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NOVEL SPIN EFFECTS IN QUANTUM CHROMODYNAMICS*

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ABSTRACT

I discuss a number of interesting hadronic spin effects which test fundamental features of perturbative and non-perturbative QCD. These include constraints on the shape and normalization of the polarized quark and gluon structure functions of the proton; the principle of hadron helicity retention in high x_F inclusive reactions; predictions based on total hadron helicity conservation in high momentum transfer exclusive reactions; the dependence of nuclear structure functions and shadowing on virtual photon polarization; and general constraints on the magnetic moment of hadrons. I also will discuss the implications of several measurements which are in striking conflict with leading-twist perturbative QCD predictions, such as the extraordinarily large spin correlation A_{NN} observed in large angle proton-proton scattering, the anomalously large $\rho\pi$ branching ratio of the J/ψ, and the rapidly changing polarization dependence of both J/ψ and continuum lepton pair hadroproduction observed at large x_F.

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1. Introduction

Polarization effects and spin correlations often provide the most sensitive tests of the underlying structure and dynamics of hadrons. The basic measures of the spin structure of the proton are its magnetic moment, which gives a global measure of the quark spin content of the proton, and the spin-dependent structure functions, which register the local distribution of the quark helicity currents as a function of their light-cone momentum fraction x [1]. The SLAC-Yale and EMC measurements show a strong positive helicity correlation between the helicity of the $u-$ and $\overline{u}$ quarks with that of the proton; the helicity of the $d-$ and $s-$quark and antiquarks are negatively correlated. Most remarkably, the net correlation of the quark plus antiquark helicity with that of the proton Δq is consistent with zero. Since the total spin projection $\frac{1}{2}\Delta q + \Delta g + L_z = \frac{1}{2}$, there must be a significant fraction of Fock states in the proton containing gluons, and there must be a non-trivial correlation of the gluon helicity with that of the proton.

Although the net correlation of the quark helicity with the proton helicity in inclusive reactions is small, the spin correlations of large angle elastic pp scattering nevertheless display a dramatic structure at the highest measured energies $\sqrt{s} \sim 5\ GeV$ [2]. These measurements are in strong conflict with the expectations of perturbative QCD which predicts a smooth power-law fall-off for exclusive helicity amplitude with increasing momentum transfer [3]. The strong polarization correlations observed in pp scattering are clearly of fundamental interest, since the microscopic QCD mechanisms that underlie the spin correlations between the incident and final hadrons must involve the coherent transfer of helicity information through their common quark and gluon constituents. The implications of the spin correlation measurements will be discussed in sections 5 and 6.

In this talk I shall emphasize a basic but non-trivial prediction of the gauge couplings of PQCD, "hadron helicity retention": a projectile hadron tends to transfer its helicity to its leading particle fragments. A particularly interesting consequence is the prediction that the J/ψ and the continuum lepton pairs produced in pion-nucleus collisions will be longitudinally polarized at large x_F. Helicity retention also provides important constraints on the shape of the gluon and quark helicity distributions. In the large x_F domain, with $Q^2(1-x)$ fixed, leading twist and multi-parton higher twist processes can be of equal importance [4]. In the case of large momentum transfer exclusive reactions, the underlying chiral structure of perturbative QCD predicts that sum of hadron helicities in the initial state must equal that of the final state [5]. Although hadron helicity conservation appears to be empirically satisfied in most reactions, the most interesting cases are its dramatic failures such as the large branching ratio for $J/\psi \rightarrow \rho\pi$. I will discuss these predictions and their experimental tests in section 7.

Although most of the topics discussed in this talk are concerned with quark or gluon helicity, there are also interesting linear polarization predicted by the theory, such as in Υ decays, or in the planar correlations of four-jet events in e^+e^- annihilation. In addition, the oblateness [6] of a gluon jet can be used to determine its axis of linear polarization.

2. The Magnetic Moment of Hadrons in QCD

Much of our understanding of the helicity structure of hadrons comes from rigorous constraints, such as the Bjorken Sum Rule for the integral of the spin dependent structure functions, and the Drell-Hearn-Gerasimov sum rule, which relates the anomalous magnetic moment of a composite system to an integral over the photoabsorption cross section. In fact Burkert and Ioffe [7] have shown that the DHG and Bjorken sum rules can be regarded as low and high Q^2 limits of the same sum rule.

One of the most interesting consequences of the DHG sum rule occurs if we take a point-like limit such that the threshold for inelastic excitation becomes infinite while the mass of the system is kept finite. Since the integral over the photoabsorption cross section vanishes in this limit, the DHG sum rule implies that the anomalous moment must also vanish. Thus in the point-like limit, the magnetic moment of a spin-half system must approach the Dirac value $\mu \rightarrow \mu_D = e/2M$ up to structure corrections of order M/Λ, [or $(M/\Lambda)^2$ if the underlying theory is chiral] [8]. Hiller and I have recently derived a generalization of the DHG sum rule for spin-one composite systems. In the point-like limit, both the magnetic moment and quadrupole moment of any spin-one system must approach the canonical values predicted by electroweak theory for the W [9].

The Drell-Hearn sum rule also has important consequences for the computation of the magnetic moments of baryons in QCD. Magnetic moments are often computed using the quark model formula $\vec{\mu} = \sum_{i=1}^{3} \vec{\mu}_i$. This formula is correct in the case of atoms where the mass of the nucleus can be taken as infinite. However, magnetic moment additivity cannot be correct in general: the DHG sum rule shows that in the limit of strong binding where the constituents become very massive and the hadron becomes point-like, its magnetic moment must equal the Dirac value, not zero as predicted by quark moment additivity. The flaw in the conventional quark model formula is that it does not take into account the fact that the moment of a system H is derived from the electron scattering amplitude $eH \rightarrow e'H'$ at non-zero momentum transfer q. The Dirac value in the point-like limit actually arises from the Wigner boost of the wavefunction from p to $p+q$. A detailed discussion of this and the resulting relativistic corrections to the moment are given Ref. [10]. On the other hand, the overlap of light-cone Fock wavefunctions does provide a general method for the evaluation of hadronic magnetic moments and form factors [8].

3. The Gluon Helicity Distribution

One of the most interesting questions in QCD spin physics is the distribution of gluon polarization in the proton. The gluon distribution of a hadron is usually assumed to be radiatively generated from the QCD evolution of the quark structure functions beginning at an initial scale Q_0^2. The evolution is incoherent; *i.e.* each quark in the hadron radiates gluons independently. However, as can be seen in the light–cone Hamiltonian approach, the higher Fock components of a bound state in QCD contain gluons at any resolution scale. Furthermore, the exchange of gluon quanta between the bound-state constituents provides an interaction potential whose energy–dependent part generates a non-trivial non-additive contribution to the full gluon distribution $G_{g/H}(x, Q_0^2)$. The physics of gluon helicity distributions clearly involves the nonperturbative structure of the proton. Nevertheless, there are constraints which we can use to limit the possible form of the helicity–aligned and anti–aligned gluon distributions: $G^+(x) = G_{g\uparrow/N\uparrow}(x)$ and

$G^-(x) = G_{g\downarrow/N\uparrow}(x)$ [11]:

1. In order to insure positivity of fragmentation functions, the distribution functions $G_{a/b}(x)$ must behave as an odd or even power of $(1-x)$ at $x \to 1$ according to the relative statistics of a and b [12]. Thus the gluon distribution of a nucleon must have the behavior: $G_{g/N}(x) \sim (1-x)^{2k}$ at $x \to 1$ to ensure correct crossing to the fragmentation function $D_{N/g}(z)$.

2. In the $x \to 1$ limit, a gluon constituent of the proton is far off–shell and the leading behavior in the hadron wavefunctions is dominated by perturbative QCD contributions to the interaction kernel. We thus may use the minimally connected tree-graphs to characterize the threshold dependence of the structure functions. We find for a three quark plus one gluon Fock state, $\lim_{x\to 1} G^+(x) \to C(1-x)^{2N_q-2} = C(1-x)^4$. The gauge theory couplings of gluons to quarks also imply $\lim_{x\to 1} G^-(x)/G^+(x) \to (1-x)^2$. Thus $G^-(x) \sim (1-x)^6$ at $x \sim 1$. QCD evolution does not change these powers appreciably since the available phase-space for secondary gluon emission is limited to $k_\perp^2 < (1-x)Q^2$.

3. In the low x domain the quarks in the hadron radiate gluons coherently. Define $\Delta G(x) = G^+(x) - G^-(x)$ and $G(x) = G^+(x) + G^-(x)$. One then finds that the asymmetry ratio $\Delta G(x)/G(x)$ vanishes linearly with x.

4. In a simplest three quark plus one-gluon Fock state model the generated gluon distribution in the nucleon at low x has the normalization [11] $\Delta G(x)/G(x) = (x/3)\langle 1/y \rangle$, where y is the quark momentum fraction in the three quark state. The factor of 1/3 is due to the fact that all of the quarks contribute positively to $G(x)$, but are proportional to the sign of their helicity in $\Delta G(x)$.

If we assume equal quark momentum partition $\langle 1/y \rangle = 3$, then the above constraints are satisfied by the simple form [11]:

$$\Delta G(x) = (N/x)[1 - (1-x)^2](1-x)^4,$$
$$G(x) = (N/x)[1 + (1-x)^2](1-x)^4 \ .$$

This gives $\Delta G / \langle x_g \rangle = 77/72 = 1.07$ for the ratio of the gluon helicity to its momentum fraction in the nucleon. Since the gluon momentum fraction is ~ 0.5, we predict the total gluon helicity correlation $\Delta G = 0.54$, which by itself saturates the proton spin sum rule. It is expected that these results should provide a good characterization of the gluon distribution at the resolution scale $Q_0^2 \simeq M_p^2$. Clearly the model could be improved by taking into account higher Fock states and QCD evolution.

A determination of the unpolarized gluon distribution of the proton at $Q^2 \sim 2\,GeV^2$ using direct photon and deep inelastic data has been given in Ref. [13]. The best fit over the interval $0.05 \leq x \leq 0.75$ assuming the form $xG(x, Q^2 = 2\,GeV^2) = A(1-x)^{\eta_g}$ gives $\eta_g = 3.9 \pm 0.11(+0.8 - 0.6)$, where the errors in parenthesis allow for systematic uncertainties. This result is compatible with the prediction $\eta_g = 4$ for the gluon distribution at the bound–state scale, allowing for the small effects due to QCD evolution.

4. Quark Helicity Distributions and Hadron Helicity Retention in Inclusive Reactions at Large x_F

Consider a general inclusive reaction $AB \to CX$ at large x_F where the helicities λ_C and λ_A are measured. To be precise, we shall use the boost-invariant light-cone momentum fraction $x_C = \frac{k_C^+}{k_A^+} = \frac{(k^0+k^z)_C}{(k^0+k^z)_A}$. Hadron helicity retention implies that the difference between λ_C and λ_A tends to a minimum at $x_C \to 1$. Hadron helicity retention follows from the helicity structure of the gauge theory interactions, and it is applicable to hadrons, quarks, gluons, leptons, or photons. For example, in QED photon radiation in lepton scattering has the well-known distribution $dN/dx \propto [1 + (1-x)^2]/x$. The first term corresponds to the case where the photon helicity has the same sign as the lepton helicity; opposite sign helicity production is suppressed by a factor $(1-x)^2$ at $x \to 1$ [14]: the projectile helicity tends to be transferred by the leading fragment at each step in perturbation theory. It is a nontrivial step to show that hadron helicity also holds for hadrons in QCD; e.g.: the structure functions of the leading quarks in the proton have the nominal power behavior: $G_{q/p}(x) \sim (1-x)^3$ for $\lambda_q = \lambda_p$ and $G_{q/p}(x) \sim (1-x)^5$ if $\lambda_q = -\lambda_p$. This result follows from the fact that at $x \to 1$ the struck quark is far off-shell and space-like: $k^2 \sim -\mu^2/(1-x)$ where μ is a typical hadron mass scale; the leading fall-off of structure functions at $x \to 1$ can thus be computed from the minimally-connected tree-graphs.

These considerations have the immediate consequence that the down and anti-down quark distribution $\Delta d(x)$ has a zero as a function of x. At large x PQCD predicts that the helicity-antiparallel distribution $d^+(x)$ is suppressed relative to the helicity-parallel distribution $d^+(x)$ by two powers of $(1-x)$. At very small x the two distributions must have equal magnitude to ensure convergence of sum rules. However, measurements imply that the integral $\Delta d = \int_0^1 dx[d^+(x) - d^-(x)]$ is negative. Thus one expects that $\Delta d(x)$ changes sign as a function of x [11].

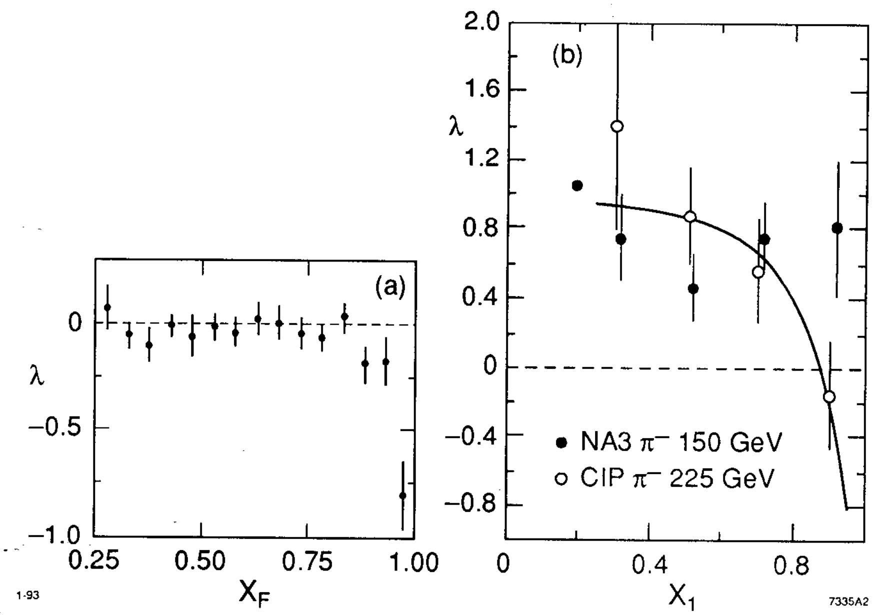

Figure 1. The x_F dependence of the polarization parameter λ for (a) J/ψ production [15] and (b) continuum lepton pair production [16] in $\pi - N$ collisions as a function of x_F.

One of the most important testing grounds for hadron helicity retention is J/ψ production in $\pi - N$ collisions. The helicity of the J/ψ can be measured from the angular distribution $1 + \lambda \cos^2 \theta_\mu$ of one of the muons in the leptonic decay of the J/ψ. At low to medium values of x_F the Chicago-Iowa-Princeton Collaboration [15] finds that $\lambda \sim 0$, which is consistent with expectations from the gluon-gluon fusion subprocess. However, at large $x_F > 0.9$ the angular distribution changes markedly to $\sin^2 \theta_\mu$; *i.e.*, the J/ψ is produced with longitudinal polarization. See Fig. 1(a). Note that the expectation of quark anti-quark fusion is $1 + \cos^2 \theta_\mu$ ($\lambda = +1$), as in the Drell-yan process. The sudden change to longitudinal polarization must mean that a new heavy quark production mechanism is present at large x_F [17]. In fact, it is easy to guess the relevant process which can produce high momentum charm quark pairs. [See Fig. 2(a).] Since nearly all of the pion's momentum is transferred to the charmonium system, one needs to consider diagrams where each valence quark in the incoming pion emits a fast gluon. The two gluons then fuse to make a fast $c\bar{c}$ pair. At large momentum fraction x, each gluon's helicity tends to be parallel to the helicity of its parent quark. Thus the angular momentum J_z of the gluon pair is transferred to the $c\bar{c}$ pair. The angular momentum tends to be preserved by any subsequent gluon radiation or gluon interaction from the heavy quarks. The J/ψ then tends to have the same helicity as the projectile at high light-cone momentum fraction.

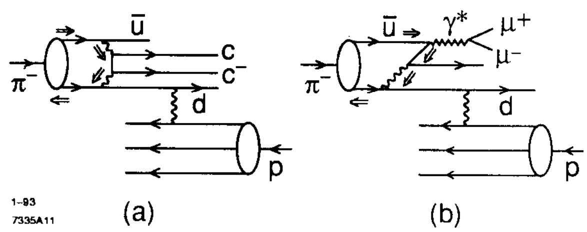

Figure 2. Higher twist mechanisms for producing (a) J/ψ and (b) massive lepton pairs at high x_F in meson-nucleon collisions.

Thus there is a natural mechanism in QCD which produces the J/ψ in the same helicity as the incoming beam hadron; the essential feature is the involvement of all of the valence quarks of the incoming hadron directly in the heavy quark production subprocess. Since such diagrams involve the correlation between the partons of the hadron, it can be classified as a higher-twist "intrinsic charm" amplitude; the production cross section is suppressed by powers of $f_\pi/M_{Q\bar{Q}}$ relative to conventional fusion processes. Although nominally higher twist, such diagrams provide an efficient way to transfer the beam momentum to the heavy quark system while stopping the valence quarks.

The intrinsic charm mechanism also can explain other features of the J/ψ hadroproduction [18,19,20]. The observed cross section persists to high x_F in excess of what is predicted from gluon fusion or quark anti–quark annihilation subprocesses; furthermore the cross section at high x_F has a strongly suppressed nuclear dependence, $A^{\alpha(x_F)} \sim 0.7$. The nuclear dependence actually depends on x_F not x_2 which rules out leading twist mechanisms. The higher-twist intrinsic charm *e.g.* $|uudc\bar{c}\rangle$ Fock state wavefunctions have maximum probability when all of the quarks have equal velocities, *i.e.* when $x_i \propto \sqrt{m^2 + k_{\perp_i}^2}$. This implies that the charm and anti–charm quarks have the majority of the momentum of the proton when they are present in the hadronic wavefunction. In a high energy proton–nucleus collision, the small transverse size, high–x intrinsic $c\bar{c}$ system can penetrate the nucleus, with minimal absorption and can coalesce to produce a charmonium state at large x_F. Since the soft quarks expand rapidly in impact space, the main interaction in the target of the intrinsic charm Fock state is with the slow valence quarks rather than the compact $c\bar{c}$ system [4]. Thus at large x_F the interaction in the nucleus should have the A-dependence of normal hadron nucleus cross sections: $\sim A^{0.7}$. Note that at high energies, the formation of the charmonium state occurs far outside the nucleus. Thus one predicts similar $A^{\alpha(x_F)}$–dependence of the J/ψ and ψ' cross sections. These predictions are in agreement with the results reported by the E–772 experiment at Fermilab [19].

In the case of continuum pair production, the lepton pair produced via the leading-twist Drell-Yan fusion mechanism $q\bar{q} \rightarrow \mu^+\mu^-$ has transverse polarization ($\lambda = 1$). However, at large x_F the muon angular distribution is observed to change to $\sin^2\theta_\mu$ [16]. See Fig. 1(b). This result was predicted [21] from the dominance of higher twist $\pi q \rightarrow \mu^+\mu^- q$ subprocess contributions at high x_F. A detailed calculation shows subprocess amplitude can be normalized to the same integral over the pion distribution amplitude $\int dx\phi(x,Q)/(1-x)$ that controls the pion form factor [22].

In the higher-twist subprocess diagram, Fig. 2(b), the lepton pair tends to have the same helicity as the beam hadron at large x_F. For example, consider $\pi^- N \to \mu^+ \mu^- X$ at high x_F. The valence d quark emits a fast gluon which in turn makes a fast-u, slow-$\overline{u}$ pair. Because of the QCD couplings, the fast u then carries the helicity of the d. The valence $\overline{u}$ then annihilates with the fast u to make the lepton pair at $x_F \sim 1$. The lepton pair thus tends to have the helicity ($J_z = 0$) of the pion, in agreement with hadron helicity retention.

5. Hadron Helicity Conservation in Hard Exclusive Reactions

There are also strong helicity constraints on form factors and other exclusive amplitudes which follow from perturbative QCD [3]. At large momentum transfer, each helicity amplitude contributing to an exclusive process at large momentum transfer can be written as a convolution of a hard quark-gluon scattering amplitude T_H which conserves quark helicity with the hadron distribution amplitudes $\phi(x_i,Q)$, which are the $L_z = 0$ projection of the hadron's valence Fock state wavefunction: $\phi(x_i,\lambda_i,Q) = \int [d^2k_\perp]\ \psi(x_i,\vec{k}_{\perp i},\lambda_i)\theta(k^2_{\perp i} < Q^2)$ where $\psi(x_i,\vec{k}_{\perp i},\lambda_i)$ is the valence wavefunction. Since ϕ only depends logarithmically on Q^2, the main dynamical dependence of $F_B(Q^2)$ is the power behavior $(Q^2)^{-2}$ derived from the scaling behavior of the elementary propagators in T_H.

As shown by Botts, Li, and Sterman [23], the virtual Sudakov form factor suppresses long distance contributions from Landshoff multiple scattering and $x \sim$

1 integration regions, so that the leading high momentum transfer behavior of hard exclusive amplitudes are generally controlled by short-distance physics. Thus quark helicity conservation of the basic QCD interactions leads to a general rule concerning the spin structure of exclusive amplitudes [5]: to leading order in $1/Q$, the total helicity of hadrons in the initial state must equal the total helicity of hadrons in the final state. This selection rule is independent of any photon or lepton spin appearing in the process. The result follows from (a) neglecting quark mass terms, (b) the vector coupling of gauge particles, and (c) the dominance of valence Fock states with zero angular momentum projection. The result is true in each order of perturbation theory in α_s.

For example, PQCD predicts that the Pauli Form factor $F_2(Q^2)$ of a baryon is suppressed relative to the helicity-conserving Dirac form factor $F_1(Q^2)$. A recent experiment at SLAC carried out by the American-University/SLAC collaboration is in fact consistent with the prediction $Q^2 F_2(Q^2)/F_1(Q^2) \to$ const. [24]. Helicity conservation holds for any baryon to baryon vector or axial vector transition amplitude at large spacelike or timelike momentum. Helicity non-conserving form factors should fall as an additional power of $1/Q^2$ [5]. Measurements [25] of the transition form factor to the $J = 3/2$ $N(1520)$ nucleon resonance are consistent with $J_z = \pm 1/2$ dominance, as predicted by the helicity conservation rule [5]. One of the most beautiful tests of perturbative QCD is in proton Compton scattering, where there are now detailed predictions available for each hadron helicity-conserving amplitude for both the spacelike and timelike processes [26]. In the case of spin-one systems such as the ρ or the deuteron, PQCD predicts that the ratio of the three form factors have the same behavior at large momentum transfer as that of the W in the electroweak theory [9].

Hadron helicity conservation in large momentum transfer exclusive reactions is a general principle of leading twist QCD. In fact, in several outstanding cases, it does not work at all, particularly in single spin asymmetries such as A_N in pp scattering, and most spectacularly in the two-body hadronic decays of the J/ψ. The inference from these failures is that non-perturbative or higher twist effects must be playing a crucial role in the kinematic range of these experiments.

The J/ψ decays into isospin-zero final states through the intermediate three-gluon channel. If $PQCD$ is applicable, then the leading contributions to the decay amplitudes preserve hadron helicity. In the case of e^+e^- annihilation into vector plus pseudoscalar mesons, Lorentz invariance requires that the vector meson will be produced transversely polarized. Since this amplitude does not conserve hadron helicity, $PQCD$ predicts that it will be dynamically suppressed at high momentum transfer. Hadron helicity conservation appears to be severely violated if one compares the exclusive decays J/ψ and $\psi' \to \rho\pi, K^*\overline{K}$ and other vector-pseudoscalar combinations. The predominant two-body hadronic decays of the J/ψ have the measured branching ratios

$$\begin{aligned} BR(J/\psi{\to}K^+K^-) &= 2.37 \pm 0.31 \times 10^{-4} \\ BR(J/\psi{\to}\rho\pi) &= 1.28 \pm 0.10 \times 10^{-2} \\ BR(J/\psi{\to}K^+K^{-*}) &= 5.0 \pm 0.4 \times 10^{-3} \ . \end{aligned}$$

Thus the vector-pseudoscalar decays are not suppressed, in striking contrast to the PQCD predictions. On the other hand, for the ψ':

$$\begin{aligned} BR(\psi'{\to}K^+K^-) &= 1.0 \pm 0.7 \times 10^{-4} \\ BR(\psi'{\to}\rho\pi) &< 8.3 \times 10^{-5} \quad (90\% \ CL) \\ BR(\psi'{\to}K^+K^{-*}) &< 1.8 \times 10^{-5} \quad (90\% \ CL) \ . \end{aligned}$$

From the standpoint of perturbative QCD, the observed suppression of ψ' to vector-pseudoscalar mesons is expected; it is the J/ψ that is anomalous [27]. What can account for the apparently strong violation of hadron helicity conservation? One possibility is that the overlap of the $c\bar{c}$ system with the wavefunctions of the ρ and π is an extremely steep function of the pair mass, as discussed by Chaichian and Tornqvist [28]. However, this seems unnatural in view of the similar size of the J/ψ and ψ' branching ratios to K^+K^-. Pinsky [29] has suggested that the ψ' decays predominantly to final states with excited vector mesons such as $\rho'\pi$, in analogy to the absence of configuration mixing in nuclear decays. However, this long-distance decay mechanism would not be expected to be important if the charmonium state decays through $c\bar{c}$ annihilation at the Compton scale $1/m_c$.

Another way in which hadron helicity conservation might fail for $J/\psi \rightarrow$ gluons $\rightarrow \pi\rho$ is if the intermediate gluons resonate to form a gluonium state $\mathcal{O}$. If such a state exists, has a mass near that of the J/ψ, and is relatively stable, then the subprocess for $J/\psi \rightarrow \pi\rho$ occurs over large distances and the helicity conservation theorem need no longer apply. This would also explain why the J/ψ decays into $\pi\rho$ and not the ψ'. Tuan, Lepage, and I [27] have thus proposed, following Hou and Soni [30], that the enhancement of $J/\psi \rightarrow K^*\overline{K}$ and $J/\psi \rightarrow \rho\pi$ decay modes is caused by a quantum mechanical mixing of the J/ψ with a $J^{PC} = 1^{--}$ vector gluonium state $\mathcal{O}$ which causes the breakdown of the QCD helicity theorem. The decay width for $J/\psi \rightarrow \rho\pi$ via the sequence $J/\psi \rightarrow \mathcal{O} \rightarrow \rho\pi$ must be substantially larger than the decay width for the (non-pole) continuum process $J/\psi \rightarrow 3$ gluons $\rightarrow \rho\pi$. In the other channels the branching ratios of the $\mathcal{O}$ must be so small that the continuum contribution governed by the QCD theorem dominates over that of the $\mathcal{O}$ pole. A gluonium state of this type was first postulated by Freund and

Nambu [31] based on OZI dynamics soon after the discovery of the J/ψ and ψ' mesons. The most direct way to search for the $\mathcal{O}$ is to scan $\overline{p}p$ or e^+e^- annihilation at $\sqrt{s}$ within $\sim$ 100 MeV of the J/ψ, triggering on vector/pseudoscalar decays such as $\pi\rho$ or $\overline{K}K^*$ and look for enhancements relative to K^+K^-. Such a search has recently been proposed for the BEPC by Chen Yu, Gu Yifan, and Wang Ping.

6. Anomalous Spin Correlations and Color Transparency Effects in Proton-Proton Scattering

The perturbative QCD analysis of exclusive amplitudes assumes that large momentum transfer exclusive scattering reactions are controlled by short distance quark-gluon subprocesses, and that corrections from quark masses and intrinsic transverse momenta can be ignored. Since hard scattering exclusive processes are dominated by valence Fock state wavefunctions of the hadrons with small impact separation and small color dipole moments, one predicts that initial and final state interactions are generally suppressed at high momentum transfer. In particular, since the formation time is long at high energies, one predicts that the attenuation of quasi-elastic processes due to Glauber inelastic scattering in a nucleus will be reduced. This is the color transparency prediction of perturbative QCD [32]. A test of color transparency in large momentum transfer quasielastic pp scattering at $\theta_{\rm cm} \simeq \pi/2$ has been carried out at BNL using several nuclear targets (C, Al, Pb) [33]. The attenuation at $p_{\rm lab} = 10$ GeV/c in the various nuclear targets was observed to be in fact much less than that predicted by traditional Glauber theory. The expectation from perturbative QCD is that the transparency effect should become even more apparent as the momentum transfer rises. However, the data at $p_{\rm lab} = 12$ GeV/c shows normal nuclear attenuation and thus a violation of color transparency.

An even more serious challenge to the PQCD predictions for exclusive scattering is the observed behavior of the normal spin-spin correlation asymmetry $A_{NN} = [d\sigma(\uparrow\uparrow) - d\sigma(\uparrow\downarrow)]/[d\sigma(\uparrow\uparrow) + d\sigma(\uparrow\downarrow)]$ measured in large momentum transfer pp elastic scattering. At $p_{\rm lab} = 11.75$ GeV/c and $\theta_{\rm cm} = \pi/2$, A_{NN} rises to $\cong 60\%$, corresponding to four times more probability for protons to scatter with their incident spins both normal to the scattering plane and parallel, rather than normal and opposite [2]. In contrast, the unpolarized data is to first approximation consistent with the fixed angle scaling law $s^{10}d\sigma/dt(pp \to pp) = f(\theta_{CM})$ expected from the perturbative analysis. The onset of new structure at $s \simeq 23$ GeV^2 suggests new degrees of freedom in the two-baryon system.

Guy De Teramond and I [34] have noted that the onset of strong spin–spin correlations, as well as the breakdown of color transparency, can be explained as the consequence of a strong threshold enhancement at the open-charm threshold for $pp \to \Lambda_c D p$ at $\sqrt{s} = 5.08$ GeV or $p_{\rm lab} \sim 12$ GeV/c. At this energy the charm quarks are produced at rest in the center of mass. Since all eight quarks have zero relative velocity, they can resonate to give a strong threshold effect in the $J = L = S = 1$ partial wave. (The orbital angular momentum of the pp state must be odd since the charm and anti–charm quarks have opposite parity.) The $J = L = S = 1$ partial wave has maximal spin correlation $A_{NN} = 1$. A charm production cross section of the order of 1 μb in the threshold region can have, by unitarity, a large effect on the large angle elastic $pp \to pp$ amplitude since the competing perturbative QCD hard-scattering amplitude at large momentum transfer is very small at $\sqrt{s} = 5$ GeV. In fact as recently shown by Manohar, Luke, and Savage [35], the QCD trace anomaly predicts that the scalar charmonium-nucleus interaction is strongly amplified at low

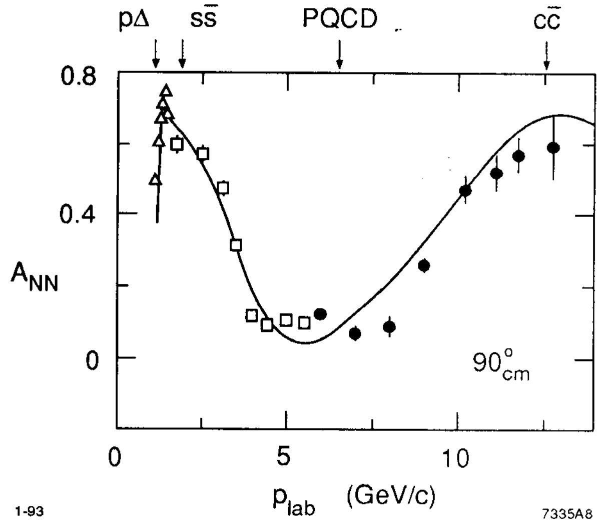

Figure 3. A_{NN} as a function of $p_{\rm lab}$ at $\theta_{cm} = \pi/2$. The data [2] are from Crosbie *et al.* (solid dots), Lin *et al.* (open squares) and Bhatia *et al.* (open triangles). The peak at $p_{\rm lab} = 1.26$ GeV/c corresponds to the $p\Delta$ threshold. The data are well reproduced by the interference of the broad resonant structures at the strange ($p_{\rm lab} = 2.35$ GeV/c) and charm ($p_{\rm lab} = 12.8$ GeV/c) thresholds, interfering with a PQCD background. The value of A_{NN} from PQCD alone is 1/3.

velocities and can lead to nuclear-bound charmonium [36].

An analytic model which contains all of these features is given in Ref. [34]. The background component of the model is the perturbative QCD amplitude with s^{-4} scaling of the $pp \to pp$ amplitude at fixed $\theta_{\rm cm}$ and the dominance of those amplitudes that conserve hadron helicity [5]. A comparison [37] of the magnitude of cross sections for different exclusive two-body scattering channels indicate that quark interchange amplitudes [38] dominate quark annihilation or gluon exchange contributions. The most striking test of the model is its prediction for the spin correlation A_{NN} shown in Fig. 3. The rise of A_{NN} to $\simeq 60\%$ at $p_{\rm lab} = 11.75$ GeV/c is correctly reproduced by the high energy J=1 resonance interfering with ϕ(PQCD). The narrow peak which appears in the data of Fig. 3 corresponds to the onset of the $pp \to p\Delta(1232)$ channel which can be interpreted as a $uuuuddq\bar{q}$ $^3F_3$ resonance. The heavy quark threshold model also provides a good description of the s and t dependence of the differential cross section, including its "oscillatory" dependence [39] in s at fixed $\theta_{\rm cm}$, and the broadening of the angular distribution near the resonances. Most important, it gives a consistent explanation for the striking behavior of both the spin-spin correlations and the anomalous energy dependence of the attenuation of quasielastic pp scattering in nuclei. A threshold enhancement or resonance couples to hadrons of conventional size. Unlike the perturbative amplitude, the protons coupling to the resonant amplitude will have normal absorption in the nucleus. Thus the nucleus acts as a filter, absorbing the non–perturbative contribution to elastic pp scattering, while allowing the hard-scattering perturbative QCD processes to occur additively throughout the nuclear volume [40]. Conversely, in the momentum range $p_{\rm lab} = 5$ to 10 GeV/c one predicts that the perturbative hard-scattering amplitude will be dominant at large angles. It is thus predicted that color transparency should reappear at higher energies ($p_{\rm lab} \geq 16$ GeV/c), and also at smaller angles ($\theta_{\rm cm} \approx 60^\circ$) at $p_{\rm lab} = 12$ GeV/c where the perturbative QCD amplitude dominates. If the resonance structures in A_{NN} are indeed associated with heavy quark degrees of freedom, then the model predicts inelastic pp cross sections of the order of 1 mb and 1μb for the production of strange and charmed hadrons near their respective thresholds. In fact, the neutral strange inclusive pp cross section measured at $p_{\rm lab} = 5.5$ GeV/c is 0.45 ± 0.04 mb [41]. Thus the crucial test of the heavy quark hypothesis for explaining A_{NN} is the observation of significant charm hadron production at $p_{\rm lab} \geq 12$ GeV/c.

Ralston and Pire [40] have suggested that the oscillations of the pp elastic cross section and the apparent breakdown of color transparency are associated with the dominance of the Landshoff pinch contributions at $\sqrt{s} \sim 5$ GeV. The oscillating behavior of $d\sigma/dt$ is then due to the energy dependence of the relative phase between the pinch and hard-scattering contributions. They assume color transparency will disappear whenever the pinch contributions are dominant since such contributions could couple to wavefunctions of large transverse size. However, the large spin correlation in A_{NN} is not readily explained in the Ralston-Pire model unless the Landshoff diagram itself has $A_{NN} \sim 1$.

7. Polarization-Dependent Nuclear Shadowing

Another interesting spin effect in QCD is the prediction that nuclear shadowing depends on the virtual photon polarization. In models where shadowing is due to the deformation of nucleon structure functions in the nucleus, one would not expect such any dependence on photon polarization. In Refs. [42] and [43] one sees that nuclear shadowing (in the target rest frame) arises from the destructive interference of the multiple scattering of a quark (or antiquark) in the nucleus. The quark comes from the upstream dissociation of the virtual photon. The $q\overline{q}$ pair is formed at a formation time (coherence length) $\tau \propto 1/x_{\rm bj}M$ before the target. In order to get significant multiple scattering and interference one needs a coherence length comparable to the nuclear size. However, Hoyer, Del Duca and I found [43] that the coherence length is significantly shorter (by a factor of $1/\sqrt{3}$) for the longitudinally polarized photon than the transverse case. The reason for this is that the internal transverse momentum and hence the virtual mass and energy of the $q\overline{q}$ pair is larger by a nearly constant factor in the longitudinal case, thus shortening its lifetime. Thus the nuclear attenuation is delayed to smaller values of $x_{\rm bj}$ in the longitudinal compared to the transverse cross section. Nikolaev [44] has also recently discussed the possibility of smaller nuclear shadowing of σ_L on the grounds that the $q\bar{q}$ system has a smaller transverse size in the case of a longitudinally polarized photon, and it is thus more color transparent. In this case diminished longitudinal shadowing would persist for all $x_{\rm bj}$.

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REFERENCES

[1] See R. Jaffe, these proceedings.

[2] For references to the data and a review, see A. Krisch, these proceedings.

[3] S. J. Brodsky, G. P. Lepage, in "Perturbative Quantum Chromodynamics", Edited by A. H. Mueller, World Scientific Publ. Co. (1989).

[4] S. J. Brodsky, P. Hoyer, A. H. Mueller, W. K. Tang, Nucl. Phys. B369, 519 (1992).

[5] S. J. Brodsky and G. P. Lepage, Phys. Rev. D24, 2848 (1981).

[6] By S. J. Brodsky, T. A. DeGrand, R. Schwitters, Phys. Lett. 79B, 255 (1978).

[7] V. D. Burkert and B. L. Ioffe, CEBAF-PR-92-018, (1992).

[8] S. J. Brodsky, S. D. Drell, Phys. Rev. D22, 2236 (1980).

[9] S. J. Brodsky, J. R. Hiller, Phys. Rev. D46, 2141 (1992).

[10] S. J. Brodsky, J. R. Primack, Annals Phys. 52, 315 (1969).

[11] S. J. Brodsky, M. Burkardt, I. A. Schmidt, in preparation; S. J. Brodsky and I. A. Schmidt, Phys. Lett. B234, 144 (1990).

[12] V. N. Gribov and L. N. Lipatov, Sov. J. Nucl. Phys. 15, 438 and 675 (1972).

[13] P. Aurenche, *et al.*, Phys. Rev. D39, 3275 (1989).

[14] J. D. Bjorken, Phys. Rev. D1, 1376 (1970).

[15] C. Biino, Phys. Rev. Lett. 58, 2523 (1987).

[16] J. G. Heinrich, *et al.*, Phys. Rev. D44, 1909 (1991); M. Gusanziroli, *et al.*, Z. Phys. C37, 545 (1988).

[17] S. J. Brodsky, W. Tang, P. Hoyer, and M. Vanttinen, in preparation.

[18] J. Badier, *et al.*, Z. Phys. C20, 101 (1983).

[19] D.M. Alde, *et al.*Phys. Rev. Lett. 66, 133 (1991).

[20] R. Vogt, S. J. Brodsky, P. Hoyer, Nucl. Phys. B360, 67 (1991).

[21] E. L. Berger and S. J. Brodsky, Phys. Rev. Lett. 42, 940 (1979).

[22] S. J. Brodsky, E. L. Berger, G. P. Lepage, SLAC-PUB-3027, published in the proceedings of the Workshop on Drell-Yan Processes, Batavia, IL, (1982).

[23] J. Botts, Phys. Rev. D44, 2768, (1991); H. Li, G. Sterman, Nucl. Phys. B381 129, (1992).

[24] P. Bosted, *et al.*, Phys. Rev. Lett. 68, 3841 (1992).

[25] V. D. Burkert, CEBAF-PR-87-006. P. Stoler, Phys. Rev. D44 73 (1991).

[26] A.S. Kronfeld B. Nizic, Phys.Rev. D44, 3445, (1991); ibid. D46, 2272 (1992). T. Hyer, SLAC-PUB-5889 (1992).

[27] S. J. Brodsky, G. P. Lepage and San Fu Tuan, Phys. Rev. Lett. 59, 621 (1987).

[28] M. Chaichian, N. A. Tornqvist, Nucl. Phys. B323, 75 (1989).

[29] S. S. Pinsky, Phys. Lett. B236, 479 (1990).

[30] Wei-Shou Hou and A. Soni, Phys. Rev. Lett. 50, 569 (1983).

[31] P. G. O. Freund and Y. Nambu, Phys. Rev. Lett. 34, 1645 (1975).

[32] S. J. Brodsky, A.H. Mueller, Phys. Lett. 206B, 685 (1988).

[33] A. S. Carroll, *et al.*, Phys. Rev. Lett. 61, 1698 (1988).

[34] S. J. Brodsky and G. de Teramond, Phys. Rev. Lett. 60, 1924 (1988).

[35] M. Luke, A. V. Manohar, M. J. Savage, Phys. Lett. B288, 355 (1992).

[36] S. J. Brodsky, I. A. Schmidt, and G. F. de Teramond, Phys. Rev. Lett. 64 1011, (1990).

[37] G. C. Blazey *et al.*, Phys. Rev. Lett. 55, 1820 (1985).

[38] J. F. Gunion, R. Blankenbecler, and S. J. Brodsky, Phys. Rev. D6, 2652 (1972).

[39] A. W. Hendry, Phys. Rev. D10, 2300 (1974).

[40] J. P. Ralston and B. Pire, Phys. Rev. Lett. 57, 2330 (1986); Phys. Lett. 117B, 233 (1982); University of Kansas preprint 5-15-92, (1992). See also G. P. Ramsey and D. Sivers, Phys. Rev. D45, 79 (1992); and C. E. Carlson, M. Chachkhunashvili, and F. Myhrer, Phys. Rev. D46 2891 (1992).

[41] G. Alexander *et al.*, Phys. Rev. 154, 1284 (1967).

[42] S. J. Brodsky and H. J. Lu, Phys. Rev. Lett. 64, 1342 (1990).

[43] V. Del Duca, S. J. Brodsky, P. Hoyer, Phys. Rev. D46, 931 (1992).

[44] N. N. Nikolaev, these proceedings.