

Collins effect in single spin asymmetries of the $p^\uparrow p \rightarrow \pi X$ process

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Received: 21 June 2004 / Revised version: 26 November 2004 /

Published online: 15 February 2005 – © Springer-Verlag / Società Italiana di Fisica 2005

Abstract. We investigate the Collins effect in single spin asymmetries (SSAs) of the $p^\uparrow p \rightarrow \pi X$ process by taking into account the transverse momentum dependence of the microscopic sub-process cross sections, with the transverse momentum in the Collins function integrated over. We find that the asymmetries due to the Collins effect can only explain the available data at best qualitatively, by using our choices of quark distributions in the quark–diquark model and a pQCD-based analysis, together with several options of the Collins function. Our results indicate the necessity to take into account contributions from other effects such as the Sivers effect or twist-3 contributions.

PACS. 13.75.Cs, 13.85.Ni, 13.87.Fh, 13.88.+e

Single spin asymmetries in hadronic reactions have become an active research topic recently since they may help us to uncover the transverse quark structure of the nucleon. Experimentally, there are data on meson electroproduction asymmetries in semi-inclusive deep inelastic scattering [1–3], whose most common explanation is to relate them to the transversity distribution of the quarks in the hadron [4, 5], convoluted with a transverse momentum dependent fragmentation function [6–10], i.e., the Collins function [6, 11]. The Collins function, which gives the distributions for a transversely polarized quark to fragment into an unpolarized hadron with non-zero transverse momentum, has aroused great interest recently since a chiral-odd structure function can be accessible together with another chiral-odd distribution/fragmentation function.

Large single spin asymmetries (SSAs) in $p^\uparrow p \rightarrow \pi X$ have been also observed by the E704 Group at Fermilab [12]. Anselmino, Boglione and Murgia [13] tried to reproduce the experimental data using a parametrization of the Collins fragmentation function, with several assumptions and in a generalized factorization scheme at the parton level. Notice that the Collins function obtained in [13] is different from the original Collins parametrization [6]. Other mechanisms such as the Sivers effect [14–17], and twist-3 contributions [18–20], have been also proposed. Very recently, Bourrely and Soffer [21] pointed out that the SSA observed several years ago at FNAL by the experiment

E704 [12] and the recent result observed at BNL-RHIC by STAR [22] are due to different phenomena. It is thus important to check whether a calculation with more detailed microscopic dynamics taken into account, and with also proper constraints on the Collins function, can indeed explain the single spin asymmetries in $p^\uparrow p \rightarrow \pi X$. We will show in the following that the Collins effect in single spin asymmetries of the $p^\uparrow p \rightarrow \pi X$ process can only explain the available experimental data at best qualitatively, but we have not been able to explain the magnitude of the data at large x_F by using several options of the Collins function and with our choices of the quark distributions in the quark–diquark model and a pQCD-based analysis. In our analysis, we work within the same simplified planar configuration as [13], using several options of Collins function with a Gaussian-type transverse momentum dependence, together with our choices of quark distributions. Notice also that measurements of SSAs in weak interaction processes can distinguish between different QCD mechanisms [16].

We first focus on the formalism of the description of single spin asymmetries at the parton level. For the inclusive production of a hadron C from the hadron A and hadron B collision process

$$A + B \rightarrow C + X, \quad (1)$$

the Mandelstam variables s , t , and u are written as

$$s = (P_A + P_B)^2 = M_A^2 + M_B^2 + 2P_A \cdot P_B, \quad (2)$$

$$t = (P_A - P_C)^2 = M_A^2 + M_C^2 - 2P_A \cdot P_C, \quad (3)$$

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$$u = (P_B - P_C)^2 = M_B^2 + M_C^2 - 2P_B \cdot P_C. \quad (4)$$

The experimental cross sections are usually expressed in terms of the experimental variables $x_F = 2P_L/\sqrt{s}$ and P_T at a given s , where $P_L = x_F\sqrt{s}/2$ and P_T are the longitudinal and transverse momenta of the produced hadron C respectively, with energy $E_C = \sqrt{M_C^2 + P_C^2} = \sqrt{M_C^2 + P_L^2 + P_T^2} = \sqrt{M_C^2 + P_T^2 + x_F^2 s/4}$ in the center of mass frame of the collision process, and s is the squared center of mass energy. The momentum and the energy of hadron A in the center of mass frame are

$$P_{AL} = |\mathbf{P}_A| = \sqrt{\frac{(s - M_A^2 - M_B^2)^2 - 4M_A^2 M_B^2}{4s}}, \quad (5)$$

$$E_A = \sqrt{M_A^2 + P_{AL}^2}, \quad (6)$$

and the momentum and the energy of hadron B are

$$\mathbf{P}_B = -\mathbf{P}_A, \quad (7)$$

$$E_B = \sqrt{M_B^2 + P_{AL}^2}. \quad (8)$$

The Mandelstam variables t and u can be expressed as

$$t = (P_A - P_C)^2 = M_A^2 + M_C^2 - 2E_A E_C + 2P_{AL} P_L, \quad (9)$$

$$u = (P_B - P_C)^2 = M_B^2 + M_C^2 - 2E_B E_C - 2P_{AL} P_L. \quad (10)$$

If the transverse momentum of the produced hadron relative to the fragmenting quark is negligible and the factorization theorem applies, the cross section can be written in terms of the subprocess at the parton level as

$$\begin{aligned} &= \sum_{a,b,c,d} \int \frac{dx_a dx_b dz}{\pi z^2} f_{a/A}(x_a) f_{b/B}(x_b) \hat{s} \delta(\hat{s} + \hat{t} + \hat{u}) \\ &\times \frac{d\hat{\sigma}^{ab \rightarrow cd}}{d\hat{t}}(x_a, x_b, z) D_{C/c}(z). \end{aligned} \quad (11)$$

However, there is no such a theorem if one considers transverse parton momenta as done sequently by taking the Collins effect into account. The Collins effect consists in taking into account the transverse momentum $\mathbf{k}_{c\perp}$ of the produced hadron C relative to the fragmenting parton c with momentum \mathbf{p}_c , so we have

$$\mathbf{P}_C = z\mathbf{p}_c + \mathbf{k}_{c\perp}, \quad (12)$$

where z is the momentum fraction of the produced hadron relative to the fragmenting parton. The cross section is then expressed as

$$\begin{aligned} d\sigma &= \frac{E_C d^3\sigma^{AB \rightarrow CX}}{d^3\mathbf{P}_C} \\ &= \sum_{a,b,c,d} \int d^2\mathbf{k}_{c\perp} \frac{dx_a dx_b dz}{\pi z^2} f_{a/A}(x_a) f_{b/B}(x_b) \\ &\times \hat{s} \delta(\hat{s} + \hat{t} + \hat{u}) \frac{d\hat{\sigma}^{ab \rightarrow cd}}{d\hat{t}}(x_a, x_b, z, \mathbf{k}_{c\perp}) D_{C/c}(z, \mathbf{k}_{c\perp}), \end{aligned} \quad (13)$$

where $D_{C/c}(z, \mathbf{k}_{c\perp})$ describes the fragmentation of hadron C with longitudinal momentum fraction z and transverse momentum $\mathbf{k}_{c\perp}$ relative to the fragmenting parton c . The denotations a , b , and c should include both the flavors and respective polarizations of the involved partons, if any or some of the incident hadrons A , B , and the produced hadron C are polarized. The cross section of the subprocess $a + b \rightarrow c + d$,

$$d\hat{\sigma} = \frac{d\hat{\sigma}^{ab \rightarrow cd}}{d\hat{t}}(x_a, x_b, z, \mathbf{k}_{c\perp}), \quad (14)$$

should be written in terms of the Mandelstam variables \hat{s} , \hat{t} , and \hat{u}

$$\hat{s} = 2p_a \cdot p_b = x_a x_b s, \quad (15)$$

$$\hat{t} = -2p_a \cdot p_c = \frac{x_a}{z} t \Phi_t(\pm k_{c\perp}), \quad (16)$$

$$\hat{u} = -2p_b \cdot p_c = \frac{x_b}{z} u \Phi_u(\pm k_{c\perp}), \quad (17)$$

where the functions $\Phi_t(\pm k_{c\perp})$ and $\Phi_u(\pm k_{c\perp})$ are given by [23]

$$\Phi_t(\pm k_{c\perp}) = g(k_{c\perp}) \quad (18)$$

$$\times \left\{ 1 \mp 2k_{c\perp} \frac{\sqrt{stu}}{t(t+u)} - [1 - g(k_{c\perp})] \frac{t-u}{2t} \right\},$$

$$\Phi_u(\pm k_{c\perp}) = g(k_{c\perp}) \quad (19)$$

$$\times \left\{ 1 \pm 2k_{c\perp} \frac{\sqrt{stu}}{u(t+u)} + [1 - g(k_{c\perp})] \frac{t-u}{2u} \right\},$$

with $g(k_{c\perp}) = \sqrt{1 - k_{c\perp}^2/P_C^2}$, and $\pm k_{c\perp}$ referring respectively to the configuration in which $\mathbf{k}_{c\perp}$ points to the left or to the right of \mathbf{p}_c . The kinematical effect from the transverse momentum $\mathbf{k}_{c\perp}$ is explicitly taken into account in $\Phi_t(\pm k_{c\perp})$ and $\Phi_u(\pm k_{c\perp})$, which become 1 for $k_{c\perp} = 0$. The four variables x_a , x_b , z , and $k_{c\perp}$ are not independent, and by exploiting the $\delta(\hat{s} + \hat{t} + \hat{u})$ -function we get

$$x_a x_b s + \frac{1}{z} [x_a t \Phi_t(\pm k_{c\perp}) + x_b u \Phi_u(\pm k_{c\perp})] = 0, \quad (20)$$

so we have

$$z = -\frac{x_a t \Phi_t(\pm k_{c\perp}) + x_b u \Phi_u(\pm k_{c\perp})}{x_a x_b s}. \quad (21)$$

After integrating the δ -function, we get the cross section

$$\begin{aligned} d\sigma &= \frac{E_C d^3\sigma^{AB \rightarrow CX}}{d^3\mathbf{P}_C} \\ &= \sum_{a,b,c,d} \int d^2\mathbf{k}_{c\perp} \\ &\times \left\{ \int_{x_a^{\min}}^1 dx_a \int_{x_b^{\min}}^1 dx_b f_{a/A}(x_a) f_{b/B}(x_b) \right. \\ &\times \left. \frac{1}{\pi z} \frac{d\hat{\sigma}^{ab \rightarrow cd}}{d\hat{t}}(x_a, x_b, z, \mathbf{k}_{c\perp}) \right\} D_{C/c}(z, \mathbf{k}_{c\perp}), \end{aligned} \quad (22)$$

where

$$x_a^{\min} = -\frac{x_b u \Phi_u(\pm k_{c\perp})}{x_b s + t \Phi_t(\pm k_{c\perp})}, \quad x_b^{\min} = -\frac{t \Phi_t(\pm k_{c\perp})}{s + u \Phi_u(\pm k_{c\perp})}. \quad (23)$$

Because of the $k_{c\perp}$ dependence of x_a^{\min} and x_b^{\min} , the integration over $\mathbf{k}_{c\perp}$ should be performed after the integrations over x_a and x_b , and $\mathbf{k}_{c\perp}$ should meet the constraint $k_{c\perp} \leq P_C$. Once we know the quark distributions $f_{a/A}(x)$ and $f_{b/B}(x)$, the cross sections $d\hat{\sigma} = \frac{d\hat{\sigma}^{ab \rightarrow cd}}{d\hat{t}}(\hat{s}, \hat{t}, \hat{u})$, and the fragmentation functions $D_{C/c}(z, \mathbf{k}_{c\perp})$, we can calculate the cross section (22) explicitly.

Now we apply the above formulas to the single spin asymmetries

$$A(x_F, P_T) = \frac{d\sigma^\uparrow - d\sigma^\downarrow}{d\sigma^\uparrow + d\sigma^\downarrow} \quad (24)$$

of $p^\uparrow p \rightarrow \pi X$ process, measured by the E704 Group at Fermilab [12]. After a calculation of the cross section in terms of the explicit polarization dependent ingredients, we find that the above asymmetries can be written as

$$A(x_F, P_T) = \frac{\int d^2\mathbf{k}_{c\perp} dx_a dx_b \delta q_a(x_a) q_b(x_b) \Delta\hat{\sigma}(\hat{s}, \hat{t}, \hat{u}) \Delta D^N(z, \mathbf{k}_{c\perp})}{\int d^2\mathbf{k}_{c\perp} dx_a dx_b q_a(x_a) q_b(x_b) \hat{\sigma}(\hat{s}, \hat{t}, \hat{u}) D(z, \mathbf{k}_{c\perp})}, \quad (25)$$

where $\delta q_a(x_a)$ is the quark transversity distribution, $q_b(x_b)$ is the usual quark distribution, $d\hat{\sigma} = d\hat{\sigma}^{\uparrow\uparrow} + d\hat{\sigma}^{\uparrow\downarrow}$, $\Delta d\hat{\sigma} = d\hat{\sigma}^{\uparrow\uparrow} - d\hat{\sigma}^{\uparrow\downarrow}$, $D(z, \mathbf{k}_{c\perp}) = [D_{\pi/c^\uparrow}(z, \mathbf{k}_{c\perp}) + D_{\pi/c^\downarrow}(z, \mathbf{k}_{c\perp})]/2$, and $\Delta D^N(z, \mathbf{k}_{c\perp}) = D_{\pi/c^\uparrow}(z, \mathbf{k}_{c\perp}) - D_{\pi/c^\downarrow}(z, \mathbf{k}_{c\perp})$ is the Collins function.

For the quark distributions, we use those in a quark–diquark model [24] and a pQCD-based analysis [25], explicitly taking into account the sea contributions based on the GRV parametrization of the parton distribution functions in [26]. The unpolarized and transversely polarized cross sections of the subprocess at the parton level, $\hat{\sigma}(\hat{s}, \hat{t}, \hat{u})$ and $\Delta\hat{\sigma}(\hat{s}, \hat{t}, \hat{u})$, can be found in [27, 28]. The detailed information on the quark transversity distributions in the quark–diquark model and the pQCD-based analysis can be found in [9]. We need to point it out here that these transversity distributions do not come from model calculations, but from relations that connect the transversity distributions with parametrized unpolarized quark distributions, so that the calculation can be performed at the same scale at the experiment. It is important to use a same set of both unpolarized and polarized quark distributions, otherwise it is not reasonable to compare the denominator with the numerator. The calculated ratio would be unreasonable if the transversity distributions were taken from a model and then use parton distributions from a parametrization in order to perform the calculation. This aspect has been carefully considered in our calculation. In a strict sense, the unpolarized quark distributions and transversity distributions evolve differently, so that we should use the relations at a specific initial scale such as $Q^2 \approx 2 \text{ GeV}^2$, and then consider the evolution of the numerator and denominator

separately. However, the effects of evolution are presumably smaller than other uncertainties in the approach, such as the neglect of unfavored fragmentation, so we neglect the explicit difference in evolution and take the energy scale the same in both the numerator and denominator. Besides, the quark–diquark model [24] has been successful in providing good descriptions of the quark helicity distributions from experiments [29], as well as data for nucleon form factors [30], so the extension of applying it to the transversity distributions is reasonable.

For the pion fragmentation functions, we take only the favored fragmentation into account, following [13], and adopt the Kretzer–Leader–Christova parametrization [31] of $D(z)$:

$$D(z) = 0.689 z^{-1.039} (1-z)^{1.241}. \quad (26)$$

The scale for $D(z)$ is $\langle Q^2 \rangle = 2.5 \text{ GeV}^2$. In our case $s = 400 \text{ GeV}^2$, but what matters for the evolutions is P_L^2 . Now $P_L = x_F \sqrt{s}/2$, which is about 5 GeV, and then P_L^2 is about 25 GeV^2 . We neglect the difference between different fragmentation functions in the evolution, as the effect due to the evolution, which is at most logarithmic, can be reasonably neglected in predicting the asymmetries when only ratios between different fragmentation functions are relevant. To take into account the $\mathbf{k}_{c\perp}$ dependence, we parameterize the fragmentation function $D(z, \mathbf{k}_{c\perp})$ as [9]

$$D(z, k_{c\perp}) = \frac{R^2}{\pi} D(z) \exp(-R^2 k_{c\perp}^2), \quad (27)$$

where $R = \frac{1}{\langle k_{c\perp}^2 \rangle^{1/2}}$ with $\langle k_{c\perp}^2 \rangle^{1/2} = 0.44 \text{ GeV}$. The Collins function $\Delta D(z, k_{c\perp})$ is $\mathbf{k}_{c\perp}$ odd, and it gives a null contribution in case $k_{c\perp} = 0$, so we adopt the Collins parametrization [6] for $\Delta D(z, k_{c\perp})$ as follows:

$$\frac{\Delta D(z, k_{c\perp})}{2D(z, k_{c\perp})} = \frac{2M_C(k_{c\perp}/z)}{(M_C^2 + k_{c\perp}^2/z^2)}, \quad (28)$$

where M_C is a typical hadronic scale around $0.3 \rightarrow 1 \text{ GeV}$ and we take its value as 0.7 GeV in our calculation. The above Collins parametrization fulfills the bound [13]:

$$\frac{|\Delta D(z, k_{c\perp})|}{2D(z, k_{c\perp})} \leq 1. \quad (29)$$

In the above formulas, all quantities needed in the calculation are clearly given except that there are still large uncertainties concerning the Collins function $\Delta D(z, k_{c\perp})$. Therefore we consider several options for the Collins function.

Option 1: The Collins parametrization (28).

Option 2: The Collins parametrization with an additional z -dependent factor

$$\frac{\Delta D(z, k_{c\perp})}{2D(z, k_{c\perp})} = \frac{2M_C(k_{c\perp}/z)}{(M_C^2 + k_{c\perp}^2/z^2)} P(z), \quad (30)$$

where $P(z) = cz^\alpha (1-z)^\beta$ with $\alpha = 15$, $\beta = 1.5$, and $c = 306.18$. This option is chosen to illustrate a case which

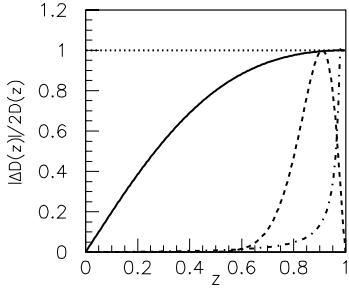


Fig. 1. The z dependence of the ratio $\frac{\Delta D(z, k_{c\perp})}{2D(z, k_{c\perp})}$ at a given value of $k_{c\perp} = k_{c\perp}^0 = M_C$ for the four options of the Collins function. The solid, dashed, dash-dotted, and dotted curves are corresponding to the results for the four options of Collins function respectively

is of similar shape and not far from that of the Anselmino–Boglionne–Murgia (ABM) parametrization (Option 3), but which can reproduce the asymmetries with magnitude around twice that of Option 3.

Option 3: The Anselmino–Boglionne–Murgia (ABM) parametrization (there is some trivial difference in details) [13]:

$$\frac{\Delta D(z, k_{c\perp})}{2D(z, k_{c\perp})} = \begin{cases} 0.007409z(1-z)^{-1.3} & z < 0.97742; \\ 1 & z \geq 0.97742, \end{cases} \quad (31)$$

at an average value of $\langle k_{c\perp}^2 \rangle^{1/2} = k_{c\perp}^0$.

Option 4: The upper limit bound

$$\frac{\Delta D(z, k_{c\perp})}{2D(z, k_{c\perp})} = 1. \quad (32)$$

The z dependence of the ratio $\frac{\Delta D(z, k_{c\perp})}{2D(z, k_{c\perp})}$ at a given $k_{c\perp} = k_{c\perp}^0 = M_C$ for the four options can be found in Fig. 1, from where we can find that the shapes of the Collins function are quite different for the four options.

Thus we have all of the ingredients for the calculation of the single spin asymmetries $A(x_F, P_T)$ in (25). The results shown in Fig. 2 are calculated at $P_T = 1.5$ GeV with the quark transversity distributions from both the quark–diquark model and the pQCD-based analysis. We notice that both models give large $A(x_F, P_T)$ at small x_F , with magnitude compatible with or larger than the data, but at large x_F the magnitude is below the data. The quark–diquark model can reproduce the trend of $A(x_F, P_T)$ for π^- qualitatively, but the pQCD-based analysis produces a result with opposite sign. The reason for the discrepancy is that the valence d -quark transversity distribution is positive at large x in the pQCD-based analysis [9]. The situation for the quark–diquark model is better, but there is still a discrepancy with the magnitude of the data. In order to reduce the discrepancy between the calculated result and the data, especially for the π^- data, we introduce an additional case (which is not a realistic case, because it breaks Soffer’s inequality [32], so this is just as an illustration) in the quark–diquark model with the valence d -quark transversity distributions enhanced by a factor of 3, i.e.,

$$\delta d_v(x) = -d_v(x)\hat{W}_v(x), \quad (33)$$

where $d_v(x)$ is the unpolarized quark distribution and $\hat{W}_v(x)$ is a Melosh–Wigner rotation factor to reflect the relativistic effect from quark transverse motions [33]. We find that the magnitude of the calculated results is more compatible with the data than both the quark–diquark model and the pQCD-based analysis, as shown in Fig. 2. This might suggest that the d -quark transversity distributions is more negatively polarized than predicted in the

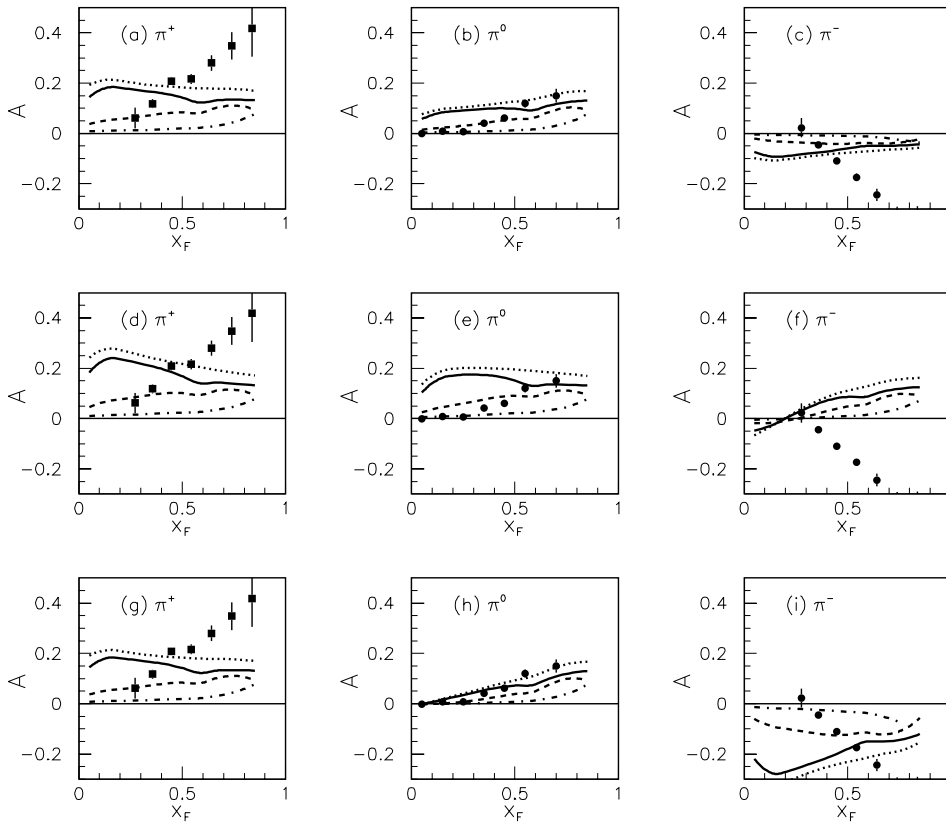


Fig. 2. The single spin asymmetries of $p^\uparrow p \rightarrow \pi X$ process. The calculated results in the quark–diquark model: **a**, **b**, and **c**, in the pQCD-based analysis: **d**, **e**, and **f**, and in the quark–diquark model with more negatively polarized valence d -quark: **g**, **h**, and **i**. The solid, dashed, dash-dotted, and dotted curves are corresponding to the results for the four options of Collins function respectively. The experimental data are given by the E704 Group [12]

quark–diquark model, as supported by sum-rule based arguments [34], if one only considers the Collins effect. From the results of the upper limit bound of the Collins function, we notice that any improvement in the parametrization of the Collins function cannot improve the fit to the data at large x_F . Inclusion of unfavored fragmentation [9] will lead the calculated results to go in the opposite direction from the data. The ABM parametrization reproduces the shape of the data, but underestimates its magnitude by a factor of around 4–6. We have neglected the gluon transverse polarization in our calculation, as the gluon transversity in a spin 1/2 hadron is strictly zero due to helicity conservation. Thus the description of the data is not obtained in our calculation by taking into account the detailed microscopic cross sections with the transverse momentum in the Collins function integrated over, together with our choices of quark distributions and Collins function.

In summary, we checked the Collins effect in single spin asymmetries of the $p^\uparrow p \rightarrow \pi X$ process by taking into account the transverse momentum dependence of the microscopic sub-process cross sections at the parton level, with the transverse momentum in the Collins function integrated over. We introduced several options for the Collins effect, and found that the single spin asymmetries of $p^\uparrow p \rightarrow \pi X$ process due to the Collins effect can only explain the available data at best qualitatively in some specific situations, by using the quark distributions in the quark–diquark model and in a pQCD-based analysis. The results suggest the necessity of taking into account contributions from other effects such as the Sivers effect and twist-3 contributions. It might be also possible that some unexpected novel behaviors of the quark distributions and Collins function need to be introduced. With our present calculation results on the asymmetries and the calculation done in [21] for the cross sections themselves, we have to conclude that one needs to introduce other mechanisms in order to understand the single spin asymmetries for π inclusive production in pp collisions. A similar conclusion with a more detailed and complete analysis has also been drawn in a recent work [35].

Acknowledgements. We acknowledge the helpful comments and suggestions from Andreas Schäfer and Mauro Anselmino. This work is partially supported by National Natural Science Foundation of China, by the Key Grant Project of Chinese Ministry of Education (No. 305001), by Fondecyt (Chile) grant 1030355, by Alexander von Humboldt-Stiftung (J.J. Yang), and by the Italian Ministry of Education, University and Research (MIUR).

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