3 December 1998

PHYSICS LETTERS B

Physics Letters B 442 (1998) 470–478

Single spin asymmetries in $p^{\uparrow} p$ and $\bar{p}^{\uparrow} p$ inclusive processes

Mauro Anselmino [a], Francesco Murgia [b]

[a] *Dipartimento di Fisica Teorica, Università di Torino, and Istituto Nazionale di Fisica Nucleare, Sezione di Torino, via P. Giuria 1, I-10125 Torino, Italy*

[b] *Istituto Nazionale di Fisica Nucleare, Sezione di Cagliari, and Dipartimento di Fisica, Università di Cagliari, C.P. 170, I-09042 Monserrato (CA), Italy*

Received 8 September 1998
Editor: R. Gatto

Abstract

We consider several single spin asymmetries in inclusive $p^{\uparrow} p$ and $\bar{p}^{\uparrow} p$ processes as higher twist QCD contributions, taking into account spin and intrinsic $\boldsymbol{k}_{\perp}$ effects in the quark distribution functions. This approach has been previously applied to the description of the single spin asymmetries observed in $p^{\uparrow} p \to \pi X$ reactions and all its parameters fixed: we give here predictions for new processes, which agree with experiments for which data are available, and suggest further possible measurements. © 1998 Elsevier Science B.V. All rights reserved.

PACS: 13.88.+e; 13.85.-t; 12.38.-t

1. Introduction

In the last years a copious theoretical and experimental activity has been devoted to the study of single spin asymmetries (SSA) in inclusive particle production at high energy and moderately large p_T. On theoretical grounds, at leading twist in massless perturbative QCD, SSA at high energies and p_T are expected to be negligible. However, the presently available experimental results seem to show that this is not the case: at least in some kinematical regions (namely, at large x_F and at moderately large p_T) sizeable SSA (up to the order of 30%-40%) have been measured [1,2]. These experimental results have prompted a renewed theoretical activity, focused at the introduction in perturbative QCD schemes of higher twist contributions, previously neglected, which could play a crucial role in this context.

Several single spin effects, relevant at different steps of the inclusive hadronic production, have been suggested in the literature. In Ref. s [3–5], for example, quark transverse momentum effects have been taken into account in the structure of the initial, transversely polarized nucleon. A similar mechanism has been proposed for the fragmentation process of a polarized quark into the final observed particles [6,7]. Quark-gluon correlation functions and gluonic pole contributions have also been considered as possible origin of SSA in various processes [8–11]. Quark orbital angular momentum and a simple (non QCD) elementary dynamics is used in Ref. [12].

In Ref. [5] we start from the QCD formalism for polarized hard processes [13], with the inclusion of the intrinsic transverse momentum of quarks inside the polarized hadrons: this allows to introduce a non-diagonal (in the helicity basis) distribution func-

PII: S0370-2693(98)01267-2

tion for these quarks, denoted by $\Delta_N f$. Such distribution would be forbidden by time reversal invariance for free quarks [6], but is allowed by initial state interactions between the incoming hadrons, or by quark-gluon correlations [8–11] or in chiral models [14]. The only assumption is that the QCD factorization theorem still holds in such cases. Our formalism, whose results agree with those of Sivers [3,4], allows then in principle to obtain sizeable values for the SSA.

A simple phenomenological parametrization of the new distributions was introduced, and all parameters were fixed by fitting the data on SSA in inclusive pion production in proton-proton collisions, $p^{\uparrow} p \to \pi X$ [1]. It was shown that a good description of the experimental data is indeed possible and the resulting features of the new distribution turn out to be physically plausible and well justified.

In this paper, equipped with a definite expression for $\Delta_N f$ obtained by fitting the data on pion production, we further apply our formalism to other processes, both obtaining genuine predictions which can be compared with data already existing and suggesting possible future measurements and tests. The systematic study of several SSA allows to evaluate the consistency and relevance of our approach, and to isolate its contribution to SSA from other possible sources of spin effects.

A preliminary, partial account of this analysis was presented in [15].

The plan of the paper is the following: in Section II we recall and summarize our formalism and repeat the fitting procedure which allows to fix all parameters. Section III, which contains the main results of the paper, will be devoted to the application of the model to several processes; numerical results will be presented and discussed in details. Finally, in Section IV we give some comments and conclusions.

2. The model

In general we shall consider the high energy, high p_T, inclusive process $h_1^{\uparrow(\downarrow)} h_2 \to h_3 X$, where h_1, h_2, and h_3 are hadrons; the apex $\uparrow(\downarrow)$ means that hadron h_1 is transversely polarized with respect to the scattering plane, in the same (opposite) direction as $\boldsymbol{p}_{h_1} \times \boldsymbol{p}_{h_3}$ Some of the hadrons may be substituted by other particles, like leptons or photons. In the remaining of this section we deal with the process $p^{\uparrow} p \to \pi X$; all results can be extended to other processes, like those considered in the following section, by simply adapting the formalism to the case of interest.

The single spin asymmetry A_N for the process under consideration is defined as follows:

$$A_N(x_F, p_T) = \frac{d\sigma^{\uparrow} - d\sigma^{\downarrow}}{d\sigma^{\uparrow} + d\sigma^{\downarrow}}, \tag{1}$$

where $d\sigma^{\uparrow,\downarrow}$ stands for the differential cross section $E_\pi\, d^3\sigma^{p^{\uparrow,\downarrow} p \to \pi X}/d^3\boldsymbol{p}_\pi$; $x_F = 2p_L/\sqrt{s}$, where p_L is the pion longitudinal momentum in the pp c.m. frame, and $\sqrt{s}$ is the total c.m. energy. Of course, other sets of kinematical variables could also be used. Single spin asymmetries for other spin directions are forbidden by parity invariance.

By allowing for spin effects in the distribution function and still assuming the QCD factorization theorem to hold, one can write [5]

$$\begin{aligned} 2\, d\sigma^{unp} A_N(x_F, p_T) &= \sum_{a,b,c,d} \int d^2\boldsymbol{k}_{\perp a}\, dx_a\, dx_b\, \frac{1}{\pi z} \\ &\times \Delta^N f_{a/p^{\uparrow}}(x_a, \boldsymbol{k}_{\perp a}) f_{b/p}(x_b) \\ &\times \frac{d\hat{\sigma}^{ab \to cd}}{d\hat{t}} D_{\pi/c}(z), \end{aligned} \tag{2}$$

where $d\sigma^{unp}$ is the unpolarized differential cross section for the process under consideration; $f_{b/p}(x_b)$ is the distribution function, inside the unpolarized proton, of partons b with a fraction x_b of the proton momentum; $D_{\pi/c}(z)$ is the unpolarized fragmentation function for parton c fragmenting into a pion carrying a fraction z of its momentum; $d\hat{\sigma}^{ab \to cd}/d\hat{t}$ is the partonic cross section for the hard process $ab \to cd$.

The r.h.s. of Eq. (2) differs from the expression of the unpolarized cross section $d\sigma^{unp}$ in that the unpolarized distribution function of parton a inside the beam proton is replaced by a new, nonperturbative function, depending on the intrinsic transverse momentum of parton a,

$$\Delta^N f_{a/p^\uparrow}(x_a,\boldsymbol{k}_{\perp a}) = \sum_{\lambda_a}\Big[\hat{f}_{a,\lambda_a/p^\uparrow}(x_a,\boldsymbol{k}_{\perp a}) - \hat{f}_{a,\lambda_a/p^\uparrow}(x_a,-\boldsymbol{k}_{\perp a})\Big] \equiv 2\, I_{+-}^{a/p}(x_a,\boldsymbol{k}_{\perp a}), \tag{3}$$

which gives the difference between the total number of partons a, with momentum fraction x_a and intrinsic transverse momentum $\boldsymbol{k}_{\perp a}$ inside a proton with spin $\uparrow$ and a proton with spin $\downarrow$ [notice that $\hat{f}_{a,\lambda_a/p^\uparrow}(x_a,-\boldsymbol{k}_{\perp a}) = \hat{f}_{a,\lambda_a/p^\downarrow}(x_a,\boldsymbol{k}_{\perp a})$]. This same function is denoted by $f_{1T}^{\perp}$ in Ref. [16]. The symbol $I_{+-}^{a/p}(x_a,\boldsymbol{k}_{\perp a})$ has been introduced to show the non-diagonal nature, in the helicity basis, of this new quantity [5].

Eq. (3) explicitely shows that $\Delta^N f_{a/p^\uparrow}(x_a,\boldsymbol{k}_{\perp a})$ vanishes when $\boldsymbol{k}_{\perp a} \to 0$. Parity invariance requires $\Delta^N f$ to vanish when $\boldsymbol{k}_{\perp a}$ is perpendicular to the scattering plane (i.e. parallel to the proton spin), so that

$$\Delta^N f_{a/p^\uparrow}(x_a,\boldsymbol{k}_{\perp a}) = \hat{f}_{a/p^\uparrow}(x_a,\boldsymbol{k}_{\perp a}) - \hat{f}_{a/p^\downarrow}(x_a,\boldsymbol{k}_{\perp a}) \sim k_{\perp a}\sin\varphi \tag{4}$$

where φ is the angle between $\boldsymbol{k}_{\perp a}$ and the $\uparrow$ direction (normal to the scattering plane).

We notice that $\Delta^N f_{a/p^\uparrow}(x_a,\boldsymbol{k}_{\perp a})$ is an odd function of $\boldsymbol{k}_{\perp a}$. This means that we cannot neglect the $\boldsymbol{k}_{\perp a}$ dependence of all the other terms in the convolution integral, Eq. (2). We have to keep into account $\boldsymbol{k}_{\perp a}$ values in the partonic cross sections and the differences of dynamical contributions from partons with an intrinsic $+\boldsymbol{k}_{\perp a}$ and those with an intrinsic $-\boldsymbol{k}_{\perp a}$, which leads to an overall value of A_N due to higher twist effects.

In fact Eq. (2) for A_N can be rewritten as:

$$\begin{aligned} 2\,d\sigma^{unp}A_N(x_F,p_T) &= \sum_{a,b,c,d}\int d^2\boldsymbol{k}_{\perp a}\,dx_a\,dx_b\,\frac{1}{\pi z} \\ &\times \Delta^N f_{a/p^\uparrow}(x_a,\boldsymbol{k}_{\perp a})\,f_{b/p}(x_b) \\ &\times\left[\frac{d\hat{\sigma}^{ab\to cd}}{d\hat{t}}(\boldsymbol{k}_{\perp a}) - \frac{d\hat{\sigma}^{ab\to cd}}{d\hat{t}}(-\boldsymbol{k}_{\perp a})\right] \\ &\times D_{\pi/c}(z)\,, \end{aligned} \tag{5}$$

where now the integration on $\boldsymbol{k}_{\perp a}$ runs only over the positive half-plane of its components. The momentum fraction z is fixed in terms of x_a, x_b and $\boldsymbol{k}_{\perp a}$ by momentum conservation in the partonic process.

The only unknown function in Eq. (5) is the distribution $\Delta^N f_{a/p^\uparrow}(x_a,\boldsymbol{k}_{\perp a})$, which, at least in principle, is then measurable via the single spin asymmetry A_N.

In order to perform numerical calculations, we make some further assumptions which, while preserving the basic physical ideas of our approach, make computations easier to handle. Our first assumption is that the dominant effect is given by the valence quarks in the polarized protons. That is, we assume $\Delta^N f_{a/p^\uparrow}(x_a,\boldsymbol{k}_{\perp a})$ to be non-zero only for valence u, d quarks: while it is natural to assume a correlation between the proton polarization and the $\boldsymbol{k}_\perp$ of the polarized valence quarks, one does not expect so for the sea quarks. Moreover, sea quarks do not contribute much to the production of large x_F pions. Secondly, we evaluate Eq. (5) by assuming that the main contribution comes from $\boldsymbol{k}_{\perp a} = \boldsymbol{k}^0_{\perp a}$ where, as suggested by Eq. (4), $\boldsymbol{k}^0_{\perp a}$ lies in the overall scattering plane and its magnitude equals the average value of $\langle \boldsymbol{k}^2_{\perp a}\rangle^{1/2} = k^0_{\perp a}(x_a)$. This average value – which sets an overall physical scale for the transverse momentum effects – will in general depend on x_a. Estimates of this dependence have been given in the literature [17]; it can be well represented by the following expression (M being the proton mass):

$$\frac{1}{M}k^0_{\perp a}(x_a) = 0.47\,x_a^{0.68}(1-x_a)^{0.48}\,. \tag{6}$$

The residual x_a dependence in $\Delta^N f_{a/p^\uparrow}$ not coming from $k^0_{\perp a}$ is taken to be of the simple form

$$N_a x_a^{\alpha_a}(1-x_a)^{\beta_a}, \tag{7}$$

where N_a, α_a and β_a are free parameters.

We then end up with the simple expression:

$$\int d^2\boldsymbol{k}_{\perp a} \Delta^N f_{a/p^\uparrow}(x_a,\boldsymbol{k}_{\perp a}) \times\left[\frac{d\hat\sigma^{ab\to cd}}{d\hat t}(\boldsymbol{k}_{\perp a}) - \frac{d\hat\sigma^{ab\to cd}}{d\hat t}(-\boldsymbol{k}_{\perp a})\right] \simeq \frac{k^0_{\perp a}(x_a)}{M} N_a x_a^{\alpha_a}(1-x_a)^{\beta_a} \times\left[\frac{d\hat\sigma^{ab\to cd}}{d\hat t}(\boldsymbol{k}^0_{\perp a}) - \frac{d\hat\sigma^{ab\to cd}}{d\hat t}(-\boldsymbol{k}^0_{\perp a})\right]. \tag{8}$$

where $k^0_{\perp a}(x_a)$ is given by Eq. (6) and, choosing xz as the scattering plane and z as the direction of the incoming polarized proton, $\boldsymbol{k}^0_{\perp a} = (k^0_{\perp a},0,0)$.

Eqs. (5), (6) and (8) can be used to explain measured single spin asymmetries and to give numerical predictions for other single spin asymmetries of interest, as we shall do in the sequel. Parametrizations for the unpolarized partonic distributions are available in the literature; similarly for the fragmentation functions, for which, due to the increasing experimental information available, a lot of progress has been recently made. Finally, the analytical expressions of the elementary cross sections for all possible $ab \to cd$ partonic processes are well known; we only need to take into account the corresponding kinematical modifications due to the inclusion of the transverse momentum for parton a.

In Ref. [5] the parameters appearing in the expression of $\Delta^N f_{a/p^\uparrow}(x_a,\boldsymbol{k}_{\perp a})$, Eqs. (7) and (8), were fixed by fitting the available data on single spin asymmetries for the $p^\uparrow p \to \pi X$ process [1]. The main result of Ref. [5] was indeed to show that it is possible within such an approach to reproduce the experimental data with physically reasonable values of the parameters.

The precise values of the parameters depend on the specific choice adopted for the distribution and fragmentation functions. We have repeated here the fitting procedure of Ref. [5] choosing the MRSG parametrization [18] for the partonic distribution functions, and the parametrization (BKK1) of Ref. [19] for pion and kaon fragmentation functions. The quality of the fit – see Fig. 1 – is quite similar to the original one, while the values of the parameters do not change significantly:

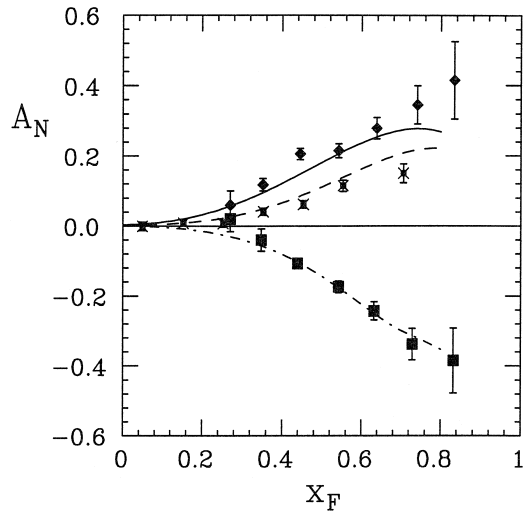

Fig. 1. Fit of the data on A_N for the process $p^\uparrow p \to \pi X$ [1], with the parameters given in Eq. (9); the upper, middle, and lower sets of data and curves refer respectively to π^+, π^0, and π^-.

	N_a	α_a	β_a	
u	3.68	1.34	3.58	(9)
d	−1.24	0.76	4.14	

Having fixed the values of the parameters we can now give predictions for several other single spin asymmetries, some of which have been measured; this is the main purpose of the paper.

3. New results

In this section we consider a number of interesting physical processes. The full set of such processes and the corresponding experimental data, either available now or, hopefully, in the near future, should allow to assess the validity of our model. Unless explicitely stated, we use the formalism, the sets of distribution and fragmentation functions discussed and the parameters derived in the previous section.

3.1. $\bar{p}^{\uparrow} p \to \pi X$

Recently the E704 Collaboration at Fermilab has presented results on SSA for inclusive production of pions in the collision of transversely polarized antiprotons off a proton target [2]. The kinematical conditions are the same as in the case of polarized protons [1]. Within our model the connection between SSA with polarized protons or antiprotons is very simple: since only valence quark contributions are taken into account, we just have to exploit charge conjugation relations for the I_{+-} distributions, i.e. $I_{+-}^{\bar{u}/\bar{p}} = I_{+-}^{u/p}$ and $I_{+-}^{\bar{d}/\bar{p}} = I_{+-}^{d/p}$. In particular, this means that

$$A_N(\bar{p}^{\uparrow} p \to \pi^{\pm}) \simeq A_N(p^{\uparrow} p \to \pi^{\mp})$$
$$A_N(\bar{p}^{\uparrow} p \to \pi^{0}) \simeq A_N(p^{\uparrow} p \to \pi^{0}). \qquad (10)$$

As it is shown in Fig. 2, this compares rather well with experimental results, although our results for π^0 and π^- are somewhat too large.

3.2. $p^{\uparrow}(\bar{p}^{\uparrow})p \to \pi X$ at fixed x_F, as a function of p_T

From presently available experimental data it is not possible to disentangle the behaviour of $A_N(x_F, p_T)$ as a function of x_F and p_T separately. This might be crucial for the refinement of theoretical models. The E704 Collaboration has published results on SSA for the processes $p^{\uparrow}(\bar{p}^{\uparrow})p \to \pi^0 X$ for $|x_F| \leq 0.15$, in the p_T range $1.0-4.5$ GeV/c [20]. No sizeable value of A_N was measured. The results of our model are in good agreement with these data: only at values of x_F greater than ~ 0.3 sizeable effects for the SSA can be obtained.

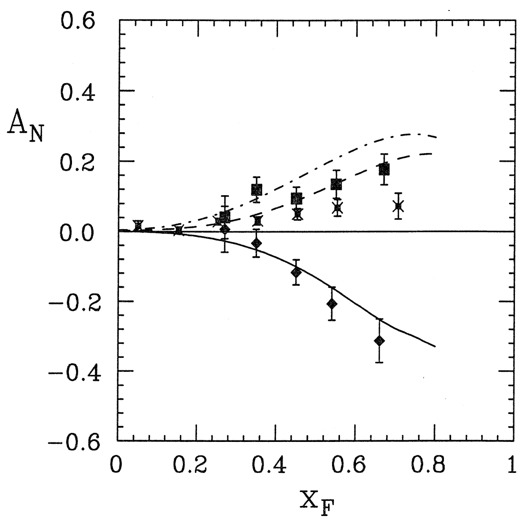

Fig. 2. Single spin asymmetries A_N for the process $\bar{p}^{\uparrow} p \to \pi X$ [2]; the lower, middle, and upper sets of data and curves refer respectively to π^+, π^0, and π^-.

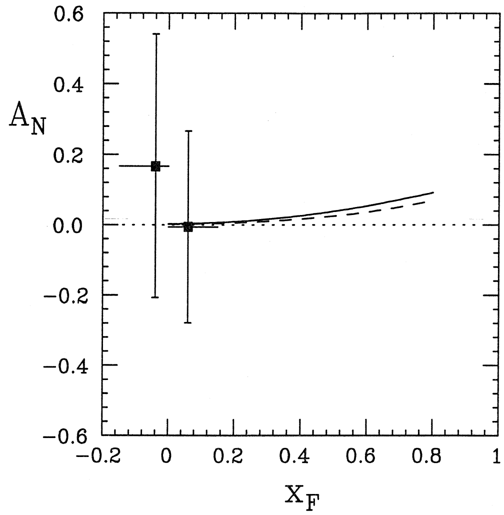

Fig. 3. Single spin asymmetry for the process $p^{\uparrow}(\bar{p}^{\uparrow})p \to \gamma X$; experimental data, at $|x_F| \leq 0.15$ and $2.5 < p_T < 3.1$ GeV/c, are from Ref. [21]. The curves show our corresponding theoretical predictions at $p_T = 2.5$ GeV/c; the solid curve refers to the $p^{\uparrow} p \to \gamma X$ process, the dashed curve to the $\bar{p}^{\uparrow} p \to \gamma X$ case.

3.3. $p^{\uparrow}(\bar{p}^{\uparrow})p \to \gamma X$

Unfortunately, only few experimental data [21], at $x_F \sim 0$ and p_T in the range $2.5-3.1$ GeV/c are available for this process at present. We stress its importance: possible origins of SSA arising in final quark fragmentation [6,7] are obviously excluded here; in principle we could then directly test those effects originating in initial state interactions and/or from quark-gluon correlations.

The predictions of our model are presented in Fig. 3, as a function of x_F. We choose the same kinematical conditions as in the available experimental data of E704 Collaboration [21]. Only two experimental points at $x_F \sim 0$, with huge errors, are available; they are in agreement with our results, but, given the large errors bars, we do not consider such an agreement as significant; further experimental data would be most helpful.

Theoretical estimates for SSA in $p^{\uparrow} p \to \gamma X$ were given also by Qiu and Sterman [8], who presented predictions for kinematical conditions similar to those of E704 collaboration [21]. Their theoretical model introduces higher twist distribution functions arising from soft quark-gluon correlations: results are given for two possible choices of these unknown higher twist distributions. These results have a very similar behaviour, as a function of x_F, when compared with ours, Fig. 3. However, they are greater by approximately a factor 2.

In Ref. [22] it has been argued that out of the two parametrizations for the higher twist distributions proposed in Ref. [8] the first one (which produces the larger result for the asymmetry) should be disregarded, since it violates some required physical constraints. The second one should be multiplied by an overall factor smaller that unity, leading to results on SSA reduced by at least a factor two. These considerations, if correct, reduce considerably the difference between the results of Qiu and Sterman and our present results.

3.4. $p^{\uparrow}(\bar{p}^{\uparrow})p \to KX$

It is interesting to investigate the inclusive production of hadrons different from the pion. On one hand, in fact, such new data would allow further tests of our model and of the parametrization of the I_{+-} distributions; on the other hand, the study of kaon production, for example, is interesting by itself, due to the presence of valence s quarks and the possibility of learning about the role of strange quarks in fragmentation processes.

To study kaon production, we need information on the corresponding fragmentation functions $D_{K/c}$. The knowledge of kaon FF is more limited than in the pion case and, to the best of our knowledge, only a few parametrizations are available in the literature. Let us briefly recall the main features of these parametrizations, which will be used in the sequel:

i) Ref. [19] not only gives the set of fragmentation functions (BKK1) used to fit the pion data, but also gives parton fragmentation functions into kaons. Although not the most recent one, it has the advantage to give separate contributions from leading (valence) and non-leading (sea) parent quarks of the produced meson, which allow to derive separate FF's for positively and negatively charged mesons;

ii) The same authors have recently published a more accurate parametrization (BKK2) [23], based on a much richer set of experimental data. However, this second set only gives FF either for the sum of charged mesons ($\pi^+ + \pi^-$, $K^+ + K^-$) or for neutral (π^0, $K^0 + \bar{K}^0$) mesons.

iii) Other sets of parametrizations for the sum of charged pion and kaon FF have been given by Greco and Rolli (GR) [24,25].

iv) Finally, some most recent fragmentation functions (IMR) can be found in Ref. [26].

Since the main contribution to the large x_F production of mesons comes from partonic processes involving valence quarks from the initial nucleon and in the final mesons, it is plausible to expect, for kaons produced from polarized protons, the following behaviour, in analogy to the pion case:

$$A_N(K^+) \sim A_N(\pi^+)\,; \quad A_N(K^0) \sim A_N(\pi^-)\,;$$
$$A_N(K^-) \sim A_N(\bar{K}^0) \sim 0\,. \tag{11}$$

This is clearly understood if we keep in mind the flavour content of pions and kaons ($\pi^+ = u\bar{d}$, $\pi^- = d\bar{u}$, $\pi^0 = (u\bar{u} - d\bar{d})/\sqrt{2}$; $K^+ = u\bar{s}$, $K^- = s\bar{u}$, $K^0 = d\bar{s}$, $\bar{K}^0 = s\bar{d}$) and remember that we are assuming that only $I_{+-}^{u,d}$ are different from zero for the incident polarized proton.

Regarding K_S^0 mesons, which are actually observed in experiments, since $\sigma(K_S^0) = \left[\sigma(K^0) + \sigma(\bar{K}^0)\right]/2$, using Eq. (1) we see that the asymmetry for K_S^0 is related to the asymmetries for K^0 and $\bar{K}^0$ by:

$$A_N(K_S^0) = \frac{A_N(K^0) + \left[\dfrac{\sigma^{unp}(\bar{K}^0)}{\sigma^{unp}(K^0)}\right] A_N(\bar{K}^0)}{1 + \dfrac{\sigma^{unp}(\bar{K}^0)}{\sigma^{unp}(K^0)}}\,. \tag{12}$$

Since we have used the BKK1 set of FF's for fitting the pion data, the first thing we can do is to use the corresponding BKK1 FF's for the kaons, keeping the same $I_{+-}^{u,d}$ functions as obtained from fitting the pion asymmetry data. Then, by inserting

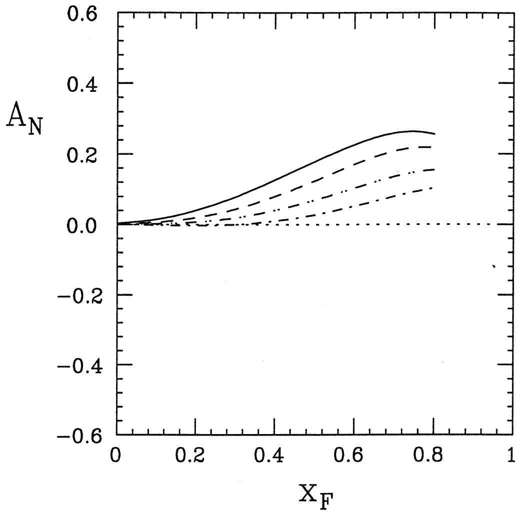

Fig. 4. Predicted single spin asymmetries for the process $p^{\uparrow} p \to KX$, with the set of kaon FF's BKK1 [19]; kinematical conditions are the same as for the pion case, at $p_T = 1.5$ GeV$/c$. The solid, dashed, dot-dashed, double dot-dashed curves refer respectively to the K^+, K^-, K^0, K^0_S cases. Results for $\bar{K}^0$ meson are very similar to those for K^- case.

the $D_{K/c}$ given in Ref. [19] into Eq. (5) and using Eqs. (6), (8) and (9), we obtain predictions for kaon SSA, in the same kinematical region as those observed for pions [1,2]: they are shown in Fig. 4.

Surprisingly, the results for K^- and K^0 are very different from what expected qualitatively, see Eq. (11) and compare Figs. 1 and 4. A look at the parametrization of the BKK1 FF's, $D_{K/q}(z)$, for the leading (valence) parent quarks (e.g. u quark inside K^+) and the non-leading (sea) ones (e.g. s quark inside K^+) helps to understand the origin of such a discrepancy: in BKK1 FF the non-leading contributions are not suppressed, for high values of z, compared to the leading ones, which is what one would have expected on general grounds, leading to Eq. (11). That the unexpected behaviour of Fig. 4 originates from this property of the BKK1 FF can be confirmed in a simple way: if we take the valence contributions of BKK1 kaon FF's, but *rescale* the sea ones by assuming the same relative behaviour as in the BKK1 pion case (that is, we redefine $D_{K/sea} \simeq D_{K/val}[D_{\pi/sea}/D_{\pi/val}]$), then the results for kaon asymmetries agree very well (see Fig. 5) with the expected qualitative behaviour, Eq. (11). Of course, this simple example has no physical motivation, but only clarifies the reason why kaon asymmetries in Fig. 4 are so different from what expected. The real, physical origin of this discrepancy is in the behaviour at high z of non-leading *vs.* leading contri-

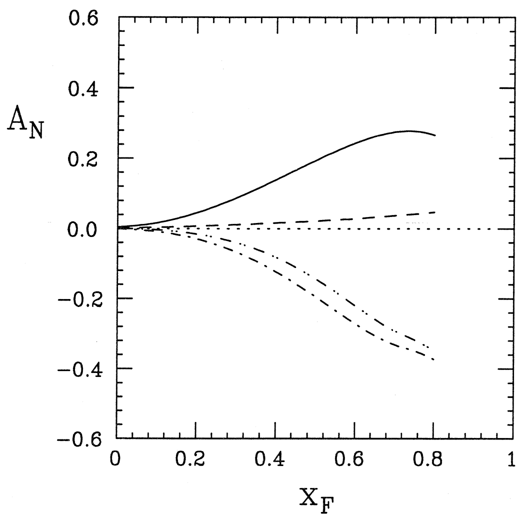

Fig. 5. The same as for Fig. 4, but using the set of kaon FF's BKK1 modified so that $D_{K/sea} \simeq D_{K/val}[D_{\pi/sea}/D_{\pi/val}]$ (see text for more details).

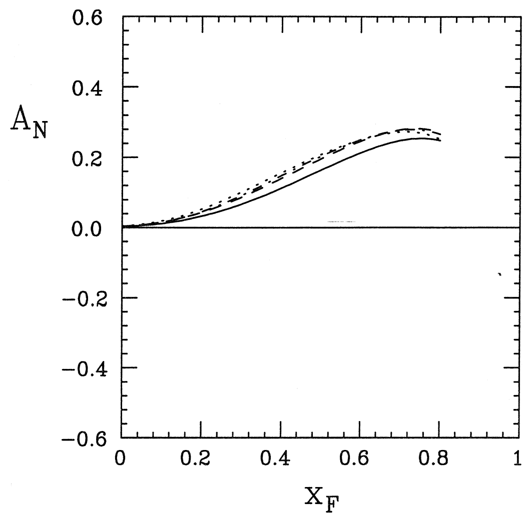

Fig. 6. Predicted single spin asymmetries for the process $p^{\uparrow} p \to (K^+ + K^-)\ X$; kinematical conditions are the same as for the pion case, at $p_T = 1.5$ GeV$/c$. The curves correspond respectively to the set of kaon FF's BKK1 [19] (solid); BKK2 [23] (dashed); GR [25] (dot-dashed); IMR [26] (dotted).

butions in the BKK1 FF's. Indeed, results consistent with Eq. (11) are found – as we shall show – when using fragmentation functions for which the sea quark contribution is suppressed with respect to the valence one.

In Fig. 6 we give and compare results obtained for $A_N(K^+ + K^-)$ with the different sets of fragmentation functions discussed above, i) to iv). This asymmetry is related to the asymmetries for the separate production of K^+ and K^- by

$$A_N(K^+ + K^-) = \frac{A_N(K^+) + \left[\frac{\sigma^{unp}(K^-)}{\sigma^{unp}(K^+)}\right] A_N(K^-)}{1 + \frac{\sigma^{unp}(K^-)}{\sigma^{unp}(K^+)}}. \qquad (13)$$

According to Eq. (11) one would expect, at large x_F, $A_N(K^+ + K^-) \simeq A_N(K^+)$. All sets of FF yield similar results.

Finally, we show in Fig. 7 the predictions obtained – again using the different sets of fragmentation functions – for $A_N(K_S^0)$. Contrary to the production of charged kaons, in this case the results strongly depend on the set of fragmentation functions: sets which enhance the role, at large z, of valence quarks give results in agreement with the expectations of Eq. (11), whereas sets which allow a large contribution from sea quarks give totally different results, even in sign. A measurement of $A_N(K_S^0)$ would then supply useful information.

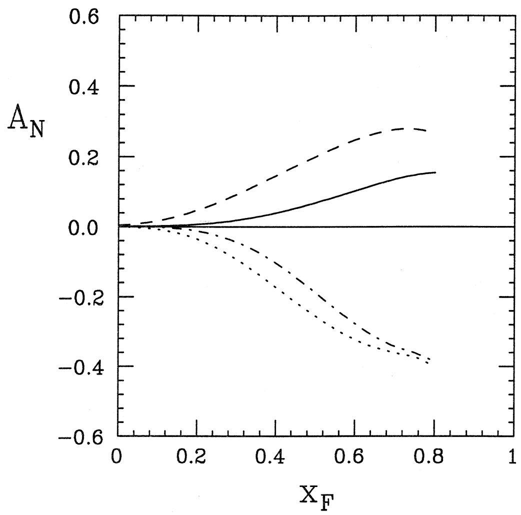

Fig. 7. Predicted single spin asymmetries for the process $p^{\uparrow} p \to K_S^0 X$; kinematical conditions are the same as for the pion case, at $p_T = 1.5$ GeV/c. Notations for the theoretical curves are the same as in Fig. 6.

We have also investigated possible contributions from a non-zero I_{+-} distribution for the strange quark inside the polarized proton. In principle, kaon asymmetries should be quite sensitive to this distribution. However, as we have checked explicitely, at least in the kinematical region considered here, the role of strange quarks is strongly suppressed. This is because strange quarks come in any case from the sea quarks in the proton, so that their (unpolarized and, to a greater extent, polarized) distributions decrease at large x much faster than the distributions of u and d quarks.

The values of the SSA for kaons produced with polarized antiprotons, $\bar{p}^{\uparrow} p \to KX$ can easily be deduced from those obtained with polarized protons by applying simple charge conjugation arguments, in analogy to Eq. (10); we have explicitely checked that this is the case.

4. Summary and conclusions

We have consistently applied a QCD hard scattering scheme, with the inclusion of some higher twist effects, to the description of single spin asymmetries (SSA) in $p^{\uparrow} p \to \pi X$, refining a previous [5] determination of a new $\boldsymbol{k}_{\perp}$ and spin dependent distribution function, introduced by several authors [3–5,16] as a possible source of single spin asymmetries.

We have then used this distribution function to compute several single spin asymmetries in other processes, namely $\bar{p}^{\uparrow} p \to \pi X$, $p^{\uparrow}(\bar{p}^{\uparrow}) p \to \gamma X$ and $p^{\uparrow}(\bar{p}^{\uparrow}) p \to KX$. In the first two cases we have no free parameters and our results are genuine predictions of the model: they turn out to agree with the experimental data, although the data on SSA in γ production are still rather qualitative with large errors.

The SSA for the $p^{\uparrow} p \to KX$ process are predicted to be large, and, for neutral kaons, their actual value strongly depend on the set of fragmentation functions used; experimental information would allow to discriminate between different sets of kaon fragmentation functions. The main feature of the quark fragmentation which influences the value of the SSA is the relative importance of sea and valence quark contributions.

Our scheme seems then to be a good phenomenological way of describing single spin asymmetries within a generalization of the QCD factorization theorem; of course, other effects [6,11], might be present and play a smaller or larger role, depending on the process considered. For example, SSA in γ production or Drell-Yan processes [11] should mostly originate from the mechanism used here or in Ref. [8]; in other cases another or several other effects might be active at the same time. A possible strategy to isolate different origins of SSA in Deep Inelastic Scattering has been discussed in Ref. [27]. More data and more phenomenological calculations are needed.

Acknowledgements

We are very grateful to M.E. Boglione, who contributed to the early stages of this work. We wish also to thank M. Greco and S. Rolli; J. Binnewies, B.A. Kniehl and G. Kramer; D. Indumathi, H.S. Mani and A. Rastogi, for kindly sending us a copy of their numerical routines, and for clarifying discussions and comments on their parametrizations of meson fragmentation functions.

References

[1] D.L. Adams et al., Phys. Lett. B 264 (1991) 462.
[2] A. Bravar et al., Phys. Rev. Lett. 77 (1996) 2626.
[3] D. Sivers, Phys. Rev. D 41 (1990) 83.
[4] D. Sivers, Phys. Rev. D 43 (1991) 261.
[5] M. Anselmino, M. Boglione, F. Murgia, Phys. Lett. B 362 (1995) 164.
[6] J. Collins, Nuc. Phys. B 396 (1993) 161.
[7] X. Artru, J. Czyzewski, H. Yabuki, Z. Phys. C 73 (1997) 527.
[8] J. Qiu, G. Sterman, Phys. Rev. Lett. 67 (1991) 2264.
[9] J. Qiu, G. Sterman, e-Print Archive: hep-ph/9806356.
[10] A.V. Efremov, V.M. Korotkiyan, O. Teryaev, Phys. Lett. B 348 (1995) 577.
[11] N. Hammon, O. Teryaev, A. Schäfer, Phys. Lett. B 390 (1997) 409.
[12] C. Boros, Liang Zuo-tang, and Meng Ta-chung, Phys. Rev. Lett. 67 (1993) 1751.
[13] J. Collins, Nucl. Phys. B 394 (1993) 169.
[14] M. Anselmino, A. Drago, F. Murgia, e-Print Archive: hep-ph/9703303.
[15] M. Anselmino, M. Boglione, F. Murgia, Proceedings of Spin96, World Scientific (1997), p. 461.
[16] D. Boer, P. Mulders, Phys. Rev. D 57 (1998) 5780.
[17] J.D. Jackson, G.G. Ross, R.G. Roberts, Phys. Lett. B 226 (1989) 159.
[18] A.D. Martin, W.J. Stirling, R.G. Roberts, Phys. Lett. B 354 (1995) 155.
[19] J. Binnewies, B.A. Kniehl, G. Kramer, Z. Phys. C 65 (1995) 471.
[20] D.L. Adams et al., Phys. Rev. D 53 (1996) 4747.
[21] D.L. Adams et al., Phys. Lett. B 345 (1995) 569.
[22] B. Ehrnsperger, A. Schäfer, W. Greiner, L. Mankiewicz, Phys. Lett. B 321 (1994) 121.
[23] J. Binnewies, B.A. Kniehl, G. Kramer, Phys. Rev. D 52 (1995) 4947; Phys. Rev. D 53 (1996) 3573.
[24] M. Greco, S. Rolli, A. Vicini, Z. Phys. C 65 (1995) 277.
[25] M. Greco, S. Rolli, Phys. Rev. D 52 (1995) 3853.
[26] D. Indumathi, H.S. Mani, A. Rastogi, e-Print Archive: hep-ph/9802324.
[27] M. Anselmino, E. Leader, F. Murgia, Phys. Rev. D 56 (1997) 0021.