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Accessing the nucleon transverse structure in inclusive deep inelastic scattering

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ABSTRACT

We revisit the standard analysis of inclusive Deep Inelastic Scattering off nucleons taking into account the fact that on-shell quarks cannot be present in the final state, but they rather decay into hadrons – a process that can be described in terms of suitable "jet" correlators. As a consequence, a spin-flip term associated with the invariant mass of the produced hadrons is generated nonperturbatively and couples to the target's transversity distribution function. In inclusive cross sections, this provides an hitherto neglected and large contribution to the twist-3 part of the g_2 structure function, that can explain the discrepancy between recent calculations and fits of this quantity. It also provides an extension of the Burkhardt–Cottingham sum rule, providing new information on the transversity function, as well as an extension of the Efremov–Teryaev–Leader sum rule, suggesting a novel way to measure the tensor charge of the proton.

1. Introduction

The tensor charge is a fundamental property of the nucleon that is at present poorly constrained but of fundamental importance, not the least because its knowledge can also be used to put constraints on searches for physics beyond the Standard Model [1–3, 61,62]. The tensor charge has been estimated in lattice QCD (see, *e.g.*, [4–8]), but only limited information is available from direct measurements. Its experimental extraction requires first of all flavor-separated measurements of the so-called transversity parton distribution function, denoted by $h_1^q(x)$ (see Ref. [9] for a review and Refs. [10–12] for the most recent extractions). Secondly, one needs to perform flavor-by-flavor integrals of these, that correspond to the contribution of a parton flavor q to the tensor charge.

The transversity distribution is notoriously difficult to access because it is a chiral-odd function and needs to be combined with a spin-flip mechanism to appear in a scattering process [13]. Usually, this spin flip is provided by another nonperturbative distribution or fragmentation function, accessible in Drell–Yan or semi-inclusive Deep Inelastic Scattering (DIS) [14–17]. The only other known way to attain spin-flip terms in Quantum Electro-Dynamics and QCD is taking into account mass corrections. In fact, it is well known that the transversity distribution gives a contribution to the structure function g_2 in inclusive DIS (see, *e.g.*, [18] and references therein), and in particular to the violation of the so-called Wandzura–Wilczek relation for g_2 [19]. However, this contribution is proportional to the current quark mass and can be expected to be negligibly small.

In this paper, we discuss a novel way of accessing the transversity parton distribution function (PDF) and measuring the proton's tensor charge in totally inclusive Deep Inelastic Scattering. We revisit the standard analysis of the DIS handbag diagram, taking into account the fact that on-shell quarks cannot, in fact, be present in the final state, but they rather decay and form (mini)jets of hadrons. This is sufficient to modify the structure of the DIS cut diagram, even if none of those hadrons is detected in the final state. For a proper description of this effect, we include "jet correlators" into the analysis, and pay particular attention to ensuring that our results are gauge invariant.

The jet correlators describe interactions of a perturbative quark with vacuum fields, that break chiral symmetry and generate a nonperturbative mass estimated in the 10-100 MeV range, potentially much larger than the current quark mass for light flavors, as also heuristically advocated in Ref. [20] for a study of transverse target single-spin asymmetries in two-photon exchange processes. Here, we formalize this idea in the context of collinear factorization, and observe that jet correlators introduce a new contribution

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already in one-photon exchange processes, and more precisely to the inclusive g_2 structure function. The new term is proportional to the transversity distribution function multiplied by a new non-perturbative "jet mass", which will be precisely defined below, and has the interesting features that: (a) it violates the Wandzura–Wilczek relation; (b) it extends the Burkhardt–Cottingham sum rule, providing new useful information on behavior of the transversity distribution; (c) it also extends the Efremov–Teryaev–Leader sum rule, providing a novel way to measure the proton's tensor charge. We estimate this new jet-mass-induced contribution based on a recent extraction of the transversity distribution, and show it can indeed be very large.

2. The quark–quark jet correlator

Motivated by mass corrections to inclusive DIS structure functions at large values of the Bjorken invariant x_B, Accardi and Qiu [21] have introduced in the LO handbag diagram a "jet correlator", also called "jet factor" by Collins, Rogers, and Stasto in Ref. [22], that accounts for invariant mass production in the current region and ensures that leading twist calculations in collinear factorization are consistent with the $x_B < 1$ requirement imposed by baryon number conservation. [21]. The jet correlator is depicted in Fig. 1(a) and is defined as

$$\Xi_{ij}(l,n_+) = \int \frac{d^4\eta}{(2\pi)^4}\, e^{il\cdot\eta}\, \langle 0|\,\mathcal{U}^{n_+}_{(+\infty,\eta)}\, \psi_i(\eta)\bar{\psi}_j(0)\,\mathcal{U}^{n_+}_{(0,+\infty)}\,|0\rangle \ , \tag{1}$$

In this definition, l is the quark's four-momentum, Ψ the quark field operator (with quark flavor index omitted for simplicity), and $|0\rangle$ is the nonperturbative vacuum state. Furthermore, the correlator's gauge invariance is explicitly guaranteed the two Wilson line operators $\mathcal{U}^{n_+}$, that run to infinity first along a light-cone plus direction determined by the vector n_+, then along the direction transverse to that vector, see [23] for details. This path choice for the Wilson line is required by QCD factorization theorems, and the vector n_+ is determined by the particular hard process to which the jet correlator contributes. For example, in the case of inclusive DIS discussed in this paper, this is determined by the four momentum transfer q and the proton's momentum p.

The correlator Ξ can be parametrized in terms of jet parton correlation functions A_i and B_i through a Lorentz covariant Dirac decomposition that utilizes the vectors l and n_+,

$$\Xi(l,n_+) = \Lambda A_1(l^2)\,\mathbf{1} + A_2(l^2)\,\not{l} + \frac{\Lambda^2}{l\cdot n_+}\not{n}_+ B_1(l^2) + \frac{i\Lambda}{2l\cdot n_+}[\,\not{l},\not{n}_+]\,B_2(l^2) \ , \tag{2}$$

where Λ is an arbitrary scale, introduced for power counting purposes. In this parametrization, no terms proportional to γ_5 enter because of parity invariance. Time reversal invariance in QCD requires $B_2 = 0$, while B_1 contributes only at twist-4 order and will not be considered further in this paper. We focus, instead, on the role of chiral odd terms in the g_2 structure function up to twist 3. At this order,

$$\Xi(l,n_+) = \Lambda A_1(l^2)\,\mathbf{1} + A_2(l^2)\,\not{l} + \mathcal{O}(\Lambda^2/Q^2) \tag{3}$$

is nothing else than the cut quark propagator; note however, that we consider here the full QCD vacuum rather than the perturbative one (or, in other words, the interacting rather than the free quark fields). The A_1 and A_2 terms can be interpreted in terms of the spectral representation of the cut quark propagator (see, e.g., Sec. 6.3 of [24] and Sec. 2.7.2 of [25]),

$$\Xi(l) = \int d\sigma^2 \big[J_1(\sigma^2)\,\sigma\,\mathbf{1} + J_2(\sigma^2)\,\not{l}\,\big]\,\delta(l^2-\sigma^2) \ , \tag{4}$$

where σ^2 can be interpreted as the invariant mass of the current jet, *i.e.*, of the particles going through the cut in the top blob of Fig. 1(a). The J_i are the spectral functions of the quark propagator, also called "jet functions" in [21], and can be interpreted as current-jet mass distributions. As a consequence of positivity constraints and CPT invariance, these satisfy [24–26]

$$J_2(\sigma^2) \ge J_1(\sigma^2) \ge 0 \quad \text{and} \quad \int d\sigma^2 J_2(\sigma^2) = 1 \ . \tag{5}$$

From a comparison of Eqs. (2) and (4), one can see that

$$A_1(l^2) = \frac{\sqrt{l^2}}{\Lambda} J_1(l^2) \qquad A_2(l^2) = J_2(l^2) \ . \tag{6}$$

When inserting the jet correlator in the handbag diagram for inclusive DIS, the integration over dl^+, or equivalently $dl^2/(2l^-)$, is kinematically coupled to the other integrations, and induces corrections of order $\mathcal{O}(1/Q^2)$ whose effect on the F_2 structure function has been studied in Ref. [21]. In this paper, where we limit our attention to effects of order $\mathcal{O}(1/Q)$, we can neglect k^- compared to q^-. As a consequence, we can extend the integration over dl^2 to infinity, with the consequence that the jet correlator decouples from the parton correlator Φ, and the inclusive structure functions only depend on the integrated jet correlator

$$\Xi(l^-,\boldsymbol{l}_T) \equiv \int \frac{dl^2}{2l^-}\,\Xi(l) = \frac{\Lambda}{2l^-}\,\xi_1\mathbf{1} + \xi_2\frac{\not{n}_-}{2} + \mathcal{O}(l_T/l^-) + \text{higher twists} \ . \tag{7}$$

The neglected $\boldsymbol{l}_T$-dependent and higher twist terms only contribute to $\mathcal{O}(1/Q^2)$ to the inclusive cross section. Note that thanks to Eq. (5) we obtain

$$\xi_1 = \int d\sigma^2 \frac{\sigma}{\Lambda} J_1(\sigma^2) \equiv \frac{M_q}{\Lambda}, \qquad \xi_2 = \int d\sigma^2 J_2(\sigma^2) = 1 \ , \tag{8}$$

where M_q can be interpreted as the average invariant mass produced in the spin-flip fragmentation processes of a quark of flavor q.

It is important to notice that,r while $\xi_2 = 1$ exactly due to CPT invariance (see Sec. 10.7 of Ref. [26]), the jet mass $M_q < \int d\sigma^2\sigma\, J_2(\sigma^2)$ is dynamically determined. From the analytic properties of the spectral functions we expect that $J_2(\sigma^2) = Z\delta(\sigma^2 - m_q^2) + \bar{J}_2(\sigma^2)\theta(\sigma^2 - m_\pi^2)$, with $Z < 1$ and the continuum starting at m_π (the mass of the pion) due to color confinement effects, indicating $M_q = O(\Lambda_{QCD})$. However, in a dynamical confinement scenario, the spectral function J_2 needs not be positive definite [27] and we therefore estimate $M_q \sim 10$–100 MeV. An experimental measurement of M_q is anyway possible, as we discuss in Section 5, and could shed some light on the confinement mechanism. We have also explicitly verified that $M_q > m_q$ in a model where quark fragmentation is simulated by a Yukawa pseudoscalar quark–meson interaction, already utilized, *e.g.*, in Ref. [28].

Although M_q is in general a nonperturbative quantity, it is interesting to notice that on the perturbative vacuum

$$\Xi^{\text{pert}}(l) = (\not{l} + m_q\mathbf{1})\,\delta(l^2 - m_q^2) + \mathcal{O}(\alpha_s) \ , \tag{9}$$

where m_q is the current quark mass; therefore $M_q^{\text{pert}} = m_q$, and one recovers the result of the calculation with the conventional handbag diagram. However, we are here considering nonperturbative effects in the quark propagation, and $M_q \gg m_q$. Therefore, differently from J_2, the J_1 function leaves an imprint on the inclusive DIS cross section even in the asymptotic $Q^2 \to \infty$ regime.