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Hadron Spin Dynamics*

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Abstract

Spin effects in exclusive and inclusive reactions provide an essential new dimension for testing QCD and unraveling hadron structure. Remarkable new experiments from SLAC, HERMES (DESY), and Jefferson Lab present many challenges to theory, including measurements at HERMES and SMC of the single spin asymmetries in $ep \to e'\pi X$ where the proton is polarized normal to the scattering plane. This type of single spin asymmetry may be due to the effects of rescattering of the outgoing quark on the spectators of the target proton, an effect usually neglected in conventional QCD analyses. Many aspects of spin, such as single-spin asymmetries and baryon magnetic moments are sensitive to the dynamics of hadrons at the amplitude level, rather than probability distributions. I will illustrate the novel features of spin dynamics for relativistic systems by examining the explicit form of the light-front wavefunctions for the two-particle Fock state of the electron in QED, thus connecting the Schwinger anomalous magnetic moment to the spin and orbital momentum carried by its Fock state constituents and providing a transparent basis for understanding the structure of relativistic composite systems and their matrix elements in hadronic physics. I also present a survey of outstanding spin puzzles in QCD, particularly A_{NN} in elastic pp scattering, the $J/\psi \to \rho\pi$ puzzle, and J/ψ polarization at the Tevatron.

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1 Introduction

One of the most important goals in quantum chromodynamics is to determine the fundamental structure of hadrons in terms of their quark and gluon degrees of freedom. Spin-dependent phenomena play a crucial role in this pursuit, dramatically increasing the sensitivity to important and subtle QCD effects [1]. Many aspects of spin phenomenology, such as the proton's magnetic moment and single-spin asymmetries, require an understanding of hadron structure at the amplitude level, rather than probability distributions. Spin phenomena include such topics as the magnetic moments of baryons, nuclear polarization effects in deuteron photo-disintegration, spin correlations in B decays, spin-spin correlations in high momentum transfer exclusive and inclusive reactions, and the spin correlations of top quarks produced in high energy colliders. The upcoming RHIC spin program will study high energy polarized proton-proton collisions, with the potential of directly measuring the spin carried by gluons in the proton [2, 3, 4]. The planned charm photoproduction studies at SLAC with polarized photons and protons will provide important checks on leading-twist QCD predictions [5]. New programs at HERMES [6] and Jefferson Laboratory [7] will study spin structure functions, spin correlations, and azimuthal polarization asymmetries in deeply virtual Compton scattering and in other exclusive channels.

A number of new experimental results from SLAC, HERMES (DESY), COMPASS at CERN, and Jefferson Lab were reported at this meeting [8], presenting many new challenges to theory. Most hadron spin experiments have been performed with protons; however there are also potentially important tests of QCD involving higher spin targets such as the deuteron [9, 10, 11] which test nuclear coherence and concepts such as hidden color [12]. Some of the most interesting confrontations with theory have arisen from the measurements by the HERMES [13] and SMC [14] collaborations of single-spin asymmetry in inclusive electroproduction $ep \to e'\pi X$ where the proton is polarized normal to the photon-to-pion scattering plane. In the case of HERMES, the proton is polarized along the incident electron beam direction, and thus the azimuthal correlation is suppressed by a kinematic factor $Q\sqrt{1-y}/\nu$, where $y = q \cdot p/p_e \cdot p$, which vanishes when the electron and photon are parallel [15, 16]. Nevertheless, the observed azimuthal spin asymmetry is large. Large single spin asymmetries have also been observed by the E704 experiment in $pp \to \pi X$ reactions at Fermilab [17]. In general, such single-spin correlations require a correlation of a particle spin with a production or scattering plane; they correspond to a process which yields the $T-odd$ triple product $i\vec{S}_A \cdot \vec{p}_B \times \vec{p}_C$ where the phase i is required by hermiticity and time-reversal invariance. In the case of deeply virtual Compton scattering, the phase can arise from the interference of the real Bethe-Heitler and largely imaginary Compton amplitude. In purely hadronic processes, such correlations can provide a window to the physics of final-state interactions, an effect usually neglected in conventional QCD analyses.

In this brief theory summary talk of SPIN2001, I will give an overview of recent progress in understanding proton spin and spin correlations in quantum chromodynamics. I will focus on results which pose challenges to QCD, such as single-spin asymmetries, the A_{NN} asymmetry in large-angle elastic pp scattering, the $J/\psi \to \rho\pi$ puzzle, and the problem of J/ψ polarization at the Tevatron. An important tool is the light-front wavefunction representation [18], a formalism in which the concepts of spin and orbital angular momentum in relativistic bound states become transparent. The wavefunctions derived from light-front quantization play a central role in QCD phenomenology, providing a frame-independent description of hadrons in terms of their quark and gluon degrees of freedom at the amplitude level [19].

2 Relativistic Spin

The spin decomposition of relativistic composite systems is particularly transparent in light-front quantized QCD in light-cone gauge $A^+ = 0$, since all non-physical degrees of freedom of the quarks and gluons are eliminated by constraints [18, 20]. Instead of ordinary time t, one quantizes the theory at fixed light-cone time [21] $\tau = t + z/c$. The generator $P^- = i\frac{d}{d\tau}$ generates light-cone time translations, and the eigen spectrum of the Lorentz scalar $H_{LC}^{QCD} = P^+P^- - \vec{P}_\perp^2$, where $P^\pm = P^0 \pm P^z$ gives the mass spectrum of the color-singlet hadron states, together with their respective frame-independent light-front wavefunctions. The total spin projection J_z, as well as the momentum generators P^+ and $\vec{P}_\perp$, are kinematical; *i.e.*, they are independent of the interactions. For example, the proton state satisfies: $H_{LC}^{QCD}\,|\,\psi_p\rangle = M_p^2\,|\,\psi_p\rangle$. The expansion of the proton eigensolution $|\,\psi_p\rangle$ on the color-singlet $B = 1$, $Q = 1$ eigenstates $\{|\,n\rangle\}$ of the free Hamiltonian $H_{LC}^{QCD}(g = 0)$ gives the light-cone Fock expansion:

$$\begin{aligned} \left|\psi_p(P^+, \vec{P}_\perp)\right\rangle &= \sum_n \prod_{i=1}^n \frac{\mathrm{d}x_i\,\mathrm{d}^2\vec{k}_{\perp i}}{\sqrt{x_i}\,16\pi^3}\; 16\pi^3\delta\left(1 - \sum_{i=1}^n x_i\right)\,\delta^{(2)}\left(\sum_{i=1}^n \vec{k}_{\perp i}\right) \\ &\quad \times \psi_n(x_i, \vec{k}_{\perp i}, \lambda_i)\left|n;\, x_iP^+, x_i\vec{P}_\perp + \vec{k}_{\perp i}, \lambda_i\right\rangle . \end{aligned} \tag{1}$$

The light-cone momentum fractions $x_i = k_i^+/P^+$ and $\vec{k}_{\perp i}$ represent the relative momentum coordinates of the QCD constituents. The physical transverse momenta are $\vec{p}_{\perp i} = x_i\vec{P}_\perp + \vec{k}_{\perp i}$. The λ_i label the light-front spin projections S^z of the quarks and gluons along the quantization direction z [22]. The corresponding spinors of the light-front formalism automatically incorporate the Melosh-Wigner transformation. The physical gluon polarization vectors $\epsilon^\mu(k,\ \lambda = \pm 1)$ are specified in light-cone gauge by the conditions $k\cdot\epsilon = 0,\ \eta\cdot\epsilon = \epsilon^+ = 0$. The light-front wavefunctions $\psi_{n/p}(x_i, \vec{k}_{\perp i}, \lambda_i)$ of a hadron in QCD are the projections of the hadronic eigenstate on the free color-singlet Fock state $|\,n\rangle$ at a given light-cone time τ. They thus represent the ensemble of quark and gluon states possible when a hadron is intercepted at the light-front.

The light-front wavefunctions are the natural extension of the Schrödinger theory for many body theory; however, they are Lorentz-invariant functions independent of the bound state's physical momentum P^+, and $P_\perp$. The parton degrees of freedom are thus all physical; there are no Faddeev-Popov ghost or other negative metric states [20].

The light-front formalism provides a remarkably transparent representation of angular momentum [23, 24, 25, 26, 27, 15, 18]. The projection J_z along the light-front quantization direction is kinematical and is conserved separately for each Fock component: each light-front Fock wavefunction satisfies the sum rule:

$$J^z = \sum_{i=1}^{n} S_i^z + \sum_{j=1}^{n-1} l_j^z \ .$$

The sum over S_i^z represents the contribution of the intrinsic spins of the n Fock state constituents. The sum over orbital angular momenta

$$l_j^z = -\mathrm{i} \left(k_j^1 \frac{\partial}{\partial k_j^2} - k_j^2 \frac{\partial}{\partial k_j^1} \right) \tag{2}$$

derives from the $n-1$ relative momenta and excludes the contribution to the orbital angular momentum due to the motion of the center of mass, which is not an intrinsic property of the hadron [25]. The numerator structure of the light-front wavefunctions in $k_\perp$ is thus largely determined by the angular momentum constraints. The spin properties of light-front wavefunctions provide a consistent basis for analyzing the role of orbital angular momentum and spin correlations and azimuthal spin asymmetries in both exclusive and inclusive reactions.

The light-front representation thus provides a rigorous underpinning for the familiar angular momentum sum rule for the proton [28, 29]:

$$J_z = \frac{1}{2} = \frac{1}{2}\Delta\Sigma + \Delta G + L_z.$$

The sum rule holds for each n particle Fock state of the proton. Here $\Delta\Sigma$ is the sum of the $S_z = \pm\frac{1}{2}$ of the quarks and antiquarks. The orbital term L_z refers to the sum over the $n-1$ internal orbital angular momenta of the quark and gluon constituents. In each case, the spin projections are referenced to the light-front quantization direction z. It is easy to check that the QCD interactions of the light-cone Hamiltonian conserve the total J_z. Thus relativistic orbital angular momentum has a physical, well-defined meaning in light-cone gauge [27].

3 Spin Dependent Local Matrix Elements

The primary measures of deeply virtual inelastic lepton scattering $\ell p \to \ell' X$ and deeply virtual Compton scattering $\gamma^* p \to \gamma p'$ are the gauge invariant matrix elements

$\langle p'|\mathcal{O}|p\rangle$ of products of quark and gluon fields at invariant separation $x^2 = \mathcal{O}(\frac{1}{Q^2})$, as generated by the operator product expansion. Comprehensive discussions of the non-forward (or skewed) parton distributions have been given by Radyushkin [30] and Ji [31]. Polarization measurements are essential for distinguishing the various operators. For example, Ji [31] has shown that there is a remarkable connection of the x-moments of the chiral-conserving and chiral-flip form factors $H(x,t,\zeta)$ and $E(x,t,\zeta)$ which appear in deeply virtual Compton scattering with the corresponding spin-conserving and spin-flip electromagnetic form factors $F_1(t)$ and $F_2(t)$ and gravitational form factors $A_{\rm q}(t)$ and $B_{\rm q}(t)$ for each quark and anti-quark constituent. Thus, in effect, one can use virtual Compton scattering to measure graviton couplings to the charged constituents of a hadron [32].

Given the light-front wavefunctions $\psi^{(\Lambda)}_{n/H}$, one can construct any spacelike electromagnetic, electroweak, or gravitational form factor or local operator product matrix element from the diagonal overlap of the LC wavefunctions [33, 34]. For example, the Pauli form factor and the anomalous magnetic moment $\kappa = \frac{e}{2M}F_2(0)$ are off-diagonal convolutions of light-front wavefunctions [34]

$$-(q^1 - {\rm i}q^2)\frac{F_2(q^2)}{2M} = \sum_a \int \frac{{\rm d}^2\vec{k}_\perp {\rm d}x}{16\pi^3} \sum_j e_j \; \psi^{\uparrow *}_a(x_i, \vec{k}'_{\perp i}, \lambda_i)\, \psi^{\downarrow}_a(x_i, \vec{k}_{\perp i}, \lambda_i) \ , \qquad (3)$$

where the summation is over all contributing Fock states a and struck constituent charges e_j. The arguments of the final-state light-front wavefunction are $\vec{k}'_{\perp i} = \vec{k}_{\perp i} + (1-x_i)\vec{q}_\perp$ for the struck constituent and $\vec{k}'_{\perp i} = \vec{k}_{\perp i} - x_i\vec{q}_\perp$ for each spectator. Notice that the anomalous magnetic moment is calculated from the spin-flip non-forward matrix element of the current where the incident and final wavefunctions have different orbital angular momenta $\Delta L_z = \pm 1$. In the ultra-relativistic limit, where the radius of the system is small compared to its Compton scale $1/M$, the anomalous magnetic moment must vanish as required by the DHG sum rule [35]. The light-cone formalism is consistent with this theorem.

The form factors of the energy-momentum tensor for a spin-$\frac{1}{2}$ system are defined by

$$\begin{aligned}\langle P'|T^{\mu\nu}(0)|P\rangle &= \overline{u}(P')\Big[A(q^2)\gamma^{(\mu}\overline{P}^{\nu)} + B(q^2)\frac{i}{2M}\overline{P}^{(\mu}\sigma^{\nu)\alpha}q_\alpha \\ &\qquad + C(q^2)\frac{1}{M}(q^\mu q^\nu - g^{\mu\nu}q^2)\Big]u(P) \ , \end{aligned} \qquad (4)$$

where $q^\mu = (P'-P)^\mu$, $\overline{P}^\mu = \frac{1}{2}(P'+P)^\mu$, $a^{(\mu}b^{\nu)} = \frac{1}{2}(a^\mu b^\nu + a^\nu b^\mu)$, and $u(P)$ is the spinor of the system. One can express the matrix elements of the energy momentum tensor $T^{\mu\nu}$ as overlap integrals of light-front wavefunctions [25]. As in the light-cone decomposition of the Dirac and Pauli form factors for the vector current, one can obtain the light-cone representation of $A(q^2)$ and $B(q^2)$ from local matrix elements of $T^{++}(0)$: An important consistency check of any relativistic formalism is to verify the vanishing of the anomalous gravito-magnetic moment $B(0)$, the spin-flip matrix element of the graviton coupling and analog of the anomalous magnetic moment $F_2(0)$. For example, at one-loop order in QED, $B_f(0) = \frac{\alpha}{3\pi}$ for the electron when the graviton interacts with the fermion line, and $B_\gamma(0) = -\frac{\alpha}{3\pi}$ when the graviton interacts with the exchanged photon. The vanishing of $B(0)$ can be shown to be exact for bound or elementary systems in the light-front formalism [25], in agreement with the equivalence principle [36, 37, 38].

In the case of the timelike matrix elements which appear in deeply virtual Compton scattering and heavy hadron decays, one also has contributions in which the initial and final Fock state number differ by $\Delta n = 2$. The skewed parton distributions which control deeply virtual Compton scattering can also be defined from diagonal and off-diagonal overlaps of light-front wavefunctions [39, 40]. The form factors and matrix elements of local currents $B|J^\mu(0)|A$ which appear in semileptonic decay amplitudes of heavy hadrons have an exact representation as overlap momentum-space overlap integrals of light-front wavefunctions [41]. The light-front Fock representation is particularly important for semi-leptonic exclusive matrix elements such as $B \to D\ell\overline{\nu}$. The Lorentz-invariant description requires both the overlap of $n' = n$ parton-number conserving wavefunctions as well as the overlap of wavefunctions with parton numbers $n' = n - 2$ which arises from the annihilation of a quark-antiquark pair in the initial wavefunction [41].

The total angular momentum projection of a state is given by [31]

$$\left\langle J^i \right\rangle = \frac{1}{2}\epsilon^{ijk} \int d^3x \left\langle T^{0k}x^j - T^{0j}x^k \right\rangle = A(0) \left\langle L^i \right\rangle + [A(0) + B(0)]\, \overline{u}(P)\frac{1}{2}\sigma^i u(P)\ . \quad (5)$$

The $\langle L^i \rangle$ term is the orbital angular momentum of the center of mass motion with respect to an arbitrary origin and can be dropped. The coefficient of the $\langle L^i \rangle$ term must be 1; $A(0) = 1$ also follows when we evaluate the four-momentum expectation value $\langle P^\mu \rangle$. Thus the total intrinsic angular momentum J^z of a nucleon can be identified with the values of the form factors $A(q^2)$ and $B(q^2)$ at $q^2 = 0$:

$$\langle J^z \rangle = \left\langle \frac{1}{2}\sigma^z \right\rangle [A(0) + B(0)]\ . \quad (6)$$

One can define individual quark and gluon contributions to the total angular momentum from the matrix elements of the energy momentum tensor [31]. However, this definition is only formal; $A_{q,g}(0)$ can be interpreted as the light-cone momentum fraction carried by the quarks or gluons $\langle x_{q,g} \rangle$. The contributions from $B_{q,g}(0)$ to J_z cancel in the sum. In fact, as noted above, the contributions to $B(0)$ vanish when summed over the constituents of each individual Fock state [25].

4 The Spin Distributions of the Proton

Given the light-front wavefunctions, one can compute the x-dependence of the quark spin and transversity distributions measured in polarized inclusive reactions. For example, the spin-polarized quark distributions at resolution (factorization scale) Λ correspond to [22, 25]

$$q_{\lambda_q/\lambda_p}(x,\Lambda) = \sum_{n,q_a} \int \prod_{j=1}^{n} dx_j d^2k_{\perp j} \sum_{\lambda_i} |\psi^{(\Lambda)}_{n/H}(x_i, \vec{k}_{\perp i}, \lambda_i)|^2 \tag{7}$$
$$\times \delta\left(1 - \sum_i^n x_i\right) \delta^{(2)}\left(\sum_i^n \vec{k}_{\perp i}\right) \delta(x - x_q)\delta_{\lambda_a,\lambda_q}\Theta(\Lambda^2 - \mathcal{M}_n^2) \ ,$$

where the sum is over all quarks q_a which match the quantum numbers, light-cone momentum fraction x, and helicity of the struck quark. We thus have $g_1(x,\Lambda) = \Delta q(x,\Lambda) = q(\lambda_q = +\frac{1}{2}, \lambda_p = +\frac{1}{2}) - q(\lambda_q = -\frac{1}{2}, \lambda_p = +\frac{1}{2})$. If the simple handbag diagram of the parton model dominates the deep inelastic cross section, then the quark distributions defined from the light-front wavefunction probabilities will give an accurate representation of the $g_1(x,Q)$ structure function in the Bjorken limit, where to leading order in α_s we can identify the factorization scale Q and resolution scale Λ. data. These distributions can also be defined as forward matrix elements of gauge invariant local quark currents [42].

The most recent determinations of the integrated quark and gluon distributions, as reported to this meeting by Ramsey, are [43]

$$\Delta u = 0.86, \Delta d = -0.40, \Delta s = -0.06, \Sigma = \Delta q = 0.40, \Delta g = 0.44, L_z = -0.1,$$

and $\Delta u_{sea} = -0.068$. The error in each case is approximately ± 0.04. The effective factorization scale is $Q^2 \simeq 1$ GeV^2. The analysis incorporates recent E154 data, the Bjorken sum rule $\Delta u - \Delta d = g_A$, and the perturbative QCD "counting rule" constraints $\Delta g(x) = xg(x)$ and $\Delta q(x)/q(x) \to 1$ at $x \to 1$ which follow from the dominance of the spin-aligned quark and gluon distributions. This will be discussed in more detail in Section 5. The results show the presence of significant orbital angular momentum L_z and gluon spin Δg, which is natural for light-cone wavefunctions. This will be illustrated in a QED example in Section 9. Applications to high p_T inclusive processes in polarized proton collisions are given by Gordon and Ramsey [44].

The gluon polarization also contributes to the quark and antiquark spin distributions: when the lepton scatters on a quark or antiquark arising from gluon splitting, the quark and antiquark contributions do not cancel. The integration over quark transverse momentum produces an anomalous contribution [45, 46, 47] $-3\frac{\alpha_s}{2\pi}\Delta G(x,Q)$ to the Ellis-Jaffe sum rule [48] from quantum fluctuations. When the quark mass is high compared to the gluon virtuality, the anomaly vanishes. This decoupling follows from an application of the Drell-Hearn Gerasimov sum rule [49].

Final-state interactions in gauge theory can affect deep inelastic scattering reactions in a profound way [50]. The rescattering of the outgoing quark leads to a leading-twist contribution to the deep inelastic cross section from diffractive channels $\gamma^* p \to q\bar{q}p'$, and the interference effects induced by these diffractive channels cause nuclear shadowing. Nuclear shadowing, such as the strong spin-dependence predicted for the deuteron spin structure function [51, 52, 53] is thus not given by the nuclear light-cone wavefunctions. The final-state interactions affect the x distributions at small x, and thus they can modify the extrapolations of the spin-dependent structure functions into the low x domain. Similar complications may also affect the distributions obtained from deeply virtual Compton scattering. The moments of the distributions which follow from the OPE are evidently not affected.

There are in fact, possible phenomenological problems with the perturbative QCD phenomenology. Recent measurements of polarized electron and proton deep inelastic scattering measurements by the $E155$ collaboration at SLAC [5] give $\int_0^1 dx g_2^p(x, Q) = -0.034 \pm 0.008$, in violation of the Burkhardt-Cottingham sum rule [54]. The second moment of g_2 also appears to be inconsistent with model predictions based on the Wandzura-Wilczek sum rule for spin-$\frac{1}{2}$ partons [55]. Future plans to extend these measurements at Jefferson laboratory were presented to this meeting by J.-P. Chen and Z.-E. Meziani. V. Guzey also has discussed complications due to the light nuclear targets used in these experiments, including possible isobar complications [56]. Measurement of diffractive dissociation $\ell A \to \ell' p(A-1)$ can help to clarify the complications of nuclear shadowing and final state interactions [50].

When leading-twist factorization holds, one can predict spin correlations for inclusive hadron production and jet measures using the standard convolution of structure functions and fragmentation functions, evolved by perturbative QCD. An excellent review of these phenomena is given by Boer and Mulders [57]. Generalizations to next-to-leading twist have been developed by Qiu and Sterman [58]. The analysis of azimuthal single-spin correlations of hadron spin with the various production planes has been pioneered by Collins [59, 16].

5 Dynamical Constraints on Polarized Structure functions and Fragmentation Distributions

One of the great challenges of nonperturbative QCD is to calculate the magnitude and shapes of the quark and gluon spin distributions from first principles in QCD. In this meeting Melnitchouk reported on lattice determinations of the distribution moments and the complications of the chiral extrapolation [60]. Another rigorous approach, the discretized light-front quantization (DLCQ) method, computes the light-front wavefunctions directly by diagonalizing the light-cone Hamiltonian on a discrete Fock basis [61]. Important progress in obtaining explicit light-front wavefunctions for model 3+1 theories has been reported by Hiller, McCartor, and myself [62]. There are also important guides obtained by QCD sum rules [63] and the meson cloud model [64, 65, 66].

Even without explicit nonperturbative solutions, one can use theoretical constraints to pin down the large and small x behavior of the spin distributions. For example, one can use the Regge behavior of high energy amplitudes to discern the power behavior of spin-dependent fragmentation functions at small x by analyzing the spin or effective spin of particles coupling in the t channel in virtual Compton scattering [67, 68, 69]. On the other hand, The behavior of structure functions where one quark has the entire momentum requires the knowledge of LC wavefunctions with $x \to 1$ for the struck quark and $x \to 0$ for the spectators. This is a highly off-shell configuration, and the fall-off the light-front wavefunctions at large $k_\perp$ and $x \to 1$ is dictated by QCD perturbation theory since the state is far-off the light-cone energy shell. Thus one can rigorously derive quark-counting and helicity-retention rules for the power-law behavior of the polarized and unpolarized quark and gluon distributions in the $x \to 1$ endpoint domain [22, 70]. Notice that $x \to 1$ corresponds to $k^z \to -\infty$ for any constituent with nonzero mass or transverse momentum.

Modulo DGLAP evolution, the counting rule for the momentum distribution for finding parton a in hadron a at large $x \sim 1$ is $G_{a/A}(x) \propto (1-x)^{2n_{\rm spect}-1+2|\Delta S_z|}$ where $n_{\rm spect}$ is the minimum number of partons left behind when parton a is removed from A, and ΔS_z is the difference of the a and A spin projections. This predicts $(1-x)^3$ behavior for valence quarks with spin aligned with the parent proton's spin projection, and $(1-x)^5$ behavior for anti-aligned quarks. Similarly, the counting rules predict that gluons with their spins aligned with the proton have a $(1-x)^4$ distribution versus $(1-x)^6$ distribution when the gluon and proton spins are anti-aligned. These constraints [70] are usually incorporated into phenomenological analyses such as the Ramsey [43] and MRST [71] approaches. DGLAP evolution is quenched in the large x limit in the fixed W^2 domain One also expects strong spin correlations in fragmentation functions. Because of the Gribov-Lipatov relation, the same counting rule behavior holds for fragmentation $D_{H\to a}(z,Q)$ distributions at $z \to 1$. This is particularly interesting in the case of the fragmentation of a quark into a Λ since its polarization is self-analyzing.

In the large x or z domain, higher twist terms involving several partons of the target can dominate inclusive rates since fewer spectators are stopped [72]. The clearest example is the Drell-Yan process $\pi A \to \ell^+\ell^- X$ at $x_1 \to 1$ where the process $\pi q \to \gamma^* q$ subprocess can dominate the leading twist $\overline{q}(x_1)q(x_2) \to \gamma^* \to \ell^+\ell^-$ amplitude. In the higher twist process the virtual photon is produced with longitudinal polarization. The muon differential cross section thus has the form

$$\frac{d\sigma}{d\cos\theta dx_1} = A(1-x_1)^2(1+\cos^2\theta) + \frac{B}{Q^2}\sin^2\theta, \tag{8}$$

in good agreement with the data. Thus at large momentum fraction x_1, the higher-twist term is amplified by a factor $\mu^2/[Q^2(1-x_1)^2]$. In addition, the coplanarity distribution of the muon pair plane with the $\pi \to \gamma^*$ production plane is consistent with these expectations [73].

6 Proton Spin Carried by Non-Valence Quarks

Since particle number is not conserved in a relativistic theory, the wavefunctions of hadrons must contain non-valence states, corresponding to additional sea quarks and gluons in flight at a fixed light-cone time τ; part of this fluctuation corresponds to the resolved substructure of the valence quarks and is associated with DGLAP evolution. However, the sea quarks and gluons which are multiply connected to the valence quarks are intrinsic to the hadron's structure and are not generated by perturbative evolution [74]. One can show rigorously from the OPE, that the momentum and spin carried by the intrinsic sea quarks in non-Abelian QCD decreases as $1/m_Q^2$ for heavy quarks like charm [75]. If one associates the intrinsic strange quarks in the proton with its $K^+\Lambda$ fluctuation, then the strange quark spin will be anti-aligned with the nucleon and the anti-strange quark will be unpolarized since it is associated with the pseudoscalar meson [65]. The momentum fraction carried by intrinsic charm quarks is of the order of 1%[76], and thus one expects the fraction of proton spin carried by carried by charm quarks to also be of this order [77, 78].

7 Heavy Quarkonium Polarization in Inclusive Reactions

There are still many unresolved problems involving of J/ψ production in high energy hadron collisions. The expectation from the color-singlet model that the J/ψ will be produced with transverse polarization at large transverse momentum in pp collisions has not been confirmed [79]. In the case of $\pi A \to J/\psi X$, the J/ψ is generally produced unpolarized. However at $x_F \sim 0.95$ the J/ψ is observed to have significant longitudinal polarization [80]. This anomalous behavior can be understood if the dominant subprocess is the diffractive excitation of the $q\bar{q}c\bar{c}$ Fock state of the pion at large x_F, since the $c\bar{c}$ pair will have the same in as the parent pion when they have a large momentum fraction $x_F = x_c + x_{\bar{c}} \to 1$ [81].

8 Transversity

The transversity distribution $\delta q(x, Q)$ gives the correlation between transversely polarized quarks and the transverse polarization of the parent hadron [82, 83, 84]. Transversity is a leading-twist distribution which can be measured in inclusive processes with two hadrons such as the Drell-Yan process. An excellent review of the theory and phenomenology of transversity was presented to this meeting by Jaffe [15].

Transversity can be determined from the light-front wavefunctions of the target. A transversely polarized spin-$\frac{1}{2}$particle polarized in the $\hat{y}$ direction can be represented as the state $\mid N\rangle = \frac{|+>-|->}{\sqrt{2}}$ where $\mid\pm\rangle$ represents the $S_z = \pm\frac{1}{2}$ state. Thus the transversity distribution is a density matrix of light-front wavefunctions:

$$\delta q(x,\Lambda) = \sum_{n,q_a}\int\prod_{j=1}^{n} dx_j d^2k_{\perp j}\sum_{\lambda_i}\Big|\psi^{(\Lambda)*}_{n/H}(x_i,\vec{k}_{\perp i},\lambda_i=\lambda_p=+\frac{1}{2}) \tag{9}$$

$$\times \qquad \psi^{(\Lambda)}_{n/H}(x_i,\vec{k}_{\perp i},\lambda_i=\lambda_p=-\frac{1}{2})\Big| \tag{10}$$

$$\times\ \delta\left(1-\sum_i^n x_i\right)\delta^{(2)}\left(\sum_i^n \vec{k}_{\perp i}\right)\delta(x-x_q)\delta_{\lambda_a,\lambda_q}\Theta(\Lambda^2-\mathcal{M}_n^2)\ ,$$

The Soffer inequality [85] $\delta q(x,\Lambda)\leq\frac{1}{2}[\Delta q(x,\Lambda)+q(x,\Lambda)]$ follows simply from the light-cone representation. The integrated transversity is the matrix element of the quark tensor charge $\overline{\psi}(0)\sigma^{+i}\gamma^5\psi(0)$ operator with quantum numbers $J^{CP}=1^{+-}$. Gronberg and Goldstein [86] have shown how one can use the measured couplings of axial vector meson to the proton in order to estimate the strength of the transversity of the u and d quarks in the proton at a factorization scale $x^2=1/\mu^2$ of the order of the size of hadrons. As a simple example of transversity, the spin decomposition of the electron in QED due to its lowest order quantum fluctuations will be analysed in the next section.

9 The Light-Front Wavefunctions and Spin Structure of Leptons in QED

The light-front wavefunctions of a lepton in QED provide an ideal system to check explicitly the intricacies of relativistic spin and angular momentum in quantum field theory [25]. Although they are derived in perturbation theory, the light-front wavefunctions of leptons and photons can be used as templates for the wavefunctions of non-perturbative composite systems resembling hadrons in QCD. For generality, one can assign a mass M to the external lepton, a different mass m to the internal fermion lines, and a mass λ to the internal photon line, thus simulating the quark-spin-1 diquark structure of a baryon. The fast-falling wavefunction of a composite state can be simulated by differentiating the wavefunction with respect to a mass parameter or employing a Pauli-Villars spectrum.

The two-particle Fock state for an electron with $J^z=+\frac{1}{2}$ has

$$\left\{\begin{array}{l}
\psi^{\uparrow}_{+\frac{1}{2}\,+1}(x,\vec{k}_{\perp})=-\sqrt{2}\frac{(-k^1+ik^2)}{x(1-x)}\,\varphi\ ,\\
\psi^{\uparrow}_{+\frac{1}{2}\,-1}(x,\vec{k}_{\perp})=-\sqrt{2}\frac{(+k^1+ik^2)}{1-x}\,\varphi\ ,\\
\psi^{\uparrow}_{-\frac{1}{2}\,+1}(x,\vec{k}_{\perp})=-\sqrt{2}(M-\frac{m}{x})\,\varphi\ ,\\
\psi^{\uparrow}_{-\frac{1}{2}\,-1}(x,\vec{k}_{\perp})=0\ ,
\end{array}\right. \tag{11}$$

where

$$\varphi=\varphi(x,\vec{k}_{\perp})=\frac{e/\sqrt{1-x}}{M^2-(\vec{k}_{\perp}^2+m^2)/x-(\vec{k}_{\perp}^2+\lambda^2)/(1-x)}\ , \tag{12}$$

with similar expressions for $\psi^{\downarrow}$. The coefficients of φ are the matrix elements of $\frac{\overline{u}(k^+,k^-,\vec{k}_\perp)}{\sqrt{k^+}}\gamma\cdot\epsilon^*\frac{u(P^+,P^-,\vec{P}_\perp)}{\sqrt{P^+}}$ which are the numerators of the wavefunctions corresponding to each constituent spin s^z configuration. The two boson polarization vectors in light-cone gauge are $\epsilon^\mu = (\epsilon^+ = 0\ ,\epsilon^- = \frac{\vec{\epsilon}_\perp\cdot\vec{k}_\perp}{2k^+},\vec{\epsilon}_\perp)$ where $\vec{\epsilon} = \vec{\epsilon_\perp}_{\uparrow,\downarrow} = \mp(1/\sqrt{2})(\widehat{x}\pm \mathrm{i}\widehat{y})$. The polarizations also satisfy the Lorentz condition $k\cdot\epsilon = 0$.

The transverse momentum dependence of the numerator of each light-front wavefunction specifies the orbital angular momentum guarantees J_z conservation for each Fock state. The fermion spin is given by the matrix element of the light-cone spin operator $\gamma^+\gamma^5$ [87]; the relative orbital angular momentum operator is $-\mathrm{i}(k^1\frac{\partial}{\partial k^2} - k^2\frac{\partial}{\partial k^1})$ [88, 89, 90, 91]. Each configuration satisfies the spin sum rule: $J^z = s^z_{\rm f} + s^z_{\rm b} + l^z$.

In the non-relativistic limit, the transverse motion of the constituents can be neglected and we have only the $\left|+\frac{1}{2}\right\rangle \to \left|-\frac{1}{2}\ +1\right\rangle$ configuration which is the non-relativistic quantum state for the spin-half system composed of a fermion and a spin-1 boson constituents. The fermion constituent has spin projection in the opposite direction to the spin J^z of the whole system. However, for ultra-relativistic binding in which the transverse motions of the constituents are large compared to the fermion masses (and at large evolution scales), the $\left|+\frac{1}{2}\right\rangle \to \left|+\frac{1}{2}\ +1\right\rangle$ and $\left|+\frac{1}{2}\right\rangle \to \left|+\frac{1}{2}\ -1\right\rangle$ configurations dominate over the $\left|+\frac{1}{2}\right\rangle \to \left|-\frac{1}{2}\ +1\right\rangle$ configuration. In this case the fermion constituent has spin projection parallel to J^z. In the case of Yukawa theory corresponding to a spin-0 diquark, the non-relativistic fermion's spin projection is aligned with the total J^z, and it is anti-aligned in the ultra-relativistic limit.

Given the light-front wavefunctions, one can explicitly calculate the helicity-flip electromagnetic and gravitational form factors for the fluctuations of the electron at one-loop, and verify the Schwinger anomalous moment and the cancellation of the sum of graviton couplings $B(q^2)$ to the constituents at $q^2 = 0$ [25]. The contribution to $B(q^2)$ where the photon couples to the photon constituent has threshold dependence $\sim\sqrt{-q^2/m^2}$ reflecting the massless two-body cut in the t channel.

The spin distributions in the QED model are also easily computed:

$$\begin{aligned}&\Delta q(x,\Lambda^2)_{\rm spin-1\ diquark} \\ &= \int\frac{{\rm d}^2\vec{k}_\perp{\rm d}x}{16\pi^3}\theta(\Lambda^2-\mathcal{M}^2)\ 2\left[\ \frac{\vec{k}^2_\perp}{x^2(1-x)^2}\ +\ \frac{\vec{k}^2_\perp}{(1-x)^2}\ -\ (M-\frac{m}{x})^2\ \right]|\varphi|^2\ .\end{aligned}\tag{13}$$

In the case of the Yukawa model, where the boson plays the role of a spin-0 diquark, one finds

$$\Delta q(x,\Lambda^2)_{\rm spin-0\ diquark} = \int\frac{{\rm d}^2\vec{k}_\perp{\rm d}x}{16\pi^3}\theta(\Lambda^2-\mathcal{M}^2)\left[\ (M+\frac{m}{x})^2\ -\ \frac{\vec{k}^2_\perp}{x^2}\ \right]|\varphi|^2\ ,\tag{14}$$

where we have regulated the integral by assuming a cutoff in the invariant mass: $\mathcal{M}^2 = \sum_i\frac{\vec{k}^2_{\perp i}+m_i^2}{x_i} < \Lambda^2$. In the spin-0 diquark model $\Delta q = 1$ in the nonrelativistic limit, and decreases toward $\Delta q = -1$ as the intrinsic transverse momentum increases.

The behavior is just opposite in the case of the spin-1 diquark. The distinct features of spin structure in the non-relativistic and ultra-relativistic limits reveals the importance of relativistic effects and supports the viewpoint [87, 92, 93] that the proton "spin puzzle" can be understood as due to the relativistic motion of quarks inside the nucleon. In particular, the spin projection of the relativistic constituent quark tends to be anti-aligned with the proton spin in a quark-diquark bound state if the diquark has spin 0. The state with orbital angular momentum $l^z = \pm 1$ in fact dominates over the states with $l^z = 0$. Thus the empirical fact that Δq is does not saturate the spin of the proton has a natural description in the light-cone Fock representation of hadrons.

In the case of transversity, we require the overlap of the $\psi^{*\uparrow}_{+\frac{1}{2}}$ and $\psi^{\downarrow}_{-\frac{1}{2}}$ states. Thus

$$\delta q(x,\Lambda^2)_{\rm spin-1\ diquark} = \int \frac{{\rm d}^2\vec{k}_\perp {\rm d}x}{16\pi^3}\theta(\Lambda^2-\mathcal{M}^2)\ 2\left[\ \frac{\vec{k}^2_\perp}{x(1-x)^2}\ \right]|\varphi|^2\ .$$

In the case of spin-0 diquarks, one finds

$$\delta q(x,\Lambda^2)_{\rm spin-0\ diquark} = \int \frac{{\rm d}^2\vec{k}_\perp {\rm d}x}{16\pi^3}\theta(\Lambda^2-\mathcal{M}^2)\left[\ (M+\frac{m}{x})^2\ \right]|\varphi|^2\ , \qquad (15)$$

In each case the result obeys the Soffer inequality [85].

10 Spin and Exclusive Processes

One of the important new areas of spin phenomenology is the exclusive decays of the B mesons [94, 95, 96]. The factorization theorems for hard exclusive processes [97] in which amplitudes factorize as products of distribution amplitudes and hard scattering quark and gluon scattering amplitudes have been generalized to exclusive B decays. The spin observables provide sensitive tests of factorization and the prediction of diminished final-state interactions (color transparency) [98]. Predictions for $B \to J/\psi K^*$ and Λ_B decay were presented at this meeting by K.-C. Yang [99]. The physics of the B distribution amplitude using constraints from heavy quark symmetry was presented by Kawamura [100]. Intrinsic charm in the B wavefunction can play an enhanced role in its weak decays because of the CKM hierarchy [101, 102].

The features of exclusive processes to leading power in the transferred momenta are well known [97]:

(1) The leading power fall-off is given by dimensional counting rules for the hard-scattering amplitude: $T_H \sim 1/Q^{n-1}$, where n is the total number of fields (quarks, leptons, or gauge fields) participating in the hard scattering [103, 104]. Thus the reaction is dominated by subprocesses and Fock states involving the minimum number of interacting fields. The hadronic amplitude follows this fall-off, modulo logarithmic corrections from the running of the QCD coupling and the evolution of the hadron

distribution amplitudes. In some cases, such as large angle $pp \to pp$ scattering, pinch contributions from multiple hard-scattering processes must also be included [105]. The general success of dimensional counting rules implies that the effective coupling $\alpha_V(Q^*)$ controlling the gluon exchange propagators in T_H are frozen in the infrared, *i.e.*, have an infrared fixed point, since the effective momentum transfers Q^* exchanged by the gluons are often a small fraction of the overall momentum transfer [106]. The pinch contributions are then suppressed by a factor decreasing faster than a fixed power [103].

(2) The leading power dependence is given by hard-scattering amplitudes T_H which conserve quark helicity [107, 108]. Since the convolution of T_H with the light-front wavefunctions projects out states with $L_z = 0$, the leading hadron amplitudes conserve the sum of light-cone spin projections s^z_i. Thus the sum of initial and final hadron helicities are conserved. Hadron helicity conservation thus follows from the underlying chiral structure of QCD.

One of the important predictions of hadron helicity conservation in perturbative QCD is the relative suppression of the Pauli versus the Dirac form factors: $F_2(Q^2)/F_1(Q^2) \to \mu^2/Q^2$ (modulo logarithmic factors) for timelike and spacelike form factors at high Q^2. However recent measurements [109] at Jefferson laboratory using the ratio of transverse and longitudinal recoil proton polarization correlated with the electron longitudinal spin gives the ratio $G^p_E(Q^2)/G^p_M(Q^2)$ directly; the results are consistent with a behavior $F_2(Q^2)/F_1(Q^2) \sim C/Q$ for $2 < Q^2 < 5$ GeV^2. The experimental results suggest that the Pauli form factor may be receiving significant logarithmic enhancement factors, possibly from the Sudakov-controlled fixed $k_\perp$, $x \to 1$ integration regime. It is also important to check the angular distribution in $p\overline{p} \to \ell^+\ell^-$ and e^+e^- annihilation to baryon pairs. Hadron helicity conservation predicts the dominance of $1 + \cos^2\theta$ distributions at large s.

Hadron-helicity conservation also predicts the suppression of vector meson states produced with $J_z = \pm 1$ in $e^+e^=$ annihilation to vector-pseudoscalar final states [107]. However, $J/\psi \to \rho\pi$ appears to occur copiously whereas $\psi' \to \rho\pi$ has never been conserved. The PQCD analysis assumes that a heavy quarkonium state such as the J/ψ always decays to light hadrons via the annihilation of its heavy quark constituents to gluons. However, as Karliner and I [110] have shown, the transition $J/\psi \to \rho\pi$ can also occur by the rearrangement of the $c\overline{c}$ from the J/ψ into the $|\, q\overline{q}c\overline{c}\rangle$ intrinsic charm Fock state of the ρ or π. On the other hand, the overlap rearrangement integral in the decay $\psi' \to \rho\pi$ will be suppressed since the intrinsic charm Fock state radial wavefunction of the light hadrons will evidently not have nodes in its radial wavefunction. This observation can provide a natural explanation of the long-standing puzzle why the J/ψ decays prominently to two-body pseudoscalar-vector final states, whereas the ψ' does not.

One of the most striking anomalies in elastic proton-proton scattering is the large spin correlation A_{NN} observed at large angles [111]. At $\sqrt{s} \simeq 5$ GeV, the rate for scattering with incident proton spins parallel and normal to the scattering plane is four times larger than that for scattering with anti-parallel polarization. This strong polarization correlation can be attributed to the onset of charm production in the intermediate state at this energy [112, 113]. A resonant intermediate state $|uuduudc\bar{c}\rangle$ has odd intrinsic parity and can thus couple to the $J = L = S = 1$ initial state, thus strongly enhancing scattering when the incident projectile and target protons have their spins parallel and normal to the scattering plane. The charm threshold can also explain the anomalous change in color transparency observed at the same energy in quasi-elastic pp scattering. A crucial test is the observation of open charm production near threshold with a cross section of order of 1μb. Analogous strong spin effects should also appear at the strangeness threshold and in exclusive photon-proton reactions such as large angle Compton scattering and pion photoproduction near the strangeness and charm thresholds. An alternate hypothesis, based on the interference of Landshoff pinch contributions has been proposed by Jain and Ralston [114].

An important new area of study is exclusive channels in electroproduction, especially for longitudinally polarized virtual photons, where factorization theorems can also be proved [115].

The application of perturbative QCD to exclusive nuclear processes such as the deuteron form factors, [116, 12, 10] photodisintegration [117], and meson photoproduction [11] on a polarized deuteron target is particularly interesting as first principle tests of the applicability of QCD to nuclear physics. For example, the dominance of helicity-conserving amplitudes in gauge theory can be shown to imply universal ratios for the charge, magnetic, and quadrupole form factors of spin-one bound state [10]. These results provide all-angle predictions for the leading power behavior of the tensor polarization $T_{20}(Q^2, \theta_{cm})$ and the invariant ratio $B(Q^2)/A(Q^2)$, although significant higher-twist corrections are expected at presently accessible momentum transfers. One can also show that the magnetic and quadrupole moments of any composite spin-one system take on the canonical values $\mu = e/M$ and $Q = -e/M^2$ in the strong binding limit of the zero bound-state radius or infinite excitation energy independent of the internal dynamics. A comprehensive review of the status of spin tests in exclusive nuclear processes and the successes and failures of the perturbative QCD approach is given by Gilman and Gross [118]. It should be emphasized that the normalization of the perturbation theory deuteron matrix elements are strongly affected by the existence of six-quark hidden color Fock states [12].

11 Single-Spin Asymmetries

Single-spin asymmetries in hadronic reactions have been among the most difficult phenomena to understand from basic principles in QCD. The problem has become more acute because of the observation by the HERMES [13] and SMC [14] collaborations of a strong correlation between the target proton spin $\vec{S}_p$ and the plane of a produced pion in semi-inclusive deep inelastic lepton scattering $\ell p^\uparrow \to \ell' \pi X$ at photon virtuality as large as $Q^2 = 6$ GeV^2. Large azimuthal single-spin asymmetries have also been seen in hadronic reactions such as $pp^\uparrow \to \pi X$ [17], where the target antiproton is polarized normal to the pion production plane, and in $pp \to \Lambda^\uparrow X$ [119], where the hyperon is polarized normal to the Λ production plane.

In the target rest frame, single-spin correlations correspond to the $T-odd$ triple product $i\vec{S}_p \cdot \vec{p}_\pi \times \vec{q}$, where the phase i is required by time-reversal invariance. In order to produce such an azimuthal correlation involving a transversely polarized proton, there are two necessary conditions: (1) There must be two proton spin amplitudes $M[\gamma^* p(J^z_p) \to F]$ with $J^z_p = \pm\frac{1}{2}$ which couple to the same final-state $|\, F\rangle$; and (2) The two amplitudes must have different, complex phases. The analysis of single-spin asymmetries thus requires an understanding of QCD at the amplitude level, well beyond the standard treatment of hard inclusive reactions based on the factorization of structure functions and fragmentation functions. Since we need the interference of two amplitudes which have different proton spin $J^z_p = \pm\frac{1}{2}$ but couple to the same final-state, the orbital angular momentum of the two proton wavefunctions must differ by $\Delta L^z = 1$. If a target is stable, its light-front wavefunction must be real. Thus the only source of a nonzero complex phase in leptoproduction in the light-front frame are final-state interactions. The rescattering corrections from final-state exchange of gauge particles produce Coulomb-like complex phases which, however, depend on the proton spin. Thus $M[\gamma^* p(J^z_p = \pm\frac{1}{2}) \to F] = |M[\gamma^* p(J^z_p = \pm\frac{1}{2}) \to F]|\, e^{i\chi_\pm}$. Each of the phases is infrared divergent; however the difference $\Delta\chi = \chi_+ - \chi_-$ is infrared finite and nonzero. The resulting single-spin asymmetry is then proportional to $\sin\Delta\chi$.

Recently, Dae Sung Hwang, Ivan Schmidt and I [120] have found that final-state interactions from gluon exchange between the outgoing quark and the target spectator system leads to single-spin asymmetries in deep inelastic lepton-proton scattering at leading twist in perturbative QCD; *i.e.*, the rescattering corrections are not power-law suppressed at large photon virtuality q^2 at fixed x_{bj}. The existence of such single-spin asymmetries requires a phase difference between two amplitudes coupling the proton target with $J^z_p = \pm\frac{1}{2}$ to the same final-state, the same amplitudes which are necessary to produce a nonzero proton anomalous magnetic moment. We find that the exchange of gauge particles between the outgoing quark and the proton spectators produces a Coulomb-like complex phase which depends on the angular momentum L^z of the proton's constituents and thus is distinct for different proton spin amplitudes. The single-spin asymmetry which arises from such final-state interactions does not factorize into a product of structure function and fragmentation function, and it is not related to the transversity distribution $\delta q(x, Q)$ which correlates transversely polarized quarks with the spin of the transversely polarized target nucleon.

We have calculated the single-spin asymmetry in semi-inclusive electroproduction $\gamma^* p \to HX$ induced by final-state interactions in a model of a spin-$\frac{1}{2}$ proton of mass M composed of charged spin-$\frac{1}{2}$- spin-0 constituents of mass m and λ, respectively, as in the QCD-motivated quark-diquark model of a nucleon. The azimuthal single-spin asymmetry $\mathcal{P}_y$ transverse to the photon-to-pion production plane decreases as

$\alpha_s(r_\perp^2)x_{bj}Mr_\perp[\ln r_\perp^2]/\bar{r}_\perp^2$ for large $r_\perp$, where $r_\perp$ is the magnitude of the momentum of the current quark jet relative to the q direction. The mass M of the physical proton mass appears here since it determines the ratio of the $L_z = 1$ and $L_z = 0$ matrix elements. The final-state interactions from gluon exchange between the outgoing quark and the target spectator system leads to single-spin asymmetries in deep inelastic lepton-proton scattering at leading twist in perturbative QCD; *i.e.*, the rescattering corrections are not power-law suppressed at large photon virtuality q^2 at fixed x_{bj}. The linear fall-off in $r_\perp$ compensates for the higher twist of the $q\overline{q}$ gluon correlation. The nominal size of the spin asymmetry is thus $C_F\alpha_s(r_\perp^2)a_p$ where a_p is the proton anomalous magnetic moment. Our analysis shows that the single-spin asymmetry which arises from final-state interactions does not factorize since the result depends on the $\left\langle p|\overline{\psi}_q A\psi|p\right\rangle$ proton correlator, not the usual quark distribution derived from $\left\langle p|\overline{\psi}_q(\xi)\psi_q(0)|p\right\rangle$ evaluated at equal light-cone time $\xi^+ = 0$. We thus predict that the single-spin asymmetry in electroproduction is independent of Q^2 at fixed $\Delta = x_{bj}$. This approach also can be applied to single-spin asymmetries in more general hadronic hard inclusive reactions such as $pp \to \Lambda X$.

In the case of deeply virtual Compton scattering $\ell p \to \ell'\gamma(k)X$, the HERMES collaboration [121] has measured a large single spin asymmetry corresponding to the azimuthal correlation $i\vec{S}_p \cdot \vec{k} \times \vec{q}$. The phase of the Bethe-Heitler amplitude is real, whereas the complex phase of the virtual Compton amplitude [69] is dictated by Regge exchange in the $t-$channel: $M_{\rm Compt} \propto (s/Q^2)^{\alpha_R(t)}[1 + e^{i\pi\alpha_R(t)}]$. This is one source of phase interference. In addition, the same final state interaction which provides a phase in $\gamma^*p \to \pi p$ will also contribute to the deeply virtual Compton amplitude.

Recently HERMES has reported a significant single-spin azimuthal asymmetry in exclusive electroproduction of π^+ mesons [122]. These new measurements will open up an important new window to the effects of rescattering in hard exclusive reactions [115] Studies of single-spin asymmetries are also of critical importance in B and B^* decays since the presence of final-state hadronic interactions have to be understood in order to interpret CP-violating parameters [123].

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