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# Single Hadronic-Spin Asymmetries in Weak Interaction Processes *

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**Abstract**

We show that measurements of single-spin asymmetries (SSAs) in charged current weak interaction processes such as deep inelastic neutrino scattering on a polarized target and inclusive $W$ production in polarized hadron-hadron collisions discriminate between the two fundamental QCD mechanisms (the Sivers and Collins effects) which have been proposed to explain such time-reversal-odd asymmetries. It has recently been shown that QCD final-state interactions due to gluon exchange between the struck quark and the proton spectators in semi-inclusive deep inelastic lepton scattering will produce non-zero Sivers-type single-spin asymmetries which survive in the Bjorken limit. We show that this QCD final-state interaction produces identical SSAs in charged and neutral current reactions. Furthermore, the contribution of each quark to the SSA from this mechanism is proportional to the contribution of that quark to the polarized baryon's anomalous magnetic moment. In contrast, the Collins effect contribution to SSAs depends on the transversity distribution of quarks in the polarized target. Since the charged current only couples to quarks of one chirality, it cannot sense the transversity distribution of the target, and thus it gives no Collins-type contribution to single-spin correlations.

# 1 Introduction

Spin correlations provide a remarkably sensitive window to hadronic structure and basic mechanisms in QCD. Among the most interesting polarization effects are single-spin azimuthal asymmetries (SSAs) in semi-inclusive deep inelastic scattering, representing the correlation of the spin of the proton target and the virtual photon to hadron production plane: $\vec{S}_p \cdot \vec{q} \times \vec{p}_H$. Such asymmetries are time-reversal odd, but they can arise in QCD through phase differences in different spin amplitudes.

The most common explanation of the pion electroproduction asymmetries in semi-inclusive deep inelastic scattering is that they are related to the transversity distribution of the quarks in the hadron $h_1$ [1, 2, 3] convoluted with the transverse momentum dependent fragmentation function $H_1^\perp$, the Collins function, which gives the distribution for a transversely polarized quark to fragment into an unpolarized hadron with non-zero transverse momentum [4, 5, 6].

Recently, an alternative physical mechanism for the azimuthal asymmetries has been proposed [7, 8, 9]. It was shown that the QCD final-state interactions (gluon exchange) between the struck quark and the proton spectators in semi-inclusive deep inelastic lepton scattering can produce single-spin asymmetries which survive in the Bjorken limit. In this case, the fragmentation of the quark into hadrons is not necessary, and one has a correlation with the production plane of the quark jet itself $\vec{S}_p \cdot \vec{q} \times \vec{p}_q$. This final-state interaction mechanism provides a physical explanation within QCD of single-spin asymmetries. The required matrix element measures the spin-orbit correlation $\vec{S} \cdot \vec{L}$ within the target hadron's wavefunction, the same matrix element which produces the anomalous magnetic moment of the proton, the Pauli form factor, and the generalized parton distribution $E$ which is measured in deeply virtual Compton scattering. Physically, the final-state interaction phase arises as the infrared-finite difference of QCD Coulomb phases for hadron wavefunctions with differing orbital angular momentum.

A related analysis also predicts that the initial-state interactions from gluon exchange between the incoming quark and the target spectator system lead to leading-twist single-spin asymmetries in the Drell-Yan process $H_1 H_2^{\updownarrow} \to \ell^+ \ell^- X$ [8, 10]. These final- and initial-state interactions can be identified as the path-ordered exponentials which are required by gauge invariance and which augment the basic light-front wavefunctions of hadrons [8, 9]. Initial-state interactions also lead to a $\cos 2\phi$ planar correlation in unpolarized Drell-Yan reactions [11, 12].

In this paper we extend the analysis of QCD final-state interactions to SSAs which can be measured in weak interaction processes. For example, consider charged current neutrino semi-inclusive deep inelastic scattering, where a hadron (pion) is measured in the final state. In this case, the transversity distribution cannot contribute to the cross section since the produced quark from the weak interaction of the $W$ boson is always left-handed. This point has also been noted by Miyama [13] and Boer [2]. On the other hand, in the final-state interaction picture the single-spin asymmetry in charged and neutral current weak interactions will be the same as in the electromagnetic case. Thus these weak interaction processes will clearly distinguish the underlying physical mechanisms which produce target single-spin asymmetries.

Measurements of SSAs in semi-inclusive neutrino deep inelastic scattering on a polarized target will be experimentally possible with the advent of a muon storage ring which can provide a high intensity well-focused neutrino beam [14]. However, there are also other weak interaction processes in which similar SSAs should be present and which can be measured experimentally at existing facilities. For example, in this paper we will also make testable predictions for other SSAs in weak interaction reactions such as the processes $pp^{\updownarrow} \to ZX$, $pp^{\updownarrow} \to WX$, which can be measured at RHIC [15], and $e^+e^- \to Z \to \pi\Lambda^{\updownarrow}X$, where the correlation of the $\Lambda$ polarization with the production plane can be measured.

## 2 SSA in Electromagnetic Interactions

The final-state interaction effects can be identified with the gauge link which is present in the gauge-invariant definition of parton distributions [8]. When the light-cone gauge is chosen, a transverse gauge link is required. Thus in any gauge the parton amplitudes need to be augmented by an additional eikonal factor incorporating the final-state interaction and its phase [9, 16]. The net effect is that it is possible to define transverse momentum dependent parton distribution functions which contain the effect of the QCD final-state interactions. We will use this description in this section. It has been shown that the same final-state interactions are responsible for the diffractive component to deep inelastic scattering, and that they play a critical role in nuclear shadowing phenomena [17].

The quark distribution in the proton is described by a correlation matrix:

$$\Phi^{\alpha\beta}(x, \boldsymbol{p}_\perp) = \int \frac{d\xi^- d^2\boldsymbol{\xi}_\perp}{(2\pi)^3} e^{ip\cdot\xi} \left\langle P,S|\overline{\psi}^\beta(0)\psi^\alpha(\xi)|P,S\right\rangle |_{\xi^+=0} \ , \tag{1}$$

where $x = p^+/P^+$. We use the convention $a^\pm = a^0 \pm a^3$, $a \cdot b = \frac{1}{2}(a^+b^- + a^-b^+) - \boldsymbol{a}_\perp \cdot \boldsymbol{b}_\perp$. The correlation matrix $\Phi$ is parameterized in terms of the transverse momentum dependent quark distribution functions [18]:

$$\begin{aligned}\Phi(x, \boldsymbol{p}_\perp) \ = \ & \frac{1}{2}\Bigg[ f_1 \not{n} + f_{1T}^\perp \frac{\epsilon_{\mu\nu\rho\sigma}\gamma^\mu n^\nu p_\perp^\rho S_\perp^\sigma}{M} + g_{1s}\gamma_5 \not{n} \\ & + h_{1T} i\gamma_5\sigma_{\mu\nu}n^\mu S_\perp^\nu + h_{1s}^\perp \frac{i\gamma_5\sigma_{\mu\nu}n^\mu p_\perp^\nu}{M} + h_1^\perp \frac{\sigma_{\mu\nu}p_\perp^\mu n^\nu}{M} \Bigg] \ ,\end{aligned} \tag{2}$$

where the distribution functions have arguments $x$ and $\boldsymbol{p}_\perp$ such as $f_1(x, \boldsymbol{p}_\perp)$, and $n^\mu = (n^+, n^-, \boldsymbol{n}_\perp) = (0, 2, \boldsymbol{0}_\perp)$. The quantity $g_{1s}$ (and similarly $h_{1s}^\perp$ and $G_{1s}$, $H_{1s}^\perp$ in (5) below) is shorthand for

$$g_{1s}(x, \boldsymbol{p}_\perp) = \lambda g_{1L}(x, \boldsymbol{p}_\perp^2) + \frac{\boldsymbol{p}_\perp \cdot \boldsymbol{S}_\perp}{M} g_{1T}(x, \boldsymbol{p}_\perp^2) \ , \tag{3}$$

with $M$ the mass, $\lambda = MS^+/P^+$ the light-cone helicity, and $\boldsymbol{S}_\perp$ the transverse spin of the target hadron. Integrating over $\boldsymbol{p}_\perp$ gives the distribution functions $f_1(x) = \int d^2\boldsymbol{p}_\perp f_1(x, \boldsymbol{p}_\perp)$, $g_1(x) = g_{1L}(x)$, and $h_1(x) = h_{1T}(x) + h_{1T}^{\perp(1)}(x)$, where $h_{1T}^{\perp(1)}(x)$ is the first $\boldsymbol{p}_\perp^2/2M^2$ moment of $h_{1T}^\perp(x, \boldsymbol{p}_\perp)$ [18].

The quark fragmentation is described by the correlation matrix given by

$$\Delta^{\alpha\beta}(z, \boldsymbol{k}_\perp) = \sum_X \int \frac{d\xi^+ d^2\boldsymbol{\xi}_\perp}{2z(2\pi)^3} e^{ik\cdot\xi} \langle 0|\psi^\alpha(\xi)|X; P_h, S_h\rangle \left\langle X; P_h, S_h|\overline{\psi}^\beta(0)|0\right\rangle |_{\xi^-=0} \,, \tag{4}$$

where $z = P_h^-/k^-$ and $\boldsymbol{k}_\perp$ is the quark transverse momentum with respect to the produced hadron and then $-z\boldsymbol{k}_\perp$ is the produced hadron transverse momentum with respect to the fragmenting quark. The correlation matrix $\Delta$ is parameterized in terms of the transverse momentum dependent quark fragmentation functions [18]:

$$\begin{aligned} \Delta(z, \boldsymbol{k}_\perp) &= \frac{1}{2}\Bigg[ D_1 \not{\overline{n}} + D_{1T}^\perp \frac{\epsilon_{\mu\nu\rho\sigma}\gamma^\mu \overline{n}^\nu k_\perp^\rho S_{h\perp}^\sigma}{M_h} + G_{1s}\gamma_5 \not{\overline{n}} \\ &\quad + H_{1T} i\gamma_5 \sigma_{\mu\nu} \overline{n}^\mu S_{h\perp}^\nu + H_{1s}^\perp \frac{i\gamma_5\sigma_{\mu\nu}\overline{n}^\mu k_\perp^\nu}{M_h} + H_1^\perp \frac{\sigma_{\mu\nu} k_\perp^\mu \overline{n}^\nu}{M_h} \Bigg] \,, \end{aligned} \tag{5}$$

where the fragmentation functions have arguments $z$ and $-z\boldsymbol{k}_\perp$ like $D_1(z, -z\boldsymbol{k}_\perp)$, and $\overline{n}^\mu = (\overline{n}^+, \overline{n}^-, \overline{\boldsymbol{n}}_\perp) = (2, 0, \boldsymbol{0}_\perp)$.

The hadronic tensor of the leptoproduction by the electromagnetic interaction in leading order in $1/Q$ is given by [18]

$$\begin{aligned} 2M\mathcal{W}^{\mu\nu}(q, P, P_h) &= \int d^2\boldsymbol{p}_\perp d^2\boldsymbol{k}_\perp \delta^2(\boldsymbol{p}_\perp + \boldsymbol{q}_\perp - \boldsymbol{k}_\perp) \\ &\quad \times \frac{1}{4}\,\mathrm{Tr}\Big[\Phi(x_B, \boldsymbol{p}_\perp)\gamma^\mu \Delta(z_h, \boldsymbol{k}_\perp)\gamma^\nu\Big] \\ &\quad + \Big(\ q \leftrightarrow -q\ ,\ \ \mu \leftrightarrow \nu\ \Big) \,, \end{aligned} \tag{6}$$

where $x_B = Q^2/2P\cdot q$ and $z_h = P\cdot P_h/P\cdot q$. The momentum $\boldsymbol{q}_\perp$ is the transverse momentum of the exchanged photon in the frame where $P$ and $P_h$ do not have transverse momenta.

The single-spin asymmetry (SSA) in semi-inclusive deep inelastic scattering (SIDIS) $ep^\updownarrow \to e'\pi X$, which is given by the correlation $\vec{S}_p \cdot \vec{q} \times \vec{p}_\pi$, is obtained from (6). For the electromagnetic interaction, there are two mechanisms for this SSA: $h_1 H_1^\perp$ and $f_{1T}^\perp D_1$. The former was first studied by Collins [4], and the latter by Sivers [19]; it was shown recently that the Sivers effect does not vanish in the Bjorken limit [7].

We have calculated [7] the single-spin asymmetry in semi-inclusive electroproduction $\gamma^* p^\updownarrow \to HX$ induced by final-state interactions in a model of a spin-$\frac{1}{2}$ proton of mass $M$ with charged spin-$\frac{1}{2}$ and spin-0 constituents of mass $m$ and $\lambda$, respectively, as in the QCD-motivated quark-diquark model of a nucleon. The basic electroproduction reaction is then $\gamma^* p \to q(qq)_0$. In fact, the asymmetry comes from the interference of two amplitudes which have different proton spin but couple to the same final quark spin state, and therefore it involves the interference of tree and one-loop diagrams with a final-state interaction. In this simple model the azimuthal target single-spin asymmetry $A_{UT}^{\sin\phi}$ is given by [7]

$$
\begin{aligned}
A_{UT}^{\sin\phi} &= C_F \alpha_s(\mu^2)\, \frac{\Big(\, \Delta\, M + m\, \Big)\; r_\perp}{\Big[\, \Big(\, \Delta\, M + m\, \Big)^2 \; + \; \boldsymbol{r}_\perp^2 \,\Big]} \\
&\times \Big[\, \boldsymbol{r}_\perp^2 + \Delta(1-\Delta)(-M^2 + \frac{m^2}{\Delta} + \frac{\lambda^2}{1-\Delta})\, \Big] \\
&\times \frac{1}{\boldsymbol{r}_\perp^2}\, \ln \frac{\boldsymbol{r}_\perp^2 + \Delta(1-\Delta)(-M^2 + \frac{m^2}{\Delta} + \frac{\lambda^2}{1-\Delta})}{\Delta(1-\Delta)(-M^2 + \frac{m^2}{\Delta} + \frac{\lambda^2}{1-\Delta})}\; . \qquad (7)
\end{aligned}
$$

Here $r_\perp$ is the magnitude of the transverse momentum of the current quark jet relative to the virtual photon direction, and $\Delta = x_{Bj}$ is the usual Bjorken variable. To obtain (7) from Eq. (21) of [7], we used the correspondence $\frac{|e_1 e_2|}{4\pi} \to C_F \alpha_s(\mu^2)$ and the fact that the sign of the charges $e_1$ and $e_2$ of the quark and diquark are opposite since they constitute a bound state. It has been shown that the result (7) corresponds to the $f_{1T}^{\perp} D_1$ mechanism [8, 9]. The formula (7) is equivalent to $-\,(r_\perp^1\, f_{1T}^{\perp}(\Delta, \boldsymbol{r}_\perp))\,/\,(M\, f_1(\Delta, \boldsymbol{r}_\perp))$.

We show in Fig. 1 the predictions of our model for the asymmetry $A_{UT}^{\sin\phi}$ of the $\vec{S}_p \cdot \vec{q} \times \vec{p}_q$ correlation based on Eq. (7). As representative parameters we take $\alpha_s = 0.3$, $M = 0.94$ GeV for the proton mass, $m = 0.3$ GeV for the fermion constituent and $\lambda = 0.8$ GeV for the spin-0 spectator. The single-spin asymmetry $A_{UT}^{\sin\phi}$ is shown as a function of $\Delta$ and $r_\perp$ (GeV) in Fig. 1(a) . The asymmetry measured at HERMES [20] $A_{UL}^{\sin\phi} = K A_{UT}^{\sin\phi}$ contains a kinematic factor $K = \frac{Q}{\nu}\sqrt{1-y} = \sqrt{\frac{2Mx}{E}}\sqrt{\frac{1-y}{y}}$ because the proton is polarized along the incident electron direction. The resulting prediction for $A_{UL}^{\sin\phi}$ is shown in Fig. 1(b) . Note that $\vec{r} = \vec{p}_q - \vec{q}$ is the momentum of the current quark jet relative to the photon momentum. The asymmetry as a function of the pion momentum $\vec{p}_\pi$ requires a convolution with the quark fragmentation function.

Since the same matrix element controls the Pauli form factor, the contribution of each quark current to the SSA is proportional to the contribution $\kappa_{q/p}$ of that quark to the proton target's anomalous magnetic moment

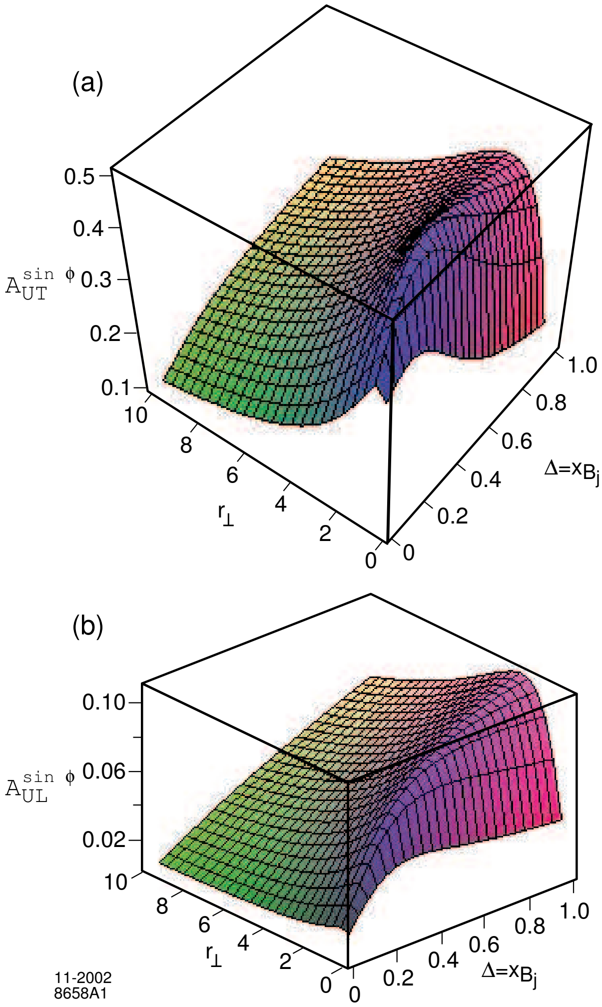


Figure 1: Model predictions for the target single-spin asymmetry $A_{UT}^{\sin\phi}$ for charged and neutral current deep inelastic scattering resulting from gluon exchange in the final state. Here $r_\perp$ is the magnitude of the transverse momentum of the outgoing quark relative to the photon or vector boson direction, and $\Delta = x_{bj}$ is the light-cone momentum fraction of the struck quark. The parameters of the model are given in the text. In (a) the target polarization is transverse to the incident lepton direction. The asymmetry in (b) $A_{UL}^{\sin\phi} = K A_{UT}^{\sin\phi}$ includes a kinematic factor $K = \frac{Q}{\nu}\sqrt{1-y}$ for the case where the target nucleon is polarized along the incident lepton direction. For illustration, we have taken $K = 0.26\sqrt{x}$, corresponding to the kinematics of the HERMES experiment [20] with $E_{lab} = 27.6$ GeV and $y = 0.5$.

$\kappa_p = \sum_q e_q \kappa_{q/p}$ [7].

The SSA of the Drell-Yan processes, such as $\pi p^\updownarrow$ (or $pp^\updownarrow$) $\to \gamma^* X \to \ell^+\ell^- X$, is related to initial-state interactions. The simplest way to get the result is applying crossing symmetry to the SIDIS processes. This was done in Ref. [10] with the result that the SSA in the Drell-Yan process is the same as that obtained in SIDIS, with the appropriate identification of variables, but with the opposite sign [8, 10]. This result corresponds to the $f_{1T}^\perp f_1$ mechanism. The SSA of Drell-Yan processes can also arise from the $h_1 h_1^\perp$ mechanism.

We can also consider the SSA of $e^+e^-$ annihilation processes such as $e^+e^- \to \gamma^* \to \pi\Lambda^\updownarrow X$. The $\Lambda$ reveals its polarization via its decay $\Lambda \to p\pi^-$. The spin of the $\Lambda$ is normal to the decay plane. Thus we can look for a SSA through the T-odd correlation $\epsilon_{\mu\nu\rho\sigma} S_\Lambda^\mu p_\Lambda^\nu q_{\gamma^*}^\rho p_\pi^\sigma$. This is related by crossing to SIDIS on a $\Lambda$ target. The SSA of this process can arise from the $H_1 H_1^\perp$ and $D_{1T}^\perp D_1$ mechanisms.

# 3 SSA in Weak Interactions

## 3.1 Charged Current

Let us consider the SSA in the charged current (CC) weak interaction process $\nu p^\updownarrow \to \ell\pi X$. For the CC weak interaction, the trace in (6) becomes

$$\mathrm{Tr}\Big[\Phi\gamma^\mu P_L \Delta\gamma^\nu P_L\Big] = \mathrm{Tr}\Big[\Phi P_R\gamma^\mu P_L \Delta P_R \gamma^\nu P_L\Big] = \mathrm{Tr}\Big[\Phi_{\mathrm{CC}}\gamma^\mu \Delta_{\mathrm{CC}}\gamma^\nu\Big] \ , \qquad (8)$$

where $P_L = (1-\gamma_5)/2$, $P_R = (1+\gamma_5)/2$, and

$$\Phi_{\mathrm{CC}} \equiv P_L \Phi P_R \ , \qquad \Delta_{\mathrm{CC}} \equiv P_L \Delta P_R \ . \qquad (9)$$

When we use $\Phi(x, \boldsymbol{p}_\perp)$ and $\Delta(z, \boldsymbol{k}_\perp)$ given in (2) and (5), (9) gives

$$\begin{aligned} \Phi_{\mathrm{CC}}(x, \boldsymbol{p}_\perp) &= \frac{1}{2}\, P_L \left( f_1 \not{n} + f_{1T}^\perp \frac{\epsilon_{\mu\nu\rho\sigma}\gamma^\mu n^\nu p_\perp^\rho S_\perp^\sigma}{M} + g_{1s}\gamma_5 \not{n} \right) , \qquad (10) \\ \Delta_{\mathrm{CC}}(z, \boldsymbol{k}_\perp) &= \frac{1}{2}\, P_L \left( D_1 \not{\overline{n}} + D_{1T}^\perp \frac{\epsilon_{\mu\nu\rho\sigma}\gamma^\mu \overline{n}^\nu k_\perp^\rho S_{h\perp}^\sigma}{M_h} + G_{1s}\gamma_5 \not{\overline{n}} \right) . \end{aligned}$$

We see from Eq. (10) that $\Phi_{\mathrm{CC}}$ does not contain the chiral-odd distribution functions which are present in (2), and $\Delta_{\mathrm{CC}}$ does not contain the chiral-odd fragmentation functions present in (5). The charged current only couples to a single quark chirality, and thus it is not sensitive to the transversity distribution. Thus SSAs can only arise in charged current weak interaction SIDIS from the Sivers FSI mechanism $f_{1T}^{\perp} D_1$ in leading order in $1/Q$; in contrast, both the Collins $h_1 H_1^{\perp}$ and Sivers $f_{1T}^{\perp} D_1$ mechanisms contribute to SSAs for the electromagnetic and neutral current (NC) weak interactions.

We can also consider the SSAs of the processes $\pi p^{\updownarrow}$ (or $pp^{\updownarrow}$) $\to WX \to \ell\nu X$. Again these SSAs arise from the $f_{1T}^{\perp} f_1$ mechanism, but not from the $h_1 h_1^{\perp}$ mechanism.

## 3.2 Neutral Current

Let us now consider the SSA in the neutral current weak interaction process $\nu p^{\updownarrow} \to \nu\pi X$. For the NC weak interaction, the interaction vertex of $Z$-f-f is given by $(-ie/\sin\theta_{\rm W}\cos\theta_{\rm W})(c_L P_L + c_R P_R)$ with the weak isospin-dependent coefficients $c_{L,R} = I_{\rm W}^3 - Q\sin^2\theta_{\rm W}$. Explicit values of $c_{L,R}$ are given by $c_L = \frac{1}{2} - \frac{2}{3}\sin^2\theta_{\rm W}$, $c_R = -\frac{2}{3}\sin^2\theta_{\rm W}$ for u, c, t quarks, and $c_L = -\frac{1}{2} + \frac{1}{3}\sin^2\theta_{\rm W}$, $c_R = \frac{1}{3}\sin^2\theta_{\rm W}$ for d, s, b quarks.

The trace in (6) becomes

$$\begin{aligned} & a\ \mathrm{Tr}\Big[\Phi\gamma^\mu(c_L P_L + c_R P_R)\Delta\gamma^\nu(c_L P_L + c_R P_R)\Big] \qquad (11)\\ = \ & a\ \mathrm{Tr}\Big[\Phi(c_L P_R + c_R P_L)\gamma^\mu\Delta\gamma^\nu(c_L P_L + c_R P_R)\Big] \ = \ a\ \mathrm{Tr}\Big[\Phi_{\rm NC}\gamma^\mu\Delta\gamma^\nu\Big] \ , \end{aligned}$$

where $a = 1/\sin^2\theta_{\rm W}\cos^2\theta_{\rm W}$ and

$$\Phi_{\rm NC} \equiv (c_L P_L + c_R P_R)\Phi(c_L P_R + c_R P_L)\ . \qquad (12)$$

When we use $\Phi(x, \boldsymbol{p}_\perp)$ given in (2), (12) gives

$$\begin{aligned} \Phi_{\rm NC}(x, \boldsymbol{p}_\perp) \ = \ & \frac{1}{2}\Bigg[\ (c_L^2 P_L + c_R^2 P_R)\left(f_1 \not{n} + f_{1T}^{\perp}\frac{\epsilon_{\mu\nu\rho\sigma}\gamma^\mu n^\nu p_\perp^\rho S_\perp^\sigma}{M} + g_{1s}\gamma_5\not{n}\right) \qquad (13)\\ & + c_L c_R \left(h_{1T} i\gamma_5\sigma_{\mu\nu}n^\mu S_\perp^\nu + h_{1s}^{\perp}\frac{i\gamma_5\sigma_{\mu\nu}n^\mu p_\perp^\nu}{M} + h_1^{\perp}\frac{\sigma_{\mu\nu}p_\perp^\mu n^\nu}{M}\right)\Bigg]\ . \end{aligned}$$

For the $f_{1T}^{\perp} D_1$ mechanism, we put the former parentheses part of (13) into (11) and then we have $\mathrm{Tr}[\Phi^{\rm f}\gamma^\mu(c_L^2 P_L + c_R^2 P_R)\Delta\gamma^\nu]$, where $\Phi^{\rm f}$ is the first three terms of $\Phi$ in (2). Then, we find that the SSA is given by that of the electromagnetic case with $f_{1T}^{\perp} D_1$ replaced by

$$a\ \frac{c_L^2 + c_R^2}{2}\ f_{1T}^{\perp} D_1\ . \qquad (14)$$

However, $f_1$ is also weighted by the same factor $a\,(c_L^2 + c_R^2)/2$, as we can see in (13). Therefore, the SSA from the final-state interaction mechanism in the NC weak interaction is the same as that in the electromagnetic interaction. This can be confirmed in the simple quark-diquark model.

For the $h_1 H_1^\perp$ mechanism, we put the latter parentheses part of (13) into (11) and find that the SSA is given by that of the electromagnetic case with $(h_1 H_1^\perp)/(f_1 D_1)$ replaced by

$$\frac{2c_L c_R}{c_L^2 + c_R^2}\ \frac{h_1 H_1^\perp}{f_1 D_1}\ . \tag{15}$$

That is, the SSAs are modified by the quark weak isospin-dependent factor $2c_L c_R/(c_L^2+c_R^2)$ in comparison with the electromagnetic case. The same factor appears in the linear $\cos\theta$ forward-backward asymmetry in the $e^+e^- \to Z \to q\bar{q}$ reaction.

The SSA of the Drell-Yan processes at the $Z^0$, such as $\pi p^\updownarrow$ (or $pp^\updownarrow$) $\to ZX \to \ell^+\ell^- X$, can arise from the $h_1 h_1^\perp$ and $f_{1T}^\perp f_1$ mechanisms. We can also consider the SSA of the $e^+e^-$ annihilation processes such as $e^+e^- \to Z \to \pi\Lambda^\updownarrow X$, which can arise from the $H_1 H_1^\perp$ and $D_{1T}^\perp D_1$ mechanisms [21]. The SSAs of these processes have the same situation as those of the above SIDIS case. The initial/final-state interaction mechanisms have the same formulas as the electromagnetic case, whereas the Collins mechanisms are weighted by the quark weak isospin-dependent factor $2c_L c_R/(c_L^2 + c_R^2)$ present in (15).

# 4 Conclusions

We have shown that target single-spin asymmetries in semi-inclusive deep inelastic scattering $\nu p^\updownarrow \to \ell\pi X$ from the charged current weak interaction can arise only from the $f_{1T}^\perp D_1$ mechanism in leading order in $1/Q$, whereas the SSAs in $ep^\updownarrow \to e'\pi X$ of electromagnetic interaction from the $h_1 H_1^\perp$ and $f_{1T}^\perp D_1$ mechanisms. Thus charged current weak interaction processes clearly distinguish the underlying physical mechanisms responsible for target single-spin asymmetries.

We have also analyzed the SSAs in semi-inclusive reactions such as $\nu p^\updownarrow \to \nu\pi X$ of the neutral current weak interaction and have found that the SSA from the Collins mechanism is dependent on the quark weak isospin. The phase from the QCD final-state interaction mechanism only depends on color; it is the same as that in the electromagnetic case and does not depend on the quark weak isospin. Furthermore, the contribution of each quark current to the SSA from this mechanism is proportional to the contribution of that quark to the polarized baryon's anomalous magnetic moment.

It is also important to study the SSAs in weak interaction reactions such as $pp^\updownarrow \to ZX$, $pp^\updownarrow \to WX$, which can be measured at RHIC, and $e^+e^- \to Z \to \pi\Lambda^\updownarrow X$ which can be measured in $e^+e^-$ colliders. In each case the contribution from each quark to the SSAs from initial/final-state interactions is identical to that of the corresponding electromagnetic process.

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