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Hard Exclusive and Diffractive Processes in QCD*

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Abstract

Exclusive and semi-exclusive processes, the diffractive dissociation of hadrons into jets, and hard diffractive processes such as vector meson leptoproduction provide new testing grounds for QCD and essential information on the structure of light-cone wavefunctions of hadrons, particularly the pion distribution amplitude. I review the basic features of the leading-twist QCD predictions and the problems and challenges of studying QCD at the amplitude level. The application of the light-cone formalism to the exclusive semi-leptonic decay of heavy hadrons is also discussed.

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1 Introduction

Exclusive hard-scattering reactions and hard diffractive reactions are now providing an invaluable window into the structure and dynamics of hadronic amplitudes. Recent measurements of the photon-to-pion transition form factor at CLEO,[1] the diffractive dissociation of pions into jets at Fermilab,[2] diffractive vector meson leptoproduction at Fermilab and HERA, and the new program of experiments on exclusive proton and deuteron processes at Jefferson Laboratory are now yielding fundamental information on hadronic wavefunctions, particularly the distribution amplitude of mesons. Such information is also critical for interpreting exclusive heavy hadron decays and the matrix elements and amplitudes entering CP-violating processes at the B factories.

The natural formalism for describing the hadronic wavefunctions which enter exclusive and diffractive amplitudes is the light-cone Fock representation obtained by quantizing the theory at fixed at fixed "light-cone" time $\tau = t + z/c$.[3] This representation is the extension of Schrödinger many-body theory to the relativistic domain. In a quantum theory a bound state cannot have a fixed number of constituents. For example, the proton state has the Fock expansion

$$\begin{aligned} |\, p\rangle &= \sum_n \langle n \,|\, p\rangle \;|\, n\rangle \\ &= \psi^{(\Lambda)}_{3q/p}(x_i, \vec{k}_{\perp i}, \lambda_i) \,|\, uud\rangle \\ &\quad + \psi^{(\Lambda)}_{3qg/p}(x_i, \vec{k}_{\perp i}, \lambda_i) \,|\, uudg\rangle + \cdots \end{aligned} \tag{1}$$

representing the expansion of the exact QCD eigenstate on a non-interacting quark and gluon basis. The probability amplitude for each such n-particle state of on-mass shell quarks and gluons in a hadron is given by a light-cone Fock state wavefunction $\psi_{n/H}(x_i, \vec{k}_{\perp i}, \lambda_i)$, where the constituents have longitudinal light-cone momentum fractions $x_i = k_i^+/p^+ = (k_i^0 + k_i^z)/(p^0 + p^z)$, $\sum_{i=1}^n x_i = 1$, relative transverse momentum $\vec{k}_{\perp i}$, $\sum_{i=1}^n \vec{k}_{\perp i} = \vec{0}_\perp$, and helicities λ_i. The effective lifetime of each configuration in the laboratory frame is $2P_{lab}/(\mathcal{M}_n^2 - M_p^2)$ where $\mathcal{M}_n^2 = \sum_{i=1}^n (k_{\perp i}^2 + m_i^2)/x_i < \Lambda^2$ is the off-shell invariant mass and Λ is a global ultraviolet regulator. A crucial feature of the light-cone formalism is the fact that the form of the $\psi^{(\Lambda)}_{n/H}(x_i, \vec{k}_{\perp i}, \lambda_i)$ is invariant under longitudinal boosts; *i.e.*, the light-cone wavefunctions expressed in the relative coordinates x_i and $k_{\perp i}$ are independent of the total momentum P^+, $\vec{P}_\perp$ of the hadron. The ensemble $\{\psi_{n/H}\}$ of such light-cone Fock wavefunctions is a key

concept for hadronic physics, providing a conceptual basis for representing physical hadrons (and also nuclei) in terms of their fundamental quark and gluon degrees of freedom. Given the $\psi^{(\Lambda)}_{n/H}$, we can construct any spacelike electromagnetic or electroweak form factor from the diagonal overlap of the LC wavefunctions.[4] Similarly, the matrix elements of the currents that define quark and gluon structure functions can be computed from the integrated squares of the LC wavefunctions.[5]

There has been much progress analyzing exclusive and diffractive reactions at large momentum transfer from first principles in QCD. Rigorous statements can be made on the basis of asymptotic freedom and factorization theorems which separate the underlying hard quark and gluon subprocess amplitude from the nonperturbative physics incorporated into the process-independent hadron distribution amplitudes $\phi_H(x_i, Q)$,[5] the valence light-cone wavefunctions integrated over $k_\perp^2 < Q^2$. An important new application is the recent analysis of hard exclusive B decays by Beneke, *et al.*[6] Key features of such analyses are: (a) evolution equations for distribution amplitudes which incorporate the operator product expansion, renormalization group invariance, and conformal symmetry; [5, 7, 8, 9, 10, 11] (b) hadron helicity conservation which follows from the underlying chiral structure of QCD;[12] (c) color transparency, which eliminates corrections to hard exclusive amplitudes from initial and final state interactions at leading power and reflects the underlying gauge theoretic basis for the strong interactions;[13, 14] and (d) hidden color degrees of freedom in nuclear wavefunctions, which reflects the color structure of hadron and nuclear wavefunctions.[15] There have also been recent advances eliminating renormalization scale ambiguities in hard-scattering amplitudes via commensurate scale relations[16, 17, 18] which connect the couplings entering exclusive amplitudes to the α_V coupling which controls the QCD heavy quark potential.[19] The postulate that the QCD coupling has an infrared fixed-point provides an understanding of the applicability of conformal scaling and dimensional counting rules to physical QCD processes.[20, 21, 19] The field of analyzable exclusive processes has recently been expanded to a new range of QCD processes, such as electroweak decay amplitudes, highly virtual diffractive processes such as $\gamma^* p \to \rho p$,[22, 23] and semi-exclusive processes such as $\gamma^* p \to \pi^+ X$ [24, 25, 26] where the π^+ is produced in isolation at large p_T.

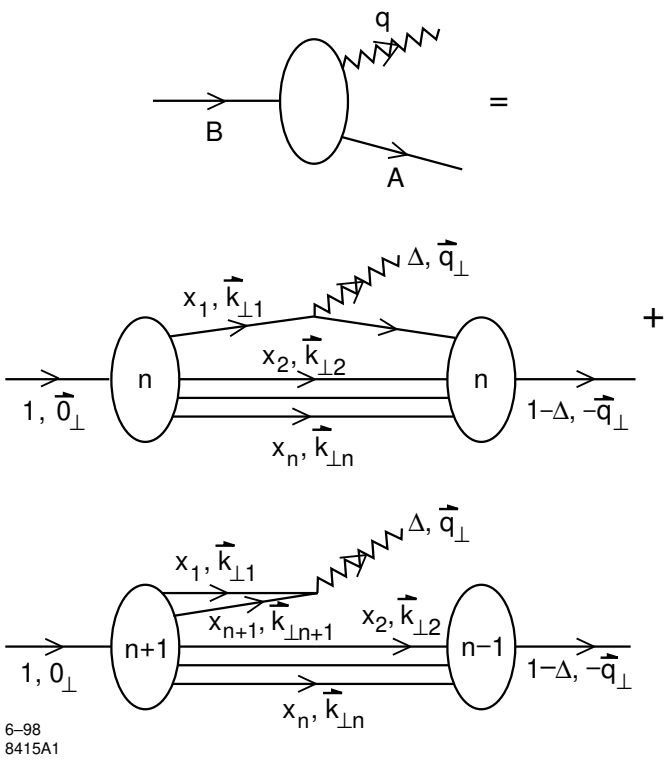

Figure 1: Exact representation of electroweak decays and time-like form factors in the light-cone Fock representation.

2 Electoweak Decays and the Light-Cone Fock Expansion

Dae Sung Hwang and I have recently shown how exclusive semi-leptonic B-decay amplitudes, such as $B \to A\ell\overline{\nu}$ can be evaluated exactly in the light-cone formalism.[27] These timelike decay matrix elements require the computation of the diagonal matrix element $n \to n$ where parton number is conserved, and the off-diagonal $n+1 \to n-1$ convolution where the current operator annihilates a $q\overline{q'}$ pair in the initial B wavefunction. See Fig. 1. This term is a consequence of the fact that the time-like decay $q^2 = (p_\ell + p_{\overline{\nu}})^2 > 0$ requires a positive light-cone momentum fraction $q^+ > 0$. Conversely for space-like currents, one can choose $q^+ = 0$, as in the Drell-Yan-West representation of the space-like electromagnetic form factors.[28, 4, 29] However, the

off-diagonal convolution can yield a nonzero q^+/q^+ limiting form as $q^+ \to 0$. This extra term appears specifically in the case of "bad" currents such as J^- in which the coupling to $q\overline{q}$ fluctuations in the light-cone wavefunctions are favored. In effect, the $q^+ \to 0$ limit generates $\delta(x)$ contributions as residues of the $n+1 \to n-1$ contributions. The necessity for this zero mode $\delta(x)$ terms has been noted by Chang, Root and Yan,[30] and Burkardt.[31]

The off-diagonal $n+1 \to n-1$ contributions provide a new perspective for the physics of B-decays. A semi-leptonic decay involves not only matrix elements where a quark changes flavor, but also a contribution where the leptonic pair is created from the annihilation of a $q\overline{q'}$ pair within the Fock states of the initial B wavefunction. The semi-leptonic decay thus can occur from the annihilation of a nonvalence quark-antiquark pair in the initial hadron. This feature will carry over to exclusive hadronic B-decays, such as $B^0 \to \pi^- D^+$. In this case the pion can be produced from the coalescence of a $d\overline{u}$ pair emerging from the initial higher particle number Fock wavefunction of the B. The D meson is then formed from the remaining quarks after the internal exchange of a W boson.

In principle, a precise evaluation of the hadronic matrix elements needed for B-decays and other exclusive electroweak decay amplitudes requires knowledge of all of the light-cone Fock wavefunctions of the initial and final state hadrons. In the case of model gauge theories such as QCD(1+1) [32] or collinear QCD [33] in one-space and one-time dimensions, the complete evaluation of the light-cone wavefunction is possible for each baryon or meson bound-state using the DLCQ method.[34, 33] It would be interesting to use such solutions as a model for physical B-decays.

3 The Transition from Soft to Hard Physics

The existence of an exact formalism provides a basis for systematic approximations and a control over neglected terms. For example, one can analyze exclusive semi-leptonic B-decays which involve hard internal momentum transfer using a perturbative QCD formalism[35, 6] patterned after the analysis of form factors at large momentum transfer.[5] The hard-scattering analysis proceeds by writing each hadronic wavefunction as a sum of soft and hard contributions

$$\psi_n = \psi_n^{\rm soft}(\mathcal{M}_n^2 < \Lambda^2) + \psi_n^{\rm hard}(\mathcal{M}_n^2 > \Lambda^2), \tag{2}$$

where $\mathcal{M}_n^2$ is the invariant mass of the partons in the n-particle Fock state and Λ is the separation scale. The high internal momentum contributions to the wavefunction $\psi_n^{\rm hard}$ can be calculated systematically from QCD perturbation theory by iterating the gluon exchange kernel. The contributions from high momentum transfer exchange to the B-decay amplitude can then be written as a convolution of a hard-scattering quark-gluon scattering amplitude T_H with the distribution amplitudes $\phi(x_i, \Lambda)$, the valence wavefunctions obtained by integrating the constituent momenta up to the separation scale $\mathcal{M}_n < \Lambda < Q$. This is the basis for the perturbative hard-scattering analyses.[35, 36, 37, 6] In the exact analysis, one can identify the hard PQCD contribution as well as the soft contribution from the convolution of the light-cone wavefunctions. Furthermore, the hard-scattering contribution can be systematically improved.

4 Hard Exclusive Processes

In general, hard exclusive hadronic amplitudes such as quarkonium decay, heavy hadron decay, and scattering amplitudes where hadrons are scattered with large momentum transfer can be factorized at leading power as a convolution of distribution amplitudes and hard-scattering quark/gluon matrix elements[5]

$$\mathcal{M}_{\rm Hadron} = \prod_H \sum_n \int \prod_{i=1}^n d^2k_\perp \prod_{i=1}^n dx\, \delta\left(1 - \sum_{i=1}^n x_i\right) \delta\left(\sum_{i=1}^n \vec{k}_{\perp i}\right) \times \psi_{n/H}^{(\Lambda)}(x_i, \vec{k}_{\perp i}, \Lambda_i)\, T_H^{(\Lambda)} \ . \tag{3}$$

Here $T_H^{(\Lambda)}$ is the underlying quark-gluon subprocess scattering amplitude in which the (incident and final) hadrons are replaced by their respective quarks and gluons with momenta $x_i p^+$, $x_i \vec{p}_\perp + \vec{k}_{\perp i}$ and invariant mass above the separation scale $\mathcal{M}_n^2 > \Lambda^2$. The essential part of the wavefunction is the hadronic distribution amplitudes, [5] defined as the integral over transverse momenta of the valence (lowest particle number) Fock wavefunction; *e.g.* for the pion

$$\phi_\pi(x_i, Q) \equiv \int d^2k_\perp\, \psi_{q\bar{q}/\pi}^{(Q)}(x_i, \vec{k}_{\perp i}, \lambda) \tag{4}$$

where the global cutoff Λ is identified with the resolution Q. The distribution amplitude controls leading-twist exclusive amplitudes at high momentum transfer, and it can be related to the gauge-invariant Bethe-Salpeter wavefunction at equal light-cone

time $\tau = x^+$. The $\log Q$ evolution of the hadron distribution amplitudes $\phi_H(x_i, Q)$ can be derived from the perturbatively-computable tail of the valence light-cone wavefunction in the high transverse momentum regime. The LC ultraviolet regulators provide a factorization scheme for elastic and inelastic scattering, separating the hard dynamical contributions with invariant mass squared $\mathcal{M}^2 > \Lambda^2_{\rm global}$ from the soft physics with $\mathcal{M}^2 \leq \Lambda^2_{\rm global}$ which is incorporated in the nonperturbative LC wavefunctions. The DGLAP evolution of quark and gluon distributions can also be derived in an analogous way by computing the variation of the Fock expansion with respect to Λ^2. The natural renormalization scheme for the QCD coupling in hard exclusive processes is $\alpha_V(Q)$, the effective charge defined from the scattering of two infinitely-heavy quark test charges. The renormalization scale can then be determined from the virtuality of the exchanged momentum of the gluons, as in the BLM and commensurate scale methods.[38, 16, 17, 18]

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

(1) The leading power fall-off is given by dimensional counting rules for the hard-scattering amplitude: $T_H \sim 1/Q^{n-1}$, where n is the total number of fields (quarks, leptons, or gauge fields) participating in the hard scattering.[20, 21] Thus the reaction is dominated by subprocesses and Fock states involving the minimum number of interacting fields. The hadronic amplitude follows this fall-off modulo logarithmic corrections from the running of the QCD coupling, and the evolution of the hadron distribution amplitudes. In some cases, such as large angle $pp \to pp$ scattering, pinch contributions from multiple hard-scattering processes must also be included.[39] The general success of dimensional counting rules implies that the effective coupling $\alpha_V(Q^*)$ controlling the gluon exchange propagators in T_H are frozen in the infrared, *i.e.*, have an infrared fixed point, since the effective momentum transfers Q^* exchanged by the gluons are often a small fraction of the overall momentum transfer.[19] The pinch contributions are suppressed by a factor decreasing faster than a fixed power.[20]

(2) The leading power dependence is given by hard-scattering amplitudes T_H which conserve quark helicity.[12, 40] Since the convolution of T_H with the light-cone wavefunctions projects out states with $L_z = 0$, the leading hadron amplitudes conserve hadron helicity; *i.e.*, the sum of initial and final hadron helicities are conserved.

(3) Since the convolution of the hard scattering amplitude T_H with the light-cone

wavefunctions projects out the valence states with small impact parameter, the essential part of the hadron wavefunction entering a hard exclusive amplitude has a small color dipole moment. This leads to the absence of initial or final state interactions among the scattering hadrons as well as the color transparency. of quasi-elastic interactions in a nuclear target.[13, 14] For example, the amplitude for diffractive vector meson photoproduction $\gamma^*(Q^2)p \to \rho p$, can be written as convolution of the virtual photon and the vector meson Fock state light-cone wavefunctions the $gp \to gp$ near-forward matrix element.[22] One can easily show that only small transverse size $b_\perp \sim 1/Q$ of the vector meson distribution amplitude is involved. The sum over the interactions of the exchanged gluons tend to cancel reflecting its small color dipole moment. Since the hadronic interactions are minimal, the $\gamma^*(Q^2)N \to \rho N$ reaction at large Q^2 can occur coherently throughout a nuclear target in reactions without absorption or final state interactions. The $\gamma^* A \to VA$ process thus provides a natural framework for testing QCD color transparency. Evidence for color transparency in such reactions has been found by Fermilab experiment E665.[41]

5 Measurement of Light-cone Wavefunctions and Tests of Color Transparency via Diffractive Dissociation.

Diffractive multi-jet production in heavy nuclei provides a novel way to measure the shape of the LC Fock state wavefunctions and test color transparency. For example, consider the reaction [42, 43, 44] $\pi A \to \mathrm{Jet}_1 + \mathrm{Jet}_2 + A'$ at high energy where the nucleus A' is left intact in its ground state. The transverse momenta of the jets have to balance so that $\vec{k}_{\perp i} + \vec{k}_{\perp 2} = \vec{q}_\perp < {R^{-1}}_A$, and the light-cone longitudinal momentum fractions have to add to $x_1 + x_2 \sim 1$ so that $\Delta p_L < R_A^{-1}$. The process can then occur coherently in the nucleus. Because of color transparency, *i.e.*, the cancelation of color interactions in a small-size color-singlet hadron, the valence wavefunction of the pion with small impact separation will penetrate the nucleus with minimal interactions, diffracting into jet pairs.[42] The $x_1 = x$, $x_2 = 1 - x$ dependence of the di-jet distributions will thus reflect the shape of the pion distribution amplitude; the $\vec{k}_{\perp 1} - \vec{k}_{\perp 2}$ relative transverse momenta of the jets also gives key information on the underlying

shape of the valence pion wavefunction.[43, 44] The QCD analysis can be confirmed by the observation that the diffractive nuclear amplitude extrapolated to $t = 0$ is linear in nuclear number A, as predicted by QCD color transparency. The integrated diffractive rate should scale as $A^2/R_A^2 \sim A^{4/3}$. A diffractive dissociation experiment of this type, E791, is now in progress at Fermilab using 500 GeV incident pions on nuclear targets.[2] The preliminary results from E791 appear to be consistent with color transparency. The momentum fraction distribution of the jets is consistent with a valence light-cone wavefunction of the pion consistent with the shape of the asymptotic distribution amplitude, $\phi_\pi^{\rm asympt}(x) = \sqrt{3} f_\pi x(1-x)$. As discussed below, data from CLEO[1] for the $\gamma\gamma^* \to \pi^0$ transition form factor also favor a form for the pion distribution amplitude close to the asymptotic solution[5] to the perturbative QCD evolution equation.[45, 46, 19, 47, 48] It will also be interesting to study diffractive tri-jet production using proton beams $pA \to {\rm Jet}_1 + {\rm Jet}_2 + {\rm Jet}_3 + A'$ to determine the fundamental shape of the 3-quark structure of the valence light-cone wavefunction of the nucleon at small transverse separation.[43] One interesting possibility is that the distribution amplitude of the $\Delta(1232)$ for $J_z = 1/2, 3/2$ is close to the asymptotic form $x_1 x_2 x_3$, but that the proton distribution amplitude is more complex. This would explain why the $p \to \Delta$ transition form factor appears to fall faster at large Q^2 than the elastic $p \to p$ and the other $p \to N^*$ transition form factors.[49] Conversely, one can use incident real and virtual photons: $\gamma^* A \to {\rm Jet}_1 + {\rm Jet}_2 + A'$ to confirm the shape of the calculable light-cone wavefunction for transversely-polarized and longitudinally-polarized virtual photons. Such experiments will open up a direct window on the amplitude structure of hadrons at short distances.

6 Leading Power Dominance in Exclusive QCD Processes

There are a large number of measured exclusive reactions in which the empirical power law fall-off predicted by dimensional counting and PQCD appears to be accurate over a large range of momentum transfer. These include processes such as the proton form factor, time-like meson pair production in e^+e^- and $\gamma\gamma$ annihilation, large-angle scattering processes such as pion photoproduction $\gamma p \to \pi^+ p$, and nuclear

processes such as the deuteron form factor at large momentum transfer and deuteron photodisintegration.[50] A spectacular example is the recent measurements at CESR of the photon to pion transition form factor in the reaction $e\gamma \to e\pi^0$.[1] As predicted by leading twist QCD[5] $Q^2 F_{\gamma\pi^0}(Q^2)$ is essentially constant for 1 GeV^2 $< Q^2 <$ 10 GeV^2. Further, the normalization is consistent with QCD at NLO if one assumes that the pion distribution amplitude takes on the form $\phi_\pi^{\rm asympt}(x) = \sqrt{3} f_\pi x(1-x)$ which is the asymptotic solution[5] to the evolution equation for the pion distribution amplitude.[45, 46, 19, 48]

The measured deuteron form factor and the deuteron photodisintegration cross section appear to follow the leading-twist QCD predictions at large momentum transfers in the few GeV region.[51, 52] The normalization of the measured deuteron form factor is large compared to model calculations [53] assuming that the deuteron's six-quark wavefunction can be represented at short distances with the color structure of two color singlet baryons. This provides indirect evidence for the presence of hidden color components as required by PQCD.[15]

There are, however, experimental exceptions to the general success of the leading twist PQCD approach, such as (a) the dominance of the $J/\psi \to \rho\pi$ decay which is forbidden by hadron helicity conservation and (b) the strong normal-normal spin asymmetry A_{NN} observed in polarized elastic $pp \to pp$ scattering and an apparent breakdown of color transparency at large CM angles and $E_{CM} \sim 5$ GeV. These conflicts with leading-twist PQCD predictions can be used to identify the presence of new physical effects. For example, It is usually assumed that a heavy quarkonium state such as the J/ψ always decays to light hadrons via the annihilation of its heavy quark constituents to gluons. However, as Karliner and I [54] have recently shown, the transition $J/\psi \to \rho\pi$ can also occur by the rearrangement of the $c\overline{c}$ from the J/ψ into the $|\, q\overline{q}c\overline{c}\rangle$ intrinsic charm Fock state of the ρ or π. On the other hand, the overlap rearrangement integral in the decay $\psi' \to \rho\pi$ will be suppressed since the intrinsic charm Fock state radial wavefunction of the light hadrons will evidently not have nodes in its radial wavefunction. This observation provides a natural explanation of the long-standing puzzle why the J/ψ decays prominently to two-body pseudoscalar-vector final states, whereas the ψ' does not. The unusual effects seen in elastic proton-proton scattering at $E_{CM} \sim 5$ GeV and large angles could be related to the charm threshold and the effect of a $|\, uuduudc\overline{c}\rangle$ resonance which would appear as in

the $J = L = S = 1$ pp partial wave.[55]

If the pion distribution amplitude is close to its asymptotic form, then one can predict the normalization of exclusive amplitudes such as the spacelike pion form factor $Q^2 F_\pi(Q^2)$. Next-to-leading order predictions are now becoming available which incorporate higher order corrections to the pion distribution amplitude as well as the hard scattering amplitude.[9, 56, 57] However, the normalization of the PQCD prediction for the pion form factor depends directly on the value of the effective coupling $\alpha_V(Q^*)$ at momenta $Q^{*2} \simeq Q^2/20$. Assuming $\alpha_V(Q^*) \simeq 0.4$, the QCD LO prediction appears to be smaller by approximately a factor of 2 compared to the presently available data extracted from the original pion electroproduction experiments from CEA.[58] A definitive comparison will require a careful extrapolation to the pion pole and extraction of the longitudinally polarized photon contribution of the $ep \to \pi^+ n$ data.

A debate has continued on whether processes such as the pion and proton form factors and elastic Compton scattering $\gamma p \to \gamma p$ might be dominated by higher twist mechanisms until very large momentum transfers.[59, 60, 61] For example, if one assumes that the light-cone wavefunction of the pion has the form $\psi_{\rm soft}(x, k_\perp) = A \exp(-b \frac{k_\perp^2}{x(1-x)})$, then the Feynman endpoint contribution to the overlap integral at small $k_\perp$ and $x \simeq 1$ will dominate the form factor compared to the hard-scattering contribution until very large Q^2. However, the above form of $\psi_{\rm soft}(x, k_\perp)$ has no suppression at $k_\perp = 0$ for any x; *i.e.*, the wavefunction in the hadron rest frame does not fall-off at all for $k_\perp = 0$ and $k_z \to -\infty$. Thus such wavefunctions do not represent well soft QCD contributions. Furthermore, such endpoint contributions will be suppressed by the QCD Sudakov form factor, reflecting the fact that a near-on-shell quark must radiate if it absorbs large momentum. If the endpoint contribution dominates proton Compton scattering, then both photons will interact on the same quark line in a local fashion and the amplitude is real, in strong contrast to the QCD predictions which have a complex phase structure. The perturbative QCD predictions[62] for the Compton amplitude phase can be tested in virtual Compton scattering by interference with Bethe-Heitler processes.[63] It should be noted that there is no apparent endpoint contribution which could explain the success of dimensional counting in large angle pion photoproduction.

It is interesting to compare the corresponding calculations of form factors of

bound states in QED. The soft wavefunction is the Schrödinger-Coulomb solution $\psi_{1s}(\vec{k}) \propto (1 + \vec{p}^2/(\alpha m_{\rm red})^2)^{-2}$, and the full wavefunction, which incorporates transversely polarized photon exchange, only differs by a factor $(1 + \vec{p}^2/m_{\rm red}^2)$. Thus the leading twist dominance of form factors in QED occurs at relativistic scales $Q^2 > m_{\rm red}^2$.[64] Furthermore, there are no extra relative factors of α in the hard-scattering contribution. If the QCD coupling α_V has an infrared fixed point, then the fall-off of the valence wavefunctions of hadrons will have analogous power-law forms, consistent with the Abelian correspondence principle.[65] If such power-law wavefunctions are indeed applicable to the soft domain of QCD then, the transition to leading-twist power law behavior will occur in the nominal hard perturbative QCD domain where $Q^2 \gg \langle k_\perp^2 \rangle, m_q^2$.

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