Instanton-induced azimuthal spin asymmetry in deep inelastic scattering

Dmitry Ostrovsky and Edward Shuryak

Department of Physics and Astronomy, State University of New York at Stony Brook, New York 11794, USA

(Received 30 September 2004; revised manuscript received 7 December 2004; published 27 January 2005)

We propose a new dynamical mechanism for spin asymmetries in semi-inclusive deep inelastic scattering, utilizing strong vacuum fluctuations of the gluon field described semiclassically by instantons. Those can generate large asymmetry because they can simultaneously generate (i) a chirality flip of the struck quark; (ii) large O(1) T-odd phase due to its final state interaction. The absolute magnitude of the effect is estimated using known parameters of the instanton ensemble in the QCD vacuum and known structure and fragmentation functions, without any new free parameters. The result agrees in sign and (roughly) in magnitude with the available data on single particle inclusive DIS. Furthermore, our predictions uniquely relate effects for longitudinally and transversely polarized targets.

DOI: 10.1103/PhysRevD.71.014037 PACS numbers: 13.88.+e

I. INTRODUCTION

Perturbative QCD is well known to account correctly for the dependence of structure and fragmentation functions on the hard scale Q^2. On the other hand, the conventional perturbative cascade of gluon emission and of the quark pair production is clearly inadequate to explain the original quark sea and glue distribution, at a nonperturbative boundary $\mu \sim 1$ GeV. Already at the level of *spin-independent* inclusive leptonic deep inelastic scattering (DIS), we learned that the quark sea is rather strongly flavor polarized. The *spin-dependent* DIS has further shown that the sea quarks are also strongly spin polarized, in the *opposite* direction to the polarization of the nucleon and the valence quarks. These puzzles about nucleon structure which were revealed by experiments, and neither their dynamical origin nor their magnitude, for hadrons other than the nucleon, are understood. The pQCD cascade (which is approximately flavor and chirality blind) obviously cannot provide their explanation.

A significant experimental program is now underway, at all experiments at HERA, at CERN (COMPASS), and Brookhaven National Laboratory (STAR and PHENIX at RHIC). The issues to be settled include the question of whether these two asymmetries mentioned are related, and also what is the degree of polarization of the glue.

The theoretical efforts to get a dynamical explanation to all these phenomena are, however, still at their infancy. A particular direction we will follow in this work is related to the nonperturbative phenomena in the QCD vacuum, described topological tunneling semiclassically by the so-called *instanton* solutions of Yang-Mills equations. There are several qualitative arguments why instantons may provide an explanation of these puzzles.

Forte and Shuryak [1] had argued that instantons provide a 100% effective mechanism of a polarization transfer from quarks to gluons. Kochelev [2] suggested that if the sea quark are produced via the instanton-induced 't Hooft vertex, they would be strongly flavor polarized. Indeed, a sea produced from, say, a valence left-handed up-quark u_L can only have $\bar{d}_L d_R$, $\bar{s}_L s_R$ pairs, but not $\bar{u}u$ ones because of Pauli principle for zero modes. Furthermore, a sea produced this way should also naturally have the opposite chirality, which is also phenomenologically welcomed. Unfortunately, any solid quantitative results on that are still lacking. (For a recent discussion of axial coupling constants of the nucleon and the "spin crisis" in an instanton framework see, e.g., a recent work by Schhafer and Zetocha [3].)

One more, although indirect, argument came from lattice calculations. Negele *et al.* (the MIT group) [4] have observed that moments of the various structure functions change very little when the true "quantum" lattice configurations are substituted by the so-called "cooled" configurations. The "cooling" procedure eliminates pQCD gluons and most of quantum fluctuations from vacuum configurations, preserving semiclassical objects (mostly instantons). If so, instanton-based fields alone may be sufficient to derive all structure functions, including the spin-dependent ones. Unfortunately, so far this idea has not been followed up to the extent to reach any quantitative conclusions.

The particular phenomenon we will discuss is single spin asymmetries (SSA). So far their theoretical discussion (see, e.g., [5]) has aimed mostly at their proper parametrization rather than explanation, but even that took time to realize the subtleties involved. One important step was an introduction of the nontrivial T-odd structure in the initial state via an appropriate structure function called the Sivers effect [6,7], while a similar effect in fragmentation function is called the Collins effect [8]. The corresponding function was introduced in [9]. In both cases the hard block remains the usual lowest-order pQCD scattering. Instantons can in principle contribute to both those effects: but this is not the intention of the present paper. One more logical alternative is that something nontrivial happens *during* the collision. For example, the next twist hard collisions were discussed, in which at the moment of hard scattering there is extra gluomagnetic field which can be also correlate with the nucleon spin; see [10]. As

 © 2005 The American Physical Society

it will become clear from the following, the dynamical process we discuss is most closely associated with this ideology, although technically we do not perform a twist expansion. The following two issues have been singled out as the necessary prerequisites to the very existence of SSA. Those are (i) the *chirality flip*; and (ii) the *final state interaction* of the outgoing quark. These issues are best explained if the state of the transversely polarized nucleon is viewed as a superposition of plus and minus chirality states. SSA can only result from the interference of these two amplitudes, or a handbaglike diagram in which the chirality is different on two ends and it thus has to be flipped in between. Obviously, the usual pQCD handbag diagram cannot do that as it conserves chirality of the quark. The spin flip issue can be in principle resolved à la Sivers, in the initial nucleon state. That was achieved by the introduction of a new component of the nucleon wave function, in which the valence quark rotates orbitally and thus has chirality *opposite* to that of the nucleon. The T-odd phase issue is related with a decade long theoretical stalemate over the very existence of the Sivers effect. Namely, Collins [8] have argued that it should be zero based on T invariance. The proof was retracted later, and the loophole is precisely the P exponents of the outgoing quarks, or their possible final state interaction.

How the Sivers effect could be incorporated into the QCD framework was shown by Brodsky, Hwang, and Schmidt (BHS) [11,12], and Collins [13]. (For an early model of a T-odd *distribution* function, see [14], as well as a bag-model calculation by Yuan [15], and a model with spin-0 and spin-1 diquarks in [16].)

BHS used a very simplified model of a nucleon, made of a valence quark plus the spin-zero diquark, thus the former carries all the angular momentum. The issue (i) is included via new p-wave wave function, and (ii) via the $O(\alpha_s)$ gluon exchange between the outgoing quark and the rest of the system (the diquark). Note that in the BHS approach there is no connection between (i) and (ii): the final state interaction is simply necessary to make the nontrivial sector of the nucleon wave function visible.

The philosophy of our approach came out of reflections about this very point. We thought it quite likely that the underlying dynamics of the quark chirality flip is related with the nonperturbative interaction producing chiral symmetry breaking. (By the way, throughout this paper we will ignore nonzero quark masses and thus treat chiral symmetry as exact.) in the QCD vacuum. Multiple arguments indicate that this phenomenon is generated by small-size instantons, see [17] for a review. And as instantons can provide the chirality flip, they also are capable to generate large $[O(1)$ rather than $O(\alpha_s)]$ phase of the P exponent, the final state interaction of the struck quark. As both are necessary for the asymmetry in question, it makes the instanton mechanism a twice more attractive candidate for its explanation.

In short, the physics of our proposal is as follows: the asymmetry is generated by events in which the point of quark-lepton collision happens to occur close to a preexisting topological vacuum fluctuation. In this sense the phenomenon is generic, due to tunneling through the topological barrier of gluonic Lagrangian. Thus we predict this effect to be universal, not related to the nucleon itself. The effect is proportional to the diluteness parameter of nonperturbative QCD vacuum (the density of instantons). Its modification inside the nucleon and other hadrons is believed to be quite small, at the level of few percent, which is ignored in what follows.

As the initial quark is moving in a strong color field of an instanton, it can disappear into a Dirac sea while instead another quark, with the *opposite* chirality, appears close by. This phenomenon, related to chiral anomaly is driven by instantons and is accounted for via the so-called 't Hooft zero-mode part of a quark propagator. As it will be seen from the calculation below, this phenomenon may occur both before or after the DIS point.

The T-odd phase of the final state interaction effect is given by another, chirality-nonflip part of a quark propagator in the instanton background. The reason it is much enhanced compared to the BHS result is simply because the gluon field of an instanton is large $A \sim 1/g$ compared to perturbative field of the exchanged gluon. As a result there would be no extra small factor $\sim\alpha_s$ in our answer.

Let us emphasize that we calculate the effect dynamically, while the previous works have mostly parametrized it. It is not a Sivers effect, which delegates the main dynamics into some unknown new structure functions, but a dynamical phenomenon with an amplitude expressed in terms of well-known parameters of the vacuum, and based on a single-instanton approximation well tested, in particular, in other applications of instanton dynamics [18].

Still this work is in a way simplistic, as we only consider the component in which the nucleon spin resigns entirely at the spin of its valence quark. In other words, in this work the chirality flip of a *nucleon* required for transverse spin asymmetry is associated with the chirality flip of a *valence quark*.

We are of course aware of the spin crises: eventually one should be able to get more realistic and include the real polarized sea and polarized gluons as well. As the authors quoted above suggested, it may also be an instanton effect: but in the formation of the nucleon wave function rather than during the DIS process, as assumed by us in this work.

HERMES experiments have reported first data on SSA in semi-inclusive deep inelastic scattering (SIDIS) for longitudinally [19–21] and transversely polarized targets [22]. The data from COMPASS [23] are for longitudinally polarized targets only. Using the dependence on the spin direction and the transverse momentum of the produced particle, it is in principle possible (and desirable) to disentangle the only kinematical structure we obtained below, without admixture of any other effects.

The plan of the paper is as follows: in Sec. II kinematics of SSA in SIDIS is introduced, in Sec. III the asymmetric part of the cross-section due to instantons is calculated, in Sec. IV an estimate of the effect and a comparison with experiment are given. In Sec. V we discuss our results and an outlook for future developments.

II. KINEMATICS OF AZIMUTHAL SSA IN SIDIS

Total cross section for deep inelastic scattering has the form

$$\frac{d\sigma}{dxdyd\phi} = \frac{\alpha_{em}^2}{Q^4} y L^{\mu\nu} W_{\mu\nu}, \tag{1}$$

where the azimuthal angle ϕ is unobservable in totally inclusive DIS.

A symmetric (spin-independent) lepton tensor is given by (see Fig. 1)

$$L^{\mu\nu} = 2(l^\mu l'^\nu + l'^\mu l^\nu) - 2g^{\mu\nu}(l \cdot l'). \tag{2}$$

A totally inclusive cross-section symmetric part of $W_{\mu\nu}$ is given by

$$W_{\mu\nu} = \left(-g_{\mu\nu} + \frac{q_\mu q_\nu}{q^2}\right) F_1(x, Q^2) + \left(P_\mu + \frac{1}{2x} q_\mu\right) \times \left(P_\nu + \frac{1}{2x} q_\nu\right) \frac{F_2(x, Q^2)}{P \cdot q}. \tag{3}$$

In SIDIS one has one more vector parameter, the momentum of the produced hadron, K_μ. This leads to the appearance of several new possible tensor structures of hadronic tensor $W_{\mu\nu}$ and new dimensionless invariants on which "structure functions" may depend. The tensor structure of $W_{\mu\nu}$ is of course limited by symmetry, $W_{\mu\nu} = W_{\nu\mu}$ (we consider only unpolarized electrons for leptonic tensor), electromagnetic gauge invariance, $q^\mu W_{\mu\nu} = 0$, and parity invariance. We are interested also in *spin-dependent* asymmetries and therefore, nucleon spin S_μ ($S^2 = -1$) must be involved in nontrivial combination with produced hadron momentum. To limit the possible structures even more, we will consider hadronic tensor only to the first power in K, assuming that it enters $W_{\mu\nu}$ in the combination K/Q, which is generally small.

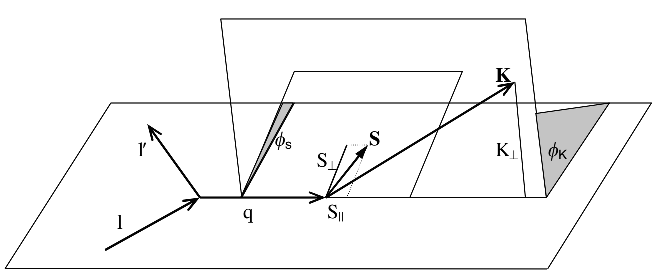

FIG. 1. Kinematics of single particle inclusive DIS in nucleon rest frame, defining all momenta and angles to be used.

Two possible tensor structures are in general possible, $[P + (1/2x)q]_{\{\mu}\epsilon_{\nu\}\rho\sigma\tau} q^\rho K^\sigma S^\tau$ and $\epsilon_{\pi\rho\sigma\tau} P^\pi q^\rho K^\sigma S^\tau$, and as the reader will see we only found that the former one contributes. Two indices with curly braces are assumed to be symmetrized.

Although the latter structure may not be excluded on general grounds, note that even if this would be the case it would lead to a smaller contribution in the parton model. Indeed, in the parton model, the dependence of $W_{\mu\nu}$ on P_μ is possible only through the momentum of struck quark. In the infinite momentum frame $p_\mu = xP_\mu$. It is also true that $K_\mu = zk_\mu$, where k_μ is the momentum of the quark after the collision, up to a relatively small correction in the fragmentation process. Momentum conservation in the interaction vertex gives $k_\mu = p_\mu + q_\mu$ up to a small correction due to the possibility of rescattering. Overall, it leads to the extra power of small transverse momentum for $\epsilon^{\pi\rho\sigma\tau} P_\pi q_\rho K_\sigma S_\tau$.

As we will see in the next section, instanton-induced contribution has the former form. Note that it contains a product of two momenta, one energy and one spin, so it is T odd. Therefore one can only find their contribution to any observable multiplied by another T-odd quantity, such as the final state interaction phase.

III. QUARK SCATTERING IN INSTANTON FIELD

For a DIS on a quark, the spin-dependent quark tensor (in analogy with hadron tensor conventionally used in DIS) is given by

$$W_{\mu\nu}^{\text{quark}} = \frac{1}{2}[\text{tr}(\rho_{\text{out}} M_{\{\mu} \rho_{\text{in}} M_{\nu\}}^*)], \tag{4}$$

where ρ_{in} and ρ_{out} are density matrices for incoming and outgoing quarks, and M_μ and M_ν are $\gamma^* qq$ vertices in the presence of the instanton field. Trace is taken over Lorentz as well as color indices. The correct normalization factor will be found later in the beginning of Sec. IV. The density matrix or an incoming quark with momentum p is $\rho_{\text{in}} = \hat{p}(1 + \gamma_5 \hat{s})$. Because we are interested only in spin-dependent part of the cross section we take only $\hat{p}\gamma_5\hat{s}$ as the quark density matrix. Because spin of outgoing quark is not measured we take $\rho_{\text{out}} = \hat{k}$. To the first order in instanton density $M_\mu = \gamma_\mu + M_\mu^1$, where M_μ^1 is a sum of amplitudes in single-instanton and anti-instanton fields. For spin-dependent quark tensor one now has

$$\Delta W_{\mu\nu}^{\text{quark}} = \text{Re}[\text{tr}(\hat{k} M_{\{\mu}^1 \hat{p} \gamma_5 \hat{s} \gamma_{\nu\}})], \tag{5}$$

Two indices with curly braces are assumed to be symmetrized.

Spin-dependent part of incoming quark density matrix is chirally odd, therefore, taking to account that all other parts of Eq. (5) are chirally even, M_μ^1 must be chirally odd. Therefore, M_μ^1 must contain propagation through zero mode in the instanton (anti-instanton) field. Calculation of M^1_μ is most easily performed in chiral basis first. We return to the Dirac fermions in the end.

The calculation of M^1_μ is in the complete analogy to the calculation of the instanton-induced chirally odd contribution to the gluon structure function made by Moch, Ringwald, and Schrempp [24], where the reader is referred for all technical details.

The quark propagator in the instanton field is known to include the so-called zero-mode part (x-space, nonamputated, left-to-right flip)

$$S_0(x, y)^{\dot{\alpha} j}_{\dot{\beta} i} = \frac{\rho^2}{\pi^2 \lambda} \frac{x_\gamma \bar{\sigma}^\gamma_{\dot{\beta}\rho} \epsilon^{\rho j} x_\delta (\sigma^\delta)^{\dot{\alpha}\pi} \epsilon_{i\pi}}{(x^2+\rho^2)^{3/2}(y^2+\rho^2)^{3/2}|x||y|}. \tag{6}$$

Here greek indices are Weyl spinor indices, i and j are color indices, $\sigma_\mu = (i, \vec{\sigma})$, $\bar{\sigma}_\mu = (-i, \vec{\sigma})$ (Euclidean space), $\epsilon^{01} = -\epsilon^{10} = -\epsilon_{01} = \epsilon_{10} = 1$ are projectors on zero-mode chiral-color states. Zero-mode propagator is normalized to λ, the lowest eigenvalue of the Dirac operator in instanton liquid. We return to the discussion of its value in Sec. IV.

Fourier transform with respect to incoming particle is given by

$$S_0(x, p)^j_{\dot{\beta}\alpha i} = \frac{2\rho^2}{\lambda} \frac{1}{(x^2+\rho^2)^{3/2}} \frac{x_\gamma \bar{\sigma}^\gamma_{\dot{\beta}\rho} \epsilon^{\rho j} \epsilon_{i\alpha}}{|x|}, \tag{7}$$

where mass-shell condition ($p^2 = 0$) is assumed and the incoming particle propagator is amputated.

Another, the so-called non-zero-mode part of the propagator does not flip chirality of a quark. For a *right-handed* quark, we will use it in the form of the Fourier transformed vertex, with an amputated line for the outgoing particle, which has the following form:

$$S_{nz}(k, x)^{\beta j}_{\alpha i} = -\frac{|x|}{\sqrt{x^2+\rho^2}} e^{ik\cdot x} \delta^\beta_\alpha \left[\delta^j_i + \frac{\rho^2}{x^2} \frac{(\bar{\tau}_\rho \tau_\sigma)^j_i k^\rho x^\sigma}{2k\cdot x} \times (1 - e^{-ik\cdot x}) \right], \tag{8}$$

where $\tau_\mu = (i, \vec{\tau})$, $\bar{\tau}_\mu = (-i, \vec{\tau})$ for color matrices.

In order to display the T-odd phase in this part of the propagator of the quark in instanton field, we (for the purpose of demonstration) will drop $\exp(-ik\cdot x)$ from the round brackets in Eq. (8) which gives the correct $x \to 0$ limit to the propagator, as well as insures the electromagnetic gauge invariance. One may also think that we are for the moment working in the $k \cdot x \gg 1$ kinematical domain. Then, Eq. (8) can be written as

$$S_{nz}(k, x) \simeq -e^{ik\cdot x} \exp\left[i \frac{\bar{\eta}^{a\mu\nu} k_\mu x_\nu \tau_a}{k\cdot x} \ln\left(\frac{x^2+\rho^2}{x^2}\right)\right], \tag{9}$$

where $\bar{\eta}^{a\mu\nu}$ is the 't Hooft symbol. This form clearly displays the $O(1)$ T-odd phase (after continuation back to the Minkowski space).

Combining two contributions to M^1_μ we get

$$M^1_\mu = \int d^4x e^{iq\cdot x} S_{nz}(k, x) \sigma_\mu S_0(x, p). \tag{10}$$

Before calculating the integral one has to perform some algebra. As chiral-color projectors of zero modes effectively mix chiral (σ) and color (τ) matrices, there is therefore not much sense in keeping the difference between the σ and τ matrices. Furthermore, we make use of the relation

$$\epsilon_{\gamma\alpha}(\sigma_\mu)^{\gamma\dot{\gamma}}(\bar{\sigma}_\nu)_{\dot{\gamma}m} = \epsilon_{\gamma m}(\sigma_\nu)^{\gamma\dot{\gamma}}(\bar{\sigma}_\mu)_{\dot{\gamma}\alpha} \tag{11}$$

to simplify the expressions.

After the x integration

$$M^1_\mu = \frac{4\pi\rho^2}{\lambda} i(\sigma_\mu \bar{\sigma}_\rho \epsilon)^{\alpha i} \epsilon_{\beta j} \Phi^\rho(k, q) \tag{12}$$

with

$$\Phi^\rho(k, q) = -\frac{p^\rho}{p^2} - \frac{k^\rho}{2k\cdot q}[1 - f(\rho|q|)], \tag{13}$$

where the "instanton form factor" $f(a) = aK_1(a)$ [$f(0) = 1$]. In derivation we took into account that $k = p + q$. Only the second term contributes to the spin-dependent part of the cross section.

So far we only included the first diagram of Fig. 2, as Eq. (12) corresponds to applying the zero-mode propagator to the incoming quark before its collision with a virtual photon. There is, of course, another diagram (see Fig. 2), which refers to the zero-mode propagator inserted in the outgoing quark line. That can easily be found by Hermitian conjugation of Eq. (12) and substitution $k \leftrightarrow -p$, which gives

$$M^{1\prime}_\mu = -\frac{4\pi\rho^2}{\lambda} i\epsilon_{i\alpha}(\epsilon\sigma_\rho\bar{\sigma}_\mu)^{j\beta}\Phi^\rho(p, q). \tag{14}$$

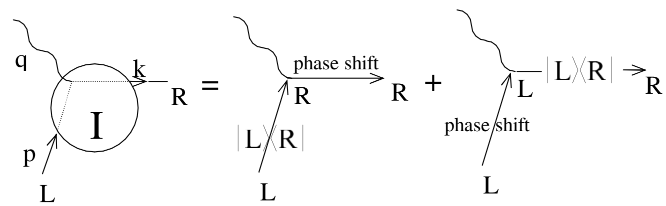

FIG. 2. The amplitude for a single quark scattering in an instanton background field (shown schematically by a circle with I) can be written as a sum of two diagrams. They differ by where the chirality flip, between a chosen left-handed (L) initial struck quark into a right-handed (R) one, which is described by the zero-mode part of the propagator. The "phase shift" subscript is a reminder of a complex phase of the nonflip part of the propagator, which has to be kept to get the nonzero answer.

Next, we take trace over color indices. Because all other parts of the diagram are trivial in color, one reduces M^1_μ to

$$M^1_\mu = \frac{4\pi\rho^2}{\lambda} i[\sigma_\mu \bar{\sigma}_\rho \Phi^\rho(k, q) - \sigma_\rho \bar{\sigma}_\mu \Phi^\rho(p, q)]. \quad (15)$$

The matrix element for propagation in the anti-instanton field is given by $\bar{\sigma} \leftrightarrow \sigma$.

Inserting M^1_μ to the Eq. (5) one has

$$\Delta W^{\rm quark}_{\mu\nu} = \frac{4\pi\rho^2}{\lambda Q^2} {\rm Im}[{\rm tr}(\hat{k}\,\hat{p}\,\gamma_{\{\mu}\hat{p}\gamma_5\hat{s}\gamma_{\nu\}}) + {\rm tr}(\hat{k}\gamma_{\{\mu}\hat{k}\,\hat{p}\,\gamma_5\hat{s}\gamma_{\nu\}})][1 - f(\rho|q|)]. \quad (16)$$

Some Dirac algebra is needed here, which finally gives

$$\begin{aligned}\Delta W^{\rm quark}_{\mu\nu} &= \frac{8\pi\rho^2}{\lambda}\frac{1}{Q^2}(k+p)_{\{\mu}{\rm Imtr}(\hat{k}\,\hat{p}\,\gamma_5\hat{s}\gamma_{\nu\}}) \times [1 - f(\rho|q|)] \\ &= -\frac{32\pi\rho^2}{\lambda}\frac{1}{Q^2}(k+p)_{\{\mu}\varepsilon_{\nu\}\rho\sigma\tau}q^\rho k^\sigma s^\tau \times [1 - f(\rho|q|)]. \end{aligned} \quad (17)$$

In the last expression we switched back to Minkowski space and made use of $\sigma^E_0 = i\sigma^M_0$.

In this calculation we have neglected the interaction between instanton and the rest of the nucleon, apart of the struck quark. This approximation is motivated by the nucleon diluteness, the fact that the typical instanton size in the QCD vacuum $\rho \approx 1/3 fm$ is small compared to nucleon size R_N. Their account would lead to corrections of the order $\rho^2/R^2_N \sim 1/10$.

In vacuum parametrized by an ensemble of instantons one has to integrate Eq. (17) over collective degrees of freedom: color rotations and instanton size.

$$\overline{\left(\frac{\rho^2}{\lambda}\right)} \to \frac{\kappa}{\bar{\rho}^2 m^*}, \quad (18)$$

where κ is the instanton diluteness factor, $\bar{\rho}^2$ is the characteristic instanton size, and $1/m^*$ is the inverse effective quark mass in the instanton-liquid model.

IV. ESTIMATE OF THE ASYMMETRY

Equation (17) constitutes the result for spin-dependent asymmetric tensor at the partonic level. To change it to the hadronic result one has to substitute $p = xP$, $k + p = 2xP + q$, and $k = K/z$.

To find out the correct kinematical normalization, we calculate the *symmetric* part of the quark tensor without instantons starting from the same Eq. (4) and compare the result to conventional parametrization of totally inclusive DIS through the distribution function. In both symmetric and asymmetric cases the normalization parameters (albeit the correct type of distribution function) must be the same.

For the spin-independent DIS cross section one has

$$W^{\rm parton}_{\mu\nu} = 2N_c[p_\mu k_\nu + p_\nu k_\mu - g_{\mu\nu}(p\cdot k)]. \quad (19)$$

Rewriting this through conventional structure functions one finds that correct normalization is given by multiplication of $W^{\rm parton}_{\mu\nu}$ on

$$\sum_q e^2_q f_q(x)/(2N_c Q^2), \quad (20)$$

which leads to Eq. (3) with $F_1(x) = F_2(x)/(2x) = \frac{1}{2}\sum_q e^2_q f_q(x)$ for conventional structure functions. Because of the spin dependence of Eq. (17) one has to use the spin-dependent quark distributions, which we tentatively will call $f_{q,s}(x)$. We can now rewrite Eq. (17) as

$$\Delta W_{\mu\nu} = \frac{32\pi\kappa}{\bar{\rho}^2 m^* N_c}\frac{x}{z}\frac{1}{Q^4}[P + (1/2x)q]_{\{\mu}\varepsilon_{\nu\}\rho\sigma\tau}q^\rho K^\sigma s^\tau \times [1 - f(\rho Q)]\sum_q e^2_q f_{q,s}(x) D_q(z). \quad (21)$$

The asymmetric part of the cross section is now

$$\begin{aligned}\frac{d\Delta\sigma}{dxdydzd\phi_K} = \frac{\alpha^2_{em}}{Q^2}\frac{32\pi\kappa}{\bar{\rho}^2 m^* N_c}\frac{|K_\perp|}{zQ}[1 - f(\rho Q)] \times \sum_q e^2_q f_{q,s}(x) D_q(z)\Big[\frac{2}{Q}\frac{1-y}{y} \\ \times \sin(\phi_K - \phi_s)|s_\perp| + \frac{(1-y/2)\sqrt{1-y}}{Mx}\sin\phi_K s_\parallel\Big]. \end{aligned} \quad (22)$$

One can now define what exactly are the spin-dependent quark distributions $f_{q,s}(x)$ we introduced in analogy to $f_q(x)$. In our simplistic model of the hadron structure, the spin-dependent quark distribution can be expressed as

$$f_{q,S,s}(x) = f_q(x) + \Delta f(x)(S_\parallel s_\parallel) + \Delta_T f(x)(\vec{S}_\perp \cdot \vec{s}_\perp), \quad (23)$$

where s and S are the (three dimensional) unit spin vectors for quark and hadron, respectively $s = (s_\parallel, \vec{s}_\perp)$, $S = (S_\parallel, \vec{S}_\perp)$. This simplification of the spin-1/2 hadron structure through twist-2 [25] is sufficient for our purposes.

Asymmetric distributions correspond to the probability to find a quark in a hadron polarized the same way as a hadron minus the probability to find a quark polarized in opposite direction then a hadron. Eliminating quark spin in (22) in favor of the nucleon spin one has

$$f_{q,s}(x)s \to \Delta f(x) S_\parallel + \Delta_T f(x) S_\perp. \quad (24)$$

Inserting it in Eq. (22) gives

$$\begin{aligned}\frac{d\Delta\sigma}{dxdydzd\phi_K} &= \frac{\alpha_{em}^2}{Q^2}\frac{32\pi\kappa}{\bar{\rho}^2 m^* N_c}\frac{|K_\perp|}{zQ}[1-f(\rho Q)]\\ &\times \sum_q e_q^2 D_q(z)\bigg[\frac{2}{Q}\frac{1-y}{y}\\ &\times \sin(\phi_K-\phi_S)|S_\perp|\Delta_T f_q(x)\\ &\times \frac{(1-y/2)\sqrt{1-y}}{Mx}\\ &\times \sin\phi_K S_\parallel \Delta f_q(x)\bigg]. \end{aligned} \quad (25)$$

To obtain relative asymmetry one has to compare Eq. (25) to totally inclusive cross section

$$\frac{d\sigma}{dxdydzd\phi} = \frac{\alpha_{em}^2}{Q^2}\frac{1+(1-y)^2}{y}\sum_q e_q^2 f_q(x) D_q(z). \quad (26)$$

From Eqs. (25) and (26) one can see that for the most simplistic model of a nucleon, where the total spin of a hadron is carried out by a single valence quark (the other two being insulated from participation in spin physics by putting them into a spin-zero diquark), relative asymmetry does not depend on distribution functions of the nucleon and is in a sense universal, applicable to *all other hadrons*.

A more realistic approximation is $\Delta f(x) = \Delta_T f(x)$. It ignores differences due to relativistic motion of the quarks inside nucleons. However, in the absence of reliable experimental data on $\Delta_T f(x)$ one can use this approximation to get a reasonable estimate of the transverse asymmetry. Model calculations also favor such an approximation.

From Eq. (25) one can readily see that if $\Delta f(x) = \Delta_T f(x)$ is assumed, the relative size of transverse and longitudinal asymmetries is purely kinematical and does not depend on any details of hadronic structure.

We will now give an estimate for prefactor in Eq. (25) from the single-instanton approximation (SIA) of the instanton-liquid model. For a general discussion of instanton phenomenology the reader can consult, e.g., [17]. We will use the usual diluteness parameter and size

$$\bar{\rho} = 1/3 \text{ fm}, \qquad \kappa = n\bar{\rho}^4 \approx 1/3^4. \quad (27)$$

As for the accuracy of SIA and the value of the (appropriately averaged) value of the Dirac eigenvalues m^*, see detailed the discussion in Ref. [18]. It is found there that if it would be simply a quantity with one zero mode, like $\langle\bar{q}q\rangle$, the accuracy of selecting one closest instanton from the ensemble and ignoring all others would be typically about 30%. In this case the definition of it (called m_{uu} in [18]) should be $m^* \equiv (\langle 1/\lambda\rangle)^{-1}$ where the angular bracket stands for real eigenvalue spectrum in the vacuum ensemble. Its numerical value changes from $m^* = 120$ MeV for random instanton-liquid model to $m^* = 170$ MeV in interacting instanton ensemble. It must be noted that in our calculation spin asymmetry depends on both chirality flip and phase shift on the same instanton. Thus, we expect that in this case SIA is more accurate and use $m^* = 170$ MeV. In summary, all instanton-related parameters appear in the following combination, which has the dimension of the energy

$$\frac{32\pi\kappa}{\bar{\rho}^2 m^* N_c} = 0.88 \text{ GeV}. \quad (28)$$

The accuracy which may be claimed for single-instanton approximation is about 30%, and thus one may say this parameter is about 1 GeV. Note, however, that it includes the instanton density which is nonperturbatively small $\kappa \sim \exp[-2\pi/\alpha_s(\rho)] \sim 1/3^4$, and the only reason its phenomenological smallness (27) is not seen is because it happens to be compensated by the large numerical factor $32\pi \sim 100$.

The relevance of the single-instanton approximation depends also on the smallness of $1/(RQ)$ parameter, where R is the typical distance between instantons in the instanton ensemble. In DIS, Q is a well-controlled parameter, which we assume is never taken smaller than 1 GeV (it helps disentangle DIS from quasielastic resonance scattering). In the instanton-liquid model [17] $R \approx 1$ fm and $1/(RQ) \sim 0.2 \ll 1$.

A. Comparison with experiment

A detailed comparison with the experiment is outside the scope of this paper. We present here only a few details to establish phenomenological relevance of our model. We consider longitudinal and transversal spin asymmetries for the production of π^+ mesons off a polarized proton target [19–21]. For simplicity, we will assume that in order to produce π^+ from the proton one has to strike a u quark. In other words, $D_q^{\pi^+}(z) = 0$ unless $q = u$. Then, from Eqs. (25) and (26) longitudinal asymmetry is

$$\begin{aligned} A_{UL}^{\sin\phi} &= 0.88 \text{ GeV}\frac{|K_\perp|}{zQ}[1-f(\rho Q)]\frac{y(1-y/2)\sqrt{1-y}}{M[1+(1-y)^2]}\\ &\times \frac{\Delta f_u(x)}{x f_u(x)}. \end{aligned} \quad (29)$$

The ratio of polarized to unpolarized distribution function is measured by the HERMES Collaboration [26–28] for the same kinematical region as spin asymmetries; this is shown in Fig. 3. It may be fitted with reasonable accuracy by simple power law $\Delta u/u = x^\alpha$ with $\alpha = 0.68 \pm 0.08$. [Note that this dependence should not be true down to very low x, or else the $\Delta u/(xu)$ blows up.] Parameters Q^2, x, and y are related by $Q^2 = xy(s - M^2)$ [here s is the Mandelstam variable, $s = 2ME$ in the proton rest frame]. $|K_\perp|$ and z can be taken as being independent from the rest of the kinematical variables as long as $|K_\perp|/z \ll Q$. Otherwise DIS separation of parallel and transversal degrees of freedom breaks down. In the HERMES experiment $\langle K_\perp\rangle = 0.44$ and $\langle z\rangle = 0.48$, while Q^2 is constrained

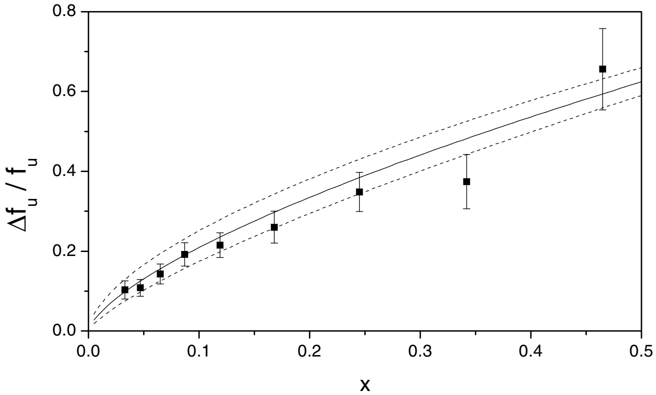

FIG. 3. Relative polarization of u-quark in the proton. The error bars combine statistical and systematic errors. Parametrization x^{α} with $\alpha = 0.68 \pm 0.08$ is shown by solid (dashed) lines.

to be >1 GeV^2. Thus, we assume $|K_\perp|/z = 0.92$ GeV in Eq. (29).

(In reality, the data do not fulfill the strong requirement $|K_\perp|/z \ll Q$ and therefore, the kinematical structures involving higher powers of $K_\perp/Q$ can also be relevant. The comparison with only the small $K_\perp/Q$ part of the data sample would be more adequate.)

Following HERMES kinematical cuts we average over $0.2 < y < 0.85$. The result for moderate values of x is shown in Fig. 4. We have excluded $x < 0.1$ because for small x our simplifying assumptions about proton structure are not applicable.

The relation of transversal to longitudinal asymmetries for the same simplified model of π^+ production we use is [see Eq. (25)]

$$\frac{A_{UT}}{A_{UL}} = \frac{2\sqrt{1-y}\sqrt{x}}{(1-y/2)y^{3/2}} \frac{M}{\sqrt{s-M^2}} \frac{\Delta_T f_u(x)}{\Delta f_u(x)}. \tag{30}$$

Recall that in our simplified model $\Delta_T f_u(x) = \Delta f_u(x)$. Taking into account HERMES kinematics as outlined above one finally has an estimate

$$\frac{A_{UT}}{A_{UL}} = 1.92\sqrt{x}, \tag{31}$$

which is compared to the available data in Fig. 5

V. CONCLUSIONS AND OUTLOOK

In this paper we have made a step toward the semiclassical theory of various spin-dependent effects in QCD, based on instantons. We repeat that we discuss neither initial state (Sivers) nor final state (Collins) effects. The instantons may well contribute to those; future work will be needed to clarify their contribution.

In this work, as the first step, we single out a dynamical process which takes time *during* the collision. Both the quark chirality flip and the T-odd phase (needed to make asymmetry visible) are associated with the instanton fluctuation, which happens in the QCD vacuum close to the DIS collision point. We emphasize that we do not introduce any new physical process or parameter, or new structure/fragmentation functions, but express the result in terms of the well-known quantities. The main one of them is the small "vacuum diluteness" parameter (27), which fortunately gets compensated by a large numerical factor 32π in the answer and provides for a result comparable to experimental value. The magnitude of the effect is thus fixed with the absolute normalization (29), based on the parameters of the instanton ensemble model known since 1982; see [17]. The result agrees in sign and magnitude with the available experimental data in suitable kinematic domain. We have argued that the asymmetry does not depend on the specific distribution functions of the nucleon, and is thus universal to all other hadrons.

Furthermore, our spin-dependent azimuthal asymmetries have a particular tensor structure in the lowest non-

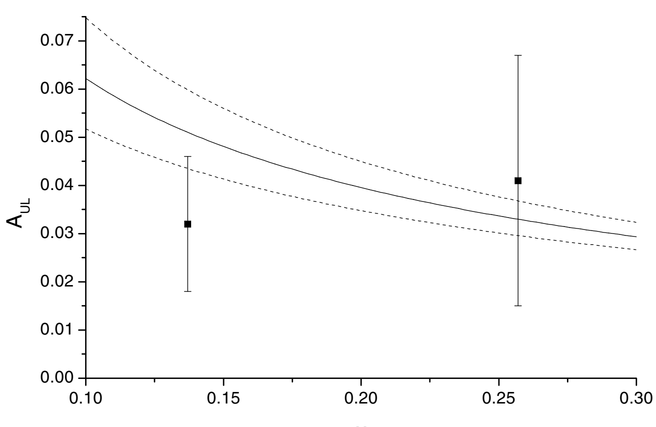

FIG. 4. Experimental values A_{UL} for moderate x are shown with comparison with the model prediction. Theoretical uncertainty is due to uncertainty in polarized distribution function.

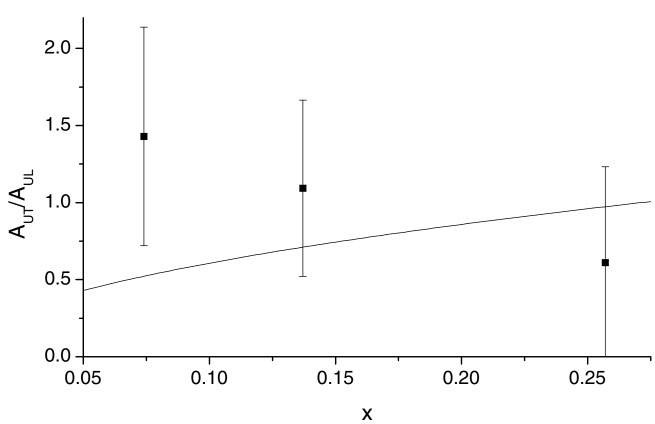

FIG. 5. Experimental values of relative transversal to longitudinal asymmetries A_{UT}/A_{UL} for moderate x are shown with comparison with the model prediction.

zero order of $K_\perp/Q$ as long as parton interpretation of hadron structure is taken into account. It leads to the specific prediction for the dependence of longitudinal and transverse asymmetries on kinematical parameters which is completely independent on the phenomenological inputs.

For the outlook, one may ask if an explanation for other known spin asymmetries, e.g., in $pp\uparrow$ collisions, can also be provided by a similar mechanisms based on instantons. In particular, well-known FERMILAB data [29–31] revealed considerable asymmetry in pion production for $x_F > 0.5$. The explanation of this asymmetry based on instanton mechanism was pioneered by Kochelev [32,33] who provided qualitative expressions for it. More quantitative calculation would, however, be needed to relate this effect to spin effects in DIS we discuss above. A promising direction for future work based on instantons may be a description of the *nonpolarized* DIS in this domain, $x_F \rightarrow 1$. It is known to be dominated by large higher twist effects, while at the same time twist expansion seems badly convergent by itself. It was speculated long ago [34] that this region may be described by instanton-based dynamics instead, but it was never demonstrated quantitatively. If that conjecture happens to be true, the instanton diluteness κ would drop out from the numerator and denominator of the spin asymmetry, resulting in a really large $\sim O(1)$ and truly universal asymmetry.

ACKNOWLEDGMENTS

One of us (E. S.) thanks S. Brodsky for useful conversation on the subject. We are also grateful to A. Bacchetta for his helpful comments on the first version of the manuscript. The work is partially supported by US-DOE Grants No. DE-FG02-88ER40388 and No. DE-FG03-97ER4014.

[1] S. Forte and E. V. Shuryak, Nucl. Phys. **B357**, 153 (1991).
[2] N. I. Kochelev, hep-ph/9511299, and subsequent works
[3] T. Schafer and V. Zetocha, Phys. Rev. D **69**, 094028 (2004).
[4] J. W. Negele *et al.*, Nucl. Phys., Proc. Suppl. **B128**, 170 (2004).
[5] P. J. Mulders and R. D. Tangeman, Nucl. Phys. **B461**, 197 (1996); **B484**, 538(E) (1997).
[6] D. W. Sivers, Phys. Rev. D **43**, 261 (1991).
[7] D. W. Sivers, Phys. Rev. D **41**, 83 (1990).
[8] J. C. Collins, Nucl. Phys. **B396**, 161 (1993).
[9] X. Artru, J. Czyzewski, and H. Yabuki, Z. Phys. C **73**, 527 (1997).
[10] J. W. Qiu and G. Sterman, Phys. Rev. D **59**, 014004 (1999).
[11] S. J. Brodsky, D. S. Hwang, and I. Schmidt, Int. J. Mod. Phys. A **18**, 1327 (2003).
[12] S. J. Brodsky, D. S. Hwang, and I. Schmidt, Phys. Lett. B **530**, 99 (2002).
[13] J. C. Collins, Phys. Lett. B **536**, 43 (2002).
[14] A. Bacchetta, R. Kundu, A. Metz, and P. J. Mulders, Phys. Lett. B **506**, 155 (2001).
[15] F. Yuan, Phys. Lett. B **575**, 45 (2003).
[16] A. Bacchetta, A. Schaefer, and J. J. Yang, Phys. Lett. B **578**, 109 (2004).
[17] T. Schafer and E. V. Shuryak, Rev. Mod. Phys. **70**, 323 (1998).
[18] P. Faccioli and E. V. Shuryak, Phys. Rev. D **64**, 114020 (2001).
[19] HERMES Collaboration, A. Airapetian *et al.*, Phys. Lett. B **562**, 182 (2003).
[20] HERMES Collaboration, A. Airapetian *et al.*, Phys. Lett. B **535**, 85 (2002).
[21] HERMES Collaboration, A. Airapetian *et al.*, Phys. Rev. Lett. **84**, 4047 (2000).
[22] HERMES Collaboration, A. Airapetian *et al.*, hep-ex/0408013.
[23] CLAS Collaboration, H. Avakian *et al.*, hep-ex/0301005.
[24] S. Moch, A. Ringwald, and F. Schrempp, Nucl. Phys. **B507**, 134 (1997)
[25] R. L. Jaffe, hep-ph/9602236.
[26] HERMES Collaboration, K. Ackerstaff *et al.*, Phys. Lett. B **464**, 123 (1999).
[27] HERMES Collaboration, A. Airapetian *et al.*, Phys. Lett. B **442**, 484 (1998).
[28] HERMES Collaboration, K. Ackerstaff *et al.*, Phys. Lett. B **404**, 383 (1997).
[29] FNAL-E704 Collaboration, D. L. Adams *et al.*, Phys. Lett. B **264**, 462 (1991).
[30] E581 Collaboration, D. L. Adams *et al.*, Phys. Lett. B **261**, 201 (1991).
[31] Fermilab E704 Collaboration, A. Bravar *et al.*, Phys. Rev. Lett. **77**, 2626 (1996).
[32] N. I. Kochelev, JETP Lett. **72**, 481 (2000) [Pis'ma Zh. Eksp. Teor. Fiz. **72**, 691 (2000)].
[33] N. I. Kochelev, hep-ph/9711274.
[34] E. V. Shuryak and A. I. Vainshtein, Nucl. Phys. **B199**, 451 (1982); **B201**, 141 (1982).