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Twist-3 fragmentation effects for A_{LT} in light hadron production from proton–proton collisions

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ABSTRACT

We compute the contribution from the twist-3 fragmentation function for light hadron production in collisions between transversely and longitudinally polarized protons, i.e., $p^\uparrow \vec{p} \to h\,X$, which can cause a double-spin asymmetry (DSA) A_{LT}. This is a naïve T-even twist-3 observable that we analyze in collinear factorization using both Feynman gauge and lightcone gauge as well as give a general proof of color gauge invariance. So far only twist-3 effects in the transversely polarized proton have been studied for A_{LT} in $p^\uparrow \vec{p} \to h\,X$. However, there are indications that the naïve T-odd transverse single-spin asymmetry (SSA) A_N in $p^\uparrow p \to h\,X$ is dominated not by such distribution effects but rather by a fragmentation mechanism. Therefore, one may expect similarly that the fragmentation contribution is important for A_{LT}. Given possible plans at RHIC to measure this observable, it is timely to provide a calculation of this term.

1. Introduction

Spin asymmetries in various hard processes have brought a new perspective to high-energy perturbative QCD theory and phenomenology. Transverse single-spin asymmetries (SSAs) A_N, which are naïve T-odd observables, were first explored in the mid-1970s in $p^\uparrow p \to \pi\,X$ at Argonne National Lab [1] and $p\,Be \to \Lambda^\uparrow X$ at FermiLab [2]. Both of these measurements led to strikingly large effects that were unexplainable in the naïve parton model [3]. Experiments continued at FermiLab in the 1990s for $p^\uparrow p \to \pi\,X$ [4] and most recently at AGS [5,6] and RHIC [7–14] for $p^\uparrow p \to \{\pi,\,K,\,jet\}\,X$. All of their measurements likewise produced substantial transverse SSAs. On the theoretical side, it was realized in the 1980s by Efremov and Teryaev that if one went beyond the simple parton model and included (collinear twist-3) quark–gluon–quark correlations in the nucleon, then there was the potential to generate these large effects [15]. A systematic approach was then developed by Qiu and Sterman in the 1990s that presented the collinear twist-3 factorization framework [16–18] with the expectation that one would be able describe transverse SSAs within this perturbative approach. Later a solid foundation was given to this formalism in [19,20] that proved the cancelation among gauge-noninvariant terms and led to an expression for the twist-3 cross section in terms of the complete set of the twist-3 quark–gluon–quark correlation functions. Over the last decade, several other analyses, including those for the extention to twist-3 fragmentation functions [21–24] and three-gluon correlation functions [25], furthered the progress of this formalism — see also [26–31] and the references therein.

For many years the main assumption was that these transverse SSAs were due to effects inside the transversely polarized proton, in particular those embodied by the so-called Qiu–Sterman function T_F [16–18,26]. However, a fit of the QS function to A_N data led to a result that was inconsistent with an extraction of the Sivers function $f_{1T}^\perp$ [32] from SIDIS, which has a model-independent relation to T_F [33], and became known as the "sign mismatch" crisis [34]. An attempt to resolve this issue through more flexible parameterizations of the Sivers function proved unsuccessful [35], and, by looking at A_N data on the target transverse SSA in inclusive DIS [36,37], it was argued in fact that the QS function could not be the main cause of A_N [38]. This led to a recent work that examined the impact of fragmentation effects from the outgoing hadron [39] based on the analytical calculation in Ref. [23]. It was determined that this fragmentation term could be the dominant source of A_N in $p^\uparrow p \to \pi\,X$ [39].

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In addition to A_N, there is another twist-3 observable in proton–proton collisions that can give insight into quark–gluon–quark correlations in the incoming protons and/or outgoing hadron. This is the longitudinal-transverse double-spin asymmetry (DSA) A_{LT}, which, unlike A_N, is a naïve T-even process. The classic reaction for which this effect has been analyzed is A_{LT} in inclusive DIS (see [40] for recent experimental results on this observable). This asymmetry has also been studied in the Drell–Yan process involving two incoming polarized hadrons [41–44]; in inclusive lepton production from W-boson decay in proton–proton scattering [45]; for jet production [46] and pion production [47] in lepton–nucleon collisions; and for direct photon production [48], jet/pion production [49], and D-meson production [50] in proton–proton collisions.

Of these works on A_{LT}, only in Ref. [47] for $\vec{\ell}\, p^\uparrow \to \pi\, X$ was the twist-3 fragmentation piece calculated (we will see the structure of that result persists in our computation), whereas the fragmentation term for $p^\uparrow \vec{p} \to \pi\, X$ has never been studied. Like with A_N, there is no reason *a priori* that this piece cannot be important or perhaps dominant in the asymmetry. Given the possible plans by the PHENIX Collaboration at RHIC to measure A_{LT} for pions [51],[1] we feel a calculation of the fragmentation term for this final state is needed at this time. Furthermore, like prior research in the literature [21–24,47,54,55], this work will continue to establish/verify the theoretical techniques used in collinear twist-3 fragmentation calculations. The remainder of the paper is organized as follows: in Section 2 we introduce the twist-3 fragmentation functions relevant for spin-0 hadron production. Next, in Section 3 we discuss the calculation of the polarized cross section formula for A_{LT}. Finally, in Section 4 we summarize our work and give an outlook. A general proof of the color gauge invariance of our result is given in the Appendix.

2. Twist-3 fragmentation functions for spin-0 hadrons

We now define the set of twist-3 fragmentation functions relevant for spin-0 hadron production. The quark–quark matrix element gives two purely real twist-3 functions, which read

$$\frac{1}{N}\sum_X \int \frac{d\lambda}{2\pi} e^{-i\frac{\lambda}{z}} \langle 0|\psi_i^q(0)|P_h; X\rangle\langle P_h; X|\bar{\psi}_j^q(\lambda w)|0\rangle = \frac{M_\chi}{z}(\mathbb{1})_{ij}\,\widehat{e}_1^{\,h/q}(z) + \frac{M_\chi}{2z}(\sigma_{\lambda\alpha} i\gamma_5)_{ij}\epsilon^{\lambda\alpha w P_h}\,\widehat{e}_{\bar{1}}^{\,h/q}(z) + \cdots, \quad (1)$$

where ψ_i is a quark field with spinor index i, and we use the simplified notation $\epsilon^{\lambda\alpha w P_h} \equiv \epsilon^{\lambda\alpha\rho\sigma} w_\rho P_{h\sigma}$ (with $\epsilon_{0123} = +1$). The color indices are summed over and divided by the number of colors $N = 3$. The scale M_χ is used to make the functions dimensionless and is on the order of the nucleon mass.[2] The vector w^μ is light-like ($w^2 = 0$) and satisfies $P_h \cdot w = 1$. We will suppress the gauge-link operators throughout for simplicity.

Next, we introduce the so-called F-type quark–gluon–quark twist-3 fragmentation functions. We can define two independent functions as

$$\frac{1}{N}\sum_X \int \frac{d\lambda}{2\pi} \int \frac{d\mu}{2\pi} e^{-i\frac{\lambda}{z_1}} e^{-i\mu\left(\frac{1}{z}-\frac{1}{z_1}\right)} \times \langle 0|\psi_i^q(0)|P_h; X\rangle\langle P_h; X|\bar{\psi}_j^q(\lambda w) g F^{\alpha w}(\mu w)|0\rangle = \frac{M_\chi}{2z}(\gamma_5 \not{P}_h \gamma_\lambda)_{ij}\epsilon^{\lambda\alpha w P_h}\,\widehat{E}_F^{\,h/q}(z_1, z) + \cdots, \quad (2)$$

$$\frac{1}{N}\sum_X \int \frac{d\lambda}{2\pi} \int \frac{d\mu}{2\pi} e^{-i\frac{\lambda}{z_1}} e^{-i\mu\left(\frac{1}{z}-\frac{1}{z_1}\right)} \times \langle 0|\bar{\psi}_j^q(\lambda w)\psi_i^q(0)|P_h; X\rangle\langle P_h; X| g F^{\alpha w}(\mu w)|0\rangle = \frac{M_\chi}{2z}(\gamma_5 \not{P}_h \gamma_\lambda)_{ij}\epsilon^{\lambda\alpha w P_h}\,\widetilde{E}_F^{\,h/q}(z_1, z) + \cdots, \quad (3)$$

where $F^{\alpha w}(\mu w)$ is the gluon field strength tensor. We note that both $\widehat{E}_F(z_1, z)$ and $\widetilde{E}_F(z_1, z)$ in general are complex functions. The correlator $\widehat{E}_F(z_1, z)$ has support on $1 > z > 0$ and $z_1 > z$, while $\widetilde{E}_F(z_1, z)$ has support on $\frac{1}{z} - \frac{1}{z_1} > 1$, $\frac{1}{z_1} < 0$, and $\frac{1}{z} > 0$ [24,56].

We can consider the so-called D-type twist-3 fragmentation functions $\widehat{E}_D(z_1, z)$ by replacing $gF^{\alpha w}(\mu w)$ in (2) with a covariant derivative $D^\alpha(\mu w) = \partial^\alpha - igA^\alpha(\mu w)$. However, $\widehat{E}_F(z_1, z)$ and $\widehat{E}_D(z_1, z)$ can be related through the identity,

$$\widehat{E}_D^{\,h/q}(z_1, z) = P\left(\frac{1}{1/z_1 - 1/z}\right)\widehat{E}_F^{\,h/q}(z_1, z) + \delta\left(\frac{1}{z_1} - \frac{1}{z}\right)\widetilde{e}^{\,h/q}(z), \quad (4)$$

where $\widetilde{e}(z)$ is another twist-3 fragmentation function that is pure imaginary and defined as

$$\frac{1}{N}\sum_X \int \frac{d\lambda}{2\pi} e^{-i\frac{\lambda}{z}} \times \langle 0|[\infty w, 0]\psi_i^q(0)|P_h; X\rangle\langle P_h; X|\bar{\psi}_j^q(\lambda w)[\lambda w, \infty w]|0\rangle \overleftarrow{\partial}^\alpha = \frac{M_\chi}{2z}(\gamma_5 \not{P}_h \gamma_\lambda)_{ij}\epsilon^{\lambda\alpha w P_h}\,\widetilde{e}^{\,h/q}(z) + \cdots. \quad (5)$$

Note that we have restored the gauge links $[a, b]$ in order to emphasize that $\overleftarrow{\partial}^\alpha$ acts on the λ dependence in both the quark field $\bar{\psi}_j^q(\lambda w)$ and the gauge link $[\lambda w, \infty w]$. The D-type function $\widehat{E}_D(z_1, z)$ has another relation associated with the QCD equation of motion,

$$z\int_z^\infty \frac{dz_1}{z_1^2}\widehat{E}_D^{\,h/q}(z_1, z) = \widehat{e}_1^{\,h/q}(z) + i\,\widehat{e}_{\bar{1}}^{\,h/q}(z). \quad (6)$$

By combining Eqs. (4), (6) we can eliminate the D-type function and obtain

$$z\int_z^\infty \frac{dz_1}{z_1^2} P\left(\frac{1}{1/z_1 - 1/z}\right)\widehat{E}_F^{\,h/q}(z_1, z) + z\,\widetilde{e}^{\,h/q}(z) = \widehat{e}_1^{\,h/q}(z) + i\,\widehat{e}_{\bar{1}}^{\,h/q}(z). \quad (7)$$

The real and imaginary parts of the above relation respectively give

$$z\int_z^\infty \frac{dz_1}{z_1^2} P\left(\frac{1}{1/z_1 - 1/z}\right)\widehat{E}_F^{\,h/q,\Re}(z_1, z) = \widehat{e}_1^{\,h/q}(z), \quad (8)$$

$$z\int_z^\infty \frac{dz_1}{z_1^2} P\left(\frac{1}{1/z_1 - 1/z}\right)\widehat{E}_F^{\,h/q,\Im}(z_1, z) + z\,\widetilde{e}^{\,h/q,\Im}(z) = \widehat{e}_{\bar{1}}^{\,h/q}(z), \quad (9)$$

where $\Re$ ($\Im$) indicates the real (imaginary) part of the function. It was shown that Eq. (9) ensures the gauge invariance of the polarized cross section formula in the case of the transverse SSA in SIDIS [24]. We will show that Eq. (8) plays the same role in the case of the longitudinal-transverse DSA in proton–proton collisions.

[1] We mention that a clear A_{LT} asymmetry has already been seen by the Hall A Collaboration at Jefferson Lab in SIDIS [52] and $\vec{\ell}\, n^\uparrow \to \pi\, X$ [53].

[2] M_χ is the scale of nonperturbative chiral-symmetry breaking (CSB), which is said to be on the order of the nucleon mass (∼1 GeV). We have chosen this scale instead of the light hadron mass M_h since twist-3 functions representing helicity flip effects are due to nonperturbative CSB whereas M_h for the pseudoscalar mesons represents the explicit CSB due to the quark mass.