$\Delta\sigma_L(pp)$ AND JET PHYSICS*

D. G. RICHARDS
High Energy Physics Division
Argonne National Laboratory
Argonne, IL 60439

We show that there is a positive contribution to $\Delta\sigma_L(pp;s) = \sigma_{\rm tot}(p(+)p(+);s) - \sigma_{\rm tot}(p(+)p(-);s)$ (where the $\pm$ refer to proton helicities) associated with the pointlike scattering of fundamental constituents. Simple arguments imply that this positive contribution would, at very high s, be larger in absolute value than the negative contribution to $\Delta\sigma_L$ predicted from the exchange of the A_1 reggeon, and furthermore may provide important insight into the shape of the spin weighted quark and gluon distributions. Measurements of $\Delta\sigma_L$ in the energy range $\sqrt{s} = 18 - 30$ GeV also should help clarify theoretical ideas associated with the observation of "minijets" and could aid in the prediction of event structure at future high energy colliders.

I. Introduction

This talk summarizes work undertaken in collaboration with Gordon Ramsey and Dennis Sivers.[1]

The recent data from the European Muon Collaboration[2] on the small-x behavior of the deep inelastic scattering asymmetry for polarized leptons and polarized protons suggests that the total spin carried by the valence quarks in a polarized proton may be approximately cancelled by a strong negative polarization of the sea of $q\bar{q}$ pairs. The spin structure of the proton that this result implies is not easily understood in terms of traditional quark-model ideas.

Various models of the spin structure of the proton have been proposed in response to the gauntlet thrown down by the EMC data.[3,4] Arguments based on the Altarelli-Parisi

*Work supported by the U.S. Department of Energy, Division of High Energy Physics, Contract W-31-109-ENG-38.

evolution equations suggest that in general the amount of angular momentum associated with polarized gluons within a polarized proton should grow as the proton is probed with increasing Q^2, and be exactly balanced by the growth of orbital angular momentum. However the "hybrid" quark-skyrme picture of the proton,[4] for example, avoids the introduction of large orbital angular momentum. Information concerning the polarized gluon structure function will be vital in discriminating between these different pictures of the proton's spin structure.

In this talk I shall show that the contribution to $\Delta\sigma_L(pp;s) = \sigma_{\text{tot}}(p(+)p(+);s) - \sigma_{\text{tot}}(p(+)p(-);s)$ (where the $\pm$ refer to proton helicities) associated with the pointlike scattering of the fundamental constituents of the proton would be substantial at very large s and could provide important information concerning the spin-weighted parton densities, and in particular that of the gluon. Furthermore such measurements would prove a useful tool in examining the interplay between hard, point-like scattering mechanisms and soft, coherent dynamics. Attention has been focussed on this issue by the observation of a large cross section for "minijets" at the CERN $S\bar{p}pS$ collider.[5,6]

Phenomenological questions concerning the impact of jet physics begin with a breakup of the total inelastic cross section

$$\sigma_{\text{inel}}(pp;s) = \sigma_{\text{soft}}(pp;p_0,s) + \sigma_{\text{jet}}(pp;p_0,s), \tag{1.1}$$

where the jet cross section, $\sigma_{\text{jet}}(pp;p_0,s)$, is the cross section for observing at least one large-p_T jet ($p_T > p_0$). The split-up in (1.1) is obviously dependent on the cutoff, p_0, used to define the jet cross section. If the cutoff is chosen to be large enough, the jet cross section should be given by the integral of the large-p_T differential cross section which is calculable within the framework of the QCD-aided parton model.[7]

Although the basic theoretical justification for using QCD perturbation theory to calculate a wide variety of hard processes has been given a boost by recent work on perturbative factorization,[8] the split-up in (1.1) is subject to a considerable amount of theoretical uncertainty. This uncertainty is conveniently discussed in the form of the renormalization prescription dependence and the factorization prescription dependence of the low-order QCD calculation.[9] As an indication of this uncertainty, theoretical estimates for the magnitude, cutoff dependence and energy dependence of $\sigma_{\text{jet}}(pp;p_0,s)$ are highly sensitive to the small-x behavior of the quark and gluon distributions in the proton. Alternate prescriptions can differ significantly concerning these distributions. The experimental jet cross section must be measured as a function of energy and cutoff and compared with the calculated values to test the validity of the overall prescription scheme. Only then can the QCD prediction for the jet cross section be connected to predictions for other hard processes.

We can use the ideas behind the decomposition (1.1) in the measurement of $\Delta\sigma_L(pp;s)$. Applying the perturbative factorization hypothesis to individual helicity cross sections, the decomposition

$$\Delta\sigma_L(pp;s) = \Delta\sigma_L^{\mathrm{soft}}(pp;p_0,s) + \Delta\sigma_L^{\mathrm{jet}}(pp;p_0,s), \tag{1.2}$$

allows an independent test of the ideas behind the hard scattering model. The differential cross sections for scattering of quarks and gluons from definite helicity states are known in the large momentum transfer limit from QCD perturbation theory.[10] Measurement of $\Delta\sigma_L^{\mathrm{jet}}(pp;p_0,s)$ for different values of the jet cutoff can therefore give important information about the nature of the helicity-weighted quark and gluon distributions in a polarized proton. In particular our numerical studies indicate that $\Delta\sigma_L^{\mathrm{jet}}(pp;p_0,s)$ can be sensitive to the behavior of $\Delta G_{g/p}(x,\mu^2) = G_{g(+)/p(+)}(x,\mu^2) - G_{g(-)/p(+)}(x,\mu^2)$ at small x. Since there have been,to date, no hard experimental results concerning $\Delta G_{g/p}(x,\mu^2)$, it would be very valuable to compare data on $\Delta\sigma_L^{\mathrm{jet}}(pp;p_0,s)$ with models of the constituent distributions based on current theoretical wisdom. It is possible, therefore, that early measurements with the proposed Fermilab polarized proton beam will force us to drastically revise basic theoretical concepts involving hadron structure.[11] It is certain that these measurements will be the first among a variety of high-p_T spin asymmetry measurements which are necessary if we are to complete our knowledge of the helicity-weighted quark and gluon distributions.[12]

As in the case of the unpolarized cross sections, the observation of a sizeable $\Delta\sigma_L^{\mathrm{jet}}(pp;p_0,s)$ raises some important issues concerning the interplay of coherent hadronic dynamics and point-like constituent scattering. Consider, for example, the question of the asymptotic behavior of $\Delta\sigma_L(pp;s)$. Unitarity relates $\Delta\sigma_L(pp;s)$ to the imaginary part of an elastic scattering amplitude with unnatural parity in the t-channel

$$\Delta\sigma_L(pp;s) = \frac{8\pi}{p}\mathrm{Im}U_0(s,0), \tag{1.3}$$

where U_0 is defined in terms of the s-channel helicity amplitudes for pp elastic scattering by

$$U_0(s,t) = \frac{1}{2}\left(\langle + + | + + \rangle - \langle + - | + - \rangle\right). \tag{1.4}$$

Conventional Regge-pole phenomenology[13] suggests that the A_1 pole is the leading singularity with unnatural parity which can couple to $U_0(s,t)$ at $t = 0$. If we assume that $U_0(s,t)$ displays coherent, Regge, behavior at high energy we get

$$\Delta\sigma_L(pp;s) \cong \beta(0)(s/s_0)^{\alpha_{A_1}(0)-1}, \tag{1.5}$$

where the intercept of the A_1 Regge trajectory is approximately

$$\alpha_{A_1}(0) \cong -0.15. \tag{1.6}$$

This implies an energy dependence

$$\Delta\sigma_L(pp;s) \sim s^{-1.15}. \tag{1.7}$$

This assumption has proven to be adequate for describing the energy dependence of $\Delta\sigma_L$ over the range measured with polarized beam and target at the Argonne ZGS.[14] The parameterization (1.5) has been used to provide an estimate for $\Delta\sigma_L$ at higher energies. Normalizing the asymptotic prediction to the measured value at $p_{lab} = 11.75$ GeV[14] of -500 ± 50 μb yields

$$\Delta\sigma_L \cong -17.5(s/s_0)^{-1.15}\text{mb}, \tag{1.8}$$

with $s_0 = 1$ GeV^2. For $p_{lab} = 200$ GeV/c, this gives an estimate

$$\Delta\sigma_L(p_{lab} = 200\ GeV/c) \cong -19\mu\text{b}. \tag{1.9}$$

Equation (1.8) gives the value to which we must compare our expectations for $\Delta\sigma_L^{\text{jet}}$. It is interesting to keep in mind that the value predicted by Regge theory is negative.

In contrast, the incoherent pointlike contribution to $\Delta\sigma_L^{\text{jet}}$ associated with quark-quark, quark-gluon and gluon-gluon scattering in (1.2) should provide a positive contribution to $\Delta\sigma_L$. The reason for the positive sign is straightforward. It depends on the fact that the dominant underlying two-spin asymmetries in perturbative QCD are positive and that, based on our current ideas, there should exist a positive correlation between the spin of the proton and the spin of its constituents. When we look more carefully at the asymptotic behavior of $\Delta\sigma_L^{\text{jet}}$ a surprise emerges. As we shall demonstrate, the asymptotic behavior of $\Delta\sigma_L^{\text{jet}}$ depends on the small-x behavior of the spin-weighted quark and gluon distributions. If the distributions have the behavior near $x = 0$

$$\lim_{x\to 0} \Delta G_{i/p}(x, \mu^2) = ax^{-J}, \tag{1.10}$$

then the asymptotic behavior of the jet cross section is

$$\Delta\sigma_L^{\text{jet}} = \beta s^{J-1} \ln(s/s_0). \tag{1.11}$$

Simple model estimates of the small-x behavior of the helicity-weighted quark and gluon distributions suggest that the value of J in (1.10) should be near $\alpha_\rho(0) \cong 1/2$ so $\Delta\sigma_L^{\text{jet}}$ may, in fact, be larger in absolute value than the Regge estimate for $\Delta\sigma_L$, (1.8), at asymptotic values of s. We will investigate this question in more detail below. In principle, the small-x behavior of the constituent distributions can be measured in a variety of processes. Measurement of the energy dependence of $\Delta\sigma_L^{\text{jet}}$ for fixed cutoff can therefore provide important new information concerning the regime where the hard-scattering approximation is valid. Our numerical estimates for the jet contribution to the cross section using the available models[15] for the spin-weighted distributions give

$$\Delta\sigma_L^{\text{jet}} \left(pp; p_0 = \sqrt{5}\ \text{GeV}, \sqrt{s} = 20\ \text{GeV}\right) \cong 1.4\mu\text{b}, \tag{1.12}$$

which can be compared to (1.9). While the jet component of $\Delta\sigma_L^{\text{jet}}$ is expected to be smaller than the coherent component in this energy range, it cannot be neglected completely. Based on current models for the constituent distributions, we estimate

$$\left|\Delta\sigma_L^{\text{jet}} \left(pp; p_0^2 = 5\ \text{GeV}^2, \sqrt{s}\right)\right| > \left|\Delta\sigma_L(\sqrt{s})\right|_{A_1} \tag{1.13}$$

for $\sqrt{s} > 40$ GeV.

The remainder of this talk will be organized as follows. Sec. II shows the calculation of $\Delta\sigma_L^{\text{jet}}$ within the framework of the simple parton model and discusses the expectations for the spin-weighted distributions. By using some simple models for these distributions, we can calculate the range of possibilities for $\Delta\sigma_L^{\text{jet}}$. Section III concludes with some discussion about the experimental situation and the interpretation of these estimates.

II. An Estimate for $\Delta\sigma_L^{\text{jet}}(pp;p_0,\sqrt{s})$

As discussed in the introduction, the theoretical framework of the QCD-based parton model can be used to estimate a contribution to $\Delta\sigma_L(pp;s)$ associated with the hard scattering of fundamental constituents. We hypothesize the split-up of $\Delta\sigma_L$ given in (1.2). In analogy with the procedures for unpolarized jet cross sections we can then obtain the formula:

$$\Delta\sigma_L^{\text{jet}}(pp;p_0,\sqrt{s}) = \frac{\pi\alpha_s^2(\mu^2)}{2s}\sum_{i,j}\int_{x_1x_2>\xi} dx_1dx_2\frac{\Delta G_{i/p}(x_1,\mu^2)}{x_1}\frac{\Delta G_{j/p}(x_2,\mu^2)}{x_2}\Delta H_{ij}(z_0). \quad (2.1)$$

where $\xi = 4p_0^2/s$ and $z_0 = (1-\xi/x_1x_2)^{1/2}$. The cross section factor, ΔH_{ij}, is

$$\Delta H_{ij}(z_0) = \int_{-z_0}^{+z_0} dz\left[\frac{2s}{\alpha^2}\frac{d\Delta\sigma_{ij}}{dx_1dx_2dz}(z)\right]. \quad (2.2)$$

The spin-weighted distribution functions $\Delta G_{i/p}(x,\mu^2)$ give the probability for finding a given constituent in the proton with its helicity aligned with that of the proton minus the probability with helicity opposite that of the proton

$$\Delta G_{i/p}(x,\mu^2) = G_{i(+)/p(+)}(x,\mu^2) - G_{i(-)/p(+)}(x,\mu^2). \quad (2.3)$$

We will be dealing with quarks and massless gluons so that the $(+)$ and $(-)$ are the only allowed helicities. For the time being we will leave the choice of factorization scale, μ^2, open. Before proceeding with numerical results, however, we want to show that there exists an interesting puzzle concerning the asymptotic behavior of $\Delta\sigma_L$.

In order to obtain a crude estimate for the asymptotic behavior of $\Delta\sigma_L^{\text{jet}}(pp;p_0,\sqrt{s})$ we will initially assume that the process is dominated by the uu scattering contribution. We will also assume a simple approximation for the small-x region of the spin-weighted distribution

$$\Delta u(x,\mu^2) \cong a_u(\mu^2)x^{-J(\mu^2)}. \quad (2.4)$$

We would expect $J \cong \alpha_\rho(0) + O(\alpha_s/\pi)$ in analogy to the association of the small-x behavior of the nonsinglet distributions with the intercept of the pomeron. Using the form of the angular factor

$$\Delta H_{uu} = \frac{8}{9}\left[c_0 + a_0 \ln\left(\frac{x_1 x_2}{\xi}\right) + \sum_{n=1}^{\infty} c_n \xi^n x_1^{-n} x_2^{-n}\right], \tag{2.5}$$

with $c_0 = \left[\frac{16}{3}\ln 4 - 4\right]$, $a_0 = \frac{16}{3}, \ldots$ we use (2.4) and (2.5) to get

$$\Delta\sigma_L^{jet}(pp;\xi,s) \cong \frac{\pi\alpha_s^2}{2s} a_u^2 \left(\frac{8}{9}\right) \xi^{-J}\left[A(J) + B(J)\ln\xi\right], \tag{2.6}$$

where

$$\begin{aligned} A(J) &= -\frac{32}{3}\frac{1}{J^3} - \left(\frac{16}{3}\ln 4 - 4\right)\frac{1}{J^2} - \sum_{n=1}^{\infty}\frac{c_n}{(J+n)^2} \\ B(J) &= -\frac{16}{3}\frac{1}{J^2} - \left(\frac{16}{3}\ln 4 - 4\right)\frac{1}{J} - \sum_{n=1}^{\infty}\frac{c_n}{(J+n)}. \end{aligned} \tag{2.7}$$

If our current ideas concerning the small-x behavior of the structure functions are found to be correct and (2.4) is valid with $J = \alpha_\rho(0) \cong 1/2$, then we have

$$\Delta\sigma_L^{jet} \cong s^{-1/2}\left[A + B\ln s\right]. \tag{2.8}$$

Note that this dominates asymptotically any contribution to $\Delta\sigma_L$ associated with the exchange of the A_1 trajectory. One possibility which must be considered to allow the smooth merger of the hard scattering contribution and the coherent contribution is that the small-x behavior of the spin-weighted quark distribution is, instead, related to the intercept of the A_1 trajectory

$$\begin{aligned} \lim_{x\to 0}\Delta u(x,\mu^2) &\cong a_u(\mu^2) x^{-\alpha_{A_1}(0)} \\ \lim_{x\to 0}\Delta d(x,\mu^2) &\cong a_d(\mu^2) x^{-\alpha_{A_1}(0)}. \end{aligned} \tag{2.9}$$

Although this behavior is not postulated in the models for these functions, it is not ruled out by the data. It is a symptom of our lack of knowledge concerning these distributions that this possibility is still viable.

The assumption that one subprocess dominates the asymptotic behavior of $\Delta\sigma_L^{jet}$ is considerably less believable than is the case for the single effective process approximation[16] for unpolarized scattering. However, the puzzle concerning the asymptotic behavior of $\Delta\sigma_L$

only becomes more complicated when many processes are considered, and the underlying contradiction remains unresolved.

Different subprocesses provide both positive and negative contributions to $\Delta\sigma_L^{\rm jet}(pp; p_0, s)$. At very large $\sqrt{s}$, the gg process is significant and the behavior of the helicity-weighted gluon distribution near $x = 0$ is particularly important. From a study of the evolution equation it has been conjectured[17]

$$\lim_{x\to 0}\frac{\Delta G(x,\mu^2)}{G(x,\mu^2)} \sim x\exp\left[c\left(\ln\ln(\mu^2/\mu_0^2)\ln(1/x)\right)^{1/2}\right]. \tag{2.10}$$

We can adopt a simple power approximation for the small-x behaviour of the gluon distribution

$$\lim_{x\to 0} G(x,\mu^2) = a(\mu^2)x^{-J(\mu^2)}. \tag{2.11}$$

Collins[18] has emphasized that the study of scaling violations of the gluon distribution leads to a leading singularity in (2.11) that is above one, and arguments based on QCD perturbation theory[19] suggest

$$\begin{aligned} J(\mu^2) &= 1 + (12\ln 2)\frac{\alpha_s(\mu^2)}{\pi} + O\left(\alpha_s^2\right) \\ &\cong 1.67 \quad (\alpha_s = 0.25). \end{aligned} \tag{2.12}$$

Combining this model of the small-x behavior of the gluon distribution with (2.10) gives the behavior of the gluon-gluon and gluon-quark processes. We then arrive at the conclusion that the contributions to $\Delta\sigma_L^{\rm jet}$ from these processes should also fall less rapidly than the Regge prediction for $\Delta\sigma_L$ from A_1 exchange.

These arguments indicate that there exists an interesting relationship between the hard-scattering mechanisms of QCD and the coherent processes which lead to $\Delta\sigma_L$. The situation is similar to that of the unpolarized cross section except that the "contradiction" concerning the asymptotic behavior of the hard component here involves the spin-weighted valence quark distributions. Unlike the gluon distribution, these can be measured in polarized electroproduction experiments and the small-x behavior, (2.4), can be checked. The shadowing corrections, multiparton processes and unitarity constraints involved in the interpretation of the unpolarized jet data[20,21] should have their place in understanding $\Delta\sigma_L^{\rm jet}$. The final resolution may also involve new concepts.

To understand the question of the importance of the hard component to $\Delta\sigma_L$ at a given energy, we must do more than the simple analytic estimates above. We require explicit representations of the polarized parton distributions. As emphasized earlier, there is no data whatsoever on the polarized gluon distribution, and models of the spin-weighted sea quark distributions may need radical revision in light of the EMC data.[2] We are now investigating[22] the range of parametrizations of the polarized structure functions consistent both with the data and our theoretical expectations. However in the following numerical estimates our approach is to calculate $\Delta\sigma_L^{\text{jet}}$ using (2.1) and the best currently available models of the constituent distributions.

We will employ the parametrizations of the structure functions obtained by Chiappetta and Soffer,[23] which we will review briefly here. The polarized valence quark distributions,

$$\Delta q(x,\mu^2) = \Delta G_{q/p}(x,\mu^2) \tag{2.13}$$

are given by

$$\Delta u_v(x,\mu^2) = \left(u_v(x,\mu^2) - \frac{2}{3}d_v(x,\mu^2)\right)\left[1 + H_0(\mu^2)\frac{(1-x)^2}{\sqrt{x}}\right]^{-1} \tag{2.14a}$$

$$\Delta d_v(x,\mu^2) = -\frac{1}{3}d_v(x,\mu^2)\left[1 + H_0(\mu^2)\frac{(1-x)^2}{\sqrt{x}}\right]^{-1} \tag{2.14b}$$

where $u_v(x,\mu^2)$ and $d_v(x,\mu^2)$ are the unpolarized up and down valence quark distributions extracted by Gluck, Hoffmann and Reya.[24] The initial polarized gluon and sea quark distributions are assumed to be generated from the polarized valence distributions by a process of bremsstrahlung and quark-antiquark pair creation, with the normalization determined by requiring that the sum of the third components of the spins of the constituents be $\frac{1}{2}$.[10] Explicitly the gluon and sea distributions are given by

$$\begin{aligned}\Delta q_s(x,\mu_0^2) &= 0.0327x(2-x)(1-x)^{6.5}\\ \Delta g(x,\mu_0^2) &= 0.141x(5-2x)(1-x)^5.\end{aligned} \tag{2.15}$$

with

$$\mu_0^2 = 5\text{GeV}^2.$$

The distributions are evolved according to the Altarelli-Parisi equations, and the energy dependence of the spin dilution parameter H_0 of Eq.(2.14) is extracted by requiring that the Bjorken sum rule be satisfied. The value used is $H_0 = .114$ at $\mu^2 = 5$ GeV^2.

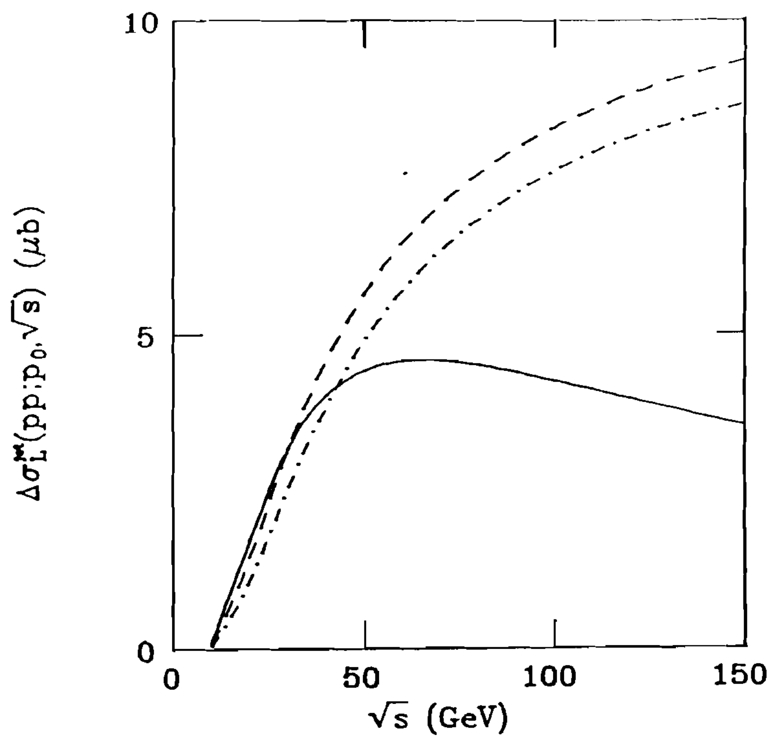

Fig. 1. The polarized jet cross section $\Delta\sigma_L^{jet}(pp;p_0,\sqrt{s})$ is shown as a function of $\sqrt{s}$ with cutoff $p_0^2 = 5$ GeV^2. The three curves correspond to $\mu^2 = 5$ GeV^2 (solid line), $\mu^2 = 9$ GeV^2 (dashed line) and $\mu^2 = 20$ GeV^2 (dot-dash)

We now address the choice of renormalization and factorization prescription. Any comprehensive treatment of this question must await the calculation of higher order corrections. However it is clear that the renormalization scale should be chosen to be of order p_0^2, the typical momentum transfer occurring in the interaction. Furthermore, if we set $\mu^2 = \lambda p_0^2$ then we wish $\Delta\sigma_L^{jet}$ to be reasonably insensitive to changes in λ. This issue is well illustrated in Fig. 1 where we show $\Delta\sigma_L^{jet}$ as a function of $\sqrt{s}$ at fixed cut-off $p_0^2 = 5$ GeV^2 for $\mu^2 = 5$, 9 and 20 GeV^2, the first value being just that of the starting distributions of Eq. (2.14). Whereas the calculations employing the higher two energy scales exhibit qualitatively similar behavior, the calculation at $\mu^2 = p_0^2$ peaks at a very much smaller value of $\sqrt{s}$. This qualitative difference in the behavior is attributable purely to the large increase in the gluon distribution at small x as the energy scale μ^2 is increased. Thus to the extent that the dichotomy into a hard and a soft cross-section exhibited in Eq. (1.2) is indeed valid at p_0^2 as low as 5 GeV^2, the choice of factorization scale $\mu^2 = 2p_0^2$ would seem to yield mere stable results than $\mu^2 = p_0^2$, and we shall use this factorization prescription

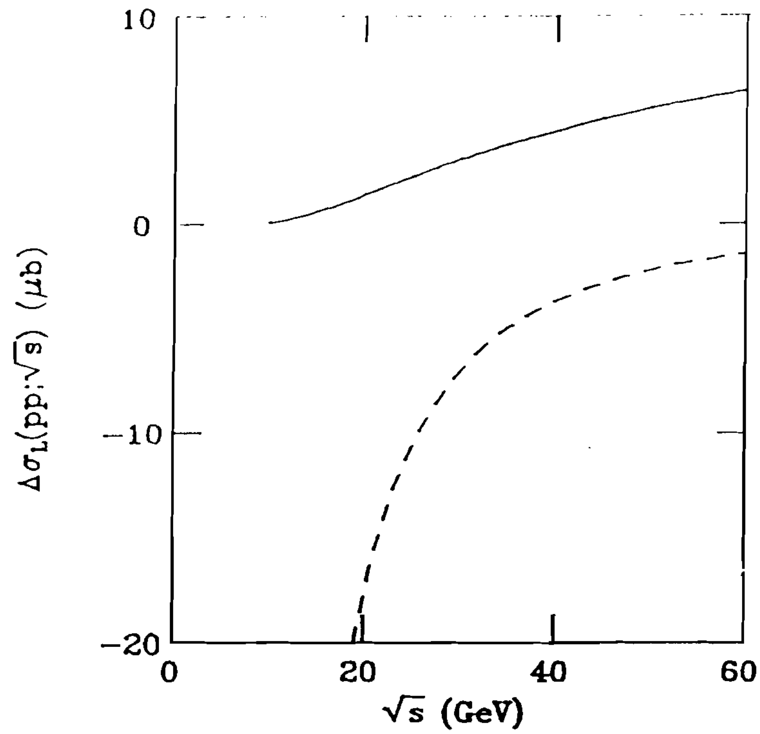

Fig. 2. The dashed line represents the expectation for $\Delta\sigma_L(pp;\sqrt{s})$ as a function of $\sqrt{s}$ using Eq. (1.8). The solid line is $\Delta\sigma_L^{\rm jet}(pp;p_0,\sqrt{s})$ at fixed cutoff $p_0^2 = 5$ GeV^2 and $\mu^2 = 10$ GeV^2 as a function of $\sqrt{s}$.

in the remainder of this section. A rather more informative, but equivalent, interpretation of this choice is that our numerical results should not depend strongly on the energy scale at which the starting distributions of Eq. (2.14) are defined.

Given this prescription choice, we plot the energy dependence of $\Delta\sigma_L^{\rm jet}(pp;p_0,\sqrt{s})$ for $p_0^2 = 5$ GeV^2 in Fig. 2 and compare it to the energy dependence for $\Delta\sigma_L(pp;\sqrt{s})$ predicted from the Regge pole fit in Eq.(1.8). As discussed above, the hard scattering contribution dominates for $\sqrt{s} > 40$ GeV. At $\sqrt{s} = 20$ GeV, where some data should be available within the next year,[11] the contributions to $\Delta\sigma_L^{\rm jet}(pp;p_0,\sqrt{s})$ are dominated by the valence quark subprocess. This is shown in Fig. 3 where the contributions from the different subprocesses are separated. In fact the valence quark contribution to the unpolarized jet cross section $\sigma_{\rm jet}(pp;p_0,\sqrt{s})$ is also large at this energy. Measurements of $\Delta\sigma_L^{\rm jet}(pp;p_0,\sqrt{s})$ in this energy range should give a good indication of the basic viability of the hard-scattering formalism involving polarized constituents. In addition, the behavior of the cross sections for a limited range of $\sqrt{s}$ and p_0 should give some idea whether the

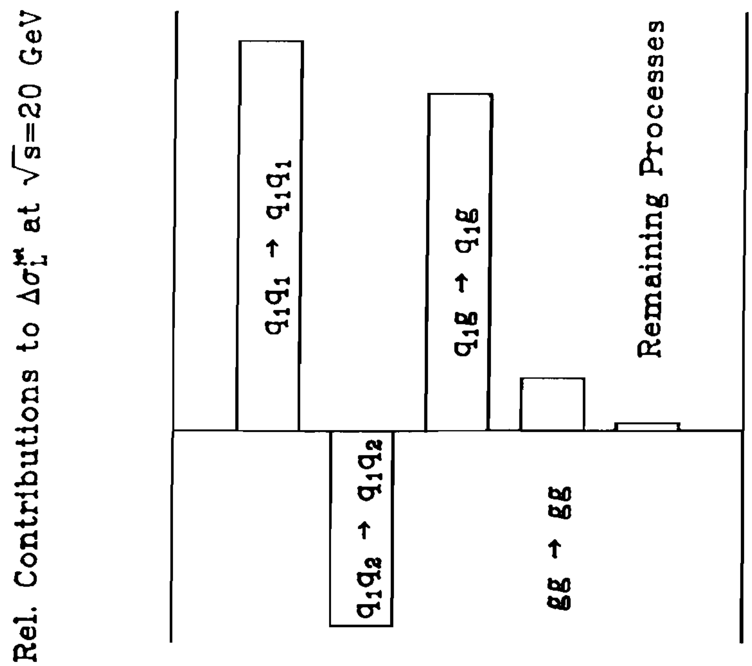

Fig. 3. The relative contributions of different processes to $\Delta\sigma_L^{\text{jet}}(pp; p_0, \sqrt{s})$ are shown at $p_0^2 = 5$ GeV^2, $\mu^2 = 10$ GeV^2 and $\sqrt{s} = 20$ GeV.

constituent distributions follow the expected pattern.

It would be very desirable if experiments with polarized hadron beams can be performed at significantly higher energies, comparable to the energies where the data[5,6] on unpolarized jet cross sections show the strong energy thresholds. At $\sqrt{s} = 100$ GeV, our estimate for $\Delta\sigma_L^{\text{jet}}$ is

$$\Delta\sigma_L^{\text{jet}}\left(pp; p_0^2 = 5 \text{ GeV}^2, \sqrt{s} = 100 \text{ GeV}\right) = 8.2\ \mu b \qquad \left[\mu^2 = 10 \text{ GeV}^2\right]. \tag{2.16}$$

This prediction, however, involves the extrapolation of the $\Delta G_{i/p}(x, \mu^2)$ into x-regions where they are not well known. The break-up of the cross section at this energy into the contribution of the various subprocesses is shown in Fig. 4. As can be seen there, most of the growth of the cross section in our calculation can be attributed to the gluon processes. Since our parameterization builds in the growth of $\langle s \rangle_x$ associated with small-x gluons expected in most models, this energy dependence is an important feature of the theory. Similarly the data for the unpolarized jet cross section is also dominated by the gluon processes at this energy.

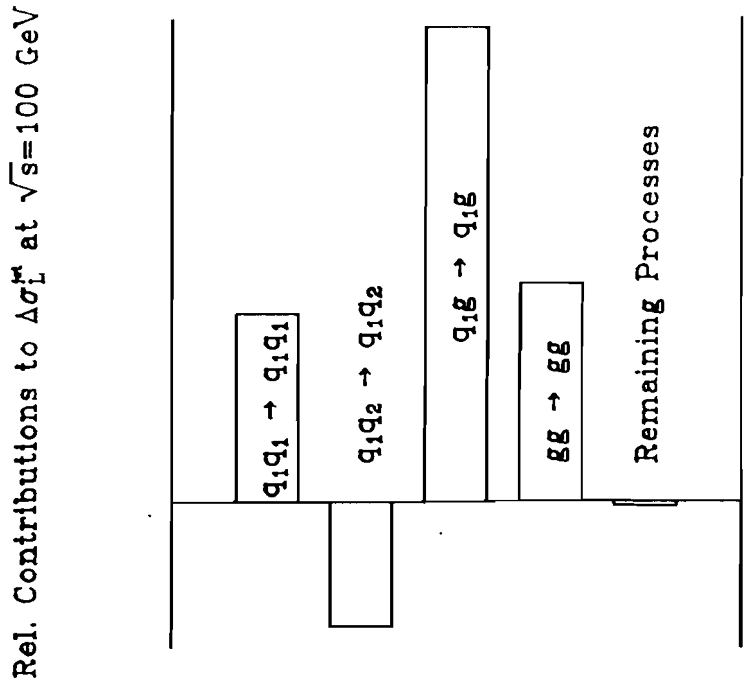

Fig. 4. The relative contributions of different processes to $\Delta\sigma_L^{\rm jet}(pp;p_0,\sqrt{s})$ are shown at $p_0^2 = 5$ GeV^2, $\mu^2 = 10$ GeV^2 and $\sqrt{s} = 100$ GeV.

III. Experimental Consequences and Discussion

Our numerical calculations above indicate that it should be possible to observe interesting structure associated with $\Delta\sigma_L^{\rm jet}(pp;p_0^2,s)$ in experiments using high energy polarized proton beams. The interpretation of our calculation involves some further discussion. Consider the hadronic multiplicity sum rule

$$\int \frac{d^3p}{E} A_{LL} \frac{Ed\sigma}{d^3p}(pp \to hX) = \int \frac{d^3p}{E}\left[\frac{Ed\sigma}{d^3p}(p(+)p(+) \to hX) - \frac{Ed\sigma}{d^3p}(p(+)p(-) \to hX)\right]$$
$$= \langle n_h^{++}\rangle \sigma\,(p(+)p(+);s) - \langle n_h^{+-}\rangle \sigma\,(p(+)p(-);s)\,. \tag{3.1}$$

We now write the average multiplicities in the form

$$\begin{aligned} \langle n_h^{++}\rangle &= \bar{n}_h + \langle \delta n_h\rangle \\ \langle n_h^{+-}\rangle &= \bar{n}_h - \langle \delta n_h\rangle, \end{aligned} \tag{3.2}$$

so that

$$\int \frac{d^3p}{E} A_{LL} \frac{Ed\sigma}{d^3p}(pp \to hX) = \bar{n}_h \Delta\sigma_L(pp;s) + \langle \delta n_h \rangle \sigma_{\rm inel}(pp;s) \tag{3.3}$$

Because of the conservation of transverse momentum, an event with a high-p_T jet will always have at least one other jet balancing it. Absorbing this constraint into the definition of jet multiplicities and using the QCD-parton model expression (2.1) gives, in analogy to (3.3)

$$\Delta\sigma_L^{\rm jet}(pp;p_0,\sqrt{s}) = \bar{n}_{\rm jet}(p_0,\sqrt{s})\Delta\sigma_L(pp;\sqrt{s}) + \left\langle \delta n_{\rm jet}(p_0,\sqrt{s})\right\rangle \sigma_{\rm inel}(pp;\sqrt{s}), \tag{3.4}$$

and

$$\Delta\sigma_L^{\rm soft}(pp;p_0,\sqrt{s}) = \left(1 - \bar{n}_{\rm jet}(p_0,\sqrt{s})\right) \Delta\sigma_L(pp;\sqrt{s}) - \left\langle \delta n_{\rm jet}(p_0;\sqrt{s})\right\rangle \sigma_{\rm inel}(pp;\sqrt{s}). \tag{3.5}$$

If we assume that our theoretical assumptions concerning *both* the small-x behavior of the distribution functions *and* the Regge behavior of $\Delta\sigma_L$ are correct, we have an energy regime where

$$\lim_{s\to\infty} \left|\Delta\sigma_L^{\rm jet}(pp;p_0,\sqrt{s})\right| \gg \left|\bar{n}_{\rm jet}(p_0,\sqrt{s})\Delta\sigma_L(pp,\sqrt{s})\right|,$$

so that

$$\begin{aligned} \lim_{s\to\infty} \Delta\sigma_L^{\rm jet}(pp;p_0,\sqrt{s}) &\cong \langle \delta n_{\rm jet}(p_0,\sqrt{s})\rangle \sigma_{\rm inel}(pp,\sqrt{s}), \\ \lim_{s\to\infty} \Delta\sigma_L^{\rm soft}(pp;p_0,\sqrt{s}) &\cong -\langle \delta n_{\rm jet}(p_0,\sqrt{s})\rangle \sigma_{\rm inel}(pp,\sqrt{s}). \end{aligned} \tag{3.6}$$

The perturbative expression, (2.6), has the form

$$\Delta\sigma_L^{\rm jet}(pp;p_0,\sqrt{s}) = C_1 \frac{s^{+J_{\rm eff}-1}}{(p_0^2)^{J_{\rm eff}}} \ln\left(C_2 \frac{s}{p_0^2}\right). \tag{3.7}$$

This quantity is positive and blows up as $p_0 \to 0$. However, for small p_0, we are not allowed to use the hard-scattering expression (2.1). For the hard scattering expression, (2.1), to be valid we require $s \gg p_0^2 \gg m^2$ where $m^2 \cong 1$ GeV^2 is a hadronic mass scale. For finite values of s, we therefore expect that

$$A_h(k_T) = \int \frac{d^3p}{E} A_{LL} \frac{Ed\sigma}{d^3p}(pp \to hX)\delta(p_T - k_T) \qquad (h = \pi, \ldots), \tag{3.8}$$

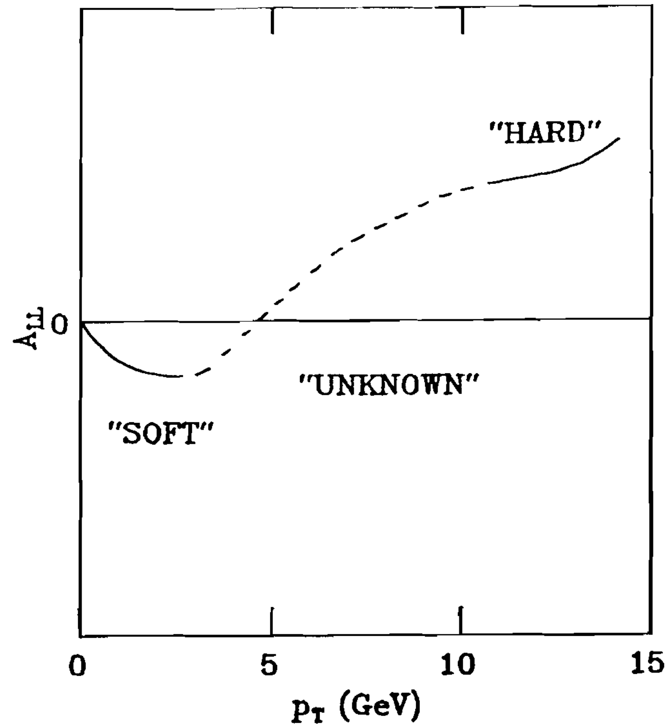

Fig. 5 The expected shape for the integrated spin asymmetry $A_{LL}(p_T)$ as a function of transverse momentum p_T.

should have the form sketched in Fig. 5. The negative region is associated with the negative value of $\Delta\sigma_L$ from A_1 exchange while the positive region corresponds to the onset of the hard scattering regime. The question of where the transition between soft and hard scattering takes place is potentially very interesting but cannot be answered at the present time.

Our discussion of $\Delta\sigma_L^{\text{jet}}$ has, so far, ignored the contribution of subasymptotic "higher-twist" contributions to the hard-scattering cross section. Since our estimates involve the extrapolation of structure functions to small x, we have to estimate the possible contribution of such terms. We can begin by writing the definite-helicity cross sections

$$\begin{aligned} \sigma\left(p(+)p(+) \to \text{jet}; p_0, \sqrt{s}\right) &= \sigma_0\left(p_0, \sqrt{s}\right)\left[1 + A_{++}\xi^{\epsilon}\left(1 + o(1)\right) + B_{++}\frac{m^2}{4p_0^2} + \cdots\right] \\ \sigma\left(p(+)p(-) \to \text{jet}; p_0, \sqrt{s}\right) &= \sigma_0\left(p_0, \sqrt{s}\right)\left[1 + A_{+-}\xi^{\epsilon}\left(1 + o(1)\right) + B_{+-}\frac{m^2}{4p_0^2} + \cdots\right], \end{aligned} \tag{3.9}$$

where $\sigma_0(p, \sqrt{s})$ is the "leading" contribution to the jet cross section. For fixed p_0 at high energy we expect that it will become approximately independent of s. The terms involving

$\xi = 4p_0^2/s$ are the spin-dependent effects we are attempting to calculate. The possibility of spin-dependent higher twist contributions is indicated by the terms proportional to $m^2/4p_0^2$. The requirement $s \gg 4p_0^2 \gg m^2$ mentioned above guarantees that both types of subleading terms are negligible in the unpolarized cross section. However, if we now consider the polarized case we have

$$\Delta\sigma_L^{\rm jet}(pp; p_0, \sqrt{s}) = \sigma_0 \left[\Delta A \xi^\epsilon + \Delta B \frac{m^2}{4p_0^2} + \cdots \right]. \tag{3.10}$$

Thus at fixed p_0, with $\epsilon > 0$, it is clear that the spin-dependent leading-twist calculation outlined above can be swamped by unknown higher-twist effects as s increases.

It is possible to avoid this problem by increasing the cutoff. For example, our expectation discussed in Sec. II would be that (3.10) is valid with $\epsilon \cong 1/2$. The condition for choosing the cutoff, is then:

$$p_0 > \frac{1}{2} \left(\frac{\Delta B}{\Delta A} \right)^{1/3} \sqrt{s}^{1/3} m^{2/3}. \tag{3.11}$$

The general condition

$$p_0 > \frac{1}{2} s^{\frac{\epsilon}{2+2\epsilon}} m^{\frac{2}{2+2\epsilon}} \left(\frac{\Delta B}{\Delta A} \right)^{\frac{1}{2+2\epsilon}}, \tag{3.12}$$

assures that we can take $\xi \to 0$ and stay away from possible higher twist effects. If we adopt the stringent condition

$$\left(4p_0^2\right)^2 = m^2 s, \tag{3.13}$$

which corresponds to $\epsilon = 1$ in (3.12) then we expect the higher twist contributions to be relatively more strongly suppressed as $\sqrt{s}$ increases. We can estimate the numerical impact. Figure 6 shows $\Delta\sigma_L^{\rm jet}(pp; p_0(\sqrt{s}), \sqrt{s})$ with $p_0 = \frac{1}{2} m^{1/2} \sqrt{s}^{1/2}$. It should be emphasized that this conservative approach may be necessary to evade problems with unknown dynamics but there is no indication that higher twist effects of the type indicated in (3.9) are necessary.

One final intriguing experimental possibility is that $\Delta\sigma_L^{\rm soft}(pp; p_0, \sqrt{s}) \ll \Delta\sigma_L^{\rm jet}(pp; p_0, \sqrt{s})$ for some choice of available p_0 and $\sqrt{s}$. This would be associated with the contribution

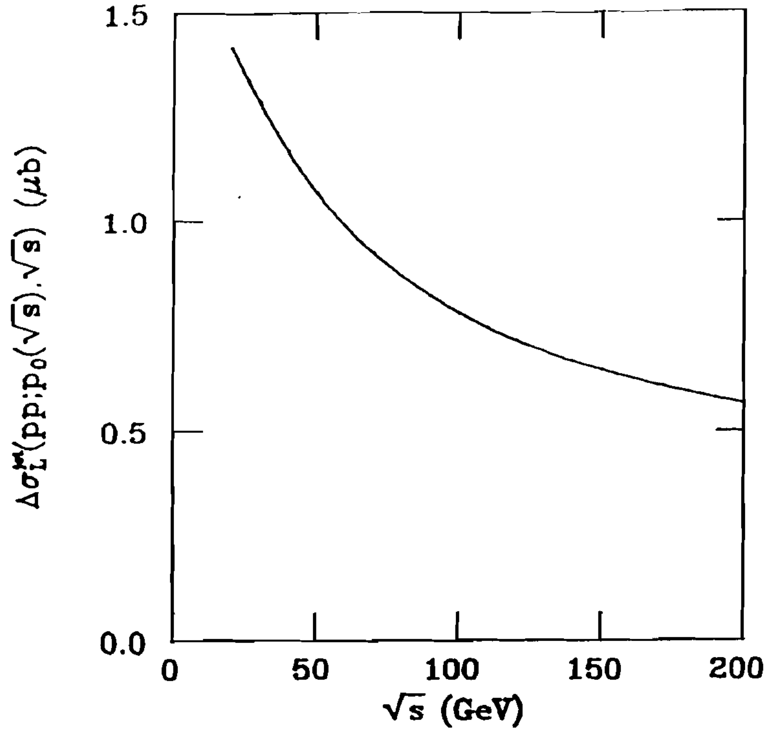

Fig. 6. We show $\Delta\sigma_L^{\rm jet}(pp; p_0(\sqrt{s}), \sqrt{s})$ as a function of $\sqrt{s}$ with p_0 given by Eq. (3.13) and $\mu^2 = 2p_0^2$.

of a "new" Regge contribution to $U_0(s,t)$. The interpretation of this singularity is uncertain. An indication that something like this might happen would be the measurement of $A_{LL}(pp \to \pi X) > 0$ for small p_T.

It should be kept in mind that all the models for reconciling the data for unpolarized jet physics with the asymptotic behavior of total cross sections, will have application to $\Delta\sigma_L$. Measurements of $\Delta\sigma_L^{\rm jet}$ can provide important checks of the various theoretical ideas which have been proposed. Furthermore this and other high p_T spin asymmetry measurements afford the best possibility of determining the helicity-weighted gluon distribution, a task that is essential if we are to unravel the spin structure of the proton. Both reasons provide a strong incentive for pursuing an experimental program with polarized beams and targets.

Acknowledgements

I am grateful to Gordon Ramsey and Dennis Sivers, with whom this work was performed. I would also like to acknowledge fruitful conversations with E. Berger, J. Collins, and Jian-Wei Qiu.

References

1. Gordon P. Ramsey, David Richards and Dennis Sivers, Phys. Rev. **D37**, 3140 (1988).

2. J. Ashman *et al.*, CERN preprint CERN-EP/87-230 EMC Collaboration.

3. S. J. Brodsky, J. Ellis and M. Karliner, Phys. Lett. **B206**, 309 (1988).

4. G. P. Ramsey, J.-W. Qiu, D. Richards and D. Sivers, Argonne preprint ANL-HEP-PR-88-16 (1988).

5. G. Arnison *et al.*, Phys. Lett. **118B**, 167 (1982); G. Lattes *et al.*, Phys. Rep.**65**, 151 (1980); A. Breakstone *et al.*, Phys. Rev. **D30**, 528 (1984); G. Bagnaia *et al.*, Phys. Lett. **B144**, 283 (1984).

6. G. Alpgard *et al.*, Phys. Lett. **121B**, 209 (1983); G. Alner *et al.*, Phys. Lett. **138B**, 304 (1984).

7. F. Halzen and F. Herzog, Phys. Rev. **D30**, 2326 (1984); T. Gaisser and F. Halzen, Phys. Rev. Lett. **54**, 1754 (1985).

8. G. Bodwin, Phys. Rev. **D32**, 2616 (1985); J. Collins, D. Soper and G. Sterman, Nucl. Phys. **B261**, 104 (1985).

9. For a discussion of the phenomenological implications of prescription dependence, see W. Celmaster and D. Sivers, Ann. Phys. (N.Y.) **143**, 1 (1982).

10. J. Babcock, E. Monsay and D. Sivers, Phys. Rev. Lett. 40, 1161 (1978); Phys. Rev. **D19**, 1483 (1979).

11. A. Yokosawa *et al.*, FNAL proposal #(E-704).

12. A. Yokosawa, ANL Preprint ANL-HEP-CP-87-110, Workshop on the Experimental Program at UNK, Protovino, USSR (1987) and references therein.

13. W. Grein and P. Kroll, Nucl. Phys. **B137**, 173 (1978); E. L. Berger, A. C. Irving, C. Sorensen, Phys. Rev. **D17**, 2971 (1978).

14. I. P. Auer *et al.*, Phys. Lett. **67B**, 113 (1977); Phys. Lett. **70B**, 475 (1977); Phys. Rev. Lett. **41**, 558 (1978).

15. See, for example, P. I. Chiapetta and J. Soffer, Phys. Rev. **D31**, 1019 (1985) and references therein; N. S. Craigie, K. Hidaka, M. Jacob and F. M. Renard, Phys. Reports **99**, 69 (1983) and references therein.

16. B. L. Combridge and C. J. Maxwell, Nucl. Phys. **B239**, 429 (1984); F. Halzen and P. Hoyer, Phys. Lett. **B130**, 326 (1983).

17. M. Einhorn and J. Soffer, Nuclear Physics **B274**, 714 (1986).

18. J. Collins, in "Proceedings of the 1984 DPF Summer Study on the Design and Utilization of the SSC", R. Donaldson and J. Morfin (eds.).

19. L. V. Gribov, E. M. Levin, M. G. Ryskin, Phys. Rep. **100**, 1 (1983).

20. L. Durand and Pi Hong, Phys. Rev. Lett. **58**, 303 (1987); P. L'Heureux, B. Margolis, P. Valin, Phys. Rev. **D32**, 1681 (1985).

21. J. Kwiecinski, Krakow preprint; L. Ametller and D. Trelani, Trieste preprint 49/87/EP.

22. J. Qiu, G. Ramsey, D. Richards and D. Sivers, in preparation.

23. P. Chiappetta, J. Ph. Guillet and J. Soffer, Nucl. Phys. **B262**, 187 (1985).

24. M. Gluck, E. Hoffman and E. Reya, Z. fur Phys. **C13**, 119 (1982).