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Exclusive Processes: Tests of Coherent QCD Phenomena and Nucleon Substructure at CEBAF*

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EXCLUSIVE PROCESSES: TESTS OF COHERENT QCD PHENOMENA AND NUCLEON SUBSTRUCTURE AT CEBAF

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ABSTRACT

Measurements of exclusive processes such as electroproduction, photoproduction, and Compton scattering are among the most sensitive probes of proton structure and coherent phenomena in quantum chromodynamics. The continuous electron beam at CEBAF, upgraded in laboratory energy to 10–12 GeV, will allow a systematic study of exclusive, semi-inclusive, and inclusive reactions in a kinematic range well-tuned to the study of fundamental nucleon and nuclear substructure. I also discuss the potential at CEBAF for studying novel QCD phenomena at the charm production threshold, including the possible production of nuclear-bound quarkonium.

1. Introduction: Electroproduction at CEBAF

Quantum Chromodynamics has become the central focus of particle and nuclear physics since it potentially can describe all strong interactions in terms of fundamental quark and gluon degrees of freedom. Although there have been many empirical successes of QCD, crucial elements of the theory remain unexplored. For example, the fundamental structure of hadronic wavefunctions, the nature of confinement, and the structure of the QCD vacuum are understood only qualitatively at best. The study of hadronic wavefunctions has now become even more critical in view of the need to understand exclusive and inclusive charm and bottom hadronic from first principles. An important challenge is to devise tests of the theory and hadron structure at the amplitude level which can be made as systematically precise as possible.

A characteristic scale of quantum chromodynamics is the mean transverse momentum of quarks within hadrons: $\langle k_T^2 \rangle^{1/2} \simeq 200 - 300$ MeV. Hadronic processes involving momentum transfer much larger than this scale $Q^2 \gg \langle k_T^2 \rangle^{1/2}$ can be traced to underlying hard scattering process $T_H(Q)$ involving the minimal number of quark and gluons. Because of asymptotic freedom, $T_H(Q)$ can be computed in perturbation theory as an expansion in powers of $\alpha_s(Q^*)$, where $Q^* = \mathcal{O}(Q)$. This observation is the basis of the QCD factorization theorems, in which color confinement and the non-perturbative effects of hadronic binding are isolated in terms of

the hadron structure functions and fragmentation functions in the case of inclusive reactions, and distribution amplitudes $\phi_H(x,Q)$ in the case of exclusive reactions.[1] Processes which involve the hard scattering of more than the minimum numbers of quarks or gluons are "higher-twist"; *i.e.* they are dynamically suppressed by powers of Λ^2/Q^2.

The extended laboratory electron energy range now being contemplated at CEBAF, up to 10-12 GeV, allows one to probe QCD effects in the transition regime between coherent and incoherent quark subprocesses. In the case of inclusive electroproduction at CEBAF, the dominant subprocess at large momentum transfer corresponds to the electron scattering on one of the quark constituents of the target nucleon or nucleus. Thus the deep inelastic cross section $d\sigma(eN \to e'X)$ is to first approximation given by the convolution of the hard scattering $d\sigma(eq \to e'q')$ cross section, multiplied by target structure functions $G_{q/N}(x,Q)$, the probability distributions describing the spin and flavor distributions of quarks in the proton, neutron, or nuclei at light-cone momentum fraction $x = k^+/p^+ = (k^0 + k^3)/(p^0 + p^3)$. On-shell kinematics allows one to identify x with the Bjorken variable $x_{bj} = Q^2/2p \cdot q$. Since the beam is continuous and the number of produced hadrons is not enormous, one can measure the complete final state in electroproduction at CEBAF, and thus follow the evolution of the produced quarks and gluons from the hard subprocess and the remnant quarks and gluons spectators from the target into final state hadronic systems. Since the momentum transfers is moderate, one can also detect the non-factorizing effects of coherence, such as the interference between subprocesses in which the electron scatters on different quarks in the target.

At very high electroproduction energies, for example in electron-proton collisions at HERA, deep inelastic lepton scattering has only minimal sensitivity to the valence parton structure of the proton. From the perspective of the proton rest frame, the high energy virtual photon fluctuates into virtual $q\bar{q}$ system which then scatters on the gluonic field of the target; *i.e.*, physics associated with photon dissociation and central rapidity region processes. Thus physics at HERA focuses more on the structure of the photon rather than resolving nucleon substructure. In contrast, at CEBAF energies, the dominant electroproduction physics is controlled by the quark structure of the target nucleon or nucleus.

Monte-Carlo and string fragmentation programs are often used to simulate the main features of the final state hadronization in electroproduction; However the goal is to acquire a fundamental understanding of hadronization at the amplitude level. Studies at CEBAF have the potential for studying the basic physical processes involved in the processes in which a confined quark or gluon turns into hadronic matter. In the case of nuclear target, one can resolve the effects of the background nuclear field such as quark energy loss transverse momentum smearing and co-mover interactions on the materialization of the final state.

Although the energy range proposed for an upgraded CEBAF is well-tuned to resolving proton structure in QCD, in general one also needs to take into account coherent effects and multiparticle subprocesses. For example, at moderate momentum

transfers, the electron will often interact with more than one target constituent; *e.g.*, $eqq \to e'q'q'$. Since the two quarks can scatter coherently as a bosonic system, they can produce a large longitudinal cross section $R = \sigma_L/\sigma_T$. The two quarks together carry a large fraction of the target momentum, and thus such higher twist contributions can actually dominate the electroproduction cross section at large $x \sim 1$.

Although higher twist terms may complicate the physical interpretation of electroproduction at CEBAF, they are important and interesting topics in their own right. For example, Brandenburg, Khoze, Müller, and I[2] have recently shown that measurements of the $\cos\phi$ and $\cos 2\phi$ azimuthal angular dependence of the lepton in the Drell-Yan process $HN \to \ell\bar{\ell}X$ with $H = \pi, K, \bar{p}$ and p at forward x_F are sensitive to the shape of the projectile's distribution amplitudes. One can analytically cross these predictions to obtain a theory of meson electroproduction including higher twist contributions where the meson interacts directly within the hard subprocess. In the case of meson electroproduction, ϕ is the azimuthal angle between the lepton scattering plane and the meson production plane. Thus measurements at CEBAF of the full azimuthal and polar angular distribution of the lepton system in meson electroproduction and lepton pair hadroproduction should provide an important measure of the structure of hadrons at the amplitude level.

The study of exclusive reactions at CEBAF such as elastic electron-proton scattering, real and virtual Compton scattering, and meson electroproduction, provides a complimentary measure of nucleon structure to the purely inclusive studies.[3] An analogy is an electron microscope, where information from both elastic and inelastic scattering are combined to generate the image of the target. In exclusive reactions in QCD all of the constituents of the scattered hadrons must be rearranged from the initial to final state. Dimensional counting rules[4] show that the leading subprocesses to order Q^{-1} involve the minimum number of incident and final constituents. Furthermore, the valence quarks exchange their hard momenta when their transverse separations are small: $b_\perp^i = \mathcal{O}(Q^{-1})$. Thus large momentum transfer exclusive reactions are controlled by the hadron distribution amplitudes $\phi_H(x_i, Q)$, the valence light-cone Fock wavefunction at small transverse separation. This leads to the remarkable "color transparency"[5] property of QCD, since such small color singlet fluctuations have only minimal initial and final state interactions as they transit through nuclei. In addition, the study of exclusive pion and kaon electroproduction over a large kinelatic range of energy and momentum transfer is necessary in order to reliably determine the spacelike meson form factors. A more detailed discussion will be presented in Section 4.

In the case of electroproduction at CEBAF one can use a nuclear target as a "color filter" to separate large and small structure events. Recently, the E665 group at Fermilab[6] has reported preliminary results on coherent and incoherent ρ leptoproduction in nuclei. At low momentum transfer, the ρ vis strongly absorbed, as in conventional Glauber theory; however, as Q^2 increases beyond a few GeV^2 the reactions tend to occur uniformly throughout the nuclear volume, as predicted by color transparency. Thus diffractive muo-production of ρ mesons occurs in a nuclear

target without final absorption of the ρ in the nucleus. More generally, by using nuclear targets, one can change the hadronic environment and study not only the shadowing and antishadowing of nuclear structure functions, but also the influence of the nuclear field on evolution of the final state hadronic system, including the induced radiation of the outgoing quarks.[7]It will be clearly interesting to trace the color transparency effects seen in vector meson leptoproduction at Fermilab to the lower CEBAF energy range where the formation times are moderate.

Figure 1. Comparison of photoproduction data with the dimensional counting power-law prediction. The data are summarized in Ref. 8.

A central prediction of perturbative QCD for exclusive electroproduction at large momentum transfer is fixed center-of-mass angle scaling:

$$\frac{d\sigma}{dt}(\gamma^* p \to MB) = \frac{f(t/s, Q^2/s)}{s^N}.$$

The nominal scaling power $N \simeq 7$ follows from dimensional counting: there are 4 incident and 5 outgoing elementary fields. One important test of this scaling is shown in Fig. 1 for pion photoproduction $\gamma p \to \pi^+ n$ at $\theta_{cm} = \pi/2$. The nominal s^{-7} predicted power law behavior is consistent with experiment over the energy range contemplated at CEBAF. This scaling behavior needs to be checked systematically in electroproduction, for example, as a function of virtual photon mass and polarization, and the angular and energy range. The leading power should correspond to helicity amplitudes which conserve the total hadron helicity from the initial and final state, independent of the photon polarization. Hadron helicity conservation is discussed in more detail in Section 4.

In general PQCD dimensional counting has been shown to be good guide to the scaling behavior of general fixed angle two-body scattering reactions.[3] A systematic study of meson-baryon reactions has recently been completed at Brookhaven.[9] The large relative normalization of large angle cross sections such as $K^+p \to K^+p$ compared to $K^-p \to K^-p$ shows that the dominant interaction controlling exclusive processes at large momentum transfer involves the interchange of the valence quarks[10] rather than multiple gluon exchange. Other important tests involve exclusive two-photon reactions such as $\gamma^*\gamma \to M^0$, $\gamma\gamma \to M\bar{M}$, and proton-proton annihilation. The general success of dimensional counting in the fixed angle domain is evidence that leading twist PQCD mechanisms dominate exclusive amplitudes and form factors at momentum transfers $Q^2 \sim 5$ GeV^2.

The structure of hadron wavefunctions in terms of their quark and gluon degrees of freedom at the amplitude level remains one of the most important frontiers in QCD studies. The natural formulation of hadron wavefunctions is the light-cone Fock expansion.[3] As noted above, the basic quantity which characterizes the part of the hadron bound state which enters exclusive hard-scattering subprocesses is the gauge and frame-independent distribution amplitude $\phi(x_i, Q)$ which in turn describes the valence quark structure of the hadrons at impact separation $b_\perp^i = \mathcal{O}(Q^{-1})$ is a function of the light-cone momentum fractions. The work of Sterman *et al.*,[11] has shown in detail how Sudakov suppression of large size configurations of the hadron wavefunctions are suppressed in large momentum transfer exclusive processes, confirming the validity of the PQCD description of these processes and the corresponding predictions of QCD color transparency. The interrelation of the Landshoff triple-gluon contributions to elastic proton-proton scattering to hard scattering PQCD mechanisms such as quark interchange has now been clarified by Sterman and Sotiropoulos.[12]

QCD sum rule methods and lattice gauge theory now supply important theoretical constraints on the form of the distribution amplitudes, although the reliability of these predictions is unknown. There is now much theoretical work exploring other non-perturbative QCD methods such as light-cone Hamiltonian diagonalization. One can also derive constraints on hadron light-cone wavefunctions from their static properties such as the baryon magnetic moments and their axial couplings.

2. The Charm Threshold in Electroproduction

One of the most interesting physics areas which can be studied at CEBAF at electron beam energies above 8 GeV will be the onset of charm electroproduction. The threshold virtual photon energy for the lowest hidden charm system $\gamma^* p \to \eta_c(2.9788\text{GeV})p(0.9383\text{GeV})$ is $\nu_{th} = q \cdot p/M = 7.707$ GeV $+(Q^2/2M_p)$. The production of open charm $\gamma^* p \to D^0(1.8645\text{GeV})\Lambda_c(2.2849\text{GeV})$ begins at $\nu_{th} = 8.7057$ GeV $+(Q^2/2M_p)$. In the threshold regime one probes extreme configurations of the proton target quark structure as it strains to produce the new heavy systems. The study of this physics also has implications for the charm structure function at large x_{bj} and the threshold production of heavy systems such as beauty, top, and supersymmetric particles.

Although the charm electroproduction cross section is inevitably suppressed at threshold by phase space, there is reason to believe that the production rate will be substantially enhanced by dynamical effects in non-perturbative QCD. Because of the disparate charmonium and proton size scales, one can classify and compute their two-gluon exchange couplings using the operator product expansion. Luke, Manohar, and Savage[13] have shown that the scalar part of the two-gluon exchange interaction is related to the trace of the energy momentum tensor and thus its coupling to nucleons or nuclei is proportional to the target mass. The two-gluon coupling to the small size charmonium state can be computed using conventional potential models. This analysis leads to the remarkable prediction that there is a strong QCD van der Waal attraction of the quarkonium state to ordinary hadrons at small relative velocity. DeTeramond, Schmidt, and I[14] have argued that the QCD van der Waal effects at low relative velocity could be sufficiently strong as to bind charmonium states to ordinary hadrons or light nuclei. Such nuclear-bound quarkonium states could show up as narrow $s-$channel resonance in electroproduction: $\gamma^* p \to (\eta_c p)$, $\gamma^* d \to (\eta_c d)$, $\gamma^* p \to (J/\psi p)$, etc. just below the charmonium production threshold. This could be a very interesting experiment CEBAF. However, it should be emphasized that even if the QCD van der Waals force is not sufficient to actually form bound states, one still expects to see strong threshold effects in the production cross section perhaps similar to the enhancements that have been observed for η production at threshold.[15]

Note that even though the rate at threshold may be small, the cross section can be enhanced by using nuclear Fermi-motion to effectively increase the available energy. For example, the anti-deuteron was first observed below the nominal threshold energy in the 1960s at Brookhaven by Ting and Lederman using heavy nuclei as targets. It is also interesting to use the charm threshold to measure the extreme limits of the Fermi momentum spectrum, since its origin involves nuclear short-range interactions.

In the case of open charm, the simplest electroproduction mechanism is quark interchange, as illustrated in Fig. 2. The interchange amplitude can be written in an

Figure 2. Quark interchange contribution to charm electroproduction.

elegant form as a convolution over valence light-cone wavefunctions:

$$\int \frac{d^2k_\perp dx}{16\pi^3}\psi^\dagger_D(x,k_T{+}xr_\perp)\psi^\dagger_{\Lambda_c}(x,k_T{+}(1{-}x)q_\perp)\Delta\psi_{\gamma^*}(x,k_T{+}xr_\perp{+}(1{-}x)q_\perp)\psi_p(x,k_T),$$

where $q_\perp^2 = -t, r_\perp^2 = -u$ and Δ is the inverse of the light-cone energy denominator. The complete analysis is given in Ref. 10.

There are a number of experiments which indicate that non-perturbative QCD mechanisms are necessary for understanding heavy quark production in the regimes where either the charm system is produced at extreme kinematic configurations such as large x_F or large x_{bj} or at small relative velocity to other quarks:

1. The anomalously high $c(x)$ distribution measured at large x_{bj} by EMC.[16] The CERN measurements disagree with photon-gluon fusion by a factor of 20 to 30 at $Q^2 = 75$ GeV^2 and $x_{bj} = 75$ GeV^2.
2. In the case of J/ψ hadroproduction from pion beams, the CERN experiment NA-3 has reported a strong excess of quarkonium at large x_F with a non-factorizing nuclear dependence. In addition, the Fermilab Chicago-Iowa Princeton group has reported an anomalously sudden change in polarization of the J/ψ at large x_F in $\pi N \to \mu^+\mu^- X$. The dramatic shift to longitudinal polarization is inconsistent with leading order QCD predictions.
3. An interesting unresolved issue is the leading particle effect in charmed hadron production. The quark structure of leading D mesons has been shown to depend strongly on the valence quantum numbers of the beam hadron in direct conflict with the factorization principle at the heart of most perturbative QCD predictions. The mechanisms in which the beam quarks and heavy quarks coalesce is at the heart of hadronization dynamics, and much more critical work will be needed especially in the production of b-quark systems.

4. There are other signals for anomalous charm baryon hadroproduction at large x_F, including the reports of Λ_c production from E-400 at Fermilab using neutron beams and the measurements of WA-62 from CERN which observed charm-strange baryons using hyperon beams. There are also measurements from NA-3 at CERN which show that double J/ψ pairs are hadro-produced only at large x_F.
5. The anomalously strong nuclear dependence of large x_F J/ψ hadroproduction, as reported by NA10 and E789 are in direct contradiction to leading-twist PQCD factorization.

Much of the above physics can be accounted for by the picture of Hoyer, Mueller, Tang and myself,[17] where the hadronization in a high energy collision occurs in the following novel way: the heavy quark system is first formed as a virtual fluctuation as a light-cone Fock state component in the incoming hadron wavefunction; a light spectator quark is then stripped away in the target leaving the $Q\bar{Q}$ system to hadronize into the final heavy hadrons. This type of intrinsic heavy quark picture also explains the excess of charm quarks seen in the EMC measurements of the charm structure function of the nucleon. This new picture of hadron formation opens up a whole new avenue for studying the far-off-shell structure of hadrons. It is thus critical that a new measurement of the charm and beauty structure functions be performed.

Measurements of charm electroproduction near threshold at CEBAF should provide new insights into the collective multi-quark mechanisms needed to understand the charm production anomalies.

3. Virtual Compton Scattering

The Compton scattering process $\gamma p \to \gamma p$ is the fundamental way to "look" at proton structure. Virtual Compton scattering $\gamma^* p \to \gamma p$ is particularly interesting to measure at CEBAF since it can be probed as a function of the photon's transverse or longitudinal polarization, the target polarization, over a large domain of kinematics s, t, u, and photon virtuality $Q^2 = -q^2$.

It should be noted that the cross section for the process $ep \to e'p\gamma$ receives contributions not only from virtual Compton scattering, but also from Bethe-Heitler bremmstrahlung from the scattered electron. The two processes lead to the same final state, and thus they interfere. The Bethe-Heitler process is completely determined from elastic ep scattering and is purely real. Thus one can use the interference between the Compton and bremmstrahlung processes to determine the real part of the Compton amplitude.[18] In the case of deep inelastic Compton scattering $ep \to e\gamma X$, one can use the same interference effect to deduce new structure functions and sum rules proportional to the sum of quark charges cubed.[19]

There are many different physics aspects of virtual Compton scattering depending on the accessed kinematical domain.

1. In the case of low energy virtual Compton scattering with $s = (q+p)^2 \simeq M_{N^*}^2$ one can study the s-channel effects of baryon resonances in the Compton amplitude and their relative coupling as a function of photon virtuality.[20]
2. In the Regge limit $s \gg -t$ and fixed Q^2 one can use the Regge pole analysis, as in the paper of Damashek and Gilman.[21] Each Compton helicity amplitude has the form of a sum over $t-$channel Regge exchange contributions: $\mathcal{M} = \Sigma_R s^{\alpha_R(t)} \beta(t, Q^2)$. In Compton scattering one can have contribution from all C-even exchanges: the diffractive Pomeron contributions,the pion and A_2 and f^0 $C = +$ Reggeon trajectories. In addition, QCD predicts a special contribution which cannot occur in hadron-hadron scattering: a $j = 0$ fixed pole, the Kronecker $\delta_{j,0}$ contribution, which can be traced to the presence of quark Z-graphs.
3. An important feature of the Regge theory which is testable in virtual Compton scattering is that the Regge trajectory $\alpha_R(t)$ must be independent of Q^2 at fixed $-t$. Only the residue $\beta(t, Q^2)$ can depend on the photon virtuality. In fact, one expects that all of the normal trajectories have decreased couplings to the virtual; Compton amplitude as Q^2 increases, leaving the $j = 0$ fixed pole as the dominant and surviving contribution to the amplitude. This special contribution to Compton scattering gives an energy independent contribution to the real part of the $\gamma^* p \to \gamma p$ amplitude.[18] The $t-$dependence of the $j = 0$ fixed pole amplitude is expected to be similar to that of the helicity-conserving Dirac form factor of the proton.
4. As the momentum transfer squared to the proton increases $-t$, the Pomeron trajectory is expected to stay at $\alpha_p(-t) \simeq 1$. The non-singlet Regge trajectories are predicted to decrease monotonically to $\alpha_R(-t) \simeq 0$. For a recent discussion and further references see Ref. 22.
5. In the large momentum transfer domain with fixed $\cos\theta_{cm}$ the virtual exclusive Compton amplitude $\gamma^* p \to \gamma p$ can be analysed using perturbative QCD factorization. Detailed QCD predictions have been made by Kronfeld and Nizic, Hyer, and Gunion *et al.*. This will be discussed in detail in Section 5. In addition at CEBAF, one may be able to test these predictions as a function of photon polarization and virtuality.

4. Exclusive Processes and the Structure of Hadrons[23]

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[24]

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

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

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.[25] 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.[26,27]

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.[26] For example, the pion's electromagnetic form factor can be written as[26,27,28]

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

Here T_H is the scattering amplitude obtained when pions replaced by collinear $q\bar{q}$ 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,[26] $\phi_B(x_i, \lambda_i, Q)$, 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:

$$\phi_\pi(x,Q) = \int \frac{d^2\vec{k}_\perp}{16\pi^3}\,\psi^{(Q)}_{q\bar{q}/\pi}(x,\vec{k}_\perp)$$

$$= P^+_\pi \int \frac{dz^-}{4\pi}\, e^{ixP^+_\pi z^-/2} \left\langle 0 \left| \bar{\psi}(0)\,\frac{\gamma^+\gamma_5}{2\sqrt{2n_c}}\,\psi(z) \right| \pi \right\rangle^{(Q)} \Bigg|_{z^+ = \vec{z}_\perp = 0}.$$

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.[3,29] If one can calculate the distribution amplitude at an initial scale Q_0 using QCD sum rules or lattice gauge theory,[29] 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.[30] 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",[5] 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.[31]

5. 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 $\gamma\gamma \to H\bar{H}$ are sensitive to the x_i dependence of the hadron distribution amplitudes.[32] Lowest order predictions for meson pair production in two photon collisions using this formalism are given in Refs. 32 and 29; 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.[33] 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. 34.

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

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

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 methods of Ref. 76and 77 as discussed in the Introduction. 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.[35,36,37,38]

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.[37] 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.[39] A careful analytic treatment of the integration over the on-shell intermediate states has been given by Kronfeld and Nizic.[37]

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

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

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[40] 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[29] by Chernyak, Oglobin, and Zhitnitskii, and King and Sachrajda, appear to give a reasonable representation of the existing data. A definitive prediction for the normalization of form factors and other exclusive amplitudes in perturbative QCD will require not only a careful analysis of the non-perturbative input for the distribution amplitudes, but also a detailed calculation of the crossed-graph and other irreducible contributions to the hard-scattering QCD kernels.

More recent QCD sum rule analyses of the proton distribution amplitude are given in Ref. 41. 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 $J/\psi \to p\bar{p}$ 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.*,[35] Millers and Gunion[36], and Hyer.[38]

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 Compton to elastic differential cross sections

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

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

6. 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.[7,42,43]

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 $\gamma^* A \to V^0 A$ 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 $d\sigma/dt \sim \exp^{b_A t}$ where $b_A \propto R_A^2$ and R_A is the nuclear size. Thus the total coherent cross section $\sigma(\gamma^* A \to V^0 A)$ is predicted to scale with nuclear number as $A^{4/3}/R_A^2 \sim A^{2/3}$.
3. The predictions for $\gamma^* A \to V^0 A'$ 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[6] 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.[44] 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 leptoproduction 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. As I have emphasized in the Introduction, these illuminating studies and tests of PQCD can be carried out in detail at CEBAF.

7. 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. [45] 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[46] and Radyushkin[47] 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 $\sum_i(\epsilon \cdot p_i/k \cdot p_i)\eta_i q_i$ 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. 26 and 48. An elegant and much more complete discussion has now been given by Botts and Li and Sterman.[11] The radiation from the colored lines in QCD have similar coherence properties as in QED:[49] 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.[11,38] 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 $b_\perp = \mathcal{O}(x_iQ)^{-1}$, 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.[26,48,50] 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.[11] This analysis has now also been applied to two-photon reactions and the timelike proton form factor by Hyer.[38]

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

$$T_H \propto \frac{\alpha_s}{(1-x)(1-y)Q^2 + (k_1^\perp + k_2^\perp)^2}.$$

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[11] 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.[5]

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 resolved, such as how to extend these calculations to baryon processes, how to set the renormalization scale in α_s,[76,77] how to implement conformal symmetry and its breaking,[30,51] and how to formulate and solve the evolution equations for the hadron distribution amplitudes to next-to-leading order.

An important question for evaluating exclusive amplitudes in the transition region between hard and soft QCD processes is how to analytically separate perturbative contributions from contributions intrinsic to the bound-state wavefunction itself. The physical amplitude of course must be independent of the choice separation scale μ. Recently Ji, Pang, and Szczepaniak have observed that the natural variable to make this separation is the light-cone energy or equivalently the invariant mass of the off-shell partonic system, rather than gluon virtuality of T_H. The PQCD contributions from the invariant mass regime $\mu > 1$ GeV can then account substantially for the empirical pion form factor at $Q^2 > 1$ GeV^2. One also expects significant contributions from PQCD from higher order contributions.

It should be emphasized that the measurements of the pion form factor from electroproduction at large Q^2 are quite uncertain since they requires extrapolation to the pion $t-$ channel pole. CEBAF measurements can thus contribute significantly to this fundamental hadronic measure.

One of the most significant problems in computing the normalization of perturbative QCD predictions for exclusive processes is the uncertainty in setting the renormalization scale μ of the QCD coupling $\alpha_s(\mu)$ in the hard scattering amplitude T_H. A related problem is the question of the corresponding scale to use in evaluating the hadron distribution amplitudes.

Given a renormalization scheme, the QCD Lagrangian $\mathcal{L}_{QCD}$ is a function of the bare parameters $\alpha_s(\mu), m_q(\mu)$, etc. In principle, the values of the bare parameters can be fixed given a set of input measurements. Thus given a finite number of empirical values, all other QCD observables should be computable order by order in perturbation theory. The relation between the input and output observables must be independent of the choice of the renormalization scale μ as well as the choice of intermediate renormalization scheme. This invariance of the predictions for observables under changes of the intermediate renormalization scheme constitutes the generalized renormalization group invariance of Peterman and Stückelberg.[79]

Recently, Hung Jung Lu and I[77] have shown how this problem can be avoided by directly relating observables through commensurate scale relations. The conventional $\overline{M}S$ scheme serves as an intermediary calculational tool, but it can be systematically eliminated when relating observables. For example, the entire radiative corrections to the annihilation cross section is expressed as the effective charge $\alpha_R(Q)$ where $Q = \sqrt{s}$:

$$R(Q) \equiv 3 \sum_f Q_f^2 \left[1 + \frac{\alpha_R(Q)}{\pi}\right].$$

Similarly, we can define the entire radiative correction to the Bjorken sum rule as the effective charge $\alpha_{g_1}(Q)$ where Q is the lepton momentum transfer:

$$\int_0^1 dx \left[g_1^{ep}(x,Q^2) - g_1^{en}(x,Q^2)\right] \equiv \frac{1}{6}\left[\frac{g_A}{g_V}\right]\left[1 - \frac{\alpha_{g_1}(Q)}{\pi}\right].$$

We now use the known expressions to three loops in $\overline{M}S$ scheme and choose the scales Q^* and Q^{**} to re-sum all quark and gluon vacuum polarization corrections into the running couplings.[76] The relative scales insure that each observable pass through the heavy quark thresholds at their commensurate physical scales. The final result is remarkably simple:

$$\frac{\alpha_{g_1}(Q)}{\pi} = \frac{\alpha_R(Q^*)}{\pi} - \left(\frac{\alpha_R(Q^{**})}{\pi}\right)^2 + \left(\frac{\alpha_R(Q^{***})}{\pi}\right)^3 + \cdots$$

The fundamental test of QCD is then to verify empirically that the observables track in both normalization and shape as given by these relations. The coefficients in the series (aside for a factor of C_F, which can be absorbed in the definition of α_s) are actually independent of color and are the same in abelian, non-abelian, and conformal gauge theory. The non-Abelian structure of the theory is reflected in the scales Q^* and Q^{**}. The commensurate scale relations thus provide fundamental tests of QCD which can be made increasingly precise and independent of any scheme or other theoretical convention.

In the case of exclusive processes, the coupling associated with each virtual gluon exchange carrying momentum transfer ℓ_i^μ in the hard-scattering subprocess tree amplitude T_H can be identified with the running coupling $\alpha_V(\ell_i^2)$ appearing in the heavy quark potential.[80] We can determine the numerical values for $\alpha_V(Q^2)$ in many ways: directly from the heavy quarkonium spectrum and heavy quark lattice gauge theory[81] or from the commensurate scale relations which connect the α_V scheme to $\alpha_{\overline{MS}}$, or effective charges such as α_R, α_{g_1}, the Gross-Llewellyn Smith sum rule, etc. at their appropriate commensurate scales. Note that higher order corrections to the hard scattering amplitude from crossed graph kernels contribute even if the theory were conformal invariant; *i.e.* even of the coupling did not run. A related method can be used to choose the separation scale which controls the evolution of the hadron distribution amplitudes. By using this procedure, one should be able to substantially reduce the uncertainty in form factors and other exclusive processes from renormalization scale and scheme ambiguities.

8. 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. 52.

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 multiquark higher twist contributions can be computed using the exclusive factorization analysis, and they contribute at the same order as the leading twist terms.[53,17] 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 $x \sim 1$.[54]

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",[55] *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.

9. Exclusive Nuclear Processes

An ultimate goal of QCD phenomenology is to describe the nuclear force and the structure of nuclei in terms of quark and gluon degrees of freedom.

One of the most elegant areas of application of QCD to nuclear physics is the domain of large momentum transfer exclusive nuclear processes. Rigorous results have been given by Lepage, Ji and myself[56] for the asymptotic properties of the deuteron form factor at large momentum transfer. In the asymptotic $Q^2 \to \infty$ limit the deuteron distribution amplitude, which controls large momentum transfer deuteron reactions, becomes fully symmetric among the five possible color-singlet combinations of the six quarks. One can also study the evolution of the "hidden color" components (orthogonal to the np and $\Delta\Delta$ degrees of freedom) from intermediate to large momentum transfer scales; the results also give constraints on the nature of the nuclear force at short distances in QCD. The existence of hidden color degrees of freedom further illustrates the complexity of nuclear systems in QCD. It is conceivable that six-quark d^* resonances corresponds to these new degrees of freedom may be found by careful searches of the $\gamma^* d \to \gamma d$ and $\gamma^* d \to \pi d$ channels.

The basic scaling law for the helicity-conserving deuteron form factor is $F_d(Q^2) \sim 1/Q^{10}$ which comes from simple quark counting rules, as well as perturbative QCD. One cannot expect this asymptotic prediction to become accurate until very large Q^2 is reached since the momentum transfer has to be shared by at least six constituents. However, one can identify the QCD physics due to the compositeness of the nucleus, with respect to its nucleon degrees of freedom by using the reduced amplitude formalism.[57] For example, consider the deuteron form factor in QCD. By definition this quantity is the probability amplitude for the deuteron to scatter from p to $p+q$ but remain intact. Note that for vanishing nuclear binding energy $\epsilon_d \to 0$, the deuteron can be regarded as two nucleons sharing the deuteron four-momentum (see Fig. 3). The momentum ℓ is limited by the binding and can thus be neglected. To first approximation the proton and neutron share the deuteron's momentum equally. Since the deuteron form factor contains the probability amplitudes for the proton and neutron to scatter from $p/2$ to $p/2+q/2$; it is natural to define the reduced deuteron form factor[57,58,59]

$$f_d(Q^2) \equiv \frac{F_d(Q^2)}{F_{1N}\left(\frac{Q^2}{4}\right)\, F_{1N}\left(\frac{Q^2}{4}\right)}.$$

The effect of nucleon compositeness is removed from the reduced form factor. QCD then predicts the scaling

$$f_d(Q^2) \sim \frac{1}{Q^2}$$

i.e. the same scaling law as a meson form factor. Diagrammatically, the extra power of $1/Q^2$ comes from the propagator of the struck quark line, the one propagator not contained in the nucleon form factors. Because of hadron helicity conservation, the prediction is for the leading helicity-conserving deuteron form factor ($\lambda = \lambda' = 0$.) As shown in Fig. 4, this scaling is consistent with experiment for $Q = p_T \gtrsim 1$ GeV.

Figure 3. Application of the reduced amplitude formalism to the deuteron form factor at large momentum transfer.

The distinction between the QCD and other treatments of nuclear amplitudes is particularly clear in the reaction $\gamma d \to np$; *i.e.* photo-disintegration of the deuteron at fixed center of mass angle. Using dimensional counting, the leading power-law prediction from QCD is simply $\frac{d\sigma}{dt}(\gamma d \to np) \sim F(\theta_{\rm cm})/s^{11}$. Again we note that the virtual momenta are partitioned among many quarks and gluons, so that finite mass corrections will be significant at low to medium energies. Nevertheless, one can test the basic QCD dynamics in these reactions taking into account much of the finite-mass, higher-twist corrections by using the "reduced amplitude" formalism.[58,59] Thus the photo-disintegration amplitude contains the probability amplitude (*i.e.* nucleon form factors) for the proton and neutron to each remain intact after absorbing momentum transfers $p_p - 1/2p_d$ and $p_n - 1/2p_d$, respectively (see Fig. 5). After the form factors are removed, the remaining "reduced" amplitude should scale as $F(\theta_{\rm cm})/p_T$. The single inverse power of transverse momentum p_T is the slowest conceivable in any theory, but it is the unique power predicted by PQCD.

Figure 4. Scaling of the deuteron reduced form factor. The data are summarized in Ref. 58.

Figure 5. Construction of the reduced nuclear amplitude for two-body inelastic deuteron reactions.[58]

The prediction that $f(\theta_{\rm cm})$ is energy dependent at high-momentum transfer is compared with experiment in Fig. 6. It is particularly striking to see the QCD prediction verified at incident photon lab energies as low as 1 GeV. A comparison with a standard nuclear physics model with exchange currents is also shown for comparison as the solid curve in Fig. 6(a). The fact that this prediction falls less fast than the data suggests that meson and nucleon compositeness are not taken to into account

correctly. An extension of these data to other angles and higher energy would clearly be very valuable.

Figure 6. Comparison of deuteron photodisintegration data with the scaling prediction which requires $f^2(\theta_{cm})$ to be at most logarithmically dependent on energy at large momentum transfer. The data in (a) are from the recent experiment of Ref. 60. The nuclear physics prediction shown in (a) is from Ref. 61. The data in (b) are from Ref. 62.

The derivation of the evolution equation for the deuteron and other multiquark states is given in Refs. 56 and 59. In the case of the deuteron, the evolution equation couples five different color singlet states composed of the six quarks. The leading anomalous dimension for the deuteron distribution amplitude and the helicity-conserving deuteron form factor at asymptotic Q^2 is given in Ref. 56.

There are a number of related tests of QCD and reduced amplitudes which require $\bar{p}$ beams[59] such as $\bar{p}d \to \gamma n$ and $\bar{p}d \to \pi p$ in the fixed $\theta_{\rm cm}$ region. These reactions are particularly interesting tests of QCD in nuclei. Dimensional counting rules predict the asymptotic behavior $\frac{d\sigma}{dt}\,(\bar{p}d \to \pi p) \sim \frac{1}{(p_T^2)^{12}}\, f(\theta_{\rm cm})$ since there are 14 initial and final quanta involved. Again one notes that the $\bar{p}d \to \pi p$ amplitude contains a factor representing the probability amplitude (*i.e.* form factor) for the proton to remain intact after absorbing momentum transfer squared $\widehat{t} = (p - 1/2p_d)^2$ and the $\bar{N}N$ time-like form factor at $\widehat{s} = (\bar{p} + 1/2p_d)^2$. Thus $\mathcal{M}_{\bar{p}d\to\pi p} \sim F_{1N}(\widehat{t})\; F_{1N}(\widehat{s})\, \mathcal{M}_r$,

where $\mathcal{M}_r$ has the same QCD scaling properties as quark meson scattering. One thus predicts

$$\frac{\frac{d\sigma}{d\Omega}\ (\bar{p}d \to \pi p)}{F_{1N}^2(\widehat{t})\, F_{1N}^2(\widehat{s})} \sim \frac{f(\Omega)}{p_T^2}\ .$$

The reduced amplitude scaling for $\gamma d \to pn$ at large angles and $p_T \gtrsim 1$ GeV (see Fig. 6). One thus expects similar precocious scaling behavior to hold for $\bar{p}d \to \pi p$ and other $\bar{p}d$ exclusive reduced amplitudes. An analysis by Kondratyuk and Sapozhnikov[63] shows that standard nuclear physics wavefunctions and interactions cannot explain the magnitude of the data for two-body anti-proton annihilation reactions such as $\bar{p}d \to \pi p$.

10. 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[64] 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[65] are based on the interference of Landshoff pinch singularities[66] with the quark interchange amplitude, but there is no understanding why the Landshoff contribution would itself have a large A_{NN}[67] or sufficient normalization[68] to explain this phenomena. Guy de Teramond and I have proposed[69] 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 in pp collisions at a rate of order of 1 microbarn. A similar enhancement of A_{NN} is seen at the open strangeness threshold regime. and is consistent with the 1 millibarn cross section observed for the production of strange hadrons just above threshold. The heavy quark explanation has received some support from the work of Luke, Savage, and Manohar,[13] 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.[14]

Why does QCD color transparency appear to break down in quasielastic pp scattering. The Brookhaven measurements[70] 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 NE-18[71] 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 $\frac{d\sigma dQ^2(eA\to e'p(A-1))}{Zd\sigma dQ^2(ep\to e'p)}$. 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[44] 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.[72] 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.[73] 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.[73] 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 $-t$.[74] 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.[75] Quark interchange[10] 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[9] 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.[10] 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.

The short-distance structure of hadrons, hadron dynamics, and hadronization is thus one of the frontier areas of study in testing quantum chromodynamics. Electroproduction at CEBAF will play an essential role in resolving this fundamental area of physics.

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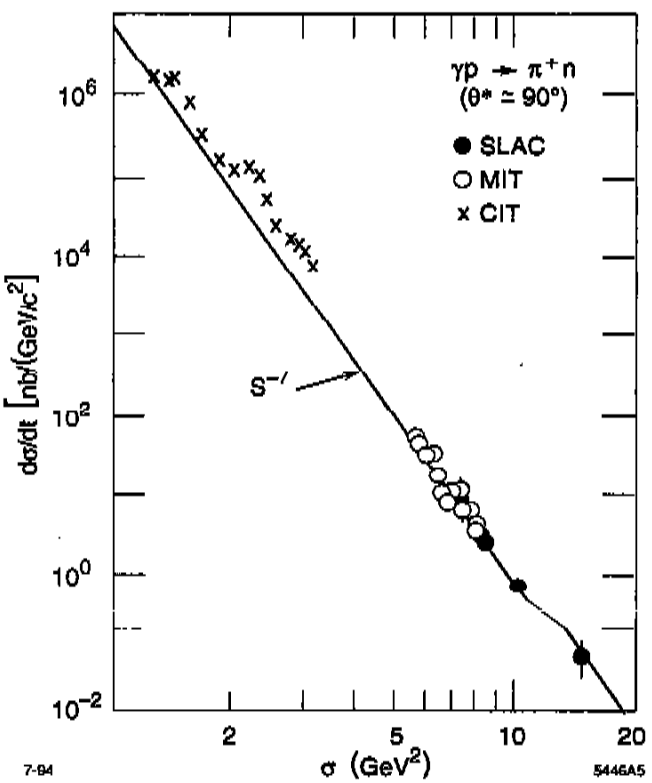
γp → π⁺n
(θ* ≃ 90°)
SLAC
MIT
CIT
S^-7
dσ/dt [nb/(GeV/c)²]
10⁶
10⁴
10²
10⁰
10⁻²
2
5
10
20
σ (GeV²)
7-94
5446A5

Figure 1

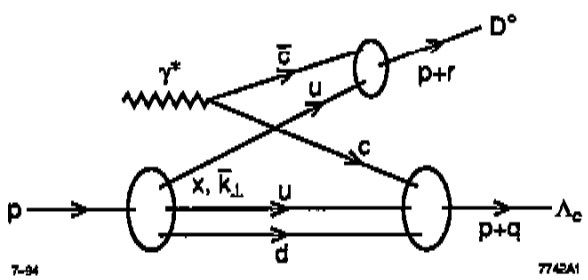

Figure 2

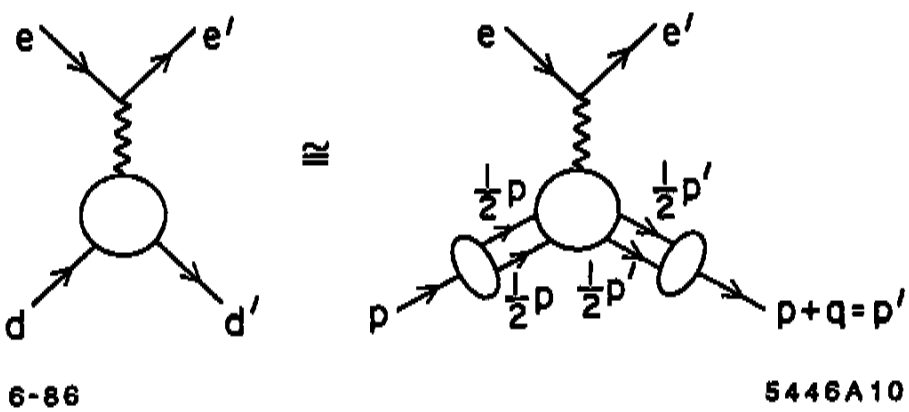

Figure 3

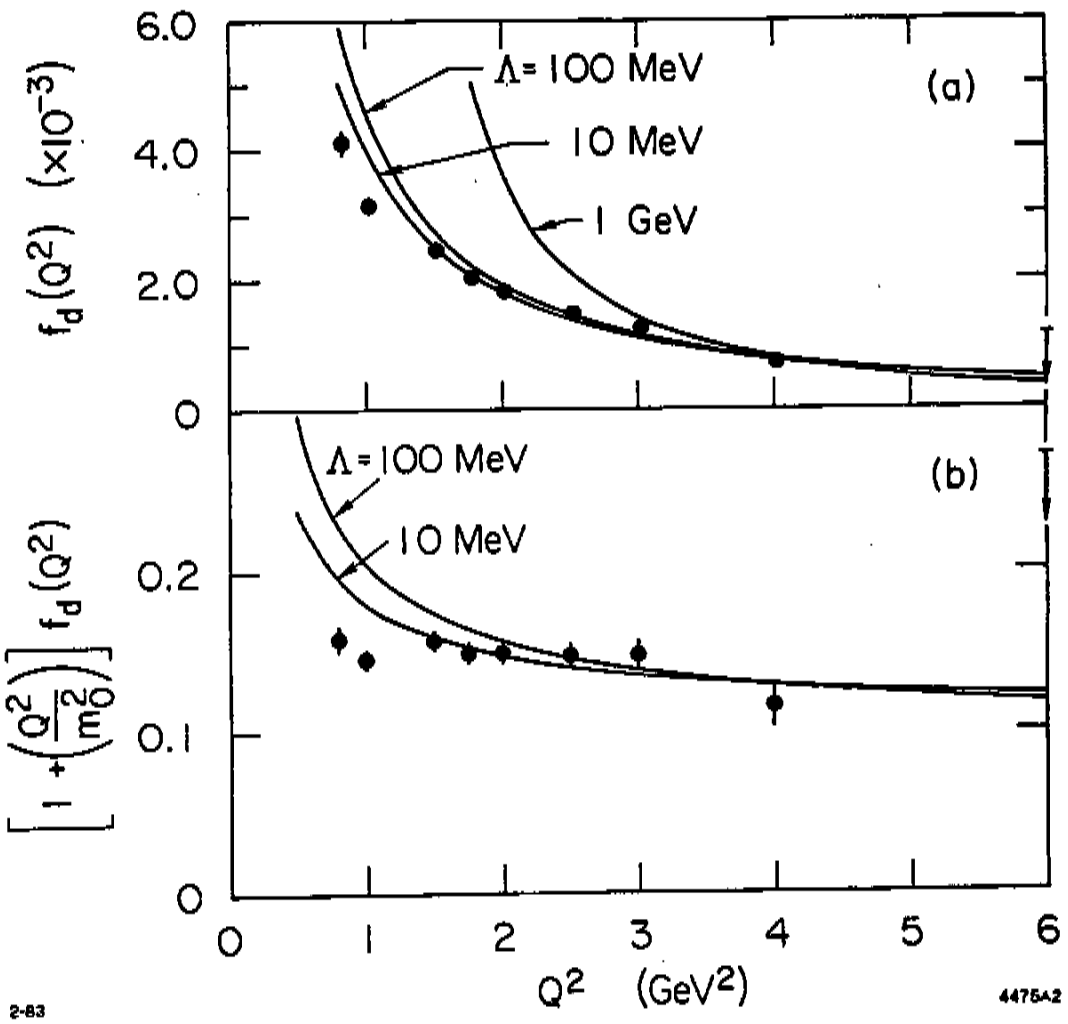
6.0
4.0
2.0
0
$f_d(Q^2)$ $(\times 10^{-3})$
Λ = 100 MeV
10 MeV
1 GeV
(a)
Λ = 100 MeV
10 MeV
(b)
0.2
0.1
0
$\left[1+\left(\frac{Q^2}{m_0^2}\right)\right] f_d(Q^2)$
0 1 2 3 4 5 6
Q^2 (GeV^2)
2-83
4475A2

Figure 4

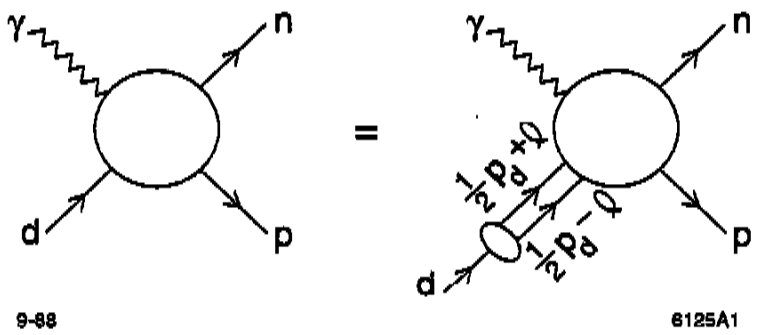

Figure 5

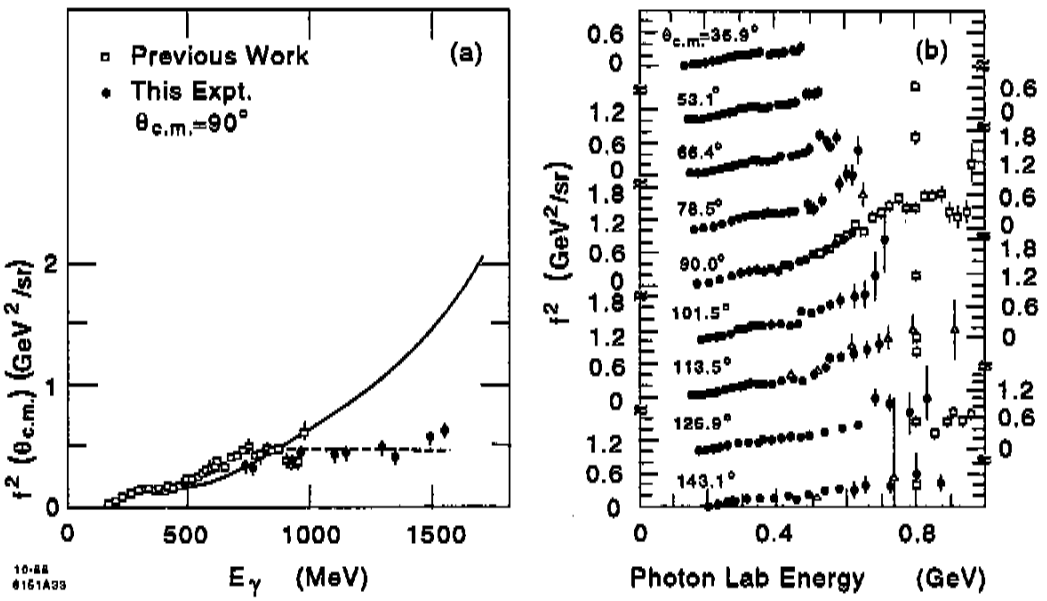
Previous Work
This Expt.
θc.m.=90°
(a)
f² (θc.m.) (GeV²/sr)
Eγ (MeV)
(b)
θc.m.=35.9°
53.1°
66.4°
78.5°
90.0°
101.5°
113.5°
125.9°
143.1°
f² (GeV²/sr)
Photon Lab Energy (GeV)

Figure 6