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Single-Spin Asymmetries and Transversity in QCD*

Stanley J. Brodsky
Stanford Linear Accelerator Center, Stanford University
Stanford, California 94309
e-mail: sjbth@slac.stanford.edu

Abstract

Initial- and final-state interactions from gluon exchange, normally neglected in the parton model, have a profound effect in QCD hard-scattering reactions, leading to leading-twist single-spin asymmetries, diffractive deep inelastic scattering, diffractive hard hadronic reactions, as well as nuclear shadowing and antishadowing–leading-twist physics not incorporated in the light-front wavefunctions of the target computed in isolation. The physics of such processes thus require the understanding of QCD at the amplitude level; in particular, the physics of spin requires an understanding of the phase structure of final-state and initial-state interactions, as well as the structure of the basic wavefunctions of hadrons themselves. I also discuss transversity in exclusive channels, including how one can use single-spin asymmetries to determine the relative phases of the timelike baryon form factors, as well as the anomalous physics of the normal-normal spin-spin correlation observed in large-angle proton-proton elastic scattering. As an illustration of the utility of light-front wavefunctions, the transversity distribution of a single electron is computed, as defined from its two-particle QED quantum fluctuations.

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1 Introduction

Spin measureables, such as single-spin asymmetries in deep inelastic lepton scattering and transversity correlations in the Drell-Yan process $\overline{p}p \to \ell^+\ell^- X$, probe the structure of hadrons at a fundamental level. Initial- and final-state interactions from gluon-exchange, normally neglected in the parton model, have a profound effect in QCD hard-scattering reactions, leading to leading-twist single-spin asymmetries [1], leading-twist diffractive deep inelastic scattering [2], diffractive hard hadronic reactions, and nuclear shadowing and antishadowing–physics not incorporated in the light-front wavefunctions of the target computed in isolation; in particular, the single-spin asymmetry in semi-inclusive deep inelastic scattering depends on the phase difference of the final-state interactions in different partial waves and the same matrix elements which produce the anomalous magnetic moment of the target nucleon. The physics of spin thus requires an understanding of the phase structure of final-state and initial state interactions as well as knowledge of the basic wavefunctions of hadrons themselves. In this talk I will review the theory of single-spin asymmetries and the utility of light-front wavefunctions in representing the fundamental structure of hadrons in terms of quark and gluon degrees of freedom. I will also show how this formalism allows the computation of the transversity distribution arising from the quantum fluctuation of a physical electron.

2 Single-Spin Asymmetries from Final-State Interactions

Spin correlations provide a remarkably sensitive window to hadronic structure and basic mechanisms in QCD. Among the most interesting polarization effects are single-spin azimuthal asymmetries in semi-inclusive deep inelastic scattering, representing the correlation of the spin of the proton target and the virtual photon to hadron production plane: $\vec{S}_p \cdot \vec{q} \times \vec{p}_H$ [3]. Such asymmetries are time-reversal odd, but they can arise in QCD through phase differences in different spin amplitudes.

The traditional explanation of pion electroproduction single-spin asymmetries in semi-inclusive deep inelastic scattering is that they are proportional to the transversity distribution of the quarks in the hadron h_1 [4, 5, 6] convoluted with the transverse momentum dependent fragmentation (Collins) function $H_1^\perp$, the distribution for a transversely polarized quark to fragment into an unpolarized hadron with non-zero transverse momentum [7, 8, 9, 10, 11].

Dae Sung Hwang, Ivan Schmidt and I have showed that an alternative physical mechanism for the azimuthal asymmetries also exists [1, 12, 13]. The same QCD final-state interactions (gluon exchange) between the struck quark and the proton spectators which lead to diffractive events also can produce single-spin asymmetries (the Sivers effect) in semi-inclusive deep inelastic lepton scattering which survive in the Bjorken limit. See Fig. 1. In contrast to the SSAs arising from transversity and the Collins fragmentation function, the fragmentation of the quark into hadrons is not necessary; one predicts a correlation with the production plane of the quark jet itself $\vec{S}_p \cdot \vec{q} \times \vec{p}_q$.

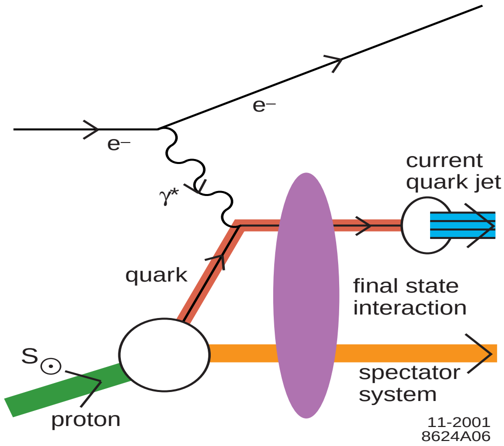

Figure 1: A final-state interaction from gluon exchange in deep inelastic lepton scattering. The difference of the QCD Coulomb-like phases in different orbital states of the proton produces a single proton spin asymmetry.

The final-state interaction mechanism provides an appealing physical explanation within QCD of single-spin asymmetries. Remarkably, the same matrix element which determines the spin-orbit correlation $\vec{S} \cdot \vec{L}$ also produces the anomalous magnetic moment of the proton, the Pauli form factor, and the generalized parton distribution E which is measured in deeply virtual Compton scattering. Physically, the final-state interaction phase arises as the infrared-finite difference of QCD Coulomb phases for hadron wave functions with differing orbital angular momentum. An elegant discussion of the Sivers effect including its sign has been given by Burkardt [14].

The final-state interaction effects can also be identified with the gauge link which is present in the gauge-invariant definition of parton distributions [15]. Even when the light-cone gauge is chosen, a transverse gauge link is required. Thus in any gauge the parton amplitudes need to be augmented by an additional eikonal factor incorporating the final-state interaction and its phase [13, 16]. The net effect is that it is possible to define transverse momentum dependent parton distribution functions which contain the effect of the QCD final-state interactions.

A related analysis also predicts that the initial-state interactions from gluon exchange between the incoming quark and the target spectator system lead to leading-twist single-spin asymmetries in the Drell-Yan process $H_1 H_2^{\updownarrow} \to \ell^+ \ell^- X$ [12, 17]. Initial-state interactions also lead to a $\cos 2\phi$ planar correlation in unpolarized Drell-Yan reactions [18].

2.1 Calculations of Single-Spin Asymmetries in QCD

Hwang, Schmidt and I calculated [1] the single-spin Sivers asymmetry in semi-inclusive electroproduction $\gamma^* p^{\updownarrow} \to HX$ induced by final-state interactions in a model of a spin-1/2 proton of mass M with charged spin-1/2 and spin-0 constituents of mass m and λ, respectively, as in the QCD-motivated quark-diquark model of a nucleon. The basic electroproduction reaction is then $\gamma^* p \to q(qq)_0$. In fact, the asymmetry comes from the interference of two amplitudes which have different proton spin, but couple to the same final quark spin state, and therefore it involves the interference of tree and one-loop diagrams with a final-state interaction. In this simple model the azimuthal target single-spin asymmetry $A_{UT}^{\sin\phi}$ is given by

$$
\begin{aligned}
A_{UT}^{\sin\phi} &= C_F \alpha_s(\mu^2) \, \frac{\Big(\Delta\, M + m \Big)\, r_\perp}{\Big[\Big(\Delta\, M + m \Big)^2 \; + \; \vec{r}_\perp^{\,2} \; \Big]} \\
&\times \; \Bigg[\vec{r}_\perp^{\,2} + \Delta(1-\Delta)(-M^2 + \frac{m^2}{\Delta} + \frac{\lambda^2}{1-\Delta}) \Bigg] \\
&\times \; \frac{1}{\vec{r}_\perp^{\,2}} \, \ln \frac{\vec{r}_\perp^{\,2} + \Delta(1-\Delta)(-M^2 + \frac{m^2}{\Delta} + \frac{\lambda^2}{1-\Delta})}{\Delta(1-\Delta)(-M^2 + \frac{m^2}{\Delta} + \frac{\lambda^2}{1-\Delta})} \; . \qquad (1)
\end{aligned}
$$

Here $r_\perp$ is the magnitude of the transverse momentum of the current quark jet relative to the virtual photon direction, and $\Delta = x_{Bj}$ is the usual Bjorken variable. To obtain (1) from Eq. (21) of [1], we used the correspondence $|e_1 e_2|/4\pi \to C_F \alpha_s(\mu^2)$ and the fact that the sign of the charges e_1 and e_2 of the quark and diquark are opposite since they constitute a bound state. The result can be tested in jet production using an observable such as thrust to define the momentum $q + r$ of the struck quark.

Since the same matrix element controls the Pauli form factor, the contribution of each quark current to the SSA is proportional to the contribution $\kappa_{q/p}$ of that quark to the proton target's anomalous magnetic moment $\kappa_p = \sum_q e_q \kappa_{q/p}$ [1, 14].

The HERMES collaboration has measured the SSA in pion electroproduction using transverse target polarization [19]. The Sivers and Collins effects can be separated using planar correlations; both contributions are observed to contribute, with values not in disagreement with theory expectations [19, 20]. A recent comparison is given by Gamberg and Goldstein [21]. A related analysis also predicts that the initial-state interactions from gluon exchange between the incoming quark and the target spectator system lead to leading-twist single-spin asymmetries in the Drell-Yan process $H_1 H_2^{\updownarrow} \to \ell^+ \ell^- X$ [12, 17]. The SSA in the Drell-Yan process is the same as that obtained in SIDIS, with the appropriate identification of variables, but with the opposite Initial-state interactions also lead to a $\cos 2\phi$ planar correlation in unpolarized Drell-Yan reactions [18]. There is no Sivers effect in charged-current reactions since the W only couples to left-handed quarks [22].

It should be emphasized that the Sivers effect occurs even for jet production;

unlike transversity, hadronization is not required. There is no Sivers effect in charged current reactions since the W only couples to left-handed quarks [22].

We can also consider the SSA of e^+e^- annihilation processes such as $e^+e^- \to \gamma^* \to \pi\Lambda^{\leftrightarrow}X$. The Λ reveals its polarization via its decay $\Lambda \to p\pi^-$. The spin of the Λ is normal to the decay plane. Thus we can look for a SSA through the T-odd correlation $\epsilon_{\mu\nu\rho\sigma}S^\mu_\Lambda p^\nu_\Lambda q^\rho_{\gamma^*} p^\sigma_\pi$. This is related by crossing to SIDIS on a Λ target.

Measurements from Jefferson Lab [23] also show significant beam single spin asymmetries in deep inelastic scattering. Afanasev and Carlson [24] have recently shown that this asymmetry is due to the interference of longitudinal and transverse photoabsorption amplitudes which have different phases induced by the final-state interaction between the struck quark and the target spectators. Their results are consistent with the experimentally observed magnitude of this effect. Thus similar FSI mechanisms involving quark orbital angular momentum appear to be responsible for both target and beam single-spin asymmetries.

3 Light-Front Wavefunctions and Spin

The light-front Fock expansion provides a convenient and rigorous wavefunction formalism for bound states in quantum field theory, analogous to the $\psi(\vec{p})$ momentum-space wavefunction description of nonrelativistic bound states of the Schrödinger theory, but allowing for particle creation and absorption. Unlike equal time theory (the instant form), boost transformations are kinematical in the front form, so that light-front wavefunctions are independent of the total four-momentum of the system. They are relativistic, and frame-independent, describing all particle number excitations n of the hadrons.

The light-front Fock expansion follows from the quantization of QCD at fixed light-front time $x^+ = x^0 + x^3$. The bound-state hadronic solutions $\mid \Psi_H\rangle$ are eigenstates of the light-front Heisenberg equation $H_{LF} \mid \Psi_H\rangle = M^2_H \mid \Psi_H\rangle$ [25]. The spectrum of QCD is given by the eigenvalues M^2_H. The projection of each hadronic eigensolution on the free Fock basis: $\langle n \mid \Psi_H\rangle \equiv \psi_{n/H}(x_i, \vec{k}_{\perp i}, \lambda_i)$ then defines the LF Fock expansion in terms of the quark and transversely polarized gluon constituents in $A^+ = 0$ light-cone gauge. The light-front wavefunctions are functions of the constituent light-cone fractions $x_i = \frac{k^+_i}{P^+} = \frac{(k^0+k^z)_i}{P^+}$, relative transverse momenta $\vec{k}_{\perp i}$, and spin projections $S^z_i = \lambda_i$. The expansion has only transversely polarized gluons.

Recently Guy de Teramond and I have shown how one can use the AdS/CFT correspondence to obtain predictions for LFWFs of hadrons in a conformal approximation of QCD [26].

The freedom to choose the light-like quantization four-vector provides an explicitly covariant formulation of light-front quantization and can be used to determine the analytic structure of light-front wave functions and define a kinematical definition of angular momentum. Angular momentum is consistently defined in the front form.

The z component of angular momentum J^z is kinematical and conserved. The front form thus provides a consistent definition of relative orbital angular momentum and J^z conservation: the total spin projection $J^z = \sum_{i=1}^n S_i^z + \sum_i^{n-1} L_i^z$ is conserved in each Fock state. The cluster decomposition theorem and the vanishing of the "anomalous gravitomagnetic moment" $B(0)$ [27] are immediate properties of the LF Fock wavefunctions [28].

Given the light-front wavefunctions $\psi_{n/H}(x_i, \vec{k}_{\perp i}, \lambda_i)$, one can compute a large range of hadron observables. For example, the valence and sea quark and gluon distributions which are measured in deep inelastic lepton scattering are defined from the squares of the LFWFS summed over all Fock states n. Form factors, exclusive weak transition amplitudes [29] such as $B \to \ell\nu\pi$. and the generalized parton distributions [30] measured in deeply virtual Compton scattering are (assuming the "handbag" approximation) overlaps of the initial and final LFWFS with $n = n'$ and $n = n' + 2$. The gauge-invariant distribution amplitude $\phi_H(x_i, Q)$ defined from the integral over the transverse momenta $\vec{k}^2_{\perp i} \leq Q^2$ of the valence (smallest n) Fock state provides a fundamental measure of the hadron at the amplitude level [31, 32]; they are the nonperturbative input to the factorized form of hard exclusive amplitudes and exclusive heavy hadron decays in perturbative QCD. The resulting distributions obey the DGLAP and ERBL evolution equations as a function of the maximal invariant mass, thus providing a physical factorization scheme [33]. In each case, the derived quantities satisfy the appropriate operator product expansions, sum rules, and evolution equations. However, at large x where the struck quark is far-off shell, DGLAP evolution is quenched [34], so that the fall-off of the DIS cross sections in Q^2 satisfies inclusive-exclusive duality at fixed W^2.

4 Transversity of an Electron in QED

The Fock structure of the electron in QED provides an excellent example of the physics of spin distributions such as the transversity correlation and its relation to orbital angular momentum. The nonzero two-particle LFWFs of the electron in light-cone gauge $A^+ = 0$ are [35]

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

where

$$\varphi = \varphi(x, \vec{k}_\perp) = \frac{e/\sqrt{1-x}}{M^2 - (\vec{k}^2_\perp + m^2)/x - (\vec{k}^2_\perp + \lambda^2)/(1-x)} \ . \tag{3}$$

$$\begin{cases} \psi^{\downarrow}_{+\frac{1}{2}\,+1}(x,\vec{k}_\perp) = 0\ , \\ \psi^{\downarrow}_{+\frac{1}{2}\,-1}(x,\vec{k}_\perp) = -\sqrt{2}(M-\frac{m}{x})\,\varphi\ , \\ \psi^{\downarrow}_{-\frac{1}{2}\,+1}(x,\vec{k}_\perp) = -\sqrt{2}\frac{(-k^1+{\rm i}k^2)}{1-x}\,\varphi\ , \\ \psi^{\downarrow}_{-\frac{1}{2}\,-1}(x,\vec{k}_\perp) = -\sqrt{2}\frac{(+k^1+{\rm i}k^2)}{x(1-x)}\,\varphi\ . \end{cases} \tag{4}$$

The photon has transverse polarization $S^z = \pm 1$. The pre-factors indicate the orbital angular momentum $L^z = 0, \pm 1$ of the two-parton state. Note that each state explicitly satisfies angular momentum conservation $J^z = \sum_{i=1,2} S^z_i + L^z$. The LF thus provides a rigorous definition of orbital angular momentum in relativistic quantum field theory.

We can generalize this form of the LFWFs to describe a quark spin-1 diquark model of the proton by giving the photon an effective mass λ and the fermionic constituent the mass m. A cut-off on the transverse momenta is introduced to regulate the distributions and provide a factorization scale. The model can be generalized further by introducing spectral weights in λ or M to mimic the fall-off of the LFWF of a composite system in QCD. The constituents can have arbitrary charges. Current matrix elements, structure functions, and generalized parton distributions are obtained from the overlap of the LFWFs. In the QED case, $m = M$ and $\lambda = 0$, one recovers the Schwinger anomalous magnetic moment $a_e = \frac{\alpha}{2\pi}$ and the standard formulae for the electron's Dirac and Pauli form factor. In addition one can verify that the anomalous gravitomagnetic moment for the electron scattering in external gravitational field vanishes identically for each Fock state, thus verifying a remarkable theorem based on the equivalence principle.

Dae Sung Hwang and I [36] have shown how the transverse polarization can be computed from the sum and difference of longitudinal LFWFs. The unpolarized, longitudinally polarized structure function and transversity distributions of the quark constituent in the quark-diquark model of the proton can then computed from the appropriate square of the LFWFs:

$$\begin{aligned} & q(x,\Lambda^2)_{\rm spin-1\ diquark} \\ = & \int \frac{{\rm d}^2\vec{k}_\perp {\rm d}x}{16\pi^3}\theta(\Lambda^2-\mathcal{M}^2)\ 2\left[\ \frac{\vec{k}_\perp^2}{x^2(1-x)^2}\ +\ \frac{\vec{k}_\perp^2}{(1-x)^2}\ +\ (M-\frac{m}{x})^2\ \right]|\varphi|^2\ , \\ & \Delta q(x,\Lambda^2)_{\rm spin-1\ diquark} \\ = & \int \frac{{\rm d}^2\vec{k}_\perp {\rm d}x}{16\pi^3}\theta(\Lambda^2-\mathcal{M}^2)\ 2\left[\ \frac{\vec{k}_\perp^2}{x^2(1-x)^2}\ +\ \frac{\vec{k}_\perp^2}{(1-x)^2}\ -\ (M-\frac{m}{x})^2\ \right]|\varphi|^2\ , \\ & \delta q(x,\Lambda^2)_{\rm spin-1\ diquark} \\ = & \int \frac{{\rm d}^2\vec{k}_\perp {\rm d}x}{16\pi^3}\theta(\Lambda^2-\mathcal{M}^2)\ 4\left[\ \frac{\vec{k}_\perp^2}{x(1-x)^2}\ \right]|\varphi|^2\ . \end{aligned} \tag{5}$$

The Soffer bound follows immediately from the LF Fock state construction. In this model $q(x) + \Delta q(x) = 2\delta q(x)$; *i.e.*, Soffer's inequality is satisfied with the equal sign.

5 Single-Spin Asymmetry and the Phase of Time-like Form Factors

As noted by Dubnickova, Dubnicka, and Rekalo [37], and by Rock [38], the existence of the $T-$odd single-spin asymmetry normal to the scattering plane in baryon pair production $e^-e^+ \to B\overline{B}$ requires a nonzero phase difference between the G_E and G_M form factors. The phase of the ratio of form factors G_E/G_M of spin-1/2 baryons in the timelike region can thus be determined from measurements of the polarization of one of the produced baryons. Carlson, Hiller, Hwang and I have shown that measurements of the proton polarization in $e^+e^- \to p\overline{p}$ strongly discriminate between the analytic forms of models which have been suggested to fit the proton G_E/G_M data in the spacelike region [39].

There are three polarization observables, corresponding to polarizations in three directions, called longitudinal, sideways, and normal but often denoted z, x, and y, respectively. Longitudinal (z) when discussing the final state means parallel to the direction of the outgoing or in going baryon. Sideways (x) means perpendicular to the direction of the baryon but in the scattering plane. Normal (y) means normal to the scattering plane, in the direction of $\vec{k} \times \vec{p}$ where $\vec{k}$ is the electron momentum and $\vec{p}$ is the baryon momentum, with x, y, and z forming a right-handed coordinate system. The polarization P_y does not require polarization of a lepton and is [39]

$$A_N = P_y = \frac{\sin 2\theta \, \mathrm{Im} G_E^* G_M}{D\sqrt{\tau}} = \frac{(\tau - 1)\sin 2\theta \, \mathrm{Im} F_2^* F_1}{D\sqrt{\tau}} \ . \tag{6}$$

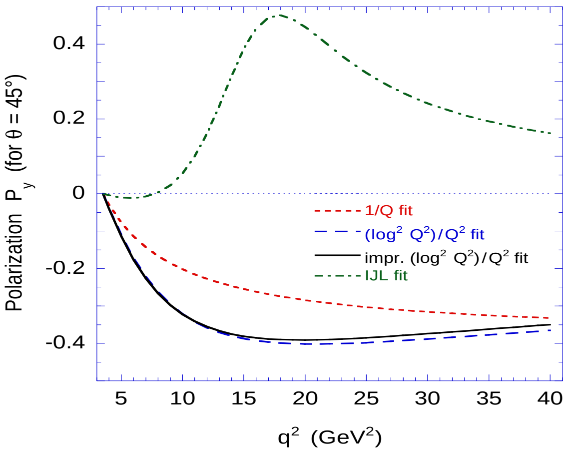

Figure 2: Predicted polarization P_y in the timelike region for selected form factor fits described in the text. The plot is for $\theta = 45°$. The four curves are for an $F_2/F_1 \propto 1/Q$ fit; the $(\log^2 Q^2)/Q^2$ fit of Belitsky *et al.* [40]; an improved $(\log^2 Q^2)/Q^2$ fit; and a fit from Iachello *et al.* [41].

The expression for polarization P_y, Eq. (6), leads to results shown in Fig. 2. The polarizations are shown for four different fits to the spacelike data as referenced in the figure. The value of P_y should be the same for $e^+e^- \to p\overline{p}$ and $\overline{p}p \to \ell^+\ell^-$ up to an overall sign. The predicted polarizations are not small. Note that a purely polynomial fit to the spacelike data gives zero P_y. The normal polarization P_y is a single-spin asymmetry and requires a phase difference between G_E and G_M. It is an example of how time-reversal-odd observables can be nonzero if final-state or initial-state interactions give interfering amplitudes with different phases.

6 Exclusive Transversity

The most remarkable spin correlation ever observed in hadron physics is the normal spin—normal spin correlation asymmetry A_{NN} measured by A. Krisch and his collaborators [42] in large CM-angle pp elastic scattering. Both the projectile and target are polarized normal to the scattering plane. The ratio r_{NN} of cross sections where protons are both parallel to antiparallel becomes as large as 4:1. Since r_{NN} rises dramatically at the strangeness and charm thresholds $\sqrt{s} = 3, 5$ GeV, Guy de Teramond and I [43] have proposed that the these correlations are related to the onset of "octoquark" $uuduuds\overline{s}, uuduudc\overline{c}$ $J = L = S = 1$ resonances; such states only couple to the normal spin parallel channel. The resonance signal interferes with the hard QCD amplitude from quark-interchange, the dominant mechanism for hard scattering in conformal QCD, as predicted from AdS/CFT duality. This mechanism correctly predicts the value of r_{NN} away from the heavy quark thresholds. Furthermore, the resonance phenomenon can account for the dramatic quenching of color transparency observed by Bunce *et al.* in quasielastic pp scattering in nuclei [44]. The octoquark model can be tested at J-PARC and GSI by measuring the rate of charm production near threshold in pp collisions. In addition one can look for the corresponding octoquark states $uud\overline{u}\overline{u}\overline{d}c\overline{c}$ state in $\overline{p}p$ scattering at GSI.

7 Diffractive Deep Inelastic Scattering

A remarkable feature of deep inelastic lepton-proton scattering at HERA is that approximately 10% events are diffractive [45, 46]: the target proton remains intact, and there is a large rapidity gap between the proton and the other hadrons in the final state. These diffractive deep inelastic scattering (DDIS) events can be understood most simply from the perspective of the color-dipole model: the $q\overline{q}$ Fock state of the high-energy virtual photon diffractively dissociates into a diffractive dijet system. The exchange of multiple gluons between the color dipole of the $q\overline{q}$ and the quarks of the target proton neutralizes the color separation and leads to the diffractive final state. The same multiple gluon exchange also controls diffractive vector meson electroproduction at large photon virtuality [47]. This observation presents a paradox: if one chooses the conventional parton model frame where the photon light-front momentum is negative $q^+ = q^0 + q^z < 0$, the virtual photon interacts with a quark constituent with light-cone momentum fraction $x = k^+/p^+ = x_{bj}$. Furthermore, the gauge link associated with the struck quark (the Wilson line) becomes unity in light-cone gauge $A^+ = 0$. Thus the struck "current" quark apparently experiences no final-state interactions. Since the light-front wavefunctions $\psi_n(x_i, k_{\perp i})$ of a stable hadron are real, it appears impossible to generate the required imaginary phase associated with pomeron exchange, let alone large rapidity gaps.

This paradox was resolved by Paul Hoyer, Nils Marchal, Stephane Peigne, Francesco Sannino and myself [2]. Consider the case where the virtual photon interacts with a strange quark—the $s\bar{s}$ pair is assumed to be produced in the target by gluon splitting. In the case of Feynman gauge, the struck s quark continues to interact in the final state via gluon exchange as described by the Wilson line. The final-state interactions occur at a light-cone time $\Delta\tau \simeq 1/\nu$ shortly after the virtual photon interacts with the struck quark. When one integrates over the nearly-on-shell intermediate state, the amplitude acquires an imaginary part. Thus the rescattering of the quark produces a separated color-singlet $s\bar{s}$ and an imaginary phase. In the case of the light-cone gauge $A^+ = \eta \cdot A = 0$, one must also consider the final-state interactions of the (unstruck) $\bar{s}$ quark. The gluon propagator in light-cone gauge $d_{LC}^{\mu\nu}(k) = (i/k^2 + i\epsilon)\left[-g^{\mu\nu} + (\eta^\mu k^\nu + k^\mu \eta^\nu / \eta \cdot k)\right]$ is singular at $k^+ = \eta \cdot k = 0$. The momentum of the exchanged gluon k^+ is of $\mathcal{O}(1/\nu)$; thus rescattering contributes at leading twist even in light-cone gauge. The net result is gauge invariant and is identical to the color dipole model calculation. The calculation of the rescattering effects on DIS in Feynman and light-cone gauge through three loops is given in detail for an Abelian model in the references [2]. The result shows that the rescattering corrections reduce the magnitude of the DIS cross section in analogy to nuclear shadowing.

A new understanding of the role of final-state interactions in deep inelastic scattering has thus emerged. The multiple scattering of the struck parton via instantaneous interactions in the target generates dominantly imaginary diffractive amplitudes, giving rise to an effective "hard pomeron" exchange. The presence of a rapidity gap between the target and diffractive system requires that the target remnant emerges in a color-singlet state; this is made possible in any gauge by the soft rescattering. The resulting diffractive contributions leave the target intact and do not resolve its quark structure; thus there are contributions to the DIS structure functions which cannot be interpreted as parton probabilities [2]; the leading-twist contribution to DIS from rescattering of a quark in the target is a coherent effect which is not included in the light-front wave functions computed in isolation. One can augment the light-front wave functions with a gauge link corresponding to an external field created by the virtual photon $q\bar{q}$ pair current [16, 15]. Such a gauge link is process dependent [12], so the resulting augmented LFWFs are not universal [2, 16, 48]. We also note that the shadowing of nuclear structure functions is due to the destructive interference between multi-nucleon amplitudes involving diffractive DIS and on-shell intermediate states with a complex phase. In contrast, the wave function of a stable target is strictly real since it does not have on-energy-shell intermediate state configurations. The physics of rescattering and shadowing is thus not included in the nuclear light-front wave functions, and a probabilistic interpretation of the nuclear DIS cross section is precluded.

Rikard Enberg, Paul Hoyer, Gunnar Ingelman and I [49] have shown that the quark structure function of the effective hard pomeron has the same form as the quark contribution of the gluon structure function. The hard pomeron is not an intrinsic part of the proton; rather it must be considered as a dynamical effect of the lepton-proton interaction. Our QCD-based picture also applies to diffraction in hadron-initiated processes. The rescattering is different in virtual photon- and hadron-induced processes due to the different color environment, which accounts for the observed non-universality of diffractive parton distributions. This framework also provides a theoretical basis for the phenomenologically successful Soft Color Interaction (SCI) model [50] which includes rescattering effects and thus generates a variety of final states with rapidity gaps.

The phase structure of hadron matrix elements is thus an essential feature of hadron dynamics. Although the LFWFs are real for a stable hadron, they acquire phases from initial state and final state interactions. In addition, the violation of CP invariance leads to a specific phase structure of the LFWFs [51]. Dae Sung Hwang, Susan Gardner and I [51] have shown that this in turn leads to the electric dipole moment of the hadron and a general relation between the electric dipole moment and anomalous magnetic moment, Fock state by Fock state.

The rescattering of the struck parton in DIS generates dominantly imaginary diffractive amplitudes, giving rise to an effective "hard pomeron" exchange and a rapidity gap between the target and diffractive system, while leaving the target intact. This Bjorken-scaling physics, which is associated with the Wilson line connecting the currents in the virtual Compton amplitude survives even in light-cone gauge. Diffractive deep inelastic scattering in turn leads to nuclear shadowing at leading twist as a result of the destructive interference of multi-step processes within the nucleus. Shadowing and antishadowing thus arise because of the $\gamma^* A$ collision and the history of the $q\overline{q}$ dipole as it propagates through the nucleus. Thus there are contributions to the DIS structure functions which are not included in the light-front wave functions computed in isolation and which cannot be interpreted as parton probabilities [2].

There are also leading-twist diffractive contributions $\gamma^* N_1 \to (q\overline{q}) N_1$ arising from Reggeon exchanges in the t-channel [52]. For example, isospin–non-singlet $C = +$ Reggeons contribute to the difference of proton and neutron structure functions, giving the characteristic Kuti-Weisskopf $F_{2p} - F_{2n} \sim x^{1-\alpha_R(0)} \sim x^{0.5}$ behavior at small x. The x dependence of the structure functions reflects the Regge behavior $\nu^{\alpha_R(0)}$ of the virtual Compton amplitude at fixed Q^2 and $t = 0$. The phase of the diffractive amplitude is determined by analyticity and crossing to be proportional to $-1 + i$ for

$\alpha_R = 0.5$, which together with the phase from the Glauber cut, leads to *constructive* interference of the diffractive and nondiffractive multi-step nuclear amplitudes. Furthermore, because of its x dependence, the nuclear structure function is enhanced precisely in the domain $0.1 < x < 0.2$ where antishadowing is empirically observed. The strength of the Reggeon amplitudes is fixed by the fits to the nucleon structure functions, so there is little model dependence. Ivan Schmidt, Jian-Jun Yang, and I [53] have applied this analysis to the shadowing and antishadowing of all of the electroweak structure functions. Quarks of different flavors will couple to different Reggeons; this leads to the remarkable prediction that nuclear antishadowing is not universal; it depends on the quantum numbers of the struck quark. This picture leads to substantially different antishadowing for charged and neutral current reactions, thus affecting the extraction of the weak-mixing angle θ_W. We find that part of the anomalous NuTeV result [54] for θ_W could be due to the non-universality of nuclear antishadowing for charged and neutral currents. Detailed measurements of the nuclear dependence of individual quark structure functions are thus needed to establish the distinctive phenomenology of shadowing and antishadowing and to make the NuTeV results definitive. Antishadowing can also depend on the target and beam polarization.

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