

Twist-3 single-spin asymmetries in semi-inclusive deep-inelastic scattering

A. Metz^a and M. Schlegel

Institut für Theoretische Physik II, Ruhr-Universität Bochum, D-44780 Bochum, Germany

Received: 17 May 2004 / Revised version: 2 July 2004 /

Published online: 11 November 2004 – © Società Italiana di Fisica / Springer-Verlag 2004

Communicated by U.-G. Meißner

Abstract. The single-spin asymmetries for a longitudinally polarized lepton beam or a longitudinally polarized nucleon target in semi-inclusive deep-inelastic scattering are twist-3 observables. We study these asymmetries in a simple diquark spectator model of the nucleon. Analogous to the case of transverse target polarization, non-vanishing asymmetries are generated by gluon exchange between the struck quark and the target system. It is pointed out that the coupling of the virtual photon to the diquark is needed in order to preserve electromagnetic gauge invariance at the twist-3 level. The calculation indicates that previous analyses of these observables are incomplete.

PACS. 13.60.Hb Total and inclusive cross-sections (including deep-inelastic processes) – 13.88.+e Polarization in interactions and scattering

1 Introduction

The measurements of A_{UL} (longitudinal target polarization) and A_{LU} (longitudinal lepton beam polarization) by the HERMES [1–3] and CLAS [4] Collaborations constitute the first clear evidence of non-vanishing single-spin asymmetries (SSA) in semi-inclusive deep-inelastic scattering (SIDIS) off the nucleon. From a theoretical point of view, SSA in hard processes are very interesting because of their relation to time-reversal odd (T-odd) correlation functions (parton distributions and fragmentation functions).

Already for more than a decade the existence of T-odd fragmentation functions has been considered to be established [5]. In the meantime, explicit model calculations including final-state interactions in the fragmentation process have provided non-vanishing results for such functions (see, *e.g.*, refs. [6, 7]). On the other hand, it has been shown only recently that non-zero T-odd parton distributions are compatible with time-reversal invariance of the strong interaction [8, 9] (see also refs. [10–12] for related work). In DIS, T-odd parton distributions arise due to the exchange of longitudinally polarized gluons between the struck quark and the target system. This rescattering effect is encoded in the gauge link appearing in the definition of parton distributions. An alternative picture according to which T-odd parton distributions can be gen-

erated without rescattering of the struck quark [13] seems to be ruled out [14].

In particular, due to the recent developments [8, 9], the T-odd and transverse-momentum-dependent (k_{\perp} -dependent) functions f_{1T}^{\perp} (Sivers function, describing the distribution of unpolarized quarks in a transversely polarized target) [15] and h_1^{\perp} (distribution of transversely polarized quarks in an unpolarized target) [16] may well exist. From a practical point of view both distributions can be considered as twist-2 functions, since they appear in observables at leading order of a $1/Q$ -expansion, where Q denotes the large scale of the hard process. For instance, f_{1T}^{\perp} enters the leading twist SSA A_{UT} in SIDIS [16].

Despite the recent progress in understanding the nature of T-odd effects, a complete formalism (including subleading T-odd parton distributions) of T-odd twist-3 observables is still missing even at tree level. So far, such effects have only been treated on the fragmentation side [17, 18]. This point may also be quite important for the description of the twist-3 asymmetries A_{UL} and A_{LU} in SIDIS. With the exception of refs. [19, 20], all present analyses/calculations of these observables are based on [18], *i.e.*, they include only T-odd fragmentation functions (see, *e.g.*, refs. [21–26]). In this scenario one finds schematically $A_{LU} \propto e H_1^{\perp}$ and $A_{UL} \propto h_L H_1^{\perp}$, with the twist-3 T-even parton distributions e and h_L , and the twist-2 T-odd Collins fragmentation function H_1^{\perp} [5]¹.

^a e-mail: andreas.metz@tp2-ruhr-uni-bochum.de

¹ In our calculation for A_{UL} the target is polarized along the direction of its momentum (and the direction of the virtual

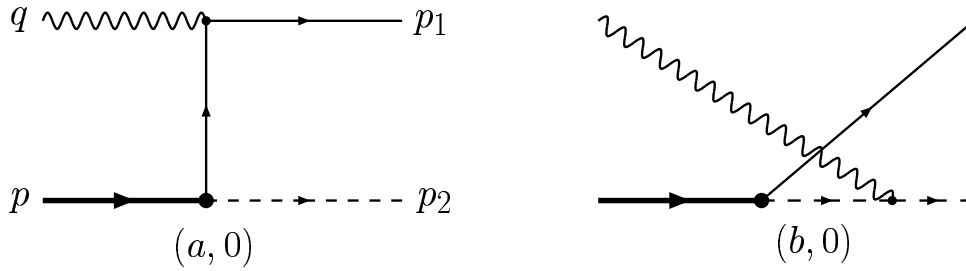


Fig. 1. Tree level diagrams of the process in (1). In order to preserve gauge invariance both diagrams have to be considered.

Analogous to the treatment of A_{UT} presented in ref. [8], we compute A_{LU} and A_{UL} in the framework of a simple diquark spectator model of the nucleon, in order to investigate whether T-odd parton distributions may be relevant in these cases. The rescattering of the struck quark, which serves as the potential source of T-odd effects, is modelled by the exchange of an Abelian gauge boson. Such a study has already been performed in ref. [27]. However, in [27] only A_{LU} has been computed explicitly. Moreover, the calculation of [27] is not gauge invariant. To preserve electromagnetic-current conservation also the coupling of the virtual photon to the diquark has to be included.

Both A_{LU} and A_{UL} turn out to be non-zero indicating that, in a factorized picture, T-odd distributions have to be taken into account. A first step in this direction has been made in [19], where it has been demonstrated that the T-odd distribution h_1^\perp appears in the description of A_{LU} through a term $h_1^\perp E$, where E is a twist-3 fragmentation function. As will be discussed below, our calculation of A_{LU} , however, cannot be identified with such a term suggesting that the formula of [19] for the beam SSA is not yet complete.

2 Tree diagrams

In order to study SIDIS off a spin- $\frac{1}{2}$ particle (for definiteness we think of a proton) in the framework of a spectator model we consider the process (compare also ref. [8])

$$\gamma^*(q) + p(p, \lambda) \rightarrow q(p_1, \lambda') + s(p_2). \quad (1)$$

In full SIDIS, both the quark and the spectator in the final state fragment into hadrons, where we are interested in the situation that one of the hadrons from the quark fragmentation is detected at low transverse momentum. However, for the study of possible T-odd effects related with parton distributions it is not necessary to include the fragmentation process in the calculation. We use the model of [8] with a scalar diquark spectator s . In this model the proton has no electromagnetic charge, and a charge e_1 is assigned to the quark. The interaction between

photon). In experiments for the longitudinal target asymmetry, however, the polarization is along the direction of the incoming lepton. Both situations differ by a kinematical twist-3 term which is given by A_{UT} .

the proton, the quark and the spectator is described by a scalar vertex with the coupling constant g .

We treat the process (1) in the Breit frame of the virtual photon. The proton has a large plus-momentum Q/x , where $x = x_{Bj} + \mathcal{O}(1/Q^2)$. The quark carries the large minus-momentum $p_1^- \approx q^-$ and a soft transverse momentum $\vec{\Delta}_\perp$. These requirements specify the kinematics:

$$\begin{aligned} q &= \left(-Q, Q, \vec{0}_\perp \right), & p &= \left(\frac{Q}{x}, \frac{xM^2}{Q}, \vec{0}_\perp \right), \\ p_1 &= \left(\frac{\vec{\Delta}_\perp^2}{Q}, Q, \vec{\Delta}_\perp \right), \\ p_2 &= \left(\frac{Q(1-x)}{x}, \frac{x(\vec{\Delta}_\perp^2 + m_s^2)}{Q(1-x)}, -\vec{\Delta}_\perp \right). \end{aligned} \quad (2)$$

The expressions for q and p are exact, while for p_1 and p_2 just the leading terms have been listed. In particular, sometimes the $1/Q^2$ corrections of p_1^- and p_2^+ are needed which can be readily obtained from 4-momentum conservation. To simplify the calculation we consider massless quarks.

The tree level diagrams of the process (1) are shown in fig. 1. Their currents, depending on the helicities of the proton and the quark, read

$$J_{(a,0)}^\mu(\lambda, \lambda') = e_1 g \frac{1}{(p_1 - q)^2} \bar{u}(p_1, \lambda') \gamma^\mu (\not{p}_1 - \not{q}) u(p, \lambda), \quad (3)$$

$$J_{(b,0)}^\mu(\lambda, \lambda') = -e_1 g \frac{1}{(p_2 - q)^2 - m_s^2} (2p_2^\mu - q^\mu) \bar{u}(p_1, \lambda') u(p, \lambda). \quad (4)$$

We have defined the current by means of the scattering amplitude according to $T = \varepsilon_\mu J^\mu$, with ε denoting the polarization vector of the virtual photon. It is easy to check that current conservation holds for the sum of the two diagrams, *i.e.*,

$$q_\mu \left(J_{(a,0)}^\mu + J_{(b,0)}^\mu \right) = 0. \quad (5)$$

As long as one is just interested in leading twist observables, it is sufficient to consider diagram $(a, 0)$. This is for instance the case in the calculation of the transverse SSA A_{UT} in ref. [8]. The specific kinematics in eq. (2) is the reason for the suppression of diagram $(b, 0)$ relative

$$J_{(a,1)}^\mu(\lambda, \lambda') = i(e_1)^3 g \int \frac{d^4 k}{(2\pi)^4} \frac{\bar{u}(p_1, \lambda') (\not{p}_1 + 2\not{p}_2 - \not{k}) \not{k} \gamma^\mu (\not{k} - \not{q}) u(p, \lambda)}{[k^2 + i\epsilon][(p-q)^2 + i\epsilon][(p+q-k)^2 - m_s^2 + i\epsilon][(k-p_1)^2 - \mu^2 + i\epsilon]}, \quad (13)$$

$$J_{(b,1)}^\mu(\lambda, \lambda') = -i(e_1)^3 g \int \frac{d^4 k}{(2\pi)^4} \frac{\bar{u}(p_1, \lambda') (\not{p}_1 + 2\not{p}_2 - \not{k}) \not{k} (2\not{p}^\mu + \not{q}^\mu - 2\not{k}^\mu) u(p, \lambda)}{[k^2 + i\epsilon][(p-k)^2 - m_s^2 + i\epsilon][(p+q-k)^2 - m_s^2 + i\epsilon][(k-p_1)^2 - \mu^2 + i\epsilon]} \quad (14)$$

to $(a, 0)$. The propagator of the diquark in $(b, 0)$ behaves like $1/Q^2$, while there is no large momentum flow through the quark propagator in $(a, 0)$. Nevertheless, as we discuss in the following, for subleading twist observables diagram $(b, 0)$ can no longer be neglected.

We compute the various components of the currents using the lightfront helicity spinors of [28]. For $J_{(a,0)}^\mu$ one obtains

$$J_{(a,0)}^1(\lambda, \lambda') = -e_1 g \frac{1-x}{\sqrt{x} |\vec{\Delta}_\perp|} \frac{Q}{\vec{\Delta}_\perp^2 + \tilde{m}^2} \times \left[Mx (\Delta^1 - i\lambda\Delta^2) \delta_{\lambda, \lambda'} - \lambda (\Delta^1 + i\lambda\Delta^2)^2 \delta_{\lambda, -\lambda'} \right], \quad (6)$$

$$J_{(a,0)}^2(\lambda, \lambda') = -e_1 g \frac{1-x}{\sqrt{x} |\vec{\Delta}_\perp|} \frac{Q}{\vec{\Delta}_\perp^2 + \tilde{m}^2} \times \left[Mx (i\lambda\Delta^1 + \Delta^2) \delta_{\lambda, \lambda'} - i (-i\lambda\Delta^1 + \Delta^2)^2 \delta_{\lambda, -\lambda'} \right], \quad (7)$$

$$J_{(a,0)}^+(\lambda, \lambda') = -2e_1 g \frac{1-x}{\sqrt{x}} \frac{|\vec{\Delta}_\perp|}{\vec{\Delta}_\perp^2 + \tilde{m}^2} \times \left[Mx \delta_{\lambda, \lambda'} - (\lambda\Delta^1 + i\Delta^2) \delta_{\lambda, -\lambda'} \right], \quad (8)$$

$$J_{(a,0)}^-(\lambda, \lambda') = -2e_1 g \frac{1-x}{\sqrt{x}} \frac{|\vec{\Delta}_\perp|}{\vec{\Delta}_\perp^2 + \tilde{m}^2} \times \left[Mx \delta_{\lambda, \lambda'} - \frac{xM^2}{\vec{\Delta}_\perp^2} \left(1 - \frac{\vec{\Delta}_\perp^2 + m_s^2}{M^2(1-x)} \right) \times (\lambda\Delta^1 + i\Delta^2) \delta_{\lambda, -\lambda'} \right], \quad (9)$$

$$\text{with } \tilde{m}^2 = x(1-x) \left(-M^2 + \frac{m_s^2}{1-x} \right).$$

One observes here the well-known result that for DIS off a spin- $\frac{1}{2}$ particle the transverse current is dominating in the Breit frame. For the second tree graph we find

$$J_{(b,0)}^+(\lambda, \lambda') = -e_1 g \frac{2-x}{\sqrt{x} |\vec{\Delta}_\perp|} (\lambda\Delta^1 + i\Delta^2) \delta_{\lambda, -\lambda'}, \quad (10)$$

$$J_{(b,0)}^-(\lambda, \lambda') = e_1 g \frac{\sqrt{x}}{|\vec{\Delta}_\perp|} (\lambda\Delta^1 + i\Delta^2) \delta_{\lambda, -\lambda'}. \quad (11)$$

The transverse components $J_{(b,0)}^i$ are proportional to $1/Q$ and, hence, indeed are suppressed by a factor $1/Q^2$ compared to $J_{(a,0)}^i$ as expected. Therefore, these terms are not relevant for the discussion of twist-3 observables. In

contrast, the plus and the minus components for both diagrams are of the same order. In this case, the suppression of $(b, 0)$ caused by the propagator of the diquark is compensated by a factor Q at the photon-diquark vertex and the fact that $\bar{u}(p_1, \lambda') u(p, \lambda) \propto Q$. Even though diagram $(b, 0)$ is not compatible with the parton model, since large momentum transfers at the proton-quark-diquark vertex are allowed, a consistent calculation of twist-3 observables in the spectator model must contain this contribution. We note that our results obey the gauge invariance constraint

$$J_{(a,0)}^+ + J_{(b,0)}^+ = J_{(a,0)}^- + J_{(b,0)}^-. \quad (12)$$

Including by hand a formfactor at the proton-quark-diquark vertex in order to suppress large momentum transfers destroys the gauge invariance.

3 One-loop diagrams

To obtain non-zero SSA in a spectator model one has to go beyond the tree level approximation and take the rescattering of the quark into account. For our purpose it is sufficient to model this effect by one-photon exchange. Since the imaginary part of the one-loop amplitude is needed for the computation of SSA, just the two diagrams in fig. 2 have to be considered. Self-energy and vertex correction diagrams are relevant for the real part of the amplitude, but cannot acquire an imaginary part, because we are dealing with either on-shell or even space-like (internal) lines. In order to avoid infrared singularities at intermediate steps of the calculation we assign a mass μ to the photon. The final results for A_{LU} and A_{UL} must be infrared finite which serves as a non-trivial check of the calculation. The currents of the diagrams in fig. 2 are given by

see eqs. (13) and (14) above.

Applying Cutkosky rules to calculate the imaginary part, one can verify the gauge invariance condition

$$q_\mu \left(\text{Im } J_{(a,1)}^\mu + \text{Im } J_{(b,1)}^\mu \right) = 0. \quad (15)$$

It is obvious that the full current (including the real part) for the sum of both diagrams is not gauge invariant.

The calculation of the imaginary parts has been performed similar to the study of T-odd fragmentation in ref. [29]. We refrain from giving any details and just quote

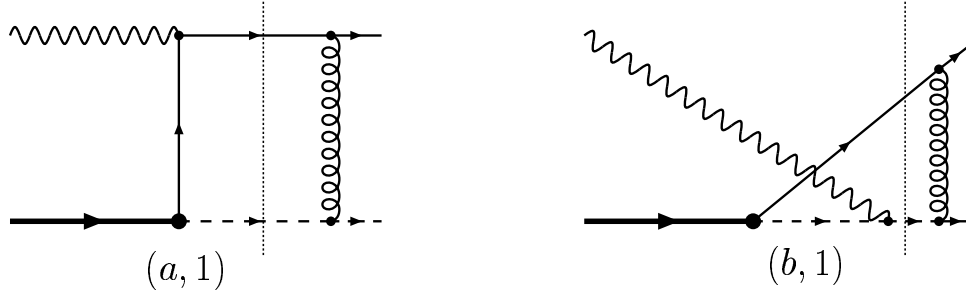


Fig. 2. One-loop diagrams for the imaginary part of the process in (1). In each diagram the possible on-shell intermediate state is indicated by a thin dotted line. The interaction between the quark and the spectator is modelled by the exchange of an Abelian gauge boson.

the final results. For diagram (a, 1) we find

$$\begin{aligned} \text{Im } J_{(a,1)}^1(\lambda, \lambda') &= -\frac{(e_1)^3 g}{8\pi} \frac{1-x}{\sqrt{x}|\vec{\Delta}_\perp|} Q \left[\frac{\ln L}{\sqrt{\Lambda}} Mx \right. \\ &\times \left(\Delta^1 - i\lambda\Delta^2 \right) \delta_{\lambda, \lambda'} - \frac{1}{2\Delta_\perp^2} \left(\left(\vec{\Delta}_\perp^2 - \tilde{m}^2 + \mu^2 \right) \frac{\ln L}{\sqrt{\Lambda}} \right. \\ &\left. \left. + \ln \frac{\tilde{m}^2}{\mu^2} \right) \lambda \left(\Delta^1 + i\lambda\Delta^2 \right)^2 \delta_{\lambda, -\lambda'} \right], \end{aligned} \quad (16)$$

$$\begin{aligned} \text{Im } J_{(a,1)}^2(\lambda, \lambda') &= -\frac{(e_1)^3 g}{8\pi} \frac{1-x}{\sqrt{x}|\vec{\Delta}_\perp|} Q \left[\frac{\ln L}{\sqrt{\Lambda}} Mx \right. \\ &\times \left(i\lambda\Delta^1 + \Delta^2 \right) \delta_{\lambda, \lambda'} - \frac{1}{2\Delta_\perp^2} \left(\left(\vec{\Delta}_\perp^2 - \tilde{m}^2 + \mu^2 \right) \frac{\ln L}{\sqrt{\Lambda}} \right. \\ &\left. \left. + \ln \frac{\tilde{m}^2}{\mu^2} \right) i \left(-i\lambda\Delta^1 + \Delta^2 \right)^2 \delta_{\lambda, -\lambda'} \right], \end{aligned} \quad (17)$$

$$\begin{aligned} \text{Im } J_{(a,1)}^+(\lambda, \lambda') &= -\frac{(e_1)^3 g}{8\pi} \frac{1-x}{\sqrt{x}|\vec{\Delta}_\perp|} \left[Mx \right. \\ &\times \left(\left(\vec{\Delta}_\perp^2 - \tilde{m}^2 + \mu^2 \right) \frac{\ln L}{\sqrt{\Lambda}} + \ln \frac{\tilde{m}^2}{\mu^2} \right) \delta_{\lambda, \lambda'} \\ &\left. + 2 \left(\tilde{m}^2 \frac{\ln L}{\sqrt{\Lambda}} - \ln \frac{Q^2(1-x)}{\mu^2 x} + 1 \right) \left(\lambda\Delta^1 + i\Delta^2 \right) \delta_{\lambda, -\lambda'} \right], \end{aligned} \quad (18)$$

$$\begin{aligned} \text{with } L &= \frac{\vec{\Delta}_\perp^2 + \tilde{m}^2 + \mu^2 + \sqrt{\Lambda}}{\vec{\Delta}_\perp^2 + \tilde{m}^2 + \mu^2 - \sqrt{\Lambda}}, \\ \Lambda &= \left(\vec{\Delta}_\perp^2 + \tilde{m}^2 - \mu^2 \right)^2 + 4\vec{\Delta}_\perp^2 \mu^2. \end{aligned}$$

The plus-component of diagram (b, 1) is given by

$$\begin{aligned} \text{Im } J_{(b,1)}^+(\lambda, \lambda') &= -\frac{(e_1)^3 g}{8\pi} \frac{1}{\sqrt{x}|\vec{\Delta}_\perp|} \\ &\times \left((2-x) \ln \frac{Q^2(1-x)}{\mu^2 x} - 2(1-x) \right) \left(\lambda\Delta^1 + i\Delta^2 \right) \delta_{\lambda, -\lambda'}. \end{aligned} \quad (19)$$

For the one-loop calculation we make use of gauge invariance to eliminate the minus component of the current. The Q behaviour of the one-loop expressions corresponds

to the one of the tree graphs. Note also that the plus component of the currents for both diagrams contains a $\ln Q^2$ -term, which is not compatible with the parton model. From our results for the transverse currents in eqs. (16), (17) we were able to reproduce the transverse target SSA computed in ref. [8] (up to an overall sign).

4 Spin asymmetries

Eventually, we proceed to the calculation of A_{LU} and A_{UL} . The full cross-section in DIS (including the leptons) in the one-photon exchange approximation can be expressed in the standard form

$$\sigma \propto L_{\mu\nu} W^{\mu\nu}, \quad (20)$$

with the lepton tensor

$$L^{\mu\nu} = 2 \left(l^\mu l'^\nu + l'^\nu l^\mu - \frac{Q^2}{2} g^{\mu\nu} + i\lambda_e \varepsilon^{\mu\nu\rho\sigma} q_\rho l_\sigma \right). \quad (21)$$

In eq. (21), the 4-momentum of the incoming (outgoing) lepton is denoted by l (l'), and $Q^2 = -(l-l')^2$. The hadron tensor is obtained from the above currents according to

$$W^{\mu\nu} = \left(J^\mu \right)^\dagger J^\nu. \quad (22)$$

Now, we exploit gauge invariance of both the lepton and hadron tensor, take q from eq. (2), choose the lepton momenta to be in the xz -plane, and ignore a contribution of W^{++} which for our calculation is at least suppressed by a factor $1/Q^2$ relative to the leading term in the cross-section. This allows us to write the cross-section as

$$\begin{aligned} \sigma \propto \frac{4Q^2}{y^2} \left[- (2-y) \sqrt{1-y} W_S^{+1} + \left(1-y + \frac{y^2}{4} \right) W_S^{11} \right. \\ \left. + \frac{y^2}{4} W_S^{22} + i\lambda_e \left(y \sqrt{1-y} W_A^{+2} - y \left(1 - \frac{y}{2} \right) W_A^{12} \right) \right], \end{aligned} \quad (23)$$

where $W_S^{\mu\nu} = (W^{\mu\nu} + W^{\nu\mu})/2$ and $W_A^{\mu\nu} = (W^{\mu\nu} - W^{\nu\mu})/2$ represent the symmetric and the antisymmetric part of the hadronic tensor, respectively. We also used the standard definition $y = p \cdot q / p \cdot l$. The first line in eq. (23)

is relevant for the target spin asymmetry, and the second one for the beam asymmetry. Actually, it turns out that the purely transverse components of the hadronic tensor (W_S^{11} , W_S^{22} , W_A^{12}) do not contribute to the spin asymmetries at twist-3 level.

In order to specify the asymmetries, we define

$$W_{UU}^{11} = \frac{1}{2} \sum_{\lambda, \lambda'} \left(J^1(\lambda, \lambda') \right)^\dagger J^1(\lambda, \lambda'), \quad (24)$$

$$W_{A,LU}^{+2} = \frac{1}{4} \sum_{\lambda, \lambda'} \left[\left(J^+(\lambda, \lambda') \right)^\dagger J^2(\lambda, \lambda') - \left(J^2(\lambda, \lambda') \right)^\dagger J^+(\lambda, \lambda') \right], \quad (25)$$

$$W_{S,UL}^{+1} = \frac{1}{4} \sum_{\lambda'} \left[\left(\left(J^+(\uparrow, \lambda') \right)^\dagger J^1(\uparrow, \lambda') - \left(J^+(\downarrow, \lambda') \right)^\dagger J^1(\downarrow, \lambda') \right) + \left(\left(J^1(\uparrow, \lambda') \right)^\dagger J^+(\uparrow, \lambda') - \left(J^1(\downarrow, \lambda') \right)^\dagger J^+(\downarrow, \lambda') \right) \right], \quad (26)$$

where polarization “ \uparrow ” in (26) means polarization along the positive z -axis, *i.e.*, along the direction of the target momentum. The element W_{UU}^{11} is given by the tree level result for diagram $(a, 0)$ in eq. (6). Non-zero contributions to $W_{A,LU}^{+2}$ and $W_{S,UL}^{+1}$ are generated by interference of the tree level amplitude with the imaginary part of the one-loop amplitude. While the transverse currents in (25) and (26) are obtained from diagrams $(a, 0)$ and $(a, 1)$, all four diagrams contribute to the plus component of the current. The final results for the asymmetries read²

$$A_{LU} = \frac{iW_{A,LU}^{+2}}{W_{UU}^{11}} = \frac{(e_1)^2}{4\pi} \frac{\vec{\Delta}_\perp^2 + \tilde{m}^2}{M^2 x^2 + \vec{\Delta}_\perp^2} \frac{\Delta^2}{Q} \times \left[\frac{1}{\vec{\Delta}_\perp^2} \left(-M^2 x^2 - \tilde{m}^2 \frac{2-x}{2(1-x)} \right) \ln \frac{\vec{\Delta}_\perp^2 + \tilde{m}^2}{\tilde{m}^2} - \frac{x}{2(1-x)} \ln \frac{Q^2(1-x)}{(\vec{\Delta}_\perp^2 + \tilde{m}^2)x} \right], \quad (27)$$

$$A_{UL} = \frac{W_{S,UL}^{+1}}{W_{UU}^{11}} = \frac{(e_1)^2}{4\pi} \frac{\vec{\Delta}_\perp^2 + \tilde{m}^2}{M^2 x^2 + \vec{\Delta}_\perp^2} \frac{\Delta^2}{Q} \times \left[\frac{1}{\vec{\Delta}_\perp^2} \left(M^2 x^2 - \tilde{m}^2 \frac{2-x}{2(1-x)} \right) \ln \frac{\vec{\Delta}_\perp^2 + \tilde{m}^2}{\tilde{m}^2} - \frac{x}{2(1-x)} \ln \frac{Q^2(1-x)}{(\vec{\Delta}_\perp^2 + \tilde{m}^2)x} \right]. \quad (28)$$

² As a reference we mention that the transverse target spin asymmetry of [8] is given by $A_{UT} = W_{UT}^{11}/W_{UU}^{11}$, with $W_{UT}^{11} = W_{S,UT}^{11}$ defined analogous to eq. (26).

We would like to add some remarks:

- An explicit non-zero result for A_{LU} in the framework of the diquark spectator model has already been obtained in ref. [27]. Our calculation shows that both A_{LU} and A_{UL} remain finite once all diagrams required by electromagnetic gauge invariance are taken into consideration.
- We believe that the effect which generates A_{LU} in our calculation is not related to a term proportional to $h_\perp^\dagger E$ discussed in ref. [19]. While h_\perp^\dagger is chirally odd, we have summed over the polarizations of the outgoing quark.
- Since the asymmetries are proportional to $\Delta^2 = |\vec{\Delta}_\perp| \sin \phi_q$ we expect in full SIDIS an effect proportional to $\sin \phi_h$, where ϕ_h is the azimuthal angle of the produced hadron. The mechanisms which have been discussed so far in the literature in connection with A_{LU} and A_{UL} [17–19] show the same $\sin \phi_h$ behaviour. In addition, the different contributions to the asymmetries have the same y -dependence ($y\sqrt{1-y}$ for A_{LU} and $(2-y)\sqrt{1-y}$ for A_{UL} , see eq. (23)).
- In the final results for the asymmetries we have performed the limit $\mu \rightarrow 0$ without encountering a divergence. We agree with the observation made in ref. [27], that the contribution from diagrams $(a, 0)$ and $(a, 1)$ to A_{LU} is separately infrared finite. This behaviour, which holds for A_{UL} as well, seems to be accidental.

5 Summary and conclusions

In summary, we have calculated the twist-3 single-spin asymmetries A_{LU} and A_{UL} for semi-inclusive DIS off a nucleon target in the framework of a simple diquark spectator model. Our study completes the previous work in ref. [27], where only A_{LU} has been computed explicitly. Moreover, the treatment in [27] is lacking electromagnetic gauge invariance.

Both A_{LU} and A_{UL} turn out to be non-zero. Although the spectator model calculation contains contributions which are not compatible with the parton model, the non-vanishing results indicate that T-odd distributions have to be included in a factorized description of the asymmetries. So far this has only been done partly in the literature. In fact, we have argued that apparently none of the present analyses/calculations of A_{LU} and A_{UL} within the parton model is complete. Mainly for two reasons we feel confident to make such a speculation: first, within our calculation non-zero asymmetries arise already from the diagram $(a, 1)$, whose kinematics is compatible with the parton model. Second, there is no reason why the asymmetries should not contain higher-order T-odd distribution functions. The status of the parton model formulae for A_{LU} and A_{UL} needs to be clarified before any definite conclusion can be extracted from the data.

Note added: After this work has been completed, a revised parton model analysis for A_{LU} and A_{UL} appeared [30]. The analysis confirms our suspicion that both asymmetries should contain an additional term with a twist-3

T-odd distribution function, which have not been taken into account in the literature before.

We are grateful to J.C. Collins and N. Kivel for discussions. The work has been partly supported by the Sofia Kovalevskaya Programme of the Alexander von Humboldt Foundation, the DFG and the COSY-Juelich project.

References

1. HERMES Collaboration (A. Airapetian *et al.*), Phys. Rev. Lett. **84**, 4047 (2000).
2. HERMES Collaboration (A. Airapetian *et al.*), Phys. Rev. D **64**, 097101 (2001).
3. HERMES Collaboration (A. Airapetian *et al.*), Phys. Lett. B **562**, 182 (2003).
4. CLAS Collaboration (H. Avakian *et al.*), hep-ex/0301005.
5. J.C. Collins, Nucl. Phys. B **396**, 161 (1993).
6. J.C. Collins, G.A. Ladinsky, hep-ph/9411444.
7. A. Bacchetta, R. Kundu, A. Metz, P.J. Mulders, Phys. Lett. B **506**, 155 (2001).
8. S.J. Brodsky, D.S. Hwang, I. Schmidt, Phys. Lett. B **530**, 99 (2002).
9. J.C. Collins, Phys. Lett. B **536**, 43 (2002).
10. S.J. Brodsky, D.S. Hwang, I. Schmidt, Nucl. Phys. B **642**, 344 (2002).
11. A.V. Belitsky, X. Ji, F. Yuan, Nucl. Phys. B **656**, 165 (2003).
12. D. Boer, P.J. Mulders, F. Pijlman, Nucl. Phys. B **667**, 201 (2003).
13. M. Anselmino, V. Barone, A. Drago, F. Murgia, hep-ph/0209073.
14. P.V. Pobylitsa, hep-ph/0212027.
15. D.W. Sivers, Phys. Rev. D **41**, 83 (1990); **43**, 261 (1991).
16. D. Boer, P.J. Mulders, Phys. Rev. D **57**, 5780 (1998).
17. J. Levelt, P.J. Mulders, Phys. Lett. B **338**, 357 (1994).
18. P.J. Mulders, R.D. Tangerman, Nucl. Phys. B **461**, 197 (1996); **484**, 538 (1997)(E).
19. F. Yuan, hep-ph/0310279.
20. L.P. Gamberg, D.S. Hwang, K.A. Oganessyan, hep-ph/0311221.
21. K.A. Oganessyan, H.R. Avakian, N. Bianchi, A.M. Kotzinian, hep-ph/9808368.
22. A.V. Efremov, K. Goeke, M.V. Polyakov, D. Urbano, Phys. Lett. B **478**, 94 (2000).
23. E. De Sanctis, W.D. Nowak, K.A. Oganessyan, Phys. Lett. B **483**, 69 (2000).
24. B.Q. Ma, I. Schmidt, J.J. Yang, Phys. Rev. D **63**, 037501 (2001).
25. A.V. Efremov, K. Goeke, P. Schweitzer, Phys. Lett. B **522**, 37 (2001); **544**, 389 (2002)(E).
26. A.V. Efremov, K. Goeke, P. Schweitzer, Phys. Rev. D **67**, 114014 (2003).
27. A. Afanasev, C.E. Carlson, hep-ph/0308163.
28. G.P. Lepage, S.J. Brodsky, Phys. Rev. D **22**, 2157 (1980).
29. A. Metz, Phys. Lett. B **549**, 139 (2002).
30. A. Bacchetta, P.J. Mulders, F. Pijlman, hep-ph/0405154.