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SPIN PHYSICS AND INCLUSIVE PROCESSES AT SHORT DISTANCES

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SPIN PHYSICS AND INCLUSIVE PROCESSES AT SHORT DISTANCES *

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TABLE OF CONTENTS

I. INTRODUCTION: WHY SPIN PHYSICS AT SHORT DISTANCES?

There are a large number of inclusive processes, each characterized by a single high frequency or short distance scale, for which there is considerable hope that they will be described within the framework of perturbative QCD and the parton model. One of the main conceptual and technical break throughs in this regard, was the so-called factorization theorem [1] for parton densities inside hadrons. However this subject is still beset with problems. One has to contend with large logarithmic corrections or even high power background mechanisms at the current range of momentum transfers. There are even some deep technical worries that something has been overlooked in the factorization proofs (see the recent works of Bodwin, Brodsky and Lepage [2] and independently Mueller [3]). With this in mind one might ask why even contemplate spin physics at short distances at this time. However there is a simple reason why it might be precisely the right thing to do. Namely, since the underlying interactions arise from a gauge theory, consequently helicity is propagated in a particularly simple way through the basic bremsstrahlung and quark or gluon pair creation process. This means typically hard processes occurring at short distances tend to have characteristically large helicity asymmetries. Further, we shall argue that the basic nature of the asymmetry (i.e. sign and magnitude) may be more reliably predicted than for example the overall normalization, thereby providing a clear test of the underlying QCD interactions.

The basic quark-gluon vector coupling is given by

$$\bar{q}\, A\cdot\gamma\, q \;=\; \bar{q}_L\, \gamma_\mu\, q_L\, A^\mu \;+\; \bar{q}_R\, \gamma_\mu\, q_R\, A^\mu \quad ,$$

i.e. the right- and left-handed fermions only can communicate through mass terms, which are proportional to $m\bar{q}_L q_R$. For u and d quarks these are small and can be neglected in leading order QCD calculations, so helicity is conserved in the basic interactions. One thus has two very simple rules, namely:

a) whenever quarks scatter through gluon exchange they conserve helicity,
b) when a quark and anti-quark annihilate through one gluon exchange (or photon exchange) they do so in opposite helicities.

In fact the situation is a little more complicated, but the basic observation is more or less correct.

In this talk we will argue that although spin physics at short distances (i.e. high momentum transfers) is difficult to do experimentally, it may still provide a valuable tool for learning how perturbative QCD works in the presence of the non-perturbative hadronic structure. It may also be a valuable way of seeing new interactions that might occur at very short distances.

Finally in this brief write-up, we have often skipped unimportant references. We therefore urge interested readers to look at the Physics Reports on "Spin Physics at Short Distances", by N.S. Craigie, M. Jacob, K. Hidaka and F.M. Renard, which will appear shortly.

II. FACTORIZATION THEOREM IN PERTURBATIVE QCD AND HELICITY ASYMMETRIES IN INCLUSIVE PROCESSES

If we consider an inclusive process involving hadrons A,B,... with momenta $p_A, p_B, \ldots$ respectively, then in the limit $p_A \cdot p_B$, $p_A \cdot p_C$, $p_B\, p_C \to \infty$ at the same rate, the process can be described by the parton model formula

$$\sigma^{AB\ldots}(p_A, p_B, \ldots) = \sum_{a,b\ldots} \int\int \cdots dx_a\, dx_b \cdots D_A^a(x_a)\, D_B^b(a_b) \cdots \sigma^{ab\ldots}(x_a p_A,\, x_b P_B, \ldots) \quad , \qquad (1)$$

where $D_A^a(x)$ is the probability density for finding a parton of type a in hadron A carrying fraction x of the latter momentum. $\sigma^{ab\ldots}$ is the corresponding cross-section at the parton level. Such a formula was proposed by Feynman, Bjorken and others [4], based on interpreting the process in terms of the underlying constituents of hadrons and assuming no interference at high momentum transfers. In a gauge theory like QCD, on the other hand, one knows that interfering diagrams like that in Fig.1(a) are in general on the same footing as non-interfering diagrams like that shown in Fig.1(b)

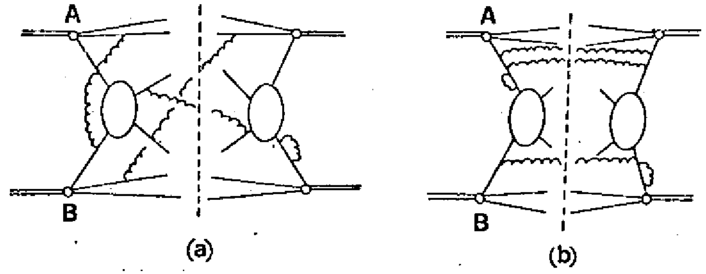


Fig.1: Typical diagrams in a gauge theory for a hard process. Diagrams of type (b) lead to a factorized parton model description of the process.

The remarkable thing that emerged between 1978-1980 was that by extensive use of the underlying Ward identities in a gauge theory or by the choice of a suitable non-covariant gauge ($\eta \cdot A = 0$) one could show that the effect of all interfering diagrams of the type in Fig.1(a) could be incorporated into factorizing diagrams like those in Fig.1(b) and that the latter lead to the parton model formula (1) with $D_A{}^a(x)$ replaced by $D_A{}^a(x,Q^2)$ and $\sigma^{ab\ldots}$ replaced by the lowest order Born cross-section evaluated at the running coupling constant $\alpha_S(Q^2) = 1/b \log Q^2/\Lambda^2$ ($b = (11-2n_f/3)/8\pi$ and Λ is the QCD mass scale parameter). Here Q^2 is some characteristic momentum transfer scale proportional to $p_A \cdot p_B$, $p_A \cdot p_C$, etc. In this form the formula holds to order $\alpha_s(Q^2)$, i.e. in the next-to-leading logarithmic order there are corrections which have been shown to be not small in some cases, depending on the renormalization scheme used and the way we define the parton densities. The latter are generally extracted from deep inelastic scattering, where it is perhaps natural to define them simply by expressing $F_2(x,Q^2)$ in the form

$$F_2(x,Q^2) = \sum_{f = u,d,\ldots} e_f^2\, D_f(x,Q^2) \quad . \qquad (2)$$

We have no room here to go into any of these matters, so we leave the discussion here, referring to the Physics Reports mentioned in the Introduction. The latter discusses these issues with particular regard to the generalization to spin asymmetries.

As regards spin asymmetries, we can summarize by saying, for states of definite helicity, all the leading order formulae apply without further assumptions and the predictions of QCD are on the same footing as for the spin averaged cross-section. For a helicity asymmetry the basic formula has as a numerator (following the notation of Eq.(1)) of the form

$$\Delta_{AB}\sigma^{AB\ldots} = \sum_{a,b} \int \ldots \Delta D_A^a \, \Delta D_B^b \ldots \Delta_{ab}\sigma^{ab\ldots} \qquad (3)$$

where $\Delta_{AB}\sigma^{AB\ldots} = \sigma^{A^+B^+} - \sigma^{A^+B^-}$ ($A^{\pm}_a$ referring to, respectively, hadron A in helicity state ±). ΔD^a_A refers to the parton flux with the same helicity as hadron A minus that with the opposite. The latter quantity can also be extracted from deep inelastic electroproduction experiments, this time using polarized electron or muon beams on a polarized proton target [5]. One can relate ΔD to the helicity structure function G_1, in fact a possible definition is

$$G_1(x,Q^2) = \sum_{f = u,d,\ldots} e_f^2 \, \Delta D^f(x,Q^2) \; .$$

(Note the Q^2 evolution of $D(x,Q^2)$ and $\Delta D(x,Q^2)$ is predicted in QCD perturbation theory through the Altarelli-Parisi-Lipatov evolution equations [6].)

These considerations do not simply carry over to transverse spin asymmetries, which is an altogether more complex issue. Part of the reason is that for helicity dependent or unpolarized processes the underlying parton process effectively involves partons which are zero mass and on-mass shell. The latter cannot transmit transverse spin. Transverse spin asymmetries are thus due to off-shell effects. We will return to the latter later on in the text.

As an illustration of the factorization in perturbative QCD theorem and its application to helicity asymmetries we end this section with two examples of immediate experimental interest:

a) Massive lepton pair production in PP and $\bar{P}P$ collisions

The basic subprocess in the Drell-Yan mechanism is $q\bar{q}$ annihilation into a lepton pair, through an intermediate massive photon. Since the quarks have to annihilate in opposite helicity states, the basic asymmetry $\Delta\sigma/\sigma = a_{LL}(\bar{q}q \rightarrow \ell^+\ell^-) = -1$. Hence if we consider for simplicity $x_F = 0$, then as a function of $x = \sqrt{Q^2/S}$, the asymmetry in $P\bar{P} \rightarrow \ell^+\ell^- + X$ is given by

$$A_{LL} = -\left[\frac{\Delta q_v(x)}{q_v(x)}\right]^2 \; , \qquad (4)$$

where $q_v = 4/9\, u(x) + 1/9\, d(x)$ (= valence quark distribution inside the proton). This leads to the prediction in Fig.2.

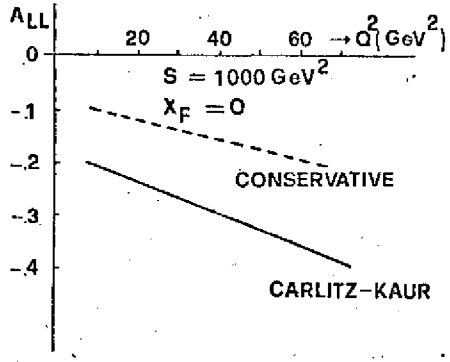


Fig.2: Predictions of Drell-Yan mechanism for asymmetry in $P\bar{P} \rightarrow \mu^+\mu^- + X$.

The model used for $\Delta q_v(x)$ was that due to Carlitz-Kauer [7], which broadly reproduces the recent SLAC measurements of $G_1(x)$ [8]. Other models can be constructed, for example, the conservative SU(6) model put forward by Sivers and collaborators [9] to represent another extreme. The main point made by the latter group was that Δq_v must be non-vanishing and in fact quite large, in order to satisfy the Bjorken sum rule [10]

$$\int_0^1 dx\, [\Delta u(x) - \Delta d(x) + \Delta\bar{u}(x) - \Delta\bar{d}(x)] = c_A/c_V \approx 1.23 \qquad (5)$$

Here we refer to the u, d,... quark distribution in the proton. One notices that the predicted asymmetry is very large indeed.

b) Lambda production at large p_T in PP collisions

This experiment offers the possibility of measuring both reflected (i.e. initial-initial) and transmitted (i.e. initial-final) spin asymmetries. The main point is that the Λ spin state can be reconstructed from its decay distribution $\Lambda \rightarrow P^+\pi^-$, i.e.

$$\Gamma(\theta) = \Gamma[1 + \beta \; \mathbb{P}\cdot\hat{n}_\Lambda]$$

where $\mathbb{P} = \Lambda$ polarization and β is a parameter associated with the decay amplitudes. Its value is of order unity, making a Λ polarization reconstruction fairly easy. For those who are not familiar with a transmitted helicity asymmetry measurement, the basic formula is given by

$$A_{LL}^{if} = \frac{\sigma(\vec{P} \rightarrow \vec{\Lambda}) - \sigma(\vec{P} \rightarrow \overleftarrow{\Lambda})}{\sigma(\vec{P} \rightarrow \vec{\Lambda}) - \sigma(\vec{P} \rightarrow \overleftarrow{\Lambda})}$$

$$= \beta_\Lambda \; \frac{\sigma(\text{forward } \pi) - \sigma(\text{backward } \pi)}{\sigma(\text{forward } \pi) + \sigma(\text{backward } \pi)} \; . \qquad (6)$$

In Eq.(6) $\sigma(\text{forward/backward } \pi)$ refers to the Λ cross-section for a proton of fixed helicity, reconstructed from Λ decays, with the pion moving forward/backward in the Λ rest frame with respect to the Λ direction in the laboratory frame.

The asymmetries for this process are given by the formula:

<u>for the reflected asymmetries</u>

$$A_{LL}^{ii} = \frac{\int_{\substack{\text{parton}\\\text{phase sp.}}} \Delta G(x_1)\ \Delta G(x_2)\ a_{LL}^{ii} \left(\frac{d\sigma}{dt}\right)_{gg \to s\bar{s}} D_{s \to \Lambda}(x_3)}{\int_{\substack{\text{parton}\\\text{phase sp.}}} G(x_1)\ G(x_2) \left(\frac{d\sigma}{dt}\right)_{gg \to s\bar{s}} D_{s \to \Lambda}(x_3)} \qquad (7)$$

<u>for the transmitted asymmetry</u>

$$A_{LL}^{if} = \frac{\int_{\substack{\text{parton}\\\text{phase sp.}}} \Delta G(x_1)\ G(x_2)\ a_{LL}^{if} \left(\frac{d\sigma}{dt}\right)_{gg \to s\bar{s}} \Delta D_{s \to \Lambda}(x_3)}{\int_{\substack{\text{parton}\\\text{phase sp.}}} G(x_1)\ G(x_2) \left(\frac{d\sigma}{dt}\right)_{gg \to s\bar{s}} D_{s \to \Lambda}(x_3)} \qquad (8)$$

The basic asymmetries are, respectively, $a_{LL}^{ii} = -1$ and $a_{LL}^{if}(gg \to s\bar{s}) = [\frac{1}{3}(u^2 + t^2)/tu - \frac{3}{4}(t-u)/s]/[\frac{1}{3}(u^2 + t^2)/tu + \frac{3}{2}(1 - tu/s^2) + \frac{3}{4}]$ which tends to +1 in both the forward and backward direction. In writing (7) and (8) we have ignored the $q\bar{q} \to s\bar{s}$ background, since this may contribute at small $x_T = 2p_T/\sqrt{s}$. We assume the Λ trigger clearly picks out the above mechanisms. However this could be improved by selecting away side jets with higher than normal strange particle content. In order to make predictions we need $\Delta G(x)$, which is not known. In fact we shall argue later that this process might be a good way of measuring it. However it is reasonable to suppose that the hard component of the gluon spectrum comes from bremsstrahlung of colinear gluons off the valence quarks. If we use the leading quark model of Carlitz and Kaur, then the helicity of the proton is completely transmitted to the gluon as $x \to 1$, i.e. $\Delta G(x)/G(x) \to \Delta u(x)/u(x) \to 1$ as $x \to 1$. We call this model the LQGB (leading-quark gluon bremsstrahlung) model. The leading order QCD predictions are shown in Fig.3,

III. EFFECT OF HIGHER ORDER RADIATIVE CORRECTIONS ON HELICITY ASYMMETRIES

From studies of radiative corrections to Drell-Yan mechanism [11] and jet production cross-section in hadron collisions [12] it has become clear that there are large order $\alpha_s(Q^2)$ corrections at currently obtainable momentum transfers. It has therefore become unimportant to reconsider the asymmetry predictions. This has recently been carried through by Ratcliffe [13] in the case of the Drell-Yan mechanism and we summarize his finding here. One defines the parton distributions to all orders in perturbation by their relations to the deep inelastic structure functions, i.e.

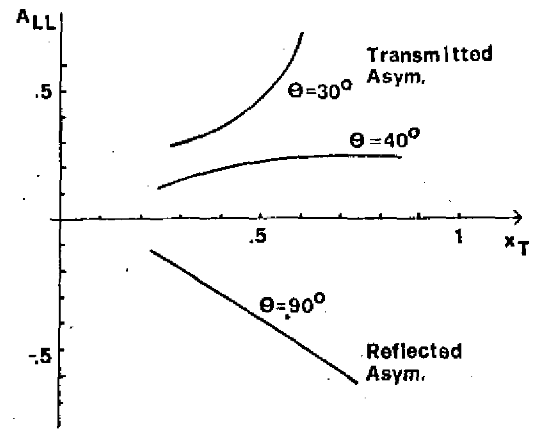


Fig.3: Helicity asymmetries for $PP \to \Lambda X$ at large p_T using LQBM model. We notice reflected asymmetry is proportional to $\Delta G(x)/G(x)$ and shows no angular dependence in contrast to the transmitted case, where the increase as θ decreases reflects helicity conservation of the fundamental interactions.

$$F_2(x,Q^2) = \sum_f e_f^2\, D^f(x,Q^2)$$

and

$$G_1(x,Q^2) = \sum_f e_f^2\, \Delta D^f(x,Q^2) \quad .$$

One therefore examines the first order radiative corrections to both deep inelastic scattering and massive lepton pair production. After factoring out the functions $D^f(x,Q^2)$, one notices that there are large order α_s corrections to the Drell-Yan mechanism and these corrections can be crudely taken into account by defining a K factor, i.e. $d\sigma/dQ^2 = K\,(d\sigma/dQ^2)_{L.O.}$ where K is between 2 and 3 and $(d\sigma/dQ^2)_{L.O.}$ is the Drell-Yan mechanism with the above definition of the parton densitites. However it is known that the large correlations leading to the factor of 2 come from ultra soft gluons, which are likely to have little effect on the helicity dependence as the analysis of Ratcliffe shows. To be more precise the K factor for $\Delta d\sigma/dQ^2$ is the same as for $d\sigma/dQ^2$ and consequently drops out of the asymmetry formulae.

We refer for details to Ref.13 and the Physics Reports mentioned in the introduction. It would be valuable to try to prove the following conjecture. Even though one has large order α_s corrections to the basic Born amplitudes of hard processes, these cancel out in the calculations of the asymmetries, consequently the leading order calculations for the latter are more reliable than the calculation of the overall normalization.

IV. HIGHER ORDER POWER MECHANISM AND SPIN ASYMMETRIES

In the Wilson operator product expansion, for example in the case of the product of two currents in deep inelastic e-p scattering, the operators are divided up according to their dimension, since higher dimensions correspond to higher inverse powers of the momentum transfer variable Q^2 in the asymptotic behaviour of the structure functions. Such power corrections exist for any inclusive process at short distances and following some observations of Politzer [14], there is hope that the corresponding cross-sections factorize in some sense. New parton densities, like the diquark distribution enter into this factorization. Whether or not these play the same universal role as the single parton densities remains to be seen. Different processes may probe the diquark structure differently. Nevertheless, it is interesting to note that, given a definite mechanism leading to a higher power scaling behaviour, then the corresponding helicity asymmetry can also be estimated, furthermore, some studies [15] have shown that the latter may be a valuable hall mark for the specific mechanism. Indeed this is the basic point we try to make in this section.

Let us again consider the reaction $PP \to \Lambda X$ and consider the Λ being formed by a gluon from one parton creating a $s\bar{s}$ pair and the strange quark picking up a diquark from the other proton to form the Λ state (Fig.4).

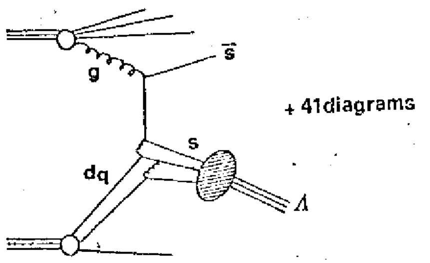


Fig.4: Higher power mechanism for hyperon production in PP collisions.

The calculation of the mechanism is too involved to enter into here, and we refer for details to Ref.15, or the Physics Reports mentioned in the introduction. We can summarize the results obtained in the reference, by saying that the initial diquark-strange quark system are prepared for the Λ bound state (as described by its asymptotic Bethe-Salpeter wave function) by a gluon exchange process. This leads to a form factor structure involving the Λ asymptotic wave function and a hard scattering subprocess $g\, dq \to \Lambda + s$, where all the particles are treated as elementary. One can write the full amplitude for $g\, dq \to \Lambda\, s$ in the form

$$T_{gdq \to \Lambda\bar{s}} = F_1\, T^{(1)}_{gdq \to \Lambda\bar{s}} + F_2\, T^{(2)}_{gdq \to \Lambda\bar{s}} \qquad (9)$$

where

$$\hat{T}^{(1)} = \frac{1}{\hat{u}}\, \bar{u}_\Lambda(k_4,h)\, [k_{1\mu}\, \epsilon_\nu(\lambda) - k_{2\nu}\, \epsilon_\mu(\lambda)]\, k_3^\mu\, \gamma^\nu\, v_{\bar{s}}(k_3)$$

$$\hat{T}^{(2)} = \frac{1}{\hat{t}}\, \bar{u}_\Lambda(k_4,h)\, \gamma^\mu\, k_4^\nu [k_{1\mu}\, \epsilon_\nu(\lambda) - k_{2\nu}\, \epsilon_\mu(\lambda)]\, v_{\bar{s}}(k_3)$$

and

$$F_1 = \frac{16\pi^2 \alpha_S^2(\hat{t})}{\hat{u}\cdot\hat{t}} \int [dy] \left\{ \phi_\Lambda^{+-+}(y_1,y_2,y_3)\, V_i^{+-}(x_1,x_2,y) + \ldots \right\} ,$$

$V_i^{hh'}$ are hard scattering kernels like those described by Lepage and Brodsky in Ref.16. The transmitted asymmetries given by the two amplitudes are ± 1. However, the form factor $F_1 \gg F_2$, for much of the kinematic region, so the overall asymmetry is negative.

If we assume the diquark system is spin zero, then there is no reflected (i.e. initial-initial) asymmetry. Further, with some further assumptions about generalized factorization in the way suggested by Politzer in Ref.14, the transmitted asymmetry is given by the formula

$$A_{LL}^{if} = - \frac{\int_{\substack{\text{parton}\\ \text{phase sp.}}} G(x)\, D_{dq}(x_1,x_2)\, \{|F_1|^2 - |F_2|^2\}\, \hat{s} + g \rightleftarrows dq}{\int_{\substack{\text{parton}\\ \text{phase sp.}}} G(x)\, D_{dq}(x_1,x_2)\, \{|F_1|^2 + |F_2|^2\}\, \hat{s} + g \rightleftarrows dq} \qquad (10)$$

where the parton phase space is given by

$$\int dx \int dx_1\, dx_2\, \theta(1-x_1-x_2)\, \delta(\hat{s}+\hat{t}+\hat{u}) \quad .$$

From Eq.(10) one sees many of the uncertainties in normalization cancel out, so that one can make a fairly straightforward asymmetry prediction. In Fig.5 we compare the asymmetry predicted by this mechanism with the leading order prediction, which is due to the mechanism $gg \to s\bar{s}$. One notices that both mechanisms predict large asymmetries, however with opposite sign and quite a diffferent shape. One common feature however is the increase of the transmitted asymmetry as we go to the forward direction (i.e. smaller trigger angles). This is simply a reflection of helicity conservation and propagation in the underlying gluon bremsstrahlung and pair creation mechanisms.

We also note that the reflected (initial-initial) asymmetry is zero for this higher power mechanism, based on the assumption that the diquark system is predominantly spin zero.

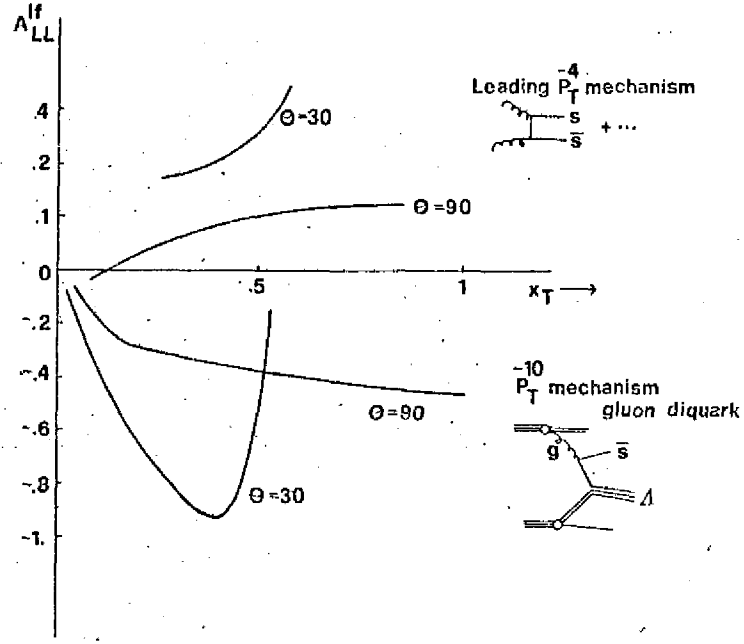


Fig.5: Transmitted asymmetries for $\vec{P}P \rightarrow \vec{\Lambda}X$ at large p_T ($s = 1000\ GeV^2$).

V. DIFFICULTIES IN UNDERSTANDING TRANSVERSE SPIN IN PERTURBATIVE QCD

Transverse spin asymmetries appear to be difficult to interpret in the parton model way, except perhaps for very massive quarks. For zero mass quarks the naive parton model would predict zero for any quantity involving quark carrying transverse spin. However, from the analysis of graphs or the Wilson OPE, one sees that even in a massless theory there can be transverse spin asymmetries due to off-shell effects. Recently, Phil Ratcliffe from Trieste has been examining the question of factorization in the chiral limit (i.e. massless quarks) and I briefly summarize his finding so far. (This item in the talk is more orientated for theorists interested in perturbative QCD.)

In the case of the leading scaling behaviour in deep inelastic scattering, which is associated with twist 2 operators of the form $O_n^{(2)} = \bar{\psi}\,\gamma_\mu (z\cdot D)^n\,\psi$, where $D_\mu = \partial_\mu - g\,A_\mu$, there is a one-to-one correspondence between the OPE renormalization group calculations and the radiative corrections to the Born diagrams. Furthermore, this

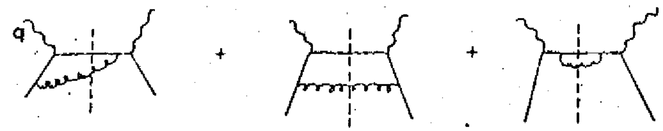


Fig.6: Radiative corrections to the Born terms in deep inelastic scattering.

correspondence persists in any gauge. In particular, the graphs in Fig.6 generate a term $\gamma_n\,\alpha_S \log q^2/\mu^2$; where γ_n is the anomalous dimension associated with the twist 2 operator above, i.e.

$$\gamma_n = \mu \frac{d}{d\mu}\ \mathrm{Ln}\ Z_n \qquad \text{where} \qquad O_n^{(2)}\Big|^{Ren} = Z_n\, O_n^{(2)}\Big|^{bare} .$$

However with the transverse spin structure function in deep inelastic scattering, which is associated with the twist 3 operator $O_n^{(3)} = \bar{\psi}\,\gamma_5\,\gamma_\lambda\,D_\sigma (z\cdot D)^n\,\psi$, the correspondence appears to be lost in the axial gauge. Indeed for the latter gauge, the Bloch-Nordsieck cancellation which is essential for the understanding of the twist 2 case, is not evident for this operator. In order to try to see what was going on and to see if spector interactions are the solution, in what after all is an effect due to virtual quarks inside a hadronic bound state, Ratcliffe considered a model involving an elementary proton-diquark coupling. On adding the set of diagrams in Fig.7

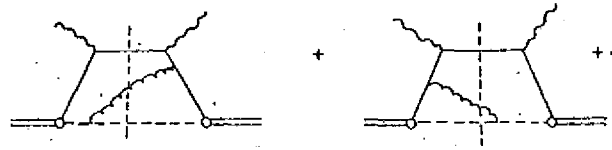


Fig.7: Spectator interactions playing a role in transverse spin asymmetries in chiral limit (i.e. zero quark mass limit).

he found that he could get the same result in all gauges. However even in the Feynman gauge these graphs generated a term $\gamma_n^{spec.}\,\alpha_S \log q^2/\mu^2$ in addition to the terms generated by graphs of the type in Fig.6. The conclusion is obvious, however I am still reluctant to draw it, since the work is still in progress and there could be cancellations. Nevertheless, it is clear from the study so far, that transverse spin is far more subtle than helicity asymmetries and it is not at all obvious that one has a factorization theorem.

This should not deplete the experimental interest in the subject. In fact I would like to end this discussion of transverse spin asymmetries by considering the remarkable results on final state hyperon polarizations in $PP \to Y + X$ at p_T now up to 4 GeV/c. [17]

. Single spin asymmetry in hyperproduction at large p_T

It appears even for unpolarized protons, that hyperons produced at transverse momentum between 1 and 4 GeV have a very large polarization asymmetry with respect to the production plane. For Λ and Ξ this asymmetry reaches + 20% and in the case of Σ it is negative with the same magnitude. In the case of Λ production this effect persists to the highest transverse momentum measured [17], namely $p_T = 4$ GeV/c. There is at present no satisfactory theoretical explanation of this effect (i.e. no one can perform a well defined computation to obtain this asymmetry). However, there are some interesting conjectures or speculations involving the notions of colour flux strings [18], Thomas-Fermi precession [19] and even resonances [20]. In order to put the matter in some kind of theoretical prospective, let us recall a simple formula for calculating single spin asymmetries

$$\sigma(\uparrow) - \sigma(\downarrow) = 2 \sum_X \mathrm{Im}\,(A^{(+)}_{PP\to YX}\, A^{(-)*}_{PP\to YX}) \;,$$

where $(\pm)$ refers to the hyperon Y helicity. Now it is very likely if one dreams up a mechanism for producing Λ's, that the same mechanism at high energies and large p_T will produce excited hyperon states with the same cross-section. The latter hyperons will decay down into Λ's. Hence the cross-section for indirect Λ production is likely to be considerably larger than that for direct production. Hence, if we write

$$A^{(\pm)}_{PP\to YX} = \sum_{h^*} A^{(h^*)}_{PP\to Y^*X}\; B^{(h^*,\pm)}$$

where $B^{(h^*,h)}$ describes decay process $B^*(h^*) \to B(h) + \ldots$

$$\sigma(\downarrow) - \sigma(\uparrow) = 2 \sum_X \mathrm{Im}\,(A^{(+)}_{PP\to Y^*X}\, A^{(-)+}_{PP\to Y^*X}) \sum_h |B^{(+h)}|^2$$
$$+ \sum_{h^* = \pm} |A^{(h^*)}_{PP\to Y^*X}|^2\; 2\cdot \mathrm{Im}(B^{(h+)}\, B^{(h-)*}) \;.$$

Hence this asymmetry can be generated by the production mechanism or the decay mechanism. Concerning the latter one has in any case to take into account all the overlapping Σ^* resonances, which decay in the channel $\Lambda\pi$. A regards the production mechanism, we need a mechanism to flip helicity on a quark line. Consider the following simple Gedanken mechanism for Λ production shown in Fig.8.

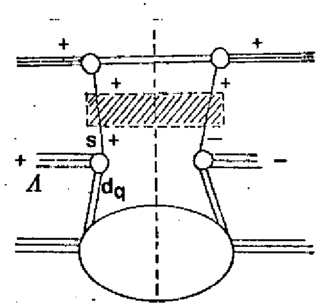
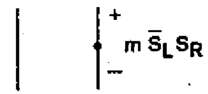
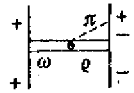


Fig.8: Gedanken mechanism for production transverse spin asymmetry in Λ production

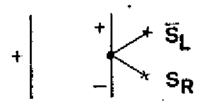
The black box in Fig.8 could involve a number of mechanisms. We list a few below:

1 

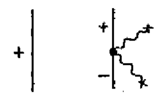
effect of quark mass

2 

$\rho\pi$-ω interference

3 

effect of chiral condensate

4 

effect of instanton gas or monopole antimonopole vacuum state of QCD .

The problem is that we do not yet know how to incorporate into a well-defined computational scheme any of these effects. It would be therefore even more challenging if the initial state polarization asymmetry is measured, with a polarized and/or target. We should in this way see if the effect is associated with the basic interaction of the partons, rather than some subtle final state interaction of decaying hyperons.

VI. HELICITY ASYMMETRIES OF SHORT DISTANCE AS A MEANS OF RECOGNIZING SUPERSYMMETRIC INTERACTIONS

Recently there has been much interest in supersymmetry (SuSy) grand unified theories (GUT's), since in the latter one can avoid the gauge hierarchy problem [21]. In particular, if a supersymmetric version of SU(5) is broken down to $SU(3)_{col} \otimes U(1)$, then we will have as a low energy theory supersymmetric QCD. In order for this to solve the gauge hierarchy problem, supersymmetry must itself be broken at relatively low masses. One can contemplate gluinos (the fermion SuSy partner of the gluon) as low as 10 GeV, which puts the excitation of this degree of freedom within the reach of present day accelerators. An ideal experiment to see SuSy strong interactions is PP collisions with polarized colliding beams. One studies the calorimeter triggered reaction

$$\vec{P}\vec{P} \longrightarrow \text{jets} + X$$

at high energies and transverse momentum. These triggers provide a valuable means of singling out interactions involving gluino production. This is illustrated in Fig.9(a).

Before turning to a particular experiment, let us briefly describe supersymmetric QCD by looking at its Lagrangian

$$\mathcal{L} = -\frac{1}{4} F^{a}_{\mu\nu} F^{\mu\nu a} \qquad \text{ordinary QCD}$$

$$= -\frac{1}{4} F^{a}_{\mu\nu} F^{\mu\nu a} + \bar{\lambda}^{a} \not{D}_{ab} \lambda^{b} \qquad N = 1 \text{ SuSy QCD}$$

$$= -\frac{1}{4} F^{a}_{\mu\nu} F^{\mu\nu a} + \sum_{i} \bar{\lambda}^{a}_{i} \not{D}_{ab} \lambda^{a}_{i} + (D^{\mu}\phi^{a}_{i})^{+} (D_{\mu}\phi^{a}_{i}) + \dots \qquad N \geqslant 2 \text{ SuSy QCD}$$

where $D_\mu = \delta_{ab}\partial_\mu - g f_{abc} A^c_\mu$ and all particles are in the adjoint representation of SU(3). The spectrum is given in the following table:

gluons	gluinos	scalars	
$g^{\pm}$	0	0	N = 0
$g^{\pm}$	$\lambda^{\pm}$	0	N = 1
$g^{\pm}$	$\lambda^{\pm}_1, \lambda^{\pm}_2$	$\phi_1, \phi^*_1, \dots$	N ⩾ 2

The basic amplitudes squared at the symmetry point $t = u$ are given by

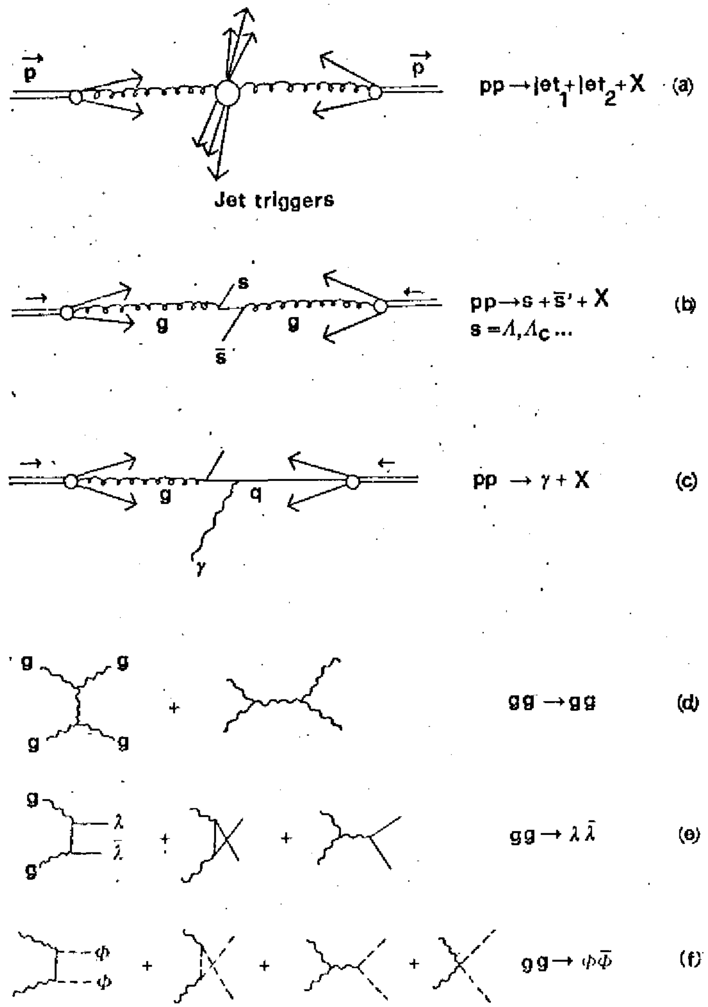


Fig.9: The fundamental hard processes in PP → hadrons at large transverse momentum (λ and φ refer to, respectively, gluinos and SuSy scalar partners in extended SuSy QCD).

$$|M_{hh'}|^2 = \frac{9}{2}\left[\left(3 - \frac{su}{t^2} - \frac{st}{u^2} - \frac{ut}{s^2}\right) + \left(-3 + \frac{2s^2}{ut} + \frac{ut}{s^2}\right) hh'\right] \qquad gg \to gg$$

$$|M_{hh'}|^2 = \frac{9}{4}\left[\frac{t^2 + u^2}{tu} - \frac{t^2 + u^2}{s^2}\right][1-hh'] \qquad gg \to \lambda\bar{\lambda}$$

From this we see that the asymmetries are significantly different and thus provide a valuable signal for supersymmetric interactions.

Before one can extract the asymmetry at the basic gluon-gluon interaction level one needs to calibrate the gluon beam, i.e. measure $\Delta G(x)/G(x)$. This can be done by studying strange or charm production, Fig.9(b) and/or prompt photon production, Fig.9(c). In the three cases, namely, jets, strangeness or photon production, the measured asymmetries are given, respectively, by

$$A_{LL}(\vec{P}\vec{P} \to \text{jets } X) = \left\langle \frac{\Delta G}{G} \frac{\Delta G}{G} a_{LL}(gg \to gg) \right\rangle$$

$$A_{LL}(\vec{P}\vec{P} \to s\bar{s}\ X) = \left\langle \frac{\Delta G}{G} \frac{\Delta G}{G} a_{LL}(gg \to s\bar{s}) \right\rangle$$

$$A_{LL}(\vec{P}\vec{P} \to \gamma\ X) = \left\langle \frac{\Delta G}{G} \frac{\Delta q}{q} a_{LL}(gq \to \gamma q) \right\rangle .$$

In each case the basic driving asymmetry is known. However, in the first case above the gluino threshold, the asymmetry will be different than that predicted by ordinary QCD. Hence the same experiment can be used to extract $\Delta G(x)/G(x)$ in the desired momentum transfer region. Now if we assume gluino jets lead to the same energy flow as for gluons and we are working at momentum transfers $Q^2 = 4 p_T^2 \gg m^2_{gluino}$, then the asymmetry changes as we go from ordinary QCD (N = 0) to the N = 1, N = 2, etc. SuSy QCD. In particular, if we take the symmetry point $\hat{t} = \hat{u}$, then

$$a_{LL}^{N=0} = 0.8, \quad a_{LL}^{N=1} = 0.6, \quad a_{LL}^{N=2} = +0.3 \quad \text{and} \quad a_{LL}^{N=4} = +0.08$$

respectively. However, the gluino must produce a new kind of stable hadron which decays only by, for example, supersymmetric electromagnetic interactions. In particular, Kane and Leveille [22] suggest the following decay mode $\lambda \to \tilde{\gamma} + q\bar{q}$, where $\tilde{\gamma}$ is a photino. The latter is the SuSy partner to the photon, which is almost inert and will surely lead to missing energy in the calorimeters. This has one immediate effect of providing a trigger bias against the gluino production in the above calorimeter experiment so the asymmetry distortion effect will be smaller. However one could use unbalanced transverse energy triggers to detect events involving gluino production as suggested by Aronson et al. from Brookhaven [23]. In this case the asymmetry should provide a clear check that one indeed has selected gluino production events. A detailed analysis of this will be published elsewhere, including a discussion on how to remove the soft KNO event background in large E_T triggers. We hope however to have illustrated the potential value of asymmetry measurements in discovering the presence of new kinds of interactions.

ACKNOWLEDGMENTS

I wish to acknowledge the hospitality of Brookhaven and Fermi National Laboratories, where this talk was written up. I would also like to thank Frank Paige and Phil Ratcliffe for discussions in preparing this talk.

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SPIN PHYSICS AND INCLUSIVE PROCESSES AT SHORT DISTANCES

N.S. Craigie

E R R A T A

On page 7 the sentence immediately after Eq.(8) should read:

The basic asymmetries are, respectively, $a_{LL}^{ii} = -1$ and $a_{LL}^{if}(gg \to s\bar{s}) = (u^2 - t^2)/(u^2 + t^2)$ which varies from +1 to -1 as we go from the forward to the backward direction........

Figs.3 and 5 should be replaced by the following:

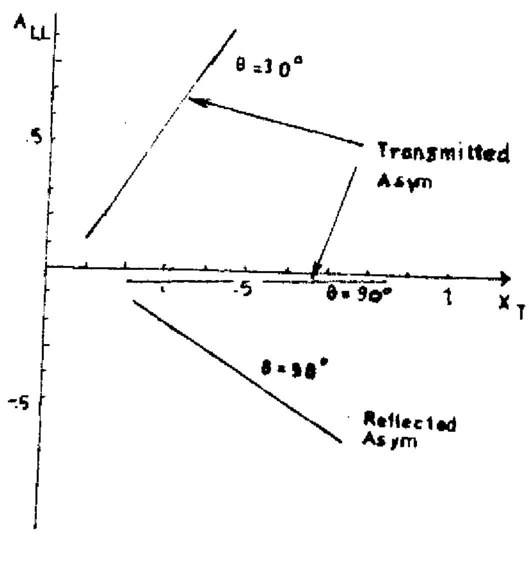

Fig.3: Helicity asymmetries for $PP \to \Lambda X$ at large p_T using LQBM model. We notice reflected asymmetry is proportional to $\Delta G(x)/G(x)$ and shows no angular dependence in contrast to the transmitted case, where the increase as θ decreases reflects helicity conservation of the fundamental interactions.

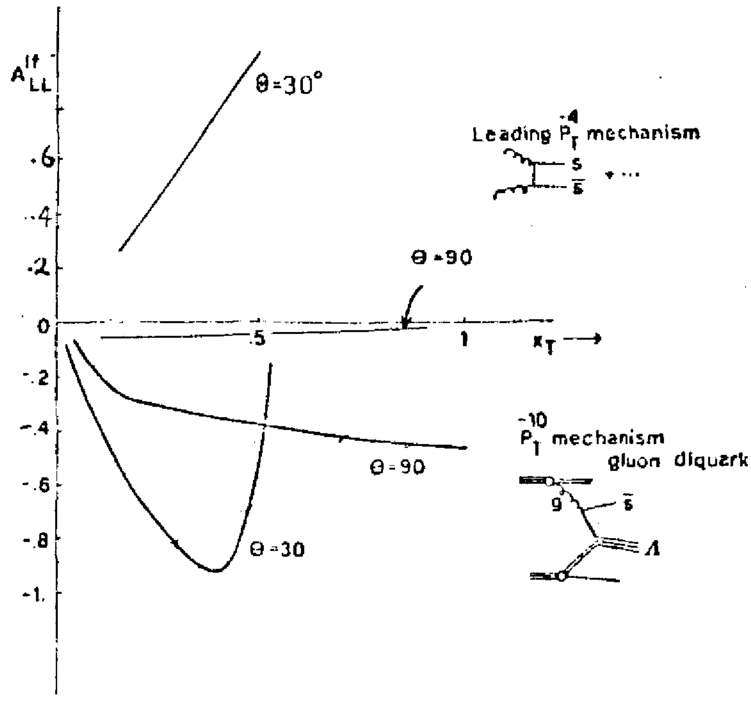

Fig.5: Transmitted asymmetries for $\vec{P}P \to \vec{\Lambda}X$ at large p_T ($s = 1000$ GeV^2).