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SPIN EFFECTS IN PERTURBATIVE QUANTUM CHROMODYNAMICS*

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ABSTRACT

The spin dependence of large momentum transfer exclusive and inclusive reactions can be used to test the gluon spin and other basic elements of QCD. In particular, exclusive processes including hadronic decays of heavy quark resonances have the potential of isolating QCD hard scattering subprocesses in situations where the helicities of all the interacting constituents are controlled. The predictions can be summarized in terms of QCD spin selection rules. We also briefly comment on the calculation of magnetic moment and other hadronic properties in QCD.

I. INTRODUCTION

Among the most difficult challenges for any dynamical theory of high energy hadronic phenomenon is the correct description of spin effects. Spin correlations and other polarization phenomenon are sensitive to the detailed helicity dependence and phase structure of hadronic amplitudes and lead to critical checks of theoretical predictions. In this talk we will discuss how the spin dependence of exclusive and inclusive charge momentum transfer reactions provide complimentary tests of some of the basic elements of quantum chromodynamics.

The predictions of QCD for large momentum transfer inclusive processes, including spin correlations, are based on the QCD factorization theorem,[1] which separates the dynamics of hard scattering quark and gluon subprocess cross sections from process-independent structure functions $G_{q/H}(x,Q)$ and $G_{g/H}(x,Q)$ -- evolved to the large momentum transfer scale Q. These predictions, though straightforward, are complicated by a number of effects which can seriously affect the results in the subasymptotic domain (see Sect. III). In addition, the correlation between constituent and parent hadron spin is only statistical and vanishes in the low x domain.

As we have discussed in a series of recent papers,[2-5] the predictions of perturbative QCD can also be extended to the domain of exclusive processes.[6] The predictions for large momentum exclusive reactions are based on a second QCD factorization theorem which separates the dynamics of hard scattering quark and gluon scattering amplitudes from process-independent "distribution amplitudes" $\phi_H(x,Q)$ evolved to the large momentum transfer scale Q. As we shall see, exclusive processes have the potential of isolating the QCD hard scattering processes in situations where the helicities of all

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the interacting constituents are controlled. In the case of the meson form factors and meson pair-production in two-photon collisions, the results are rigorous predictions of QCD which test the scaling and spin properties of quark and gluon interactions at large momentum transfer to all orders in $\alpha_s(Q^2)$, as well as the structure and helicity dependence of hadronic wave functions at short distances. In the case of other reactions such as the baryon form factor and fixed angle hadron-hadron scattering, the factorized hard scattering contributions are again predicted to give the asymptotically dominant contribution at large momentum transfer. Non-factorizable contributions from Landshoff pinch contributions[9] and the x near 1 kinematic region,[5,7] although asymptotically suppressed by Sudakov form factors,[5,10] could still play an important phenomenological role at non-asymptotic momentum. In this talk we will review the perturbative QCD predictions for many types of spin and polarization correlations and contrast the sensitivity of large momentum transfer inclusive versus exclusive reactions as basic probes of hadron dynamics.

II. QCD AND QUARK HELICITY[5,11]

A central feature of perturbative quantum chromodynamics is that quark helicity is conserved (up to terms of order m_q/Q) by the vector gluon interactions. Hard subprocesses in which all the routings of the large momentum transfer $Q^2 \gg m_q^2$ involve far-off shell intermediate states thus must conserve total quark helicity: $h_I = h_F$. The crucial step for deriving predictions for the spin-dependence of hadron reactions is to understand how the helicity of each interacting hadron is correlated with the helicity of its constituents. There is a striking difference between exclusive and inclusive reactions in this regard. In the case of inclusive reactions, the inevitable presence of quark and gluon (non-valence) spectators as well as non-zero relative orbital angular momentum strongly reduces the spin correlation between the interacting constituent and the parent hadron, expect at the kinematical limit $x \to 1$. In contrast, exclusive reactions involving large momentum transfer Q are dominated (to leading order in m_q/Q) by the simplest valence state wave function with zero relative angular momentum $L_z = 0$ for each interacting hadron. Thus, in this case, the sum of the valence constituent helicity equals the hadron helicity:

$$\sum_{\text{valence}} S_i^z = S_{\text{hadron}}^z \tag{2.1}$$

to all orders in $\alpha_s(Q^2)$ and leading order in m_q/Q. The combination of this "spin additivity" property with helicity conservation for hard subprocesses then leads to "QCD selection rules" for exclusive processes which directly reflect the spin properties of the basic quark and gluon interactions. We give a detailed discussion in Sect. IV.

The underlying link between exclusive and inclusive processes and the spin dependence of hadronic reactions in QCD is the Fock state hadronic wave function.[12] An important feature of QCD is that the wave function for hadrons can be expanded as a sum over states of definite quark and gluon number. Such renormalized Fock states can be rigorously defined because of the cancellation of all infrared divergences for color singlet bound states.[5] We will define the states at equal time $\tau = t + z$ on light-cone in the light-cone gauge $A^+ = A^0 + A^3 = 0$. The amplitude to find n (on-mass-shell) quarks and gluons in a hadron with 4-momentum P directed along the z-direction and spin projection S_z is defined to be ($k^{\pm} = k^0 \pm k^3$) (see Fig. 1).

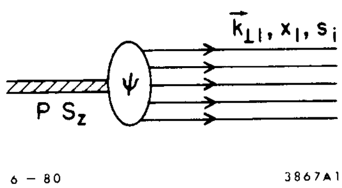

Fig. 1. A representative Fock state amplitude $\psi_{S_z}(k_{\perp i}, x_i, s_i)$ for a hadron at equal time on the light-cone.

$$\psi_{S_z}^{(n)}\left(x_i, \vec{k}_{\perp i}, s_i\right) \quad , \quad x_i \equiv \frac{k_i^+}{P^+} \quad ,$$

where by momentum conservation $\sum_{i=1}^{n} x_i = 1$ and $\sum_{i=1}^{n} \vec{k}_{\perp i} = 0$. The s_i specify the spin-projection of the constituents. The state is off the light-cone energy shell,

$$P^- - \sum_{i=1}^{n} k_i^- = \frac{M^2 - \sum_{i=1}^{n} \frac{\vec{k}_{\perp i}^2 + m_i^2}{x_i}}{P^+} < 0 \quad . \qquad (2.2)$$

The "valence" Fock states (which turn out to dominate large momentum transfer exclusive reactions) are the $|q\bar{q}\rangle$ ($n = 2$) and $|qqq\rangle$ ($n = 3$) components of the meson and baryon. For each fermion or anti-fermion constituent $\psi_{S_z}^{(n)}(k_{\perp i}, x_i, s_i)$ multiplies the spin factor $u(\vec{k}_i)/\sqrt{k_i^+}$ or $v(\vec{k}_i)/\sqrt{k_i^+}$. The wave function normalization condition is

$$\sum_{(n)(s_i)} \int \left|\psi_{S_z}^{(n)}\left(k_{\perp i}, x_i, s_i\right)\right|^2 \left[d^2k_\perp\right]\left[dx\right] = 1 \ , \qquad (2.3)$$

where

$$\left[d^2k_\perp\right] \equiv 16\pi^3 \, \delta^{(2)}\left(\sum_i k_{\perp i}\right) \prod_{i=1}^{n} \frac{d^2k_{\perp i}}{16\pi^3} \quad ,$$

and

$$\left[dx\right] = \delta\left(1 - \sum_i x_i\right) \prod_{i=1}^{n} dx_i \quad .$$

The wave functions ψ can now be used to define the coherent distribution amplitudes $\varphi(x_i, Q) \sim \int^{Q} \left[d^2k_\perp\right] \psi_V(x_i, k_\perp)$ which control high momentum transfer exclusive processes, and the probabilistic quark and gluon momentum distributions $G(x,Q) \sim \int^{Q} \left[d^2k_\perp\right] |\psi(x_i, k_\perp)|^2$ which control large momentum transfer inclusive reactions, as well as multiparticle longitudinal and transverse momentum distribution. More precisely, the quark and gluon distribution functions for large momentum transfer inclusive reactions at the scale Q^2 are

$$G_{a/H}(x_a,s_a,Q) = d_a^{-1}(Q) \sum_{n,s_{i\neq a}} \int^{k_{\perp a}^2 < Q^2} \left|\psi_{S_z}^{(n)}(k_{\perp i},x_i,s_i)\right|^2 \left[d^2k_\perp\right]\left[dx\right]\delta(x-x_a) \tag{2.4}$$

where $d_a^{-1}(Q^2)$ is due to the wave function renormalization of the constituent a. Notice that only terms which fall-off as $|\psi|^2 \sim (k_{\perp a}^2)^{-1}$ (modulo logs) contribute to the Q^2 dependence of the integral. These contributions are analyzable by the renormalization group and correspond in perturbative QCD to quark or gluon pair production or fragmentation processes associated with the struck constituent a. In general, unless x is close to 1, all Fock states in the hadron contribute to $G_{a/H}$. Multiparticle probability distributions are simple generalization of Eq. (2.4).

Inclusive cross sections in QCD are then obtained by a summation over (spin-dependent) incoherent hard scattering subprocess cross sections:[1,13]

$$d\sigma_{AB\to CX} = \sum_{\ell} \int_0^1 dx_a \int_0^1 dx_b \; G_{a/A}(x_a,s_a,\tilde{Q})\, G_{b/B}(x_b,s_b,\tilde{Q})\, d\hat{\sigma}^{\ell}_{ab\to X}(Q) \tag{2.5}$$

where each subprocess $d\hat{\sigma}$ is computed for on-shell constituents a and b which are collinear with A and B. Equation (2.5) gives the standard QCD factorization of the high momentum transfer subprocesses from the non-perturbative bound state dynamics. The distribution functions $G(x_a,\tilde{Q})$ give the probability distributions at relative impact distances of order $b_\perp \sim \mathcal{O}(1/\tilde{Q})$; the actual maximum transverse momentum scale where the factorization occurs is approximately given by $\tilde{Q}^2 \sim (1-x_a)Q^2$. [A detailed discussion is given in Ref. 14.] The subprocesses $d\sigma_{ab}$ include (high-twist) reactions where a and b are clusters of quarks and gluons in the initial state hadrons. Such terms incorporate[15] the dynamical effects of large transverse momentum quark and gluon components in the incident wave function as well as from multiple scattering effects. The hard scattering summation handles the off-shell kinematics of the constituents correctly, and can performed in a well-defined gauge-invariant manner.[15] The naive procedure of smearing the leading twist cross section leads to infinite results in the case of gluon-exchange processes and cannot be justified in QCD.

III. INCLUSIVE REACTIONS AND SPIN EFFECTS

Given forms for the $G_{a/A}(x_a,s_a,Q)$, it is, in principle, possible to calculate the dynamical and spin dependence for each hard scattering subprocess $d\hat{\sigma}$ (as a power series in $\alpha_s(Q^2)$, and obtain predictions for weak, electromagnetic, or purely hadronic large momentum transfer inclusive reactions from Eq. (2.5). For example, spin asymmetries and correlations have been calculated for the production of hadrons, jets, prompt photons, lepton pairs, weak bosons at large transverse momentum in hadron-hadron, photon-hadron, and e^+e^- collisions.[16] The basic spin correlation structure function for quarks in nucleons which is required for these predictions can be measured in polarized nucleon-polarized lepton deep inelastic reactions.[17] However, there are a number of complications and difficulties which detract from the utility of inclusive reactions as definitive probes of helicity structure in QCD:

(1) Although the Q^2-evolution of the spin-dependent quark and gluon probability distributions in hadrons is determined by perturbative QCD, the complete forms of the structure functions depends on unknown non-perturbative physics. (An important exception to this is photon-induced reactions.)

(2) The probability distributions $G_{a/H}$ receive contributions from all Fock states, including any number of spectator quark and gluons. In addition, states with any value of orbital angular momentum contribute in the square of the wave function; in particular, the spectator gluons are produced by QCD evolution with $L_Z \neq 0$. Thus (unlike exclusive reactions) the sum of constituent helicities does not generally equal the hadron helicity. (An important constraint on quark-hadron helicity correlations is however given by the Bjorken sum rule.[18])

(3) The use of Eq. (2.5) for spin correlations beyond leading order in $\alpha_s(Q^2)$ becomes very complicated because of the interrelation between many different subprocesses. For example, the $gq \to \gamma^* q$ subprocess must be explicitly taken into account in Drell-Yan massive lepton pair production beyond leading order in $\alpha_s(Q^2)$. An even more serious problem is that virtual longitudinal-scalar gluon exclusive contributions also must be taken into account beyond leading order. For example, in Born approximation the process $\gamma g \to q\bar{q}$ has zero correlation between the photon linear polarization and the $q\bar{q}$ production plane.[19] Contributions to the correlation to first order in $\alpha_s(Q^2)$ arise not only from one-loop virtual corrections[19] to $\gamma g \to q\bar{q}$, but also from the virtual gluon contributions in the higher particle subprocesses $\gamma q \to q + q\bar{q}$, and $\gamma g \to g + q\bar{q}$ (see Fig. 2). The analysis of these "coulombic" gluon contributions will be given elsewhere.

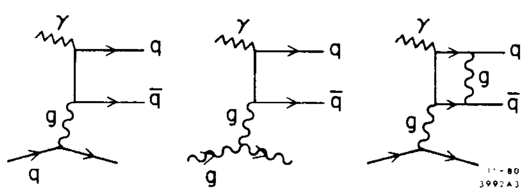

Fig. 2. Order $\alpha_s(Q^2)$ corrections to the $\gamma g \to q\bar{q}$ subprocess. The longitudinal-scalar polarization components of the exchanged gluon propagator in (a) and (b) give a non-zero contribution to the correlation of the $q\bar{q}$ production plane with the photon linear polarization. The virtual loop corrections to the $\gamma g \to q\bar{q}$ subprocess such as (c) contribute to this correlation in the same order.

(4) In the case of single-hadron production at transverse momentum, leading-twist subprocesses are unlikely to dominate the cross section below $p_T < 8$ GeV because of the well-known trigger bias effect.[21] The complete set of higher-twist subprocesses, $gq \to Mq$, $q\bar{q} \to M\bar{M}$, $qM \to qM$, etc. must be taken into account.[22] As noted above, higher-twist contributions incorporate and replace the naive $k_\perp$ smearing procedure. In general, the higher-twist and leading-twist terms have a completely unrelated helicity and spin structure.

(5) Inclusive reactions involving incident hadrons also suffer from the fact that the initiating constituents can (Glauber) scatter elastically or inelastically on spectator constituents before the hard scattering process.[13] This effect can further reduce and complicate polarization correlations at low energies.

It is clear that the net gluon or quark polarization in inclusive reactions is always less than the hadron polarization, except possibly at the kinematic limit $x \to 1$. Conversely, we expect that the correlation between constituent and hadron helicity vanish as $x \to 0$, because of all the depolarizing mechanisms.

Although the spin correlations in the valence quark, sea quark and gluon inclusive distribution functions are generally controlled by non-perturbative bound state dynamics, it is still possible to make perturbative QCD predictions for the spin correlations of fast constituents in the limit $x \to 1$.[14] Note that at the edge of phase space only the valence Fock state contributes to the leading behavior in (1-x) since this requires the fewest number of spectators to stop. Furthermore, for $x \sim 1$, the struck quark is kinematically far off-shell and spacelike: $k^2 \sim -(k_\perp^2 + m^2)/(1-x) \to -\infty$ as long as the spectator masses are non-zero. The (1-x) power behavior of the structure functions can thus be computed in leading order from the simplest QCD tree diagrams; one finds

$$G_{q/N}(x) \underset{x\to 1}{\sim} \begin{cases} (1-x)^3 & \text{parallel } q,N \text{ helicity} \\ (1-x)^5 & \text{anti-parallel } q,N \text{ helicity} \end{cases} \quad . \qquad (3.1)$$

Equation (3.1) implies that the leading quark at $x \to 1$ always carries the helicity of the nucleon. This effect, in fact, seems to be consistent with the trend of the large x_{Bj} data obtained by the SLAC-Yale polarized-electron polarized-proton deep inelastic scattering experiment.[17] Since flavor and spin are correlated in the baryon valence wave function, perturbative QCD also then predicts $G_{u/p} \neq 2G_{d/p}$ at $x \to 1$. In fact, if we assume SU(6) symmetry, we have $G_{u/p} \Rightarrow 5G_{d/p}$ for $x \to 1$.[24]

In the case of gluon distribution in the nucleon, we recall that simple $q \to gq$ bremsstrahlung has a $[(1-x)^2 + 1]/x$ dependence where the two terms correspond respectively to parallel and anti-parallel gluon and initial quark helicity. At $x \sim 0$, the gluon helicity becomes uncorrelated. Convoluting this result with $G_{q/N}$ gives perturbative contributions

$$G_{g/N}(x) \underset{x\to 1}{\sim} \begin{cases} (1-x)^4 & \text{parallel } g,N \text{ helicity} \\ (1-x)^6 & \text{anti-parallel } q,N \text{ helicity} \end{cases} \quad . \qquad (3.2)$$

i.e., in general, the leading q or g constituents at $x \sim 1$ have helicities parallel to the nucleon helicity.

The analysis of meson structure functions at $x \sim 1$ is similar to that of the baryon, with two striking differences: (1) The controlling power behavior[23] of the leading-twist contribution is $(1-x)^2$ from perturbative QCD. The extra factor of (1-x) -- compared to what would have been expected from spectator counting -- can be attributed to the mismatch between the quark spin and that of the meson. (2) The longitudinal meson structure function has an anomalous non-scaling component[24] which is finite at $x \to 1$: $F_L(x,Q) \sim Cx^2/Q^2$. This higher twist term, which comes from the lepton scattering off an instantaneous fermion-line in light-cone perturbation theory, can be rigorously computed and normalized in perturbative QCD.[14,25] The crucial fact is that the wave function evolution and spectator transverse momentum integrations can be written directly in terms of a corresponding calculation of the meson form factor. A simplified result for the pion structure function in leading order is (in analogy to the Born driving term in the Witten structure function)[14]

$$F_L^{\pi}(x,Q) \cong \frac{2x^2}{Q^2} C_F \int_{m^2/(1-x)}^{Q^2} dk^2 \, \alpha_s(k^2) \, F_{\pi}(k^2) \qquad (3.3)$$

which numerically is $F_L \sim x^2/Q^2$ (GeV^2 units).

The dominance of the longitudinal structure functions in the fixed W limit for mesons is an essential prediction of perturbative QCD and is a special type of spin test. Perhaps the most dramatic consequence is in the Drell-Yan process $\pi p \to \ell^+\ell^- X$; one predicts[26] that for fixed pair mass Q, the angular distribution of the ℓ^+ (in the pair rest frame) will change from the conventional $(1+\cos^2\theta_+)$ distribution to $\sin^2(\theta_+)$ for pairs produced at large x_L. A recent analysis of the Chicago-Illinois-Princeton experiment[27] at FNAL appears to confirm the QCD high-twist prediction with about the expected normalization. It will be very important to check whether this effect is associated with the predicted C/Q^2 behavior. Striking evidence for a higher-twist component has also been reported in a Gargamelle[28] analysis of the quark fragmentation functions in $\nu p \to \pi^+\mu^- X$. The results yield a quark fragmentation distribution into positive charged hadrons which is consistent with the predicted form:[29] $dN^+/dzdy \sim B(1-z)^2 + (C/Q^2)(1-y)$ where the (1-y) behavior corresponds to a longitudinal structure function. It is also crucial to check that the $e^+e^- \to MX$ cross section becomes purely longitudinal ($\sin^2\theta$) at large z at moderate Q^2. The implications of this higher-twist contribution for meson production at large p_T will be discussed elsewhere.[30]

IV. SPIN EFFECTS IN EXCLUSIVE REACTIONS AND QCD SELECTION RULES

As we have seen in the previous section, test of spin effects in inclusive reactions are often complicated by a number of corrections and depolarizating mechanisms. In the case of large momentum transfer exclusive reactions, such as $e^+e^- \to M\bar{M}$ and $\gamma\gamma \to M\bar{M}$, only the minimal $|q\bar{q}\rangle$ Fock state of the meson contribute to leading order in m/Q, and one can obtain direct, rigorous checks of quark and gluon dynamics at short distance.[2-8] We will focus here on the use of exclusive reactions to experimentally determine the gluon spin and interactions in QCD. As we have discussed in the introduction, exclusive reactions involving large momentum transfer can be written in a form which factorizes the dynamics of the hard scattering quark and gluon processes from the physics of the hadronic wave functions. For example, the leading contribution to the nucleon form factor is given by the product of three factors: (a) the distribution amplitude, ϕ, for finding the three-quark valence state in the incoming proton; (b) the amplitude, T_H, for this quark state to scatter with the photon producing three quarks in the final state whose momenta are roughly collinear; and (c) the amplitude, ϕ^*, for this final quark state to reform into a hadron. Thus the magnetic form factor can be written (see Fig. 3a)[3,5,31]

$$G_M(Q^2) = \int_0^1 [dx] \int_0^1 [dy] \, \phi^*(y_i,\tilde{Q}_y) \, T_H(x_i,y_i,Q) \, \phi(x_i,\tilde{Q}_x) \left[1 + \mathcal{O}(m/Q)\right] \qquad (4.1)$$

where $\tilde{Q}_x \equiv \min_i (x_i Q)$.

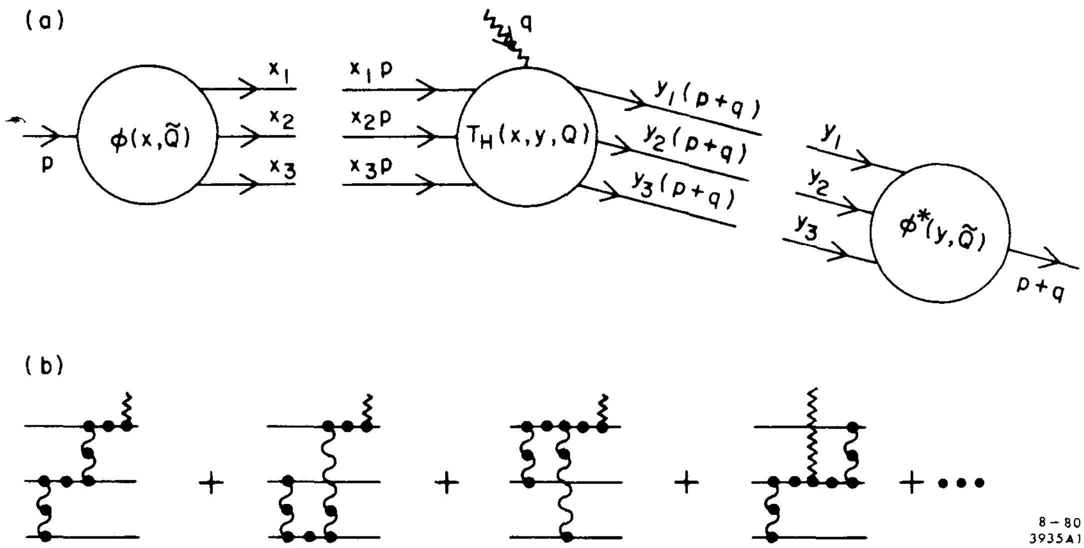

Fig. 3. (a) Factorized structure for the dominant QCD contribution to the baryon form factors. (b) Leading order (in $\alpha_s(Q^2)$) contributions to the hard scattering amplitude $(qqq)+\gamma^* \to (qqq)$. The black dots indicate renormalized vertex and self-energy insertions.

To leading order in $\alpha_s(Q^2)$, the "hard scattering amplitude" T_H is the sum of all Born diagrams for $\gamma^* + 3q \to 3q$ in perturbative QCD (see Fig. 3b). The transverse momentum fluctuations of the quarks in the initial and final protons are negligible relative to $q_\perp$, as are all particle masses. These can be ignored in T_H so that in effect each hadron is replaced by collinear on-shell valence partons. Since the final quarks are collinear, momentum of $\mathcal{O}(q_\perp) \to \infty$ must be transferred from quark line to quark line (via gluons) in T_H. This justifies the use of perturbation theory in computing T_H, since all internal propagators in the Born diagrams must then be off-shell by $\mathcal{O}(Q^2)$. Furthermore, the most important dynamical feature of the form factor -- its power-law fall-off -- can then be traced to the behavior of T_H, which falls for increasing Q^2 with a factor $\left(\alpha_s(Q^2)/Q^2\right)$ for each constituent, after the first, scattered from the incident to the final direction: i.e.,

$$T_H\left(x_i,y_i,Q\right) = \left(\frac{\alpha_s(Q^2)}{Q^2}\right)^2 T\left(x_i,y_i\right)\left[1+\mathcal{O}\left(\alpha_s(Q^2)\right)\right] \qquad (4.2)$$

where $\alpha_s(Q^2) = (4\pi/\beta)(\ell n\, Q^2/\Lambda^2)^{-1}$ is the running coupling constant (see Fig. 3b).

It is now clear that non-valence Fock states in the proton cannot contribute since all such states contain four or more constituents, each of which must be turned to the final direction. Thus T_H for these states falls as $\left(\alpha_s(Q^2)/Q^2\right)^3$ or faster and is negligible relative to (4.2) as $Q^2 \to \infty$. [This observation, while strictly true in light-cone gauge ($\eta\cdot A = A^+ = 0$), has a different interpretation in covariant gauges.] Thus non-valence ("sea") quarks and gluons in the proton do not contribute. The quantity $\phi(x,Q)$ is the "distribution amplitude" for finding the valence quark with light-cone fraction x_i in the hadron at relative separation $b_\perp \sim \mathcal{O}(1/Q)$. In fact,

$$\phi(x_i, s_i, Q) \equiv \prod_{i=1}^{n} \left[d_i^{-1}(Q^2)\right]^{\frac{1}{2}} \int^{k_{\perp i}^2 < Q^2} \left[d^2k_\perp\right] \psi^{(n)}(k_{i\perp}, x_i, s_i) \quad . \tag{4.3}$$

This amplitude is obviously process independent. It contains the essential physics of that part of the hadronic wave function which affects exclusive processes with large momentum transfer. The distribution amplitude is only weakly dependent on Q^2, and this dependence is completely specified by an evolution equation of the form (in leading order)[2-5]

$$Q^2 \frac{\partial}{\partial Q^2} \phi(x_i, Q) = \frac{\alpha_s(Q^2)}{4\pi} \int_0^1 \left[dy\right] V(x_i, y_i) \phi(y_i, Q) \tag{4.4}$$

where V can be computed from a single gluon exchange kernel. The general solution of this equation is

$$\phi(x_i, Q) = x_1 x_2 x_3 \sum_{n=0}^{\infty} a_n \left(\ell n \frac{Q^2}{\Lambda^2}\right)^{-\gamma_n} \tilde{\phi}_n(x_i) \quad . \tag{4.5}$$

Combining this expansion with Eqs. (1.5) and (1.6), we obtain the general form of G_M:

$$G_M(Q^2) = \left(\frac{\alpha_s(Q^2)}{Q^2}\right)^2 \sum_{n,m} b_{nm} \left(\ell n \frac{Q^2}{\Lambda^2}\right)^{-\gamma_n - \gamma_m} \quad . \tag{4.6}$$

The factorized form of Eq. (4.1) implies a simple space-time picture. The exchange of large transverse momentum in the hard scattering amplitude T_H occurs only when the relative separation of the constituents approaches the light-cone -- i.e., $-(z^{(i)} - z^{(j)})^2 \sim (z_\perp^{(i)} - z_\perp^{(j)})^2 \to \mathcal{O}(1/Q^2)$. The distribution amplitude ϕ is the probability amplitude for finding the valence quarks sufficiently near the light-cone; by the uncertainty principle, this corresponds to a momentum space wave function smeared over all $k^2 \lesssim 1/z_\perp^2 \sim Q^2$ as in Eq. (4.3). Each (polynomial) eigensolution $\tilde{\phi}_n(x_i)$ of the evolution equation is directly related to a term in the operator product expansion of the wave function evaluated near the light-cone. The eigenvalues γ_n are the corresponding anomalous dimensions.

Beyond leading order, the hard scattering amplitude T_H and the kernel for the distribution amplitude can be expanded in power series in $\alpha_s(Q^2)$. We note that the anomalous region $1-x_i < m/Q$ in Eq. (4.1) is suppressed in the baryon form factor by two powers of $\alpha_s(Q)$ relative to the leading hard scattering domain. The contribution of this region is further suppressed by a Sudakov quark form factor $S(Q^2)$ since for $1-x_i \sim m/Q$ a nearly on-shell struck quark must absorb the full momentum transfer Q without radiating gluons. Detailed discussions are given in Refs. 5 and 7. In the case of meson form factors, $F_{\pi\gamma}(Q^2)$, $\gamma\gamma \to M\bar{M}$, etc., the endpoint region $(1-x) \lesssim m/Q$ is suppressed by a kinematic factor of m/Q. This allows a direct proof of short distance dominance using operator product and renormalization group methods for these processes.[8]

We can generalize the above results to other exclusive processes $H_1H_2 \to H_3H_4$, $e^+e^- \to H_1 \ldots H_N$,[32] etc. where all invariants $p_i \cdot p_j$ scale with s, by computing the hard scattering amplitude $T_H(x_i, Q, \theta_{c.m.})$ -- calculated by replacing each hadron by collinear, on-shell valence quarks (with the approximate helicity) -- convoluted with the distribution amplitudes $\phi_H(x_i, Q)$ for finding the constituents in each hadron H with light-cone momentum fractions x_i at transverse separation $b_\perp \sim \mathcal{O}(1/\tilde{Q})$ with $\tilde{Q} = (\min x_i)Q$ (as in Fig. 4). By definition all the logarithmic collinear divergences are summed in the distribution amplitudes rather than T_H, which is collinear irreducible. In processes such as hadron-hadron scattering at large momentum transfer, the anomalous multiple-scattering Landshoff pinch contributions[9] which appear in tree-graph approximation are asymptotically suppressed by Sudakov form factors[5,10] for near on-shell quasi-exclusive quark-quark scattering. The possible phenomenological contribution of pinch contributions at sub-asymptotic momentum transfer is still uncertain; for simplicity we will ignore these contributions here.

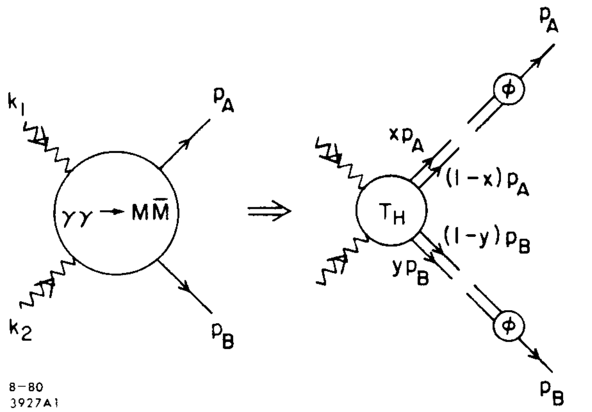

Fig. 4. Factorized structure for the process $\gamma\gamma \to M\bar{M}$ in perturbative QCD.

Note that the axial symmetric $d^2k_\perp$ integration in the definition of $\phi(x_i, Q)$ projects out only $L_z = 0$ components for the valence wave function $\psi^V(k_{\perp i}, x_i, s_i)$.[5,11] Thus the sum of the quark helicities in ϕ equals the hadron helicity. On the other hand, T_H conserves total quark helicities to leading order in m/Q because of the $\bar{\psi}\gamma_\mu\psi$ vector gluon interaction. The most important dynamical features of the hadronic amplitudes at large momentum transfer -- their power-law fall-off in Q^2, their angular dependence and their helicity dependence -- are thus determined by the Born contribution to $T_H(x_i, Q, \theta_{c.m.})$. We are thus led to a large number of detailed, experimentally testable, predictions of QCD which critically reflect its elementary scaling and spin properties at short distances. In particular there are two sets of universal predictions of QCD which follow from the properties of $T_H(x, Q^2, \theta_{c.m.})$ to leading order in 1/Q and to all orders in $\alpha_s(Q^2)$ (5.11):

(A) The dimensional counting rules[33] for the power-law behavior of exclusive processes: $\mathcal{M} \sim Q^{4-n}$, where n is the minimum number of external elementary fields (leptons, quarks, transversely-polarized gluons or photons) participating in T_H.

(B) The QCD helicity selection rule

$$\Delta h = h^{tot}_{initial} - h^{tot}_{final} = 0 \quad ; \qquad (4.7)$$

i.e., total hadron helicity is conserved. In the case of space-like electromagnetic or weak form factors, the fact that the current can only change the helicity by $\Delta J_z = h_I + h_F \le 1$ in the Breit frame leads to an even more restrictive rule:

$$h_{initial} = h_{final} = \pm\frac{1}{2} \text{ or } 0 \qquad (4.8)$$

i.e., minimal hadron helicity. We emphasize that these helicity selection rules are special features of vector gluon gauge theories and the fact that the valence Fock states dominate the amplitude at large momentum transfer. Ignoring the Sudakov-suppressed contributions, the results are true to all orders in $\alpha_s(Q^2)$. The QCD dimensional counting rule (A) for the power behavior of fixed angle scattering amplitudes and form factors appear to be consistent with experiment; detailed reviews are given in Refs. 5 and 34.

Let us now discuss some of the consequences of the QCD helicity rules. For example, space-like form factors for processes in which the hadron's helicity is changed, or in which the initial or final hadron has helicity $|h| \geq 1$ are suppressed by powers of m/Q where m is an effective quark mass. The QCD selection rules thus imply power-law suppression of $F_2^N(Q^2)/F_1^N(Q^2)$ and $\gamma^* p \to \Delta(h = 3/2)$. In the case of the deuteron, dimensional counting and helicity conservation predicts[33,35] $F_D(Q^2) \sim (1/Q^2)^5$ at large Q^2 (modulo logarithmic factors) for the dominant helicity zero → helicity zero transition form factor.

An important feature of the hard scattering perturbative QCD predictions is that all of the helicity-conserving electroweak baryon form factors can be expressed[31] as linear combinations of just two basic form factors -- $G_\parallel(Q^2)$ and $G_{\overline{\parallel}}(Q^2)$ -- corresponding to amplitudes in which the current interacts with a valence quark with helicity parallel or anti-parallel to the helicity of the nucleons, respectively. The coefficients are determined by the corresponding $SU(2)_L \times U(1)$ quark charges. Thus the nucleon magnetic form factors $G_M^n(Q^2)$ and $G_M^p(Q^2)$ are sufficient to predict all of the electroweak nucleon form factors. The assumption of the standard helicity-flavor symmetry for the baryon wave functions at short distances then leads to the specification of all the leading electroweak octet and decouplet form factors. The spatial wave functions can be assumed to be symmetrical with respect to the quarks having the same helicity, a feature which is preserved under perturbative QCD evolution. At $Q^2 \to \infty$, the spatial wave function becomes totally symmetric, $\phi_B(x_i,Q) \to x_1x_2x_3(\log Q^2/\Lambda^2)^{-\gamma_B}$, and thus the helicity-flavor structure of the baryon states satisfies exact SU(6) symmetry. The detailed results are given in Ref. 31. The ratio $G_M^p(Q^2)/G_M^n(Q^2)$ is particularly sensitive to the shape of the distribution amplitude and the spin of the gluons. Specific productions are given in Ref. 5.

For the case of timelike processes at large Q^2, $e^+e^- \to \gamma^* \to H_A + H_B$, the hadrons are predicted to be dominantly produced with opposite helicity $h_A = -h_B = 0$ or $\pm\frac{1}{2}$, since total hadron helicity is conserved. All two-body angular distributions have center-of-mass angular distributions given by $d^J_{J_z,\lambda}(\theta)$ where $\lambda = h_A - h_B$ and $J_z^\gamma = \pm 1$ $(m_e^2/s \to 0)$. The angular distirbution of the cross sections must then be proportional to $(1 + \cos^2\theta)$ for $\lambda = \pm 1$ or $\sin^2\theta$ for $\lambda = 0$; i.e., QCD predicts (modulo calculable log Q^2 factors)[11]

$$\frac{\frac{d\sigma}{d\Omega}\left(e^+e^- \to H_AH_B\right)}{\sigma\left(e^+e^- \to \mu^+\mu^-\right)} \propto \begin{cases} \dfrac{1+\cos^2\theta}{(Q^2)^4} \; , & h_A = -h_B = \pm\dfrac{1}{2} \\[2ex] \dfrac{\sin^2\theta}{(Q^2)^2} \; , & h_A = -h_B = 0 \end{cases} \tag{4.9}$$

for the leading power behavior in all orders of perturbation theory for baryon and meson pairs, respectively. In particular, we predict power-law suppression of $e^+e^- \to \pi\rho$, KK^* (since the vector meson must be produced with non-zero helicity) as well as $e^+e^- \rightarrow \rho(h=0)\bar{\rho}(h=1)$, $\rho(h=1)\bar{\rho}(h=\pm 1)$, $\Delta(h=3/2) \cdot \bar{\Delta}(h = \text{any helicity})$.

It is important to note that all of these results hold for heavy quarkonium decay $\psi,\psi',T \rightarrow H_A + H_B$; the annihilation of heavy quarks via an aribtrary number of vector gluons (see Fig. 5) into the light quarks ($m_u^2, m_d^2 << Q^2$) again conserves total hadron helicity, and the quarkonium state produced in e^+e^- must have spin ±1 along the beam direction. In fact, there is already considerable experimental data for ψ and ψ' decays which can be used to test the QCD predictions. The SPEAR Mark II data[36] for $\psi \rightarrow p\bar{p}$ are consistent with the predicted $1 + \cos^2\theta$ angular distribution. In contrast, scalar, pseudoscalar, or tensor gluon theories predict a $\sin^2\theta$ distribution in leading order! The power-law behavior predicted by QCD can be checked by comparing ψ and ψ' branching ratios into baryon pairs. The theory predicts the leading power behavior

$$\frac{BR(\psi \rightarrow p\bar{p})}{BR(\psi' \rightarrow p\bar{p})} \sim \left(\frac{M_{\psi'}}{M_\psi}\right)^8 \tag{4.10}$$

where $BR(\psi \rightarrow p\bar{p}) = \Gamma(\psi \rightarrow p\bar{p})/\Gamma(\psi \rightarrow \text{light-quark hadrons})$ removes the dependence on the ψ and ψ' wave functions. The data[36] is consistent with a ratio $(M_{\psi'}/M_\psi)^n$ with $n = 10 \pm 3$. The $\psi \rightarrow B\bar{B}$ data can also be used to normalize the baryon distribution amplitudes to leading order in $\alpha_s(M_\psi^2)$, and check predictions for relative and absolute magnitudes of the decay rate.

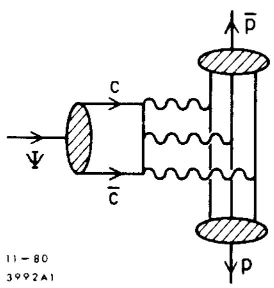

Fig. 5. The leading $\alpha_s^3(Q^2)$ contribution to the $\psi \rightarrow B\bar{B}$ amplitude. The calculation of this amplitude involves a convolution with the baryon distribution amplitude $\phi_B(x_i,Q)$.

The reactions $\psi,\psi' \rightarrow \pi\rho, KK^*, \ldots$ are suppressed in QCD because they violate hadron helicity conservation. One expects $BR(\psi \rightarrow \pi\rho)/BR(\psi' \rightarrow \pi\rho) \sim (M_{\psi'}/M_\psi)^n$ with $n \geq 6$ in QCD, whereas $n=4$ is possible in scalar or tensor gluon theories. The existing data shows that $n \geq 10$, an even stronger suppression than is expected. It is also curious that the $\pi\rho$ and KK^* rates are roughly comparable, since helicity flip amplitude are usually associated with factors of the (running) quark mass.

We can also analyze heavy pseudoscalar ($J^P = 0^-$) quarkonia decays in perturbative QCD. The decays $\chi \rightarrow B\bar{B}$ should be suppressed since quark helicity conservation requires $h_B = -h_{\bar{B}}$, in conflict with total angular momentum conservation. This is an important test of QCD since this suppression would not be present in theories with scalar or tensor glue. We would also expect suppression of χ production in $p\bar{p}$ annihilation.[37]

V. HADRON HELICITY CONSERVATION AND FIXED ANGLE SCATTERING

As we have discussed in Sect. IV, total hadron helicity conservation applies to any large momentum transfer amplitudes which factorizes into a hard scattering subprocess matrix element $T_H(x_i,Q)$ convoluted with the process-independent amplitudes $\phi_H(x,Q)$. This factorization (see Fig. 4) can be rigorously demonstrated in perturbative QCD for processes such as the fixed angle two-photon reactions[8] $\gamma\gamma \to M\bar{M}$, where $M = \pi,\rho,K,K^*$, etc. These reactions are now being using e^+e^- collisions. Landshoff pinch contributions are suppressed by a power of Q^2 ($= p_T^2$) at the Born level, even before taking into account the Sudakov form factor suppression of such amplitudes. In addition, there is no anomalous contribution from quark currents at the edge of phase space $(1-x) \lesssim m/Q$. Total helicity conservation implies that to all orders in $\alpha_s(Q^2)$ the leading power contribution for $\gamma\gamma \to M\bar{M}$ at fixed $\theta_{c.m.}$ produce mesons with equal and opposite helicities; $\gamma\gamma \to \rho(h=0)\bar{\rho}(h=\pm 1)$, $\gamma\gamma \to \rho(h=+1)\bar{\rho}(h=-1)$ are suppressed. Note that $\gamma\gamma \to \rho^+(h=+1)\rho^-(h=-1)$ is not suppressed even though $\gamma^* \to \rho^+(h=+1)\rho^-(h=-1)$ is nonleading. Complete predictions (to leading order in $\alpha_s(Q^2)$) for the $\gamma\gamma \to M\bar{M}$ cross sections, including scaling behavior, normalization, angular distributions and helicity dependence are given in Ref. 8. These processes can also be used as a sensitive probe of the structure of the distribution amplitudes. We emphasize that the $\gamma\gamma \to M\bar{M}$ predictions provide a detailed, definitive check of the basic dynamics of perturbative QCD. The corrections are rigorously of higher order in m/Q.

We also predict that hadron helicity conservation will be satisfied at asymptotic momentum transfer in fixed-angle scattering reactions such as $\pi p \to \pi p$, $\gamma p \to \gamma p$, $pp \to pp$. However, unlike the two-photon $\to \gamma\gamma \to M\bar{M}$ amplitudes, there are subasymptotic contributions from near-on-shell multiple scattering amplitudes (the Landshoff pinch contribution) and $x \sim 1$ quark scattering contributions which are not necessarily dominated by short distance dynamics. Although these contributions are suppressed at asymptotic Q^2 by Sudakov form factors, they could play a role at moderate momentum transfer. A careful analysis of helicity effects in these reactions may lead to a check on the importance of these contributions.

In the case of pseudoscalar meson-nucleon scattering, hadron helicity conservation implies that the backward peak is power-law suppressed relative to the forward peak since angular momentum cannot be conserved at 180°. This is consistent with πp and Kp scattering data.[34]

At present, the most detailed check of the helicity dependence of hadron amplitudes is provided in $pp \to pp$ scattering. In general there are five independent parity-conserving and time-reversal-invariant helicity amplitudes: $\mathcal{M}(++ \to ++)$, $\mathcal{M}(+- \to +-)$, $\mathcal{M}(-+ \to +-)$, $\mathcal{M}(++ \to +-)$, $\mathcal{M}(-- \to ++)$. The QCD selection rule $h_{initial} = h_{final}$ implies that $\mathcal{M}(++ \to +-)$ and $\mathcal{M}(-- \to ++)$ are power-law suppressed. The helicity-conserving amplitudes then are predicted to scale (modulo logarithms) as a square of the nucleon form factor, yielding the dimensional-counting prediction[5]

$$\frac{d\sigma}{dt}(pp \to pp) \sim \left[\frac{\alpha_s(p_T^2)}{p_T^2}\right]^{10} f^2(\theta_{c.m.}) \qquad (5.1)$$

(modulo logarithmic corrections from the evolution of the baryon distribution amplitudes). The pp → pp data is in fact consistent within a factor of two with fixed angle scaling $(p_T^2)^{9.7\pm0.5}\, d\sigma/dt \sim$ const. as the cross section falls than five decades in the range $4 < p_T^2 < 12$ GeV^2, $38^\circ < \theta_{c.m.} < 90^\circ$.[38] This implies that the variation of $\alpha_s(p_T^2)$ is very slow in this domain. On the other hand, the data appear to be systematically oscillating[39] about the $s^{10}\, d\sigma/dt \sim$ const. prediction, possibly suggesting the presence of an interfering subasymptotic amplitude.

The most sensitive tests of the hard scattering QCD prediction involve the polarization effects. The spin asymmetry A_{NN} is defined as

$$A_{NN} = \frac{\frac{d\sigma}{dt}(\uparrow\uparrow) + \frac{d\sigma}{dt}(\downarrow\downarrow) - \frac{d\sigma}{dt}(\uparrow\downarrow) - \frac{d\sigma}{dt}(\downarrow\uparrow)}{\frac{d\sigma}{dt}(\uparrow\uparrow) + \frac{d\sigma}{dt}(\downarrow\downarrow) + \frac{d\sigma}{dt}(\uparrow\downarrow) - \frac{d\sigma}{dt}(\downarrow\uparrow)} \quad , \qquad (5.2)$$

which measures the difference of cross sections when both nucleons are polarized parallel to the normal ($\hat{x}$) of the scattering plane or are anti-parallel. Similarly A_{LL} refers to the polarization asymmetry where the initial spins are polarized along the laboratory beam direction ($\hat{z}$) versus anti-parallel spins, and A_{SS} refers to initial spins polarized (sideways) along the third direction ($\hat{y}$).

For the scattering of identical particles at 90° all amplitudes involving a single helicity flip vanish, e.g., (++ → +−). This implies the sum rule[40]

$$A_{NN} - A_{LL} - A_{SS} = 1 \qquad \left(\theta_{c.m.} = 90^\circ\right) \quad . \qquad (5.3)$$

If in addition the double-flip amplitude (++ → −−) vanishes, as in the case of the perturbative QCD predictions, then we have $A_{NN} = -A_{SS}$ (all angles) and the above sum rule becomes[41,42]

$$2A_{NN} - A_{LL} = 1 \qquad \left(\theta_{c.m.} = 90^\circ\right) \quad . \qquad (5.4)$$

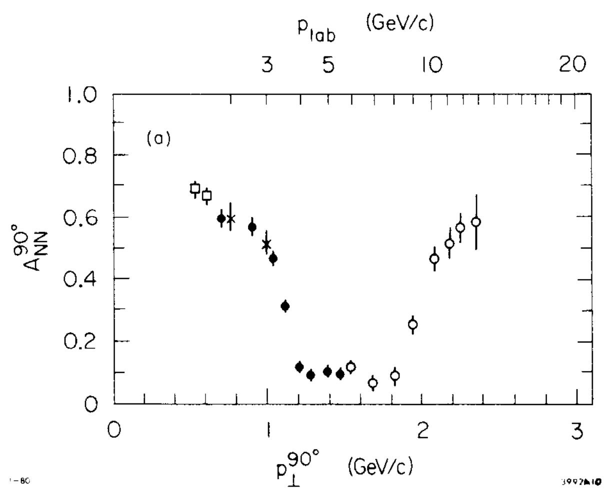

Fig. 6. Data for the spin asymmetry A_{NN} (normal to the scattering plane) for pp scattering at 90° as a function of p_{lab} and p_T.

The striking Crabb et al.[43] Argonne measurements for A_{NN} (see Fig. 6) can now be combined with new (preliminary results) for A_{LL} at 90° and $p_{lab} = 11.75$ GeV ($p_T \cong 2.4$ GeV) reported by Yokasawa at this meeting: $2A_{NN} - A_{LL} \cong 2(0.58 \pm 0.04) - (0.18 \pm 0.09) = 0.98 \pm 0.17$, which is consistent with helicity conservation. On the other hand, it should be noted that the change of A_{NN} is very rapid: $A_{NN} \cong 0.05$ at $\theta_{c.m.} \leq 60^\circ$ to $A_{NN} \cong 0.60$ at $\theta_{c.m.} \geq 70^\circ$, which is in marked contrast to the generally smooth behavior predicted from calculations of T_H for proton-proton scattering. For example, the set of hard scattering diagrams (see Fig. 7a) with only quark inter-

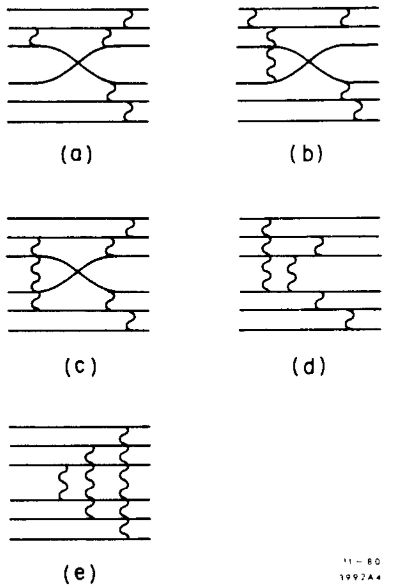

Fig. 7. Representative quark and gluon exchange contributions to the pp → pp scattering amplitude. Diagram (e) represents a Landshoff pinch contribution, corresponding to 3 sequential quark-quark scattering with $\theta_{qq} \cong \theta_{pp}^{c.m.}$.

change between the nucleons (which gives a good representation of the pp → pp angular distribution and crossing to $p\bar{p} \to p\bar{p}$) leads to the simple prediction[41,42]

$$A_{NN} = -A_{LL} = -A_{SS} = 1/3 \quad \left(\theta_{c.m.} = 90^{o}\right) \quad , \qquad (5.5)$$

with a very slow variation (≤ 2%) over all $\theta_{c.m.}$. Diagrams with quark interchange plus gluon exchange between nucleons (as in Fig. 7b) give a smaller value for A_{NN}.[45] The angular distribution predicted for diagrams with only gluon exchange (Fig. 7d,e) is incompatible with the large angle data; furthermore, if these amplitudes are normalized to the small angle regime then they are negligible at 90°.[5]

At this stage, there does not seem to be a convincing explanation of the nucleon-nucleon polarization effects at large angle.[46] It seems possible that whatever interference of amplitudes causes the oscillation of $d\sigma/dt$ around the smooth s^{-10} behavior can also lead to striking interference effects in the polarization correlations.[39,41] One possibility is that the quark interchange amplitude is asymptotically dominant, but that in the present experimental range there is significant interference with multi-Regge exchange contributions.[41] An important point is that the Landshoff pinch contribution for pp → pp scattering includes three sequential qq → qq scatterings each at approximately the same momentum transfer $\hat{t} \sim 1/9$. Since $|\hat{t}| < 1.1$ GeV^2 is not very large, ordinary Reggeon exchange could still be playing a role in the quark-quark scattering amplitude. Unfortunately, the introduction of such contributions necessary includes extra parameters and considerable model-dependence. Nevertheless, a simple estimate of the rotating phase associated with triple Regge exchange is consistent with the interference pattern indicated by the pp → pp large angle data.

VI. OBLATE JETS AS A TEST OF GLUON SPIN

The tests of gluon spin which we have discussed thus far have necessarily all been indirect. The most dramatic confirmation of confined gluon spin would be the direct measurement of gluon jet polarization. As shown in Ref. 47, the linear polarization of a produced gluon could be reflected in terms of the orientation of the axies of an oblate jet distribution, as defined by momentum weighted hadron fragmentation distributions

$$\sum_H z\, D_{H/g}(z,\phi,\, Q) \qquad .$$

A search for jet oblateness would be particularly interesting in events such as $e^+e^- \to q\bar{q}g$ (charmed $q,\bar{q}$) where the gluon jet can be specified, and the gluon linear polarization can be predicted. These tests, are complicated, however, by the fact that the subprocesses $g \to q\bar{q}$ and $g \to gg$ in the gluon jet evolution tends to give cancelling correlations with the gluon linear polarization.[47]

VII. THE MAGNETIC MOMENTS OF HADRONS

Another important area of spin effects in QCD is the prediction of the magnetic moment of the baryons. An analysis of this problem for arbitrary composite systems is given in Ref. 48 using the Fock state description of the hadron state as described in Sect. II. The general result for the anomalous moment is

$$\frac{a}{M} = -\sum_{n,j} e_j \int \left[dx\right] \int \left[d^2k_\perp\right] \sum_{i \neq j} \psi_{p\uparrow}^{(n)}\; x_i \left(\frac{\partial}{\partial k_{1i}} + i\,\frac{\partial}{\partial k_{2i}}\right) \psi_{p\downarrow}^{(n)} \qquad (7.1)$$

where $\psi_{p\downarrow}^{(n)}(x_i,\vec{k}_{\perp i},s_i)$ is the Fock state wave function for the nth state with initial spin opposite to the z-direction. The sum is over all Fock states (n) and charged constituents j.

A typical contribution to the anomalous moment is of order $\delta a^{(n)} \sim MR^{(n)}$ where $R^{(n)}$ is the mean radius of an internal Fock state (n) which contributes to the nucleon positive and negative helicity. Thus the magnetic moment of a nucleon approaches the Dirac value as the binding energy of the quark and gluon Fock states becomes arbitrary large.[48,49,50] This result has also been derived using the Drell-Hearn-Gerasimov sum rule[48,49] or sidewise dispersion relations.[50] It should be noted that the standard non-relativistic formula[51] $\vec{\mu} = \sum_{i=1}^{n} \vec{\mu}_i$ predicts $\mu \to 0$ in the strong binding limit; its validity is restricted to weak binding situations and thus it has doubtful applicability to quark model calculation. In particular, the non-relativistic formula, while Eq. (7.1), neglects the effects of the Lorentz boost in the interaction with the external field.

The complete calculation of magnetic moments, charge radii, and general form factors of hadrons will have to take into account the full relativistic and Fock state structure of QCD.[48,52] The systematic consideration of these quantities plus others such as G_A/G_V, $J=0$ matrix elements may lead to tight constraints on the details of hadron structure. We also note that the QCD analysis applies to nuclear bound states. In particular, deuteron form factors at large momentum transfer, and the parity-violating photon polarization seen in $np \to d\gamma$ capture evidentally require consideration of the quark and gluon degrees of freedom of the deuteron state at short distances, including color-polarized six quark states. Further discussion of these problems can be found in Ref. 53.

VIII. CONCLUSIONS

Let us briefly summarize some of the main conclusions concerning the testing of QCD spin effects:

(1) Tests involving inclusive processes, though usually straightforward in concept, become definitive only in the asymptotic large momentum transfer limit where the leading twist contributions of lowest order in $\alpha_s(Q^2)$ become dominant. An important example is that longitudinal/scalar gluon contributions must be taken into account at the next to leading order in $\alpha_s(Q^2)$. Among, the most promising inclusive tests of gluon spin are the distributions[54] for $e^+e^- \to q\bar{q}g$ and processes involving direct photon production. It is also important to study predictions for spin correlations in structure functions at large x, especially the longitudinal C/Q^2 contribution to the meson structure and fragmentation functions.[24,26,14] We also note that many of the complications involving hadron structure functions are elegantly circumvented in photon-induced reactions since

(i) the photon can enter directly into the leading twist subprocess; and
(ii) the photon structure function[55] for any polarization is determined by perturbative QCD.

(2) Rigorous, "first class" tests of perturbative QCD dynamics to all orders in $\alpha_s(Q^2)$ are possible in large momentum transfer exclusive reactions such as $e^+e^- \to M\bar{M}$ and $\gamma\gamma \to M\bar{M}$. QCD spin selection rules which govern the leading hard scattering contributions to exclusive reactions are given in Sect. IV. Among the most important tests of vector gluon spin are the predicted power-law suppression of $\gamma^*p \to N^*(h=3/2)$, $e^+e^- \to \pi\rho$, $e^+e^- \to \psi \to \pi\rho$, and $\chi \to p\bar{p}$. The e^+e^- annihilation processes are particularly useful for studying the helicity structure of amplitudes because the intermediate virtual photon or resonance is always polarized along the beam axis, allowing hadronic helicity conservation to be verified simply by measuring the angular distribution of the final states. In the case of $\psi \to p\bar{p}$, experiments[36] clearly indicate an angular distribution proportional to $1+\cos^2\theta$. This is strong evidence favoring a vector gluon since scalar or tensor gluon exchange models would predict a $\sin^2\theta$ distribution to leading order in the coupling constant. The predicted suppression of $\psi \to \pi\rho$ is also evident in the data. Large transverse momentum exclusive processes are particularly well suited to the study of the spin structure of hadron dyanmics at short distances because the hadron helicity equals the sum of the helicities of its valence quark constituents in all dominant amplitudes. The spin correlations are much stronger in exclusive than in inclusive reactions because of the absence of depolarizing effects such as non-interacting spectators, non-zero constituent orbital angular momentum and non-valence Fock state contributions. On the other hand, exclusive reactions such as $pp \to pp$, at the present accessible momentum transfers, may be complicated by subasymptotic contributions not controlled by short distance dynamics.

(3) Finally, we note that the calculation of the magnetic moment and other spin-dependent hadronic parameters will require consideration of the full relativistic Fock state structure of the hadronic wave functions. In particular, unless quark binding effects are negligible, the spin structure of the nucleon wave function -- Lorentz-boosted from rest -- is required to compute the nucleon magnetic moments.

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REFERENCES

1. R. K. Ellis, H. Georgi, M. Machacek, H. D. Politzer and G. Ross, Phys. Lett 78B, 281 (1978); Nucl. Phys. B152, 285 (1979). S. Gupta, A. H. Mueller, Phys. Rev. D20, 118 (1979). D. Amati, R. Petronzio and G. Veneziano, Nucl. Phys. B140, 54 and B146, 29 (1978). C. Llewellyn Smith, Proc. of the XVII Int. Universitätswochen für Kernphysik, Schladming, Austria (1978) and Oxford preprint 47/48. The expansion of QCD factorization to polarized processes is discussed by I. Antoniadis and C. Kounnas, Ecole Polytechnique preprint (1980). A comprehensive review of perturbative QCD for inclusive reactions is given in A. J. Buras, Rev. Mod. Phys., 52, #1, (1980).
2. S. J. Brodsky and G. P. Lepage (SLAC-PUBs-2294 and 2298), published in "Quantum Chromodynamics," Proceedings of the 1978 La Jolla Institute, W. Frazer and F. Henyey, eds., AIP (1979), and SLAC-PUB-2447, published in the 1979 Summer Institute on Particle Physics at SLAC, Stanford, CA. G. P. Lepage and S. J. Brodsky, Phys. Lett. 87B, 359 (1979). S. J. Brodsky and G. P. Lepage, SLAC-PUB-2605, published in the Proceedings of the XI Int. Symposium on Multiparticle Dynamics, Bruges, Belgium (1980).
3. The analysis of hadronic wavefunctions at short distances in QCD in terms of the operator product on the light-cone is given in S. J. Brodsky, Y. Frishman, G. P. Lepage and C. Sachrajda, Phys. Lett. 91B, 239 (1980).
4. Baryon form factor calculations in QCD are given in G. P. Lepage and S. J. Brodsky, Phys. Rev. Lett. 43, 545, 1625(E) (1979); M. Peskin, Phys. Lett. 88B, 128 (1979).
5. S. J. Brodsky and G. P. Lepage, Phys. Rev. D22, 2157 (1980).
6. Related analyses of the meson form factors in perturbative QCD are given in G. R. Farrar and D. R. Jackson, Phys. Rev. Lett. 43, 246 (1979); V. L. Chernyak and A. R. Vhitnishii, JETP Letters 25, 11 (1977); A. V. Efremov and A. V. Radyushkin, Dubna preprints JINR-E2-1135, 11983 and 12389; G. Parisi, Phys. Lett. 84B, 225 (1979); and M. Chase, Nucl. Phys. B167, 125 (1980). A complete renormalization group analysis has been given by A. Duncan and A. H. Mueller, Phys. Lett. 90B, 159 (1980); Phys. Rev. D21, 1636 (1980).
7. A. Duncan and A. H. Mueller, ref. 6; A. Duncan, to be published in the Proceedings of the Symposium on Topical Questions in QCD, Copenhagen, 1980, and A. H. Meuller, published in the Proceedings of the XIth Int. Symposium on Multiparticle Dynamics, Bruges, Belgium, 1980.

8. S. J. Brodsky and G. P. Lepage, SLAC-PUB-2587, to be published in the Proceedings of the XXthe Int. Conf. on High Energy Physics, Madison, Wisconsin, 1980.
9. P. V. Landshoff, Phys. Rev. D10, 1024 (1974); P. Cvitanovic, Phys. Rev. D10, 338 (1974); S. J. Brodsky and G. Farrar, Phys. Rev. D11, 1309 (1975).
10. P. V. Landshoff and D. J. Pritchard, Cambridge University preprint DAMTP 80/04.
11. S. J. Brodsky and G. P. Lepage, in preparation.
12. S. J. Brodsky, T. Huang and G. P. Lepage, SLAC-PUB-2540, contributed to the 20th Int. Conf. on High Energy Physics, Madison, Wisconsin, 1980. T. Huang, SLAC-PUB-2580, to be published in the Proc. of the 20th Int. Conf. on High Energy Physics, Madison, Wisconsin, 1980.
13. An analysis of initial state and Glauber corrections in QCD will be given in G. Bodwin, S. J. Brodsky and G. P. Lepage (in preparation).
14. S. J. Brodsky and G. P. Lepage, SLAC-PUB-2447, published in the 1979 Summer Inst. on Particle Physics at SLAC, Stanford, California, and SLAC-PUB-2601, to be published in the Proc. of the XXth Int. Conf on High Energy Physics, Madison, Wisconsin, 1980. M. Ciafaloni, Pisa preprint S.N.S. 4/80 to be published in the Proc. of the XXth Int. Conf. on High Energy Physics, Madison, Wisconson, 1980. D. Amati, A. Bassetto, M. Ciafaloni, G. Marchesini, G. Veneziano, Nucl. Phys. B173, 429 (1980).
15. W. E. Caswell, R. R. Horgan, S. J. Brodsky, Phys. Rev. D18, 2415 (1978); H. Politzer, CALT 68-789, to be published in the Proc. of the Symposium on Topical Questions in QCD, Copenhagen, Denmark., and Nucl. Phys. B172, 349 (1980).
16. The basic discussion for spin effects for leading twist QCD contributions is given in J. Babcock, E. Monsay and D. Sivers, Phys. Rev. Lett. 40, 1161 (1978); Phys. Rev. D19, 1483 (1979) and K. Hidaka, E. Monsay and D. Sivers, Phys. Rev. D5, 1503 (1979). A comprehensive survey of polarization symmetries provided by QCD is given in F. Baldracchini, N. S. Craigie, V. Roberto and M. Socolovsky, Trieste preprint IC/80/148. Other recent papers include K. Hidaka, Westfield preprint (1980) and this conference; M. Fontannaz. B. Pire, D. Schiff, Orsay preprint LPTHE 80/25 (1980).
17. V. Hughes, this conference. G. Baum, et al., SLAC-PUB-2557 (1980) and contributed to the XXth Int. Conf. on High Energy Physics, Madison, Wisconsin, 1980. M. J. Alguard, et al., Phys. Rev. Lett. 41, 70 (1978), 37, 1258 (1976). The data are consistent with the predictions of G. Farrar and D. Jackson, Phys. Rev. Lett 35, 1416 (1975) that the quark at $x \to 1$ carries to nucleon helicity. See also R. Carlitz and J. Kaur, Phys. Rev. Lett. 38, 673, 1102(E) (1967); J. Kaur, Nucl. Phys. B128, 219 (1977).
18. J. D. Bjorken, Phys. Rev. 148, 1467 (1966), D1, 1376 (1970). The basic sum-rule for the probability distributions is

$$\sum_q e_q^2 \int_o^1 dx \left[G_{q\uparrow/p\uparrow} - G_{q\downarrow/p\uparrow}\right] - \left[G_{q\uparrow/n\uparrow} - G_{q\downarrow/n\uparrow}\right] = \frac{1}{3}\,\frac{G_A}{G_V}$$

where $G_A/G_V \cong 1.23$ is the ratio of axial vector and vector coupling constant. See R. P. Feynman, Photon-Hadron Interactions (Benjamin, Reading, Mass., 1972). The order α_s corrections to the sum-rule for the spin-dependent nonsinglet structure function $g_1(x,Q)$ is given in J. Kodaira, S. Matsuda, K. Sasaki and T. Uematsu, Nucl. Phys. B159, 99, (1979), and Phys. Rev. D20, 627 (1979), and I. Antoniadis and C. Kounnas, ref. 1.

19. A. Devoto, J. Pumplin, W. Repko, G. L. Kane, Phys. Rev. Lett. 43, 1062, 1540(E) (1979); D. W. Duke and J. F. Owens, Phys. Rev. Lett. 44, 1173 (1980).
20. S. J. Brodsky and T. DeGrand, in preparation.
21. S. D. Ellis, M. Jacob, P. V. Landshoff, Nucl. Phys. B108, 93 (1976). R. Blankenbecler, S. J. Brodsky and J. F. Gunion, Phys. Rev. D18, 900 (1978); D. Sivers, S. J. Brodsky and R. Blankenbecler, Phys. Reports 23C, 1 (1976). M. K. Chase and W. J. Stirling, Nucl. Phys. B133, 157 (1978).
22. See, e.g., E. L. Berger, T. Gottschalk, D. Sivers, Argone preprint ANL-HEP-CP-80-58/59 to appear in the Proc. of the 10th Int. Conf on High Energy Physics, Madison, Wisconsin, 1980.
23. Z. F. Ezawa, Nuovo Cimento 23A, 271 (1974).
24. G. R. Farrar and D. R. Jackson, ref. 17. A. I. Vainshtain and V. I. Zakharov, Phys. Rev. Lett. 72B, 368 (1978).
25. A. Duncan and A. H. Mueller, Phys. Lett. 90B, 159 (1980).
26. S. J. Brodsky and E. L. Berger, Phys. Rev. Lett. 42, 940 (1979).
27. K. J. Anderson, et al., Phys. Rev. Lett. 73, 1219 (1979).
28. [2605-40]. M. Haguenauer, et al., CERN/EP 80-202 (1980).
29. E. L. Berger, Phys. Lett. 89B, 241 (1980). See also refs. 24 and 26.
30. E. L. Berger, S. J. Brodsky and G. P. Lepage, in preparation.
31. S. J. Brodsky, G. P. Lepage, and S. A. A. Zaidi, SLAC-PUB-2588 (1980).
32. Annihilation into multihadrons has been analyzed by S. Gupta (private communication).
33. S. J. Brodsky and G. R. Farrar, Phys. Rev. Lett. 31, 1153 (1973) and Phys. Rev. D4, 1309 (1975). V. A. Matveev, R. M. Muradyan and A. V. Tavkhelidze, Lett. Nuovo Cimento 7, 719 (1973).
34. D. Sivers, et al., ref. 21.
35. S. J. Brodsky and B. T. Chertok, Phys. Rev. Lett. 37, 269 (1976); Phys. Rev. D14, 3003 (1976). S. J. Brodsky and G. P. Lepage, SLAC-PUB-2595, to appear in the Proc. of the 9th Int. Conf. on the Few Body Problem, Eugene, Oregon, 1980.
36. I. Peruzzi, et al., Phys. Rev. D17, 2901 (1978).
37. We wish to thank M. Jacob for discussions on this point.
38. P. V. Landshoff and J. C. Polkinghorne, Phys. Lett. 44B, 293 (1973).
39. A. W. Hendry, Phys. Rev. D10, 2300 (1974); Indiana University preprint IUK-IT-60 (1980). There is, however, some question whether the oscillation is systematically observed when all experiments are combined (B. Chertok, private communication).
40. G. H. Thomas, this conference, and GRAZ Conf. 1978 viz:86.
41. S. J. Brodsky, C. E. Carlson and H. Lipkin, Phys. Rev. D20, 2278 (1978).
42. G. R. Farrar, S. Gottlieb, D. Sivers and G. H. Thomas, Phys. Rev. D20, 202 (1979).
43. D. G. Crabb, et al., Phys. Rev. Lett. 41, 1257 (1978); E. A. Crosbie, et al., Univ. of Michigan preprint UM HE 80-2 (1980). J. R. O'Fallon, et al., ibid. 39, 733 (1977).
44. A. Yokosawa, this conference. In review of the nucleon-nucleon scattering experiments see A. Yokosawa, ANL-HEP-CP-80-07-MC (1980) to be published in Phys. Reports (1980).
45. J. Szwed, Max Planck Inst. preprints MPI-PAE/PTH 33/80 and 34/80 (1980), and Phys. Lett. 93B, 485 (1980).

46. For additional discussions of the theory of pp polarization effects see C. Avilez, G. Coche and M. Moreno, preprint IFUNAM 80-17 (1980), G. Preparata and J. Soffer, Phys. Lett. 86B, 304 (1979), G. A. Walters, Phys. Rev. Lett. 45, 776 (1980), M. C. Daniel, Barcelona preprint UAB-FT-65 (1980), and the contributions of M. Doncel, G. F. Walters, G. Preparata, J. Soffer, J. Szwed, and M. Moreno to this conference.
47. S. J. Brodsky, T. A. DeGrand and R. Schwitters, Phys. Lett. 79B, 255 (1978). A. V. Smilga and H. I. Vysotsky, ITEP-12 (1979); L. Clavelli, Nucl. Phys. B167, 157 (1980).
48. S. J. Brodsky and S. D. Drell, Phys. Rev. D11, 2236 (1980).
49. S. J. Brodsky, SLAC-PUB-1699 (1975), published in the LAMPF Users Group Newsletter, Vol. 8-1, (1976).
50. G. L. Shaw, D. Silverman and R. Slansky, Phys. Lett. 94B, 57 (1980).
51. See also the discussion of M. Lipkin, Phys. Lett. 89B, 358 (1980), and I. Cohen and M. Lipkin, Phys. Lett. 93B, 56 (1980).
52. Calibrations of this type are being carried out in collaboration with T. Huang.
53. S. J. Brodsky and G. P. Lepage, ref. 35.
54. See, e.g., P. Hoyer, P. Osland, H. G. Sander and T. F. Walsh, Nucl. Phys. B161, 349 (1979), and references therein.
55. E. Witten, Nucl. Phys. B120, 189 (1977).