 [32] M. Anselmino, M. Boglione, U. D'Alesio, E. Leader, and F. Murgia, Phys. Rev. D **70**, 074025 (2004).
 [33] C. Boros, Z. T. Liang, and T. C. Meng, Phys. Rev. Lett. **70**, 1751 (1993).
 [34] N. Hammon, O. Teryaev, and A. Schäfer, Phys. Lett. B **390**, 409 (1997).
 [35] J. C. Collins *et al.*, hep-ph/0510342.
 [36] M. Anselmino *et al.*, hep-ph/0511017.
 [37] J. C. Collins, Nucl. Phys. B **396**, 161 (1993).
 [38] A. V. Efremov and O. V. Teryaev, Yad. Fiz. **39**, 1517 (1984); Phys. Lett. B **150**, 383 (1985); J. W. Qiu and G. Sterman, Phys. Rev. Lett. **67**, 2264 (1991); Nucl. Phys. B **378**, 52 (1992).
 [39] Y. Kanazawa and Y. Koike, Phys. Lett. B **478**, 121 (2000); **490**, 99 (2000).
 [40] M. Anselmino, M. Boglione, and F. Murgia, Phys. Lett. B **362**, 164 (1995).
 [41] M. Anselmino and F. Murgia, Phys. Lett. B **442**, 470 (1998).
 [42] U. D'Alesio and F. Murgia, Phys. Rev. D **70**, 074009 (2004).
 [43] M. Anselmino, M. Boglione, U. D'Alesio, E. Leader, and F. Murgia, Phys. Rev. D **71**, 014002 (2005).
 [44] B. Q. Ma, I. Schmidt, and J. J. Yang, Eur. Phys. J. C **40**, 63 (2005).
 [45] C. J. Bomhof, P. J. Mulders, and F. Pijlman, Phys. Lett. B **596**, 277 (2004); A. Bacchetta, C. J. Bomhof, P. J. Mulders, and F. Pijlman, Phys. Rev. D **72**, 034030 (2005).

- [46] A. Afanasev and C. E. Carlson, hep-ph/0308163; A. Metz and M. Schlegel, Eur. Phys. J. A **22**, 489 (2004); Ann. Phys. (Berlin) **13**, 699 (2004); K. Goetze, A. Metz, and M. Schlegel, Phys. Lett. B **618**, 90 (2005).
- [47] A. V. Efremov, K. Goetze, and P. Schweitzer, Phys. Lett. B **522**, 37 (2001); **544**, 389(E) (2002); Eur. Phys. J. C **24**, 407 (2002); Nucl. Phys. A **711**, 84 (2002); Acta Phys. Pol. B **33**, 3755 (2002); P. Schweitzer and A. Bacchetta, Nucl. Phys. A **732**, 106 (2004).
- [48] E. De Sanctis, W. D. Nowak, and K. A. Oganessian, Phys. Lett. B **483**, 69 (2000); K. A. Oganessian, N. Bianchi, E. De Sanctis, and W. D. Nowak, Nucl. Phys. A **689**, 784 (2001).
- [49] B. Q. Ma, I. Schmidt, and J. J. Yang, Phys. Rev. D **66**, 094001 (2002); **65**, 034010 (2002).
- [50] A. V. Efremov, K. Goetze, and P. Schweitzer, Phys. Lett. B **568**, 63 (2003).
- [51] M. Anselmino, M. Boglione, U. D'Alesio, A. Kotzinian, F. Murgia, and A. Prokudin, hep-ph/0412316.
- [52] A. V. Efremov, K. Goetze, and P. Schweitzer, Czech. J. Phys. **55**, A189 (2005).
- [53] A. Airapetian *et al.* (HERMES Collaboration), Phys. Lett. B **622**, 14 (2005).
- [54] J. C. Collins, Acta Phys. Pol. B **34**, 3103 (2003).
- [55] J. C. Collins, D. E. Soper, and G. Sterman, Nucl. Phys. B **250**, 199 (1985).
- [56] E. Witten, Nucl. Phys. B **160**, 57 (1979).
- [57] A. P. Balachandran, V. P. Nair, S. G. Rajeev, and A. Stern, Phys. Rev. D **27**, 1153, 2772E (1983); E. Witten, Nucl. Phys. B **223**, 433 (1983); J. L. Gervais and B. Sakita, Phys. Rev. Lett. **52**, 87 (1984); A. V. Manohar, Nucl. Phys. B **248**, 19 (1984); D. Diakonov, V. Y. Petrov, and P. V. Pobylitsa, Nucl. Phys. B **306**, 809 (1988); R. F. Dashen, E. Jenkins, and A. V. Manohar, Phys. Rev. D **49**, 4713 (1994); **51**, 2489(E) (1995).
- [58] P. V. Pobylitsa, hep-ph/0301236.
- [59] M. Anselmino, V. Barone, A. Drago, and F. Murgia, Nucl. Phys. B, Proc. Suppl. **105**, 132 (2002); hep-ph/0209073.
- [60] P. V. Pobylitsa, hep-ph/0212027.
- [61] A. Drago, Phys. Rev. D **71**, 057501 (2005).
- [62] M. Glück, E. Reya, and A. Vogt, Eur. Phys. J. C **5**, 461 (1998).
- [63] S. Kretzer, E. Leader, and E. Christova, Eur. Phys. J. C **22**, 269 (2001).
- [64] A. Bacchetta, M. Boglione, A. Henneman, and P. J. Mulders, Phys. Rev. Lett. **85**, 712 (2000).
- [65] J. Ball *et al.*, Nucl. Instrum. Methods Phys. Res., Sect. A **498**, 101 (2003).
- [66] S. Goertz, W. Meyer, and G. Reicherz, Prog. Part. Nucl. Phys. **49**, 403 (2002); **51**, 309(E) (2003).
- [67] M. Dieffenthaler, hep-ex/0507013.
- [68] M. Burkardt, Phys. Rev. D **66**, 114005 (2002); Nucl. Phys. A **735**, 185 (2004).
- [69] D. Boer, Nucl. Phys. B **603**, 195 (2001).
- [70] M. Burkardt, Phys. Rev. D **69**, 057501 (2004); **69**, 091501 (2004).
- [71] F. Landry, R. Brock, P. M. Nadolsky, and C. P. Yuan, Phys. Rev. D **67**, 073016 (2003).
- [72] K. J. Anderson *et al.*, Phys. Rev. Lett. **42**, 944 (1979).
- [73] M. Glück, E. Reya, and I. Schienbein, Eur. Phys. J. C **10**, 313 (1999).
- [74] A. D. Martin, R. G. Roberts, W. J. Stirling, and R. S. Thorne, Phys. Lett. B **531**, 216 (2002).
- [75] P. J. Sutton, A. D. Martin, R. G. Roberts, and W. J. Stirling, Phys. Rev. D **45**, 2349 (1992).
- [76] P. G. Ratcliffe, Eur. Phys. J. C **41**, 319 (2005).
- [77] H. Shimizu, G. Sterman, W. Vogelsang, and H. Yokoya, Phys. Rev. D **71**, 114007 (2005).