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The Challenges of Exclusive Processes in QCD*

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ABSTRACT

I review some of the outstanding theoretical and experimental issues confronting the application of perturbative quantum chromodynamics to large momentum transfer exclusive reactions.

1. Introduction

The analysis of exclusive hadronic amplitudes such as form factors, electroweak transition matrix elements, and two-body scattering amplitudes has remained among the most challenging computational problems in quantum chromodynamics. The physics of exclusive amplitudes clearly depends on the fundamental relativistic structure of the hadrons as well as the dynamics governing quark and gluon propagation, QCD vacuum structure, Regge behavior, and color confinement. Numerical predictions for exclusive processes involving low momentum transfer are beginning to be obtained from lattice gauge theory and QCD sum rules. However, the most interesting insights into hadron structure at the amplitude level and the most transparent connections to the underlying QCD physics emerges at high momentum transfer where perturbative analyses for the leading twist contributions to exclusive processes can be combined with non-perturbative hadron wavefunction information.

The least-complicated exclusive amplitudes to analyze from first principles in QCD are the space-like electromagnetic form factors of hadrons. An elastic form factor is the probability amplitude for a hadron to remain intact after absorbing momentum q by its local quark current. If one uses light-cone quantization in the $q^+ = q^0 + q^z = 0$ frame with $\vec{q}_\perp^{\,2} = -q^2 = Q^2$, then vacuum fluctuation contributions to the j^+ current can be avoided. Nevertheless, the computation of an elastic form factor requires knowledge of all of the hadron's light-cone Fock state wavefunctions. For example, the helicity-conserving form factor has the form[1]

$$F(Q^2) = \langle p+q|j^+|p\rangle /2p^+$$

$$= \sum_{n,\lambda_i} \sum_a e_a \int \overline{\prod_i} \frac{dx_i\, d^2\vec{k}_{\perp i}}{16\pi^3}\, \psi_n^{(\Lambda)*}(x_i, \vec{\ell}_{\perp i}, \lambda_i)\, \psi_n^{(\Lambda)}(x_i, \vec{k}_{\perp i}, \lambda_i).$$

The constituents in the initial state have longitudinal light-cone momentum fractions $x_i = (k^0 + k^z)_i/(p^0 + p^z)$, relative transverse momentum, $\vec{k}_{\perp i}$, and helicities λ_i. Here e_a is the charge of the struck quark, $\Lambda^2 \gg \vec{q}_\perp^{\,2}$, and the transverse momenta in the final state are

$$\vec{\ell}_{\perp i} \equiv \begin{cases} \vec{k}_{\perp i} - x_i \vec{q}_\perp + \vec{q}_\perp & \text{for the struck quark} \\ \vec{k}_{\perp i} - x_i \vec{q}_\perp & \text{for all other partons.} \end{cases}$$

In principle, one can obtain all of the required Fock State wavefunctions by diagonalizing the light-cone QCD Hamiltonian.[2] This has in fact been done for meson and baryon wavefunctions in the case of QCD in one-space and one-time dimensions, but the corresponding task appears to be formidable for QCD(3+1).

Fortunately, because of asymptotic freedom and the point-like behavior of quark and gluon interactions at short distances, the computation of exclusive amplitudes in QCD becomes much simpler at large momentum transfer. The primary ingredient in the analysis is factorization: the non-perturbative dynamics of the bound states can be isolated in terms of process-independent distribution amplitudes, and the dynamics of the momentum transfer to the hadrons can be isolated in terms of perturbatively-calculable hard-scattering quark and gluon subprocesses. Thus general properties of exclusive reactions at large momentum transfer can be derived without explicit knowledge of the non-perturbative structure of the theory.[3]

The most characteristic feature of an exclusive amplitude in QCD is that it falls off slowly with momentum transfer, not as an exponential or a Gaussian, but as an inverse power of $Q = p_T$ which is directly related to the degree of complexity of the scattering hadrons. The nominal power-law fall-off[4] $\mathcal{M} \sim Q^{4-n}$ of an exclusive amplitude at large momentum transfer reflects the elementary scaling of the lowest-order connected quark and gluon tree graphs obtained by replacing each of the external hadrons by its respective collinear quarks. Here n is the total number of initial state and final state lepton, photon, or quark fields entering or leaving the hard scattering subprocess. The empirical success of the dimensional counting rules for the power-law fall-off of form factors and general fixed center-of-mass angle scattering amplitudes gave early and important evidence for the scale-invariance of quark and gluon interactions at short distances.

Thus only the valence-quark Fock components of the hadron wavefunctions contribute to the leading power-law fall-off of an exclusive amplitude. In particular, since the internal momentum transfer at the quark level is required to be large, one can obtain the basic scaling and helicity structure of the hadron amplitude by simply iterating the gluon-exchange term in the effective potential for the light-cone wavefunctions. The result is that exclusive amplitudes at high momentum transfer Q^2 can be written in a factorized form as a convolution of process-independent "distribution amplitudes" $\phi(x_i, Q)$, one for each hadron involved in the amplitude, with a hard-scattering amplitude T_H describing the scattering of the valence quarks from the initial to final state.[5,6]

The distribution amplitude is the fundamental gauge invariant wavefunction which describes the fractional longitudinal momentum distributions of the valence quarks in a hadron integrated over transverse momentum up to the scale Q.[5] For example, the pion's electromagnetic form factor can be written as[5,6,7]

$$F_\pi(Q^2) = \int_0^1 dx \int_0^1 dy\, \phi_\pi^*(y,Q)\, T_H(x,y,Q)\, \phi_\pi(x,Q) \left(1 + \mathcal{O}\left(\frac{1}{Q}\right)\right).$$

Here T_H is the scattering amplitude obtained when pions replaced by collinear $q\overline{q}$ pairs. This factorized form is the prototype for the factorization of general exclusive amplitudes in QCD at high momentum transfer. All of the non-perturbative dynamics is factorized into the distribution amplitudes,[5] $\phi_B(x_i, \lambda_i, Q)$, for the baryons with $x_1 + x_2 + x_3 = 1$, and $\phi_M(x_i, \lambda_i, Q)$, for the mesons with $x_1 + x_2 = 1$ which sum all internal momentum transfers up to the scale Q^2. On the other hand, all momentum transfers higher than Q^2 appear in T_H, which can be computed perturbatively in powers of the QCD running coupling constant $\alpha_s(Q^2)$. The distribution amplitudes are thus the process-independent hadron wavefunctions which interpolate between the QCD bound state and their valence quarks at transverse separation $b_\perp \simeq 1/Q$. The pion's distribution amplitude, for example, is directly related to its valence light-cone wavefunction:

$$\begin{aligned}\phi_\pi(x,Q) &= \int \frac{d^2\vec{k}_\perp}{16\pi^3}\, \psi^{(Q)}_{q\overline{q}/\pi}(x, \vec{k}_\perp) \\ &= P_\pi^+ \int \frac{dz^-}{4\pi}\, e^{ixP_\pi^+ z^-/2} \left\langle 0 \left| \overline{\psi}(0)\, \frac{\gamma^+\gamma_5}{2\sqrt{2n_c}}\, \psi(z) \right| \pi \right\rangle^{(Q)} \Bigg|_{z^+ = \vec{z}_\perp = 0}.\end{aligned}$$

The $\vec{k}_\perp$ integration is cut off by the ultraviolet cutoff $\Lambda = Q$ implicit in the wavefunction; thus only valence Fock states with invariant mass squared $\mathcal{M}^2 \leq Q^2$ contribute.

Given the factorized structure of exclusive amplitudes at large momentum transfer, one can read off a number of general features of the PQCD predictions: the dimensional counting rules, hadron helicity conservation, and color transparency.[3] QCD also predicts calculable corrections to the nominal dimensional counting power-law behavior due to the running of the strong coupling constant, higher order corrections to the hard scattering amplitude, Sudakov effects, pinch singularities, as well as the evolution of the hadron distribution amplitudes, $\phi_H(x_i, Q)$.

Evolution equations for the meson and baryon distribution amplitudes can be derived and employed in analogy to the evolution of structure functions.[3,8] If one can calculate the distribution amplitude at an initial scale Q_0 using QCD sum rules or lattice gauge theory,[8] then one can determine $\phi(x_i, Q)$ at higher momentum scales via evolution equations in $\log Q^2$ or equivalently, the operator product expansion.[9] Empirical constraints on the hadron distribution amplitudes can be obtained from the normalization and scaling of form factors at large momentum transfer and the angular dependence of two body scattering amplitudes.

Perhaps the most surprising feature of the QCD predictions for exclusive processes in QCD is "color transparency",[10] which reflects the fact that only the small transverse separation $b_\perp \sim 1/Q$ valence wavefunction can contribute to exclusive amplitude at large momentum transfer. Since these color-singlet states have small color-dipole moments, they will have small initial and final state interactions. In particular if the large momentum transfer occurs as a quasi-elastic process within a nucleus, there will be minimal initial state or final state absorption—in striking contrast to the standard picture of strong absorption predicted in Glauber theory. A careful treatment of color transparency requires consideration of the expansion time and coherence length of the small size configurations.[11]

2. A Detailed Example: Compton Scattering in Perturbative QCD

Exclusive reactions involving two real or virtual photons provide a particularly interesting testing ground for QCD because of the relative simplicity of the couplings of the photons to the underlying quark currents, and the absence of significant initial state interactions—any remnant of vector-meson dominance contributions is suppressed at large momentum transfer, and the photon enters the amplitude as a direct point-like coupling.

The simplest example of a two-photon exclusive process is the $\gamma^*(q)\gamma \to M^0$ process which is measurable in tagged $e^+e^- \to e^+e^-M^0$ reactions. The photon to neutral meson transition form factor $F_{\gamma\to M^0}(Q^2)$ is predicted to fall as $1/Q^2$—modulo calculable logarithmic corrections from the evolution of the meson distribution amplitude. This QCD prediction reflects the elementary scaling of the quark propagator at high momentum transfer, the same scale-free behavior which leads to Bjorken scaling of the deep inelastic lepton-nucleon cross sections. The existing data from TPC/$\gamma\gamma$ are consistent with the predicted scaling and normalization of the transition form factors for the π^0, η_0, and η'.

The angular distributions for the hadron pair production processes $\gamma\gamma \to H\overline{H}$ are sensitive to the x_i dependence of the hadron distribution amplitudes.[12] Lowest order predictions for meson pair production in two photon collisions using this formalism are given in Refs. 12 and 8; the analysis of the $\gamma\gamma$ to meson pair process has been carried out to next-to-leading order in $\alpha_s(Q^2)$ by Nizic.[13] The Mark II and TPC/$\gamma\gamma$ measurements of $\gamma\gamma \to \pi^+\pi^-$ and $\gamma\gamma \to K^+K^-$ reactions are also consistent with PQCD expectations. A review of this work is given in Ref. 14.

Compton scattering $\gamma p \to \gamma p$ at large momentum transfer and its s-channel crossed reactions $\gamma\gamma \to \overline{p}p$ and $\overline{p}p \to \gamma\gamma$ are classic tests of the perturbative QCD formalism for exclusive reactions. At leading twist, each helicity amplitude has the factorized form,[3]

$$\mathcal{M}_{hh'}^{\lambda\lambda'}(s,t) = \sum_{d,i} \int [dx][dy]\phi_i(x_1,x_2,x_3,\widetilde{Q})T_i^{(d)}(x,h,\lambda;y,h',\lambda';s,t)\phi_i(y_1,y_2,y_3;\widetilde{Q}) \ .$$

The index i labels the three contributing valence Fock amplitudes at the renormalization scale $\widetilde{Q}$. The index d labels the 378 connected Feynman diagrams which contribute to the eight-point hard scattering amplitude $qqq\gamma \to qqq\gamma$ at the tree level; *i.e.* at order $\alpha\alpha_s^2(\widehat{Q})$. The arguments $\widehat{Q}$ of the QCD running coupling constant can be evaluated amplitude by amplitude using the method of Ref. 15. The evaluation of the hard scattering amplitudes $T_i^{(d)}(x,h,\lambda;y,h',\lambda';s,t)$ has now been done by several groups.[16,17,18,19]

An important simplification of Compton scattering in PQCD is the fact that pinch singularities are readily integrable and do not change the nominal power-law behavior of the basic amplitudes.[18] Physically, the pinch singularities correspond to the existence of potentially on-shell intermediate states in the hard scattering amplitudes. This leads to a non-trivial phase structure of the Compton amplitude. Such phases can in principle be measured by interfering the virtual Compton process in $e^{\pm}p \to e^{\pm}p\gamma$ with the purely real Bethe-Heitler bremsstrahlung amplitude.[20] A careful analytic treatment of the integration over the on-shell intermediate states has been given by Kronfeld and Nizic.[18]

The most characteristic feature of the PQCD predictions is the scaling of the differential Compton cross section at fixed t/s or θ_{CM}

$$s^6\frac{d\sigma}{dt}(\gamma p \to \gamma p) = F(t/s).$$

The power s^6 reflects the fact that 8 elementary fields enter or leave the hard scattering subprocess.[4] The scaling of the existing data[21] is remarkably consistent with the PQCD power-law prediction, but measurements at higher energies and momentum transfer are needed to test the predicted logarithmic corrections to this scaling behavior and determine the angular distribution of the scaled cross section over as large a range as possible.

The predictions for the normalization of the Compton cross section and the shape of its angular distribution are sensitive to the shape of the proton distribution amplitude $\phi_p(x_i,Q)$. The forms predicted for the proton distribution amplitude from QCD sum-rule constraints[8] by Chernyak, Oglobin, and Zhitnitskii, and King and Sachrajda, appear to give a reasonable representation of the existing data. More recent QCD sum rule analyses of the proton distribution amplitude are given in Ref. 22. These distributions, which predict that approximately 65% of the proton's momentum is carried by the u quark with helicity parallel to the proton's helicity also provide empirically consistent predictions for the normalization of the proton's form factor and the $J/\psi \rightarrow p\overline{p}$ decay rate. The crossing behavior from spacelike Compton scattering to the timelike annihilation channels will also provide important tests and constraints on the PQCD formalism and the shape of the proton distribution amplitudes. Predictions for the time-like processes have been made by Farrar *et al.*,[16] Millers and Gunion[17], and Hyer.[19]

The theoretical uncertainties from finite nucleon mass corrections, the magnitude of the QCD running coupling constant, and the normalization of the proton distribution amplitude largely cancel out in the ratio of differential cross sections

$$R_{\gamma\gamma/e^+e^-}(s,\theta_{cm}) = \frac{d\sigma(\overline{p}p \rightarrow \gamma\gamma)/dt}{d\sigma(\overline{p}p \rightarrow e^+e^-)/dt},$$

which is predicted by QCD to be essentially independent of s at large momentum transfer. If this scaling is confirmed, then the center-of-mass angular dependence of $R_{\gamma\gamma/e^+e^-}(s,\theta_{cm})$ will be one of the best ways to determine the shape of $\phi_p(x_i,Q)$.

2. Lepto-Production of Vector Mesons as a Test of PQCD and Color Transparency

The study of real and virtual photoproduction of vector mesons on protons and nuclei provides an elegant illustration of the emergence of perturbative QCD features in the large momentum transfer domain.[23,24]

1. At small momentum transfer and high energy where the coherence length $2\nu/(\mathcal{M}^2+Q^2)$ is large compared to the target size, the incident photon is expected to act as a coherent sum of vector mesons with mass squared $\mathcal{M}^2 \leq \mathcal{O}(Q^2)$. This is the generalized vector meson dominance picture of photon interactions. In addition, $s-$channel helicity conservation predicts that the vector meson will be dominantly produced with transverse polarization equal to that of the incident photon.
2. At small momentum transfer where photon interactions are dominantly hadron-like, the cross section for vector meson photoproduction on a nucleus should have the same nuclear properties as meson-nucleon scattering. Due to the optical theorem, the forward high energy coherent nuclear amplitude $\gamma^*A \rightarrow V^0A$ must then scale with the nuclear size the same as the total hadron-nucleus cross section; *i.e.* $A^{2/3}$. The $t-$dependence of the coherent nuclear cross section is of the form $d\sigma/dt \sim \exp^{b_At}$ where $b_A \propto R_A^2$ and R_A is the nuclear size. Thus the total coherent cross section $\sigma(\gamma^*A \rightarrow V^0A)$ is predicted to scale with nuclear number as $A^{4/3}/R_A^2 \sim A^{2/3}$.
3. The predictions for $\gamma^*A \rightarrow V^0A'$ are in striking contrast to the above results when Q^2 becomes large compared to Λ^2_{QCD}. The virtual quark loop connecting the photon to the vector meson is now highly virtual, and only the point-like

piece of the photon and the small transverse size of the valence $q\bar{q}$ light-cone wavefunction of the vector meson enter the exclusive amplitude. Thus at high Q^2 the nuclear absorption in the initial and final state should vanish, and the nuclear amplitude becomes additive: $M(\gamma^* A \to V^0 A') = A^1 M(\gamma^* N \to V^0 N')$. The integrated coherent cross section $\sigma(\gamma^* A \to V^0 A)$ is thus predicted to scale with nuclear number as $A^2/R_A^2 \sim A^{4/3}$. This contrasting nuclear dependence of the virtual photoproduction cross section provides a dramatic test of color transparency. Preliminary results from E665[25] for ρ lepto-production at Fermilab appear to confirm these QCD predictions.

4. Another important prediction of PQCD in the large Q^2 domain is that the vector meson should be produced with zero helicity since it is formed from a quark and antiquark with equal and opposite helicities.[26] The change-over from transverse to longitudinal vector meson polarization with increasing Q^2 also appears to be confirmed by the E665 data.
5. At large photon virtuality Q^2 the photon and vector meson will act as point-like systems, and thus the $t-$ dependence of the differential cross section $d\sigma/dt(\gamma^* p \to V^0 p')$ should only reflect the finite size of the scattered nucleon. At large t the form factors should reflect the underlying two-gluon exchange structure of the PQCD Pomeron.
6. At large momentum transfer $-t \gg \Lambda^2_{QCD}$, $-u \gg \Lambda^2_{QCD}$, PQCD predicts that the photoproduction cross section has the nominal fixed CM angle scaling: $d\sigma/dt(\gamma p \to V^0 p') \sim f(\theta_{CM})/s^7$. The dominant amplitudes will conserve hadron helicity: $\lambda_{p'} + \lambda_V = \lambda_p$.
7. At larger momentum transfers $-t > R_A^2$, one can study quasi-elastic lepto-production in the nucleus; $d\sigma/dt(\gamma^* A \to V^0 N' X)$ where X represents a sum over excited nuclear states, but without extra particle production. When $p_T^2 \gg \Lambda^2_{QCD}$, color transparency predicts the absence of initial or final state absorption of the incident photon and the outgoing meson and nucleon. Thus the quasi-elastic cross section should approach additivity in nuclear number at large momentum transfer.

4. When Do Leading-Twist Predictions for Exclusive Processes Become Applicable?

The factorized predictions for exclusive amplitudes are evidently rigorous predictions of QCD at large momentum transfer. However, it is important to understand the kinematic domain where the leading twist predictions become valid. The basic scales of QCD are set by the quark masses and the scale Λ_{QCD} which parameterizes the QCD running coupling constant. Thus one normally would expect that the leading power-law predictions should become dominant at momentum transfers exceeding these parameters. In the case of inclusive reactions, Bjorken scaling is already apparent at momentum transfers $Q \sim 1$ GeV or less.

In fact, the data for hadron form factors is consistent with the onset of PQCD scaling at momentum transfers of a few GeV. Stoler[27] has shown that the measurements of the transition form factors of the proton to the $N(1535)$ and $N(1680)$ resonances are consistent with the predicted PQCD Q^{-4} scaling to beyond $Q^2 = 20\ GeV^2$. The normalization is also in reasonable agreement with that predicted from QCD sum rule constraints on the nucleon distribution amplitudes, allowing for uncertainties from higher order QCD corrections. In the case of the proton to $\Delta(1232)$ transition, the form factor falls faster that Q^{-4}. This anomalous behavior is, in fact, predicted by QCD sum rule constraints, since unlike the proton, the Δ has a highly symmetric distribution amplitude which results in a small net coupling to the QCD hard scattering amplitude. The observed scaling pattern of the transition form factors gives strong support to the QCD sum rule predictions and PQCD factorization.

Isgur and Llewellyn Smith[28] and Radyushkin[29] have raised the concern that important contributions to exclusive processes could arise from the endpoint regions $x_i \to 1$; such behavior would imply the breakdown of PQCD factorization. For example, the denominator of the hard scattering amplitudes, *e.g.*, $T_H \propto \alpha_s/[(1-x)(1-y)Q^2]$ for the meson form factor becomes singular in the endpoint integration region at $x \sim 1$ and $y \sim 1$. Such endpoint regions are even further emphasized when one assumes the strongly asymmetric forms for the hadron distribution amplitudes derived from QCD sum rules. However, it is important to note that these endpoint regimes correspond to scattering processes where one quark carries nearly all of the proton's momentum and is at a fixed transverse separation $b_\perp$ from the spectator quarks.

When a quark which is isolated in space receives a large momentum transfer x_iQ, it will normally strongly radiate gluons into the final state due to the displacement of both its initial and final self-field, which is contrary to the requirements of exclusive scattering. For example, in QED the radiation from the initial and final state charged lines is controlled by the coherent sum $\sum_i \frac{\epsilon\cdot p_i}{k\cdot p_i}\eta_i q_i$ where q_i and p_i are the charges four-momenta of the charged lines, ϵ and k are polarization and four-momentum of the radiation, and $\eta_i = \pm 1$ for initial and final state particles, respectively. Radiation will occur for any finite momentum transfer scattering as long as the photon's wavelength is less than the size of the initial and final neutral bound states. The probability amplitude that radiation does not occur is given by rapidly falling Sudakov form factor, as first discussed by in Refs. 5 and 30. An elegant and much more complete discussion has now been given by Botts and Li and Sterman.[31] The radiation from the colored lines in QCD have similar coherence properties as in QED:[32] because of the destructive color interference of the radiators, the momentum of the radiated gluon in a QCD hard scattering process only ranges from k of order $1/b_\perp$, where color screening occurs, up to the momentum transfer x_iQ of the scattered quarks. This analysis and unitarity allows one to compute the probability that no radiation occurs during the hard scattering.[31,19] It is given by a rapidly falling exponentiated Sudakov form factor $S = S(x_iQ, b_\perp, \Lambda_{QCD})$; thus at large Q and fixed impact separation, the Sudakov factor strongly suppresses the endpoint contribution.

On the other hand, when $b_\perp = \mathcal{O}(x_i Q)^{-1}$, the Sudakov form factor is of order 1, and the radiation leads to logarithmic evolution and contributions of higher order in $\alpha_s(Q^2)$, the corrections already contained in the PQCD predictions.[5,30,33] This is the starting point of the detailed analysis of the suppression of endpoint contributions to meson and baryon form factors and its quantitative effect on the PQCD predictions recently presented by Li and Sterman.[31] This analysis has now also been applied to two-photon reactions and the timelike proton form factor by Hyer.[19]

Thus the leading PQCD contributions to large momentum transfer exclusive reactions derive from wavefunction configurations where the valence quarks are at small transverse separation $b_\perp = \mathcal{O}(1/k_\perp) = \mathcal{O}(1/Q)$, the regime where there is no Sudakov suppression. Furthermore, as noted by Li and Sterman, the hard scattering amplitude loses its singular endpoint structure if one explicitly retains the valence quark transverse momenta in the denominators. For example, in the case of the pion form factor, the hard scattering amplitude is effectively modified to the form

$$T_H \propto \frac{\alpha_s}{(1-x)(1-y)Q^2 + (\mathbf{k}_1^\perp + \mathbf{k}_2^\perp)^2}.$$

The Sudakov effect thus ensures that the denominators are always protected at large momentum transfers. In their numerical studies, Li and Sterman find that the pion form factor becomes relatively insensitive to soft gluon exchange at momentum transfers beyond 20 Λ_{QCD}. In the case of the proton Dirac form factor, the corresponding analysis by Li[31] is in good agreement with experiment at momentum transfers greater than 3 GeV. Thus the leading twist QCD predictions based on the factorization of long and short distance physics appear to be self-consistent and valid for momentum transfers as low as a few GeV, thus accounting for the empirical success of quark counting rules in exclusive process phenomenology. The Sudakov effect suppression also enhances the QCD "color transparency" phenomena, since only small color singlet wavefunction configurations can scatter at large momentum transfer without radiation.[10]

The extension of the leading order PQCD analysis to higher orders including Sudakov effects is technically very challenging. Thus far, the next-to-leading $\alpha_s(Q^2)$ corrections to the hard scattering amplitudes T_H have been computed for only a few exclusive processes: the meson form factor, the photon-to-meson transition form factors, and $\gamma\gamma$ to meson pairs. There are many outstanding theoretical issues which are being resolves, such as how to extend these calculations to baryon processes, how to set the renormalization scale in α_s,[15] how to implement conformal symmetry and its breaking,[9,34] and how to formulate and solve the evolution equations for the hadron distribution amplitudes to next-to-leading order.

6. Other Applications of Large Momentum Transfer Exclusive QCD.

The factorization techniques used to derive the leading-twist behavior of exclusive amplitudes have general applicability to processes where hadron wavefunctions have to be evaluated at far off-shell configurations. In each of these applications, one can separate the perturbative quark and gluon dynamics from momentum transfer higher than a scale Q from the non-perturbative long-distance physics contained in the distribution amplitudes $\phi(x_i, Q)$. For example at $x \sim 1$ the struck quark in deep inelastic lepton-hadron scattering is kinematically far off shell and space-like. Thus the leading power law fall off in $(1-x)$ is determined by iterating the gluon exchange kernel in the valence Fock state wavefunction. In this way one derives "spectator" counting rules for the nominal power law behavior [e.g. $G_{q/p}(x) \sim (1-x)^3$] and helicity-retention rules at $x \to 1$. The resulting structure functions connect smoothly to the behavior of large momentum transfer elastic and inelastic transition form factors at fixed $\mathcal{M}^2$. In fact, when $(1-x)Q^2$ is fixed, the usual evolution of the structure functions breaks down and there is no increase in the effective power beyond that given by the spectator counting rules. Further discussion may be found in Ref. 35.

Higher-twist corrections to inclusive reactions are of two types: coherent corrections which depend on the multiparticle structure of hadrons, and single particle corrections, such as mass and condensate insertions, which affect single quark or single gluon propagators. Exclusive processes represent the completely coherent limit of dynamical higher twist terms in inclusive reactions. At fixed $(1-x)Q^2$, the multi-quark higher twist contributions can be computed using the exclusive factorization analysis, and they contribute at the same order as the leading twist terms.[36,37] Strong higher-twist corrections are in fact observed in the angular and Q^2−dependence of Drell-Yan processes and in deep inelastic lepton scattering at $x \sim 1$.[38]

The factorization techniques used to derive the leading twist contributions to form factors can also be applied to the exclusive decays of heavy hadrons when large momentum transfers are involved. An interesting example of this analysis is "atomic alchemy",[39] *i.e.* the exclusive decays of muonic atoms to electronic atoms plus neutrinos. In this case the calculation requires the high momentum tail of the atomic wavefunctions, which in turn can be obtained via the iteration of the relativistic atomic bound-state equations. Again one obtains a factorization theorem for exclusive atomic transitions where the atomic wavefunction at the origin plays the role of the distribution amplitude.

7. Outstanding Phenomenological Issues in Exclusive Processes.

Although most large momentum transfer exclusive reactions appears to be empirically consistent with perturbative QCD expectations, there are a number of glaring exceptions where theory and experiment diverge. If one accepts that the underlying formalism for the leading twist behavior of exclusive reactions is reliable, then these exceptions provide important insights into new physical mechanisms within QCD.

What accounts for the structure in the spin correlations in *pp* elastic scattering at large momentum transfer? Measurements[40] of large angle *pp* elastic scattering at Argonne and Brookhaven show a dramatic spin-spin correlation

A_{NN} which reaches ~ 0.6 at $\sqrt{s} \sim 5$ GeV: *i.e.* the spin-analyzed cross section is four times larger if the protons scatter with their spins parallel and normal to the scattering plane compared to antiparallel. The explanation for this phenomena is far from settled. The most popular explanations[41] are based on the interference of Landshoff pinch singularities[42] with the quark interchange amplitude, but there is no understanding why the Landshoff contribution would itself have a large A_{NN}[43] or sufficient normalization[44] to explain this phenomena. Guy de Teramond and I have proposed[45] that the large spin correlations reflects inelastic channels corresponding to the production of charm at threshold. This effect leads to enhancement in the $J = L = S = 1$ $pp \to pp$ partial wave which implies a large value of A_{NN} at the energies sufficient to produce open charm. This explanation would be confirmed by the observation of a sizeable charm production rate of order $1 \mu bn$. A similar enhancement of A_{NN} is seen at the open strangeness threshold regime. The heavy quark explanation has received some support from the work of Luke, Savage, and Manohar,[46] who have shown that the interactions of $c\bar{c}$ systems at low relative velocity with hadrons is enhanced due to the QCD scale anomaly; in fact, the scalar exchange interaction is predicted to be strong enough to bind charmonium to heavy nuclei.[47]

Why does QCD color transparency appear to break down in quasielastic *pp* scattering? The Brookhaven measurements[48] of the transparency ratio for large angle quasi-elastic *pp* scattering increases with momentum transfer, as predicted by PQCD, but the ratio then appears to revert to normal absorption at $\sqrt{s} \sim 5$ GeV. This suggests that whatever is causing the structure in A_{NN} at the same energies and angles involves large transverse sizes and is far from perturbative in origin. The charm threshold effect is a candidate for this type of explanation.

The preliminary results for the SLAC color transparency experiment NE18 reported at this meeting[49] indicate that color transparency in quasi-elastic *ep* scattering is not a strong effect up to the accessible momentum transfers. Higher momentum transfers exceeding 5 GeV are needed for a decisive test. A sensitive test of color transparency is provided by measuring the sign of the derivative of the transparency ratio $d/dQ^2 \sigma(eA \to e'p(A-1))/Z\sigma(ep \to e'p)$. Perturbative QCD predicts a positive slope, whereas conventional Glauber theory predicts a negative derivative in the low Q^2 domain.

Why does the J/ψ decay copiously to $\rho\pi$? According to the principle of hadron helicity conservation[26] in exclusive decays, the J/ψ produced with $J_z = \pm 1$ in e^+e^- annihilation should not decay to vector plus pseudoscalar meson pairs. In fact, this is true for the ψ' and other S-state charmonium states, but in the case of the J/ψ, the $\rho\pi$ and KK^* psuedoscalar-vector meson channels are actually the dominant two-body hadronic decays. A possible explanation is that the J/ψ mixes with a nearby gluonic or hybrid $J = 1$ state $\mathcal{O}$ that favors vector plus pseudoscalar meson pair decay.[50] One can search for the $\mathcal{O}$ by looking for a $\rho\pi$ mass peak near the J/ψ in the decay $\psi' \to \pi\pi\mathcal{O} \to \pi\pi\rho\pi$.

Why do effective Reggeon trajectories flatten to values below $\alpha_R(t) = 0$ at large momentum transfer? A fundamental prediction of perturbative QCD is that the Reggeon trajectories $\alpha_\rho(t)$ and $\alpha_{A_2}(t)$ governing charge exchange reactions at high energies $s \gg -t$ monotonically approach zero at large spacelike momentum transfer.[51] More generally, the leading Reggeon in an exclusive process will reflect the minimal particle number exchange quantum numbers: two gluons in the case of the Pomeron, three gluons in the case of the Odderon, and quark plus anti-quark in the case of meson exchange trajectories. Because of asymptotic freedom the leading trajectory at large momentum transfer is thus simply $j_1 + j_2 - 1$ with corrections of order $\sqrt{\alpha_s}(-t)$. The asymptotic prediction $\lim_{-t\to\infty} \alpha_R(t) = 0$ reflects the fact that a weakly interacting quark-antiquark pair is exchanged in the $t-$channel.[51] Thus one expects that the effective ρ Reggeon should asymptote at $\alpha_\rho(t) \to 0$ at large $-t$. However, measurements of the inclusive processes $\pi^- p \to \pi^0 X$ at $s \simeq 300$ GeV^2 and $8 > -t > 2$GeV^2 indicate that the effective non-singlet ρ trajectory becomes negative at large $-t$.[52] Thorn, Tang and I have recently shown that the hard QCD part of the trajectory is weakly coupled and that its contribution may well be hidden until much higher energy.[53] Quark interchange[54] may thus be the dominant subprocess at presently accessible kinematic ranges. We also show that Reggeon contributions to exclusive and semi-inclusive mesonic exchange hadron reactions can be systematically studied in perturbative QCD.

Why is quark interchange the dominant mechanism for large-angle hadron-hadron scattering? The comprehensive measurements at BNL[55] of the relative normalization and angular dependence of a large set of exclusive hadron scattering channels strongly suggests that the dominant mechanism for scattering hadrons at large momentum transfer is quark interchange.[54] For example, if gluon exchange were the dominant mechanism, then the differential cross sections for $K^+ p \to K^+ p$ and $K^- p \to K^- p$ at large p_T would be roughly equal in magnitude and angular shape. In fact they have grossly different magnitudes and shapes. The $K^+ p \to K^+ p$ cross section has the approximate form predicted by the exchange of their common u quark. A possible explanation of this fact is that quark interchange involves the least number of large momentum exchanges within the hadron scattering amplitude.

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