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FACTORIZATION AND OTHER NOVEL EFFECTS IN QCD*

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Abstract

Recent progress in proving the validity of factorization for inclusive reactions in QCD is reviewed. A new necessary condition involving the target length is emphasized. We also discuss a number of novel effects in gauge theory including null zone phenomena, color transparency, formation zone conditions, and possible heavy quark Fock states components in ordinary hadrons.

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1. Introduction

An important way to check quantum chromodynamics is to test its novel predictions — especially effects unique to local gauge theory. In this talk I will discuss a number of unusual or unexpected aspects of QCD. These include: "null zone" phenomena — zeroes in the cross section for photon emission specific to gauge theories; "color transparency" phenomena — the small value of interaction cross sections for specific components of hadronic wavefunctions; "formation zone" phenomena — the suppression of inelastic interactions at high energies in targets of fixed length; and "intrinsic charm" — the unusual kinematical effect of virtual heavy quark components in the wavefunctions of ordinary hadrons. I will also discuss progress in proving the standard factorization ansatz for high momentum transfer inclusive processes. As we shall see, factorization for the Drell-Yan process and the absence of color correlations — now verified to two loops in perturbation theory — is itself a novel aspect of local gauge theory.

2. Radiation Null Zones[1]

A surprising feature of the subprocess cross section $\frac{d\sigma}{d\Omega}(u\,\bar{d} \to W^+\gamma)$, calculated in tree graph approximation in $SU(2) \times U(1)$ gauge theory, is the fact that each of the contributing tree graph helicity amplitudes vanishes near $\cos\theta_{CM}^{\gamma d} = 1/3$ (see Fig. 1).[2] In fact, this is a special case of a general theorem[1,3] for gauge theories applicable to any photon emission process: every tree-graph helicity amplitude $\mathcal{M}^{\mu}_{\lambda_i \ldots \lambda_n}$ for radiation produced by the scattering of n incident and final particles vanishes at the kinematic domain such that all the ratios

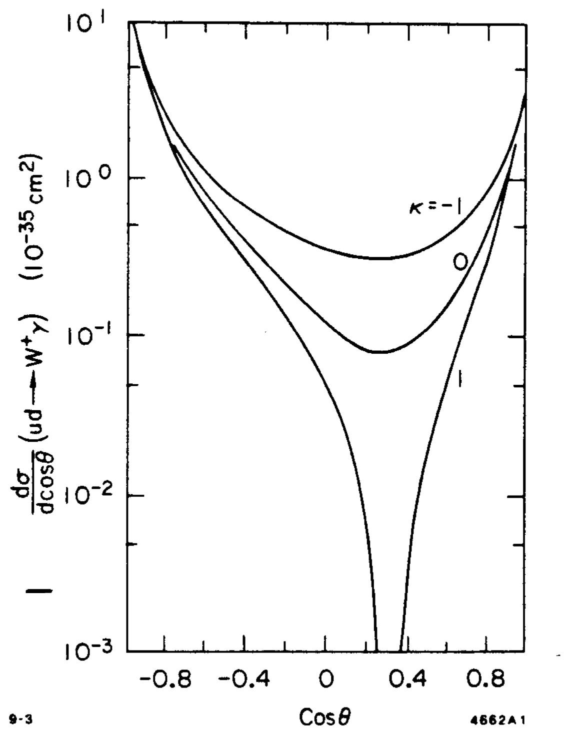

Fig. 1. The differential cross-section for $u\,\bar{d} \to W^+\gamma$ for $SU(2) \times U(1)$ gauge theory in Born approximation. The subprocess cross section vanishes identically near $\cos\theta^{\gamma d} = 1/3$.

$Q_i/p_i \cdot k$ are equal:

$$\frac{Q_i}{p_i \cdot k} = \frac{Q_1}{p_1 \cdot k} \qquad i = 2, \ldots, n \quad .$$

(Because of charge and four-momentum conservation this is actually only $n-2$ independent conditions.) The Born cross section is thus identically zero in this kinematic domain, which we will refer to as the "null zone". Notice that the photon energy is essentially unrestricted, not limited to the usual infrared regime of soft photon theorems. The general proof of this result for gauge theories for processes with any number of charged spin 0, $\frac{1}{2}$ or 1 particles with minimal electromagnetic coupling is given in Ref. 1. The essential elements of the proof are:

(1) In the null zone the radiation from the classical (convection) currents of the external lines destructively interfere.

(2) The spin currents of charged spin $\frac{1}{2}$ and spin 1 particles in gauge theory tree graph amplitudes can be represented in the form of an infinitesimal psuedo Lorentz transformation since in each case $g = 2$. In the null zone the radiation due to the spin currents then cancel, because of Lorentz invariance. The cancellation of spin current contributions is also related to the fact that for any spin $g = 2$ implies that the spin precession and Larmor frequencies are identical.[4] The radiation associated with derivative couplings and seagull contributions also cancel, again because they can be related to infinitesimal Lorentz transformations.

(3) The radiation from internal lines can be rewritten using Ward-type identities in the form of a sum of external line emission processes, each of which again give vanishing contribution in the null zone.

The null zone cancellations only hold for tree graph amplitudes — quantum corrections due to diagrams with internal loops lead to $g \neq 2$ and break the exact destructive interference.[5] The null zone phenomena, being a general result of local gauge theory, is interesting from several points of view:

(a) In the case of $u\bar{d} \to W^+\gamma$, verification of a dip in the cross section at the null zone point $\cos\theta_{cm}^{\gamma d} = \frac{1}{\beta_q}\,(Q_u - Q_{\bar{d}})/Q_W = \frac{1}{3} + O(m_q^2/m_W^2)$ tests not only that the gyromagnetic ratios g_W and g_q have the Dirac value $g = 2$ in gauge theories, but it also measures the fractional charge of the quark. The corrections from QCD higher order loop diagrams are $O(\alpha_s(m_W^2))$. In addition there are k_T-smearing and off-shell corrections due to quark transverse momentum of order $\langle k_\perp^2 \rangle / m_W^2$.

(b) We have emphasized that the null zone cancellation is a general property for all scattering amplitudes involving photon emission calculated from tree graphs in gauge theory. Most often the null zone region lies outside the physical kinematic regime unless the conditions described in Refs. 1 and 6 are met; for example: all the incident and final particles have to have the same sign of charge. Other measurable examples include the QED process $d\sigma(e^-e^- \to$

$e^-e^-\gamma$) which vanishes to leading order in α in the two-dimensional region illustrated in Fig. 2. Another interesting process is $e^+e^+ \to Q\bar{Q}\gamma$; as shown by Passarino,[6] the null zone for this process can be used to measure the heavy quark mass.

(c) Independent of whether the null cone for a given process lies in the physical region, any tree graph radiation gauge theory amplitude can be written in the compact representation[1,3]

$$M^\mu = \sum_{i=2}^{n-1}\left(\frac{Q_i}{p_i\cdot k} - \frac{Q_1}{p_1\cdot k}\right)\hat{M}_i^\mu$$

independent of helicity.

(d) The vanishing of the cross section at a specific point in momentum space is consistent with the uncertainty principle since the null zone condition only depends on the external lines which have unspecified position. The results are also consistent with the correspondence principle: the classical ($\hbar = 0$) tree graph limit of gauge theories is consistent with classical radiation. These results imply that the only consistent possiblity for electromagnetic spin couplings at the tree graph level is $g = 2$ for any spin; $g \neq 2$ must come from quantum corrections.[1] This conclusion is also consistent with the Drell-Hearn Gerasimov sum rule. Conversely, effective local field theories of nucleon and mesons which have $g_N \neq 2$ at the tree graph level are inconsistent with the correspondence principle. The null zone phenomena thus provide another criteria for constructing acceptable fundamental theories.

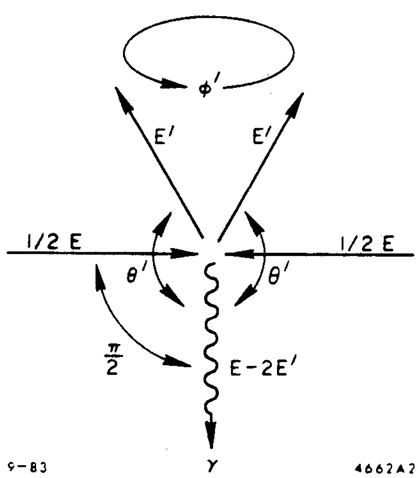

Fig. 2. Kinematics for the null zone for $e^-e^- \to e^-e^-\gamma$.

3. Factorization for High Momentum Transfer Inclusive Reactions[7]

One of the most important problems in perturbative QCD in the last two years has been to understand the validity of the standard factorization ansatz for hadron-hadron induced inclusive reactions. Although factorization is an implicit property of parton models, the existence of diagrams with color exchanging initial state interactions at the leading twist level has made the general proof of factorization in QCD highly problematical.

To see the main difficulties from a physical perspective, consider the usual form assumed for massive lepton pair production [see Fig. 3(a)]

$$\frac{d\sigma}{dx_1 dx_2}(H_A H_B \to \ell\bar{\ell}X) = \frac{1}{3}\frac{4\pi\alpha^2}{3Q^2}\sum_i Q_i^2\left[q_A^{(1)}(x_i, Q)\,\bar{q}_B^{(2)}(x_2, Q) + (1 \to 2)\right] \tag{3.1}$$

The factorization ansatz identifies the Q^2-evolved quark distributions q_A and $\bar{q}_A$ with those measured in deep inelastic lepton scattering on H_A and H_B. However, for very long targets the initial-state hadronic interactions occurring before the $q\bar{q} \to \ell\bar{\ell}$ annihilation certainly lead to induced radiation and energy loss, secondary beam production, transverse momentum fluctuations, etc. – i.e.: a profound modification of the incoming hadronic state [see Fig. 3(b)]. Since the structure functions associated with deep inelastic neutrino scattering are essentially additive in quark number even for macroscopic targets, Eq. (3.1) can obviously not be valid in general. At the least, an explicit condition related to target length must occur. The original proofs of factorization in QCD for the Drell-Yan process ignored the (Glauber) singularities associated with initial state interactions and thus had no length condition.

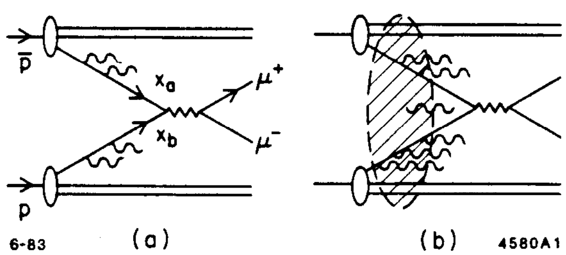

Fig. 3. (a) Gluon emission associated with QCD evolution of structure functions for the Drell-Yan process, $p\bar{p} \to \mu^+\mu^- X$. (b) Gluon emission associated with initial state interactions for the Drell-Yan process. The shaded area represents elastic and inelastic scattering of the incident quarks.

The potential problems and complications associated with "wee parton" exchange in the initial state were first mentioned by Drell and Yan[8] in their original work. Collins and Soper[9] have noted that proofs of factorization for hadron pair production in $e^+e^- \to H_A H_B X$ could not be readily extended to $H_A H_B \to \ell\bar{\ell}X$ because of the complications of initial state effects. Possible complications associated with nonperturbative interaction effects were also discussed by Ellis *et al.*[10] More recently Bodwin, Lepage, and I[7] considered the effects of initial state interactions as given by perturbative QCD and showed that specific graphs such as those in Fig. 4 lead to color exchange correlations as well as $k_\perp$ fluctuations. We also showed that induced hard collinear gluon radiation is indeed suppressed for incident energies large compared to a scale proportional to the length of the target. More recently, the question of the existence of color correlations on perturbative QCD has now been addressed systematically to two loop order by Lindsay *et al.*[11] and by Bodwin *et al.*[7] One finds that because of unitarity and local gauge invariance to two loop order the factorization theorem for $d\sigma/dQ^2\,dx_L$ is correct when applied at high energies to color singlet incident hadrons; more general proofs beyond two loop order await further work.[12] We discuss the progress in this area at the end of this section.

In addition to the above initial state interaction there are additional potential infrared problems in the non-Abelian theory associated with the breakdown of the usual Bloch-Nordsiek cancellation for soft gluon radiation. The work of Ref. 13 showed that any observable effect is suppressed by powers of s at high energies, again to at least two loop order.

In addition to these problems the high transverse momentum virtual gluon corrections to the $q\bar{q} \rightarrow \ell\bar{\ell}$ vertex lead to relatively large radiative corrections of relative order $\pi^2 C_F(\alpha_s(Q^2)/\pi)$.[14] It is usually assumed that such corrections exponentiate. As in the case of the $\Upsilon \rightarrow 3g$ problem, these corrections spoil the convergence of the perturbation theory and cannot be eliminated by choice of scale or scheme.

A remarkable feature of the QCD calculation is the fact that factorization is not destroyed by induced radiation in the target for sufficiently high energy beams. This can be understood in terms of the "formation zone" principle of Landau and Pomeranchuk:[15] a system does not alter its state for times short compared to its natural scale in its rest frame. More specifically for QCD (in the Glauber/classical scattering region), consider the diagrams for induced radiation for quark-pion scattering shown in Fig. 4(b). Here $\ell^{\pm} = \ell^0 \pm \ell^3$, $y = \ell^+/p_B^+$, $x_a = p_a^-/p_A^-$ are the usual light-cone variables. The Feynman propagators of the line before and after radiation are proportional to $y - y_1 + i\epsilon$ and $y - y_2 + i\epsilon$, where the difference of the pole contributions is $y_1 - y_2 = \mathcal{M}^2/x_a s$, and $\mathcal{M}^2$ is the mass of the quark-gluon pair after bremsstrahlung. Using partial fractions, the gluon emission amplitude is then proportional to

$$\int_0^1 dy\ \psi[(x_b - y)M_n L] \left[\frac{1}{y - y_1 + i\epsilon} - \frac{1}{y - y_2 + i\epsilon}\right] \qquad (3.2)$$

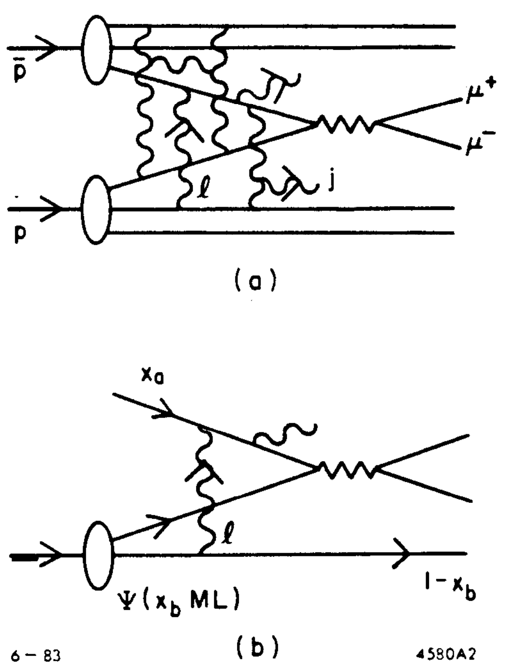

Fig. 4. (a) Representation of initial state interactions in perturbative QCD. (b) Simplest example of induced radiation by initial state interactions in $q\pi \rightarrow \ell\bar{\ell}X$. Two different physical radiation process are included in this Feynman amplitude depending on whether the intermediate state before or after the gluon emission is on-shell. The two bremsstrahlung processes destructively interfere at energies large compared to a scale proportional to the target length L.

where we have indicated the dependence on the target wave function on target length. The two poles thus cancel in the amplitude if $(\mathcal{M}^2/x_a s)\, M_N L \ll 1$; i.e. the radiation from the two Glauber processes destructively interfere and cancel for quark energies large compared to the target length. If we take $\mathcal{M}^2 \sim \mu^2$ finite, then since $Q^2 = x_a x_b s$, the condition for no induced radiation translates to

$$Q^2 \gg x_b M_N L \mu^2 \,. \tag{3.3}$$

Taking $\mu^2 \sim 0.1\, GeV^2$, this is $Q^2 \gg x_b (0.25\, GeV^2)\, A^{2/3}$; thus one requires $Q^2 \gg x_b (10\, GeV^2)$ to eliminate induced radiation in Uranium targets.

Equation (3.3) is a new necessary condition for QCD factorization; it is also a prediction that a new type of nuclear shadowing occurs for low Q^2 lepton-pair production. If this condition is not met then the cancellations found in Ref. 7, for example, fail. The same length condition affects all sources of hard collinear radiation induced by initial or final state interactions of the hadrons or quarks in a nucleus; i.e., effectively hard collinear radiation occurs outside the target at high energies. In particular, fast hadron production from jet fragmentation in $\ell p \to \ell H X$ occurs outside the target. In the case of very long or macroscopic targets, induced radiation destroys any semblance of factorization.

Although induced hard collinear radiation cancels at high energies, the basic processes of $k_\perp$ fluctuations from elastic collisions and induced central radiation [e.g. Fig. 4(a) with $j_z \sim m^2/\sqrt{s}$ in the CM] do remain. One expects that the main effects of initial state interactions can be represented by an eikonal picture where the hadronic wave functions are modified by a phase in impact space (see Fig. 5):

$$\psi_A(x_s, \vec{z}_{a\perp})\psi_B(x_b, \vec{z}_{\perp b}) \to \psi_A(x_a, \vec{z}_{a\perp})\psi_B(x_a, \vec{z}_{b\perp}) U(\vec{z}_{\perp i}) \,. \tag{3.4}$$

Here

$$U(\vec{z}_{\perp b}) = P_T \exp\left\{ -i \int_{-\infty}^{0} d\tau\, H_I(z_\perp, \tau) \right\} \tag{3.5}$$

includes elastic and soft inelastic collisions which occur up to the time $\tau = 0$ of the $q\bar{q}$ annihilation. The eikonal leads to an increased transverse smearing of the lepton pair and increased associated radiation in the central region proportional to the number of collisions

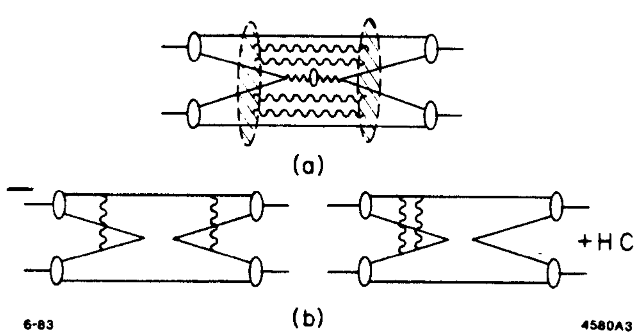

Fig. 5. (a) Representation of initial state interactions in the Drell-Yan cross section $d\sigma/dQ^2dx$. (b) Example of two-loop initial state interactions which cancel by unitarity in an Abelian gauge theory. In QCD these two contributions have different color factors.

($A^{1/3}$) of the quark in the target. For a nucleus we thus predict

$$\Delta\langle Q_\perp^2\rangle \propto A^{1/3}\ , \qquad \Delta \frac{dN}{dy} \propto A^{1/3} \tag{3.6}$$

In the case of an Abelian gauge theory the integrated cross section

$$\int \frac{d\sigma}{dQ^2\, dx_L\, d^2Q_\perp} = \frac{d\sigma}{dQ^2 dx_L} \tag{3.7}$$

in unchanged because of unitarity, $U^\dagger(z_\perp)U(z_\perp) = 1$. See Fig. 5(b). Thus for an Abelian theory, the increased production at large $Q_\perp$ from initial state-interactions must be compensated by a depletion at low $Q_\perp$.

In general, initial state interactions will have a strong modifying effect on all hadron-hadron cross sections which produce particles at large transverse momentum simply because of the $k_\perp$ smearing of very rapidly falling distributions. The initial state exchange interactions combine with the quark and gluon $k_\perp$ distributions intrinsic to the hadron wave functions as well as that induced by the radiation associated with QCD evolution to yield the total $k_\perp$ smearing effect. The unitarity structure of the initial state eikonal interactions provides a finite theory of $k_\perp$ fluctuations even when the hard scattering amplitude is singular at zero momentum transfer.

In a non-Abelian theory the eikonal unitary matrix $U(z_\perp)$ associated with the initial state interactions is a path-color-ordered exponential integrated over the paths of the incident constituents. Since U is a color matrix it would not be expected to commute with the Drell-Yan $q\bar{q} \to \ell\bar{\ell}$ matrix element

$$U^\dagger \mathcal{M}_{DY}^\dagger \mathcal{M}_{DY} U \neq \mathcal{M}_{DY}^\dagger \mathcal{M}_{DY}\ .$$

Thus unless U is effectively diagonal in color, the usual color factor $1/n_c$ in $d\sigma(q\bar{q} \to \ell\bar{\ell})$ would be expected to be modified. In principle, this effect could change the usual color factor $1/n_c$ to n_c or event to 0 without violating unitarity, although, as shown by Mueller,[16] the deviation from $1/n_c$ will be dynamically suppressed; hard gluon radiation at the subprocess vertex leads to asymptotic Sudakov form factor suppression of the color correlation effect.

Despite these general possibilities, it has now been shown that such color correlation effects actually cancel in QCD at least through two loop order, although it is present in individual diagrams. The cancellation through two loops was first demonstrated in perturbation theory by Lindsay, Ross, and Sachrajda[11] for scalar quark QCD interactions in both Feynman and light-cone gauge, and was subsequently confirmed in Feynman gauge by Bodwin *et al.*[7] A detailed physical explanation of the two-loop cancellation is not known; it seems to be a consequence of both causality at high energies and local gauge invariance; neither by itself is sufficient. We also find that the cancellation breaks down at low energies or for long targets when condition (3.3) is not satisfied. It also fails in the case of spontaneous broken gauge theories with heavy gauge boson exchange because the trigluon graph is suppressed.

An example of the nature of the color correlation cancellations is shown in Fig. 6 for $\pi\pi \to \ell\bar{\ell}X$. The diagrams shown are a gauge-invariant distinct class which have a non-trivial non-Abelian color factor and involve interactions with each of the incident spectators. The generality of the pion wave function precludes shifting of the transverse momentum interactions to other graphs. The various virtual two-gluon exchange amplitudes interfering with the zero gluon exchange amplitude each produces a $C_F C_A$ contribution which cancel in the sum. On the other hand, the imaginary part of the virtual graphs gives a non-zero contribution which potentially could lead to a color correlation at four loops. However, we find that even the potentially troublesome imaginary part is cancelled when one includes the real emission diagrams of Figs. 6(d) and 6(e). Explicitly the sum of all the virtual and real emission amplitudes is proportional to

$$\left(C_F^2 - \frac{C_F C_A}{2}\right) \frac{2\vec{\ell}_{1\perp}\cdot\vec{\ell}_{1\perp}}{\ell_{1\perp}^2 \ell_{2\perp}^2} \frac{1}{\ell_1^+ + i\epsilon} \frac{1}{\ell_2^- + i\epsilon} \frac{1}{\ell_1^+\ell_2^- - (\vec{\ell}_{1\perp} + \vec{\ell}_{2\perp})^2 - i\mathcal{E}(-\ell_1^+)} \tag{3.9}$$

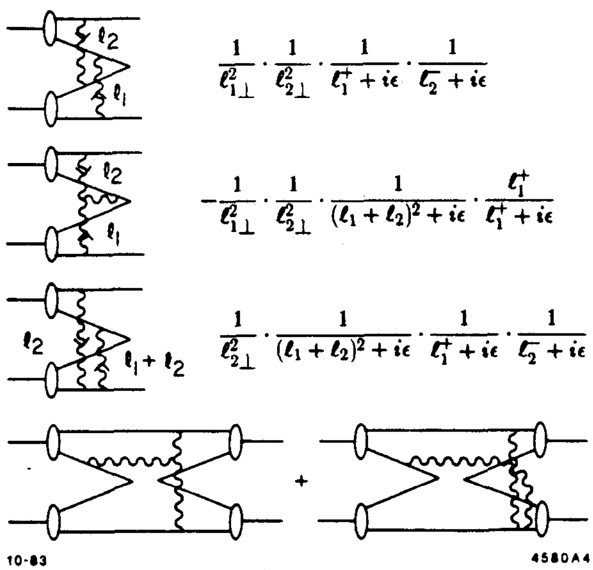

Fig. 6. Representative active spectator initial state interactions for $\pi\pi \to \ell\bar{\ell}X$ in QCD involving $C_F C_A$ evaluated in Feynman gauge. The real part of the two loop contributions represented by (a),(b), (c) (including mirror diagrams) vanishes at high energies. The imaginary parts cancel against the gluon emission contribution represented in (d) and (e).

The integration over ℓ_2^- then leads to zero contributions for the leading power behavior. More generally, the proof of factorization of the Drell-Yan cross section can be divided into two distinct steps, as indicated in Fig. 7. The first step is to prove that every contribution to initial state interactions in hadron-hadron scattering can be written as the convolution of two "eikonal-extended" structure functions as indicated in Fig. 7(a). This is the "weak-factorization" ansatz proposed by Collins, Soper, and Sterman[17] where each structure function has a eikonal factor attached which includes all of the elastic and inelastic initial state interactions of the corresponding incident annihilating quark or anti-quark. Explicitly, the eikonal-extended structure function of the target system A is defined as[18]

$$P_{q/A}(x,k_\perp) = \frac{1}{2(2\pi)^3}\int dy^- \int d^2y_\perp \, e^{i(xP_A^+ y^- - \vec{k}_\perp\cdot\vec{y}_\perp)} \times \langle A|\,\bar{\psi}_{DY}(0,y^-,\vec{y}_\perp)\gamma^+\psi_{DY}(0,0,\vec{0}_\perp)|A\rangle \tag{3.10}$$

where

$$\Psi_{DY}(y^\nu) = P \exp -ig \int_{-\infty}^{0} d\lambda\, n\cdot A(y^\nu+\lambda n^\nu)\psi(y^\nu)$$

and n^μ is chosen such that $n \cdot \ell = 2\ell^3$ in the center-of-mass frame. The path-ordered exponential contains all of the interactions of the eikonal anti-quark line with the color gauge field along the incident $\hat{z}$ direction up to the point of annihilation.

Recently, we have verified[7] that the weak factorization ansatz is correct through two loops in perturbation theory for $\mathcal{M}(A+B \to \ell\bar{\ell}X)$ despite the complicated color-topological structure of the contributing diagrams. The proof relies on splitting each Feynman amplitude into separate structure functions using identities of the form

$$\frac{1}{A\ell^+ + i\epsilon}\,\frac{1}{-B\ell^- + i\epsilon} = \left(\frac{1}{A\ell^+ + i\epsilon}\,\frac{1}{B} + \frac{1}{-B\ell^- + i\epsilon}\,\frac{1}{A}\right)\frac{1}{\ell^+ - \ell^- + i\epsilon} \tag{3.11}$$

and then analytically continuing each contribution out of the Glauber regime to either large ℓ^- or large ℓ^+, corresponding to exchange gluons collinear with the beam or target, respectively. Finally, the use of collinear Ward identities allows one to organize gauge-related diagrams into the desired weak factorization form. We are continuing efforts to try to extend the proof beyond two loop order in QCD.

The second step required to prove factorization is to show that the structure function (3.10) is actually identical to the corresponding eikonal-extended structure function for deep inelastic-lepton-hadron scattering which includes a post-factor for the final state interactions of the struck quark [see Fig. 7(b)]. This becomes intuitively obvious when one examines moments

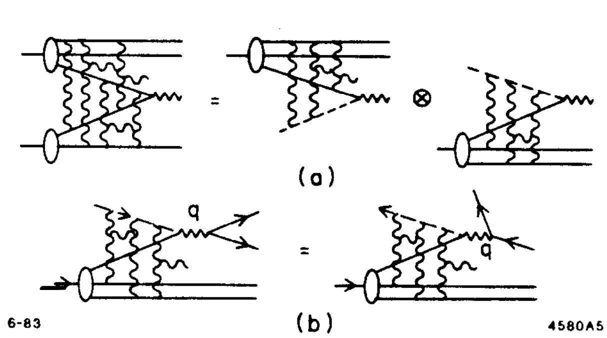

Fig. 7. (a) Schematic representation of the general decomposition required to prove weak factorization to general orders in QCD. The dotted line corresponds to the eikonal line integral of Eq. (3.10). Vertex corrections which modify the hard scattering amplitude are not shown. These provide a separate factor on the right hand side of 7(a). (b) The relationship between Drell-Yan and deep inelastic lepton scattering eikonal-extended structure functions required to prove factorization.

of the two structure functions. These moments differ only by terms proportional to powers of the integral $\int_{-\infty}^{0} dz\, E_z(z)$, where E_z is the longitudinal component of the chromo-electric field along the eikonal line. In the center of momentum frame the hadron has ultrarelativistic momentum along the z axis, and consequently the Lorentz-transformed longitudinal electric fields in the hadron are vanishing small. Thus all the moments, and therefore the structure functions themselves, become identical as $Q \to \infty$. Physically, the effective equality of the structure functions implies that the color fluctuations generated by initial and final interactions at high energies in massive lepton pair production and deep inelastic lepton scattering are basically equivalent.[18]

At this point there is no convincing counterexample to standard QCD factorization for hadron-induced large momentum transfer reactions; on the other hand, there is no proof beyond two loop order for non-Abelian theories. Clearly if factorization is a general feature of gauge theories, then it is a novel and profound feature which demands explanation in fundamental terms.[19] In any event, the initial state interactions lead to new physical phenomena for the $Q_\perp$ distributions, e.g. the nuclear number dependence of $k_\perp$ fluctuations and associated particle production (see below). Furthermore, color correlations and breakdown of factorization do explicitly occur for power-law suppressed contributions which are sensitive to the length scale of the target. Such effects should be measureable for heavy nuclear targets at moderate Q^2.

Although our analysis is based on QCD perturbation theory (to all orders) our conclusions can be expressed in terms of rather general principles:

(1) Critical Momentum Scale. The characteristic momentum of each hard subprocess must be large compared to a scale set by the length of the target (or beam), as in Eq. (3.3); otherwise the constituents, in passing through the target can lose a significant fraction of their longitudinal momentum to radiation, completely destroying any connection between the hadronic reaction and the distributions measured in deep inelastic scattering. This is related to the more general concept of the "formation zone".

(2) Formation Zone. The state of a hadronic system cannot be modified significantly in a time (in its rest system) less than its intrinsic scale. Thus, a high energy quark cannot radiate a collinear gluon $q \to q + g$ inside of a target of length L if $s \gg \Delta(\mathcal{M}^2)LM$ where $\Delta(\mathcal{M}^2)$ is the change in the square of the invariant mass, and LM/s is proportional to the Lorentz-contracted length of the target in the quark rest frame. Similarly, the fragmentation of a quark into collinear hadrons (or vice versa) occurs *outside* of the target volume at high energies. We also note that interactions between quark or gluon constituents of the same hadron do not occur (to leading order in $1/s$) during the transit through the target volume. Thus high energy interactions of hadrons within nuclei are correctly described in terms of constituent quark and gluon propagation.

(3) Large Longitudinal Range. The change of longitudinal momentum (in the CM) due to initial or final state interactions is so small that longitudinal structure in the target cannot be resolved in a target of length $L < \sqrt{s}/\langle \ell_\perp^2 \rangle$ (as measured in the CM frame).

(4) Color Singlet Cancellations. Large momentum transfer *exclusive* reactions are controlled by the Fock states with the minimum number of constituents at transverse distances $b_\perp^2 \sim (1/Q^2)$.[20] The initial and final state collisions can probe transverse distances no smaller than $1/\lambda$. Thus, such interactions cannot resolve the internal structure of the hadrons in exclusive reactions, and they do not couple to these color neutral objects. Formally, the initial and final state interactions cancel to leading order in $1/Q^2$ if one adds the contributions coming from all constituents of a color neutral hadron. This also implies that large momentum transfer quasi-elastic reactions such as $eA \to ep(A-1)$ and $\pi A \to \pi p(A-1)$ can occur deep inside a nuclear target without multiple scattering or bremsstrahlung in the target.[21] Color singlet cancellations also eliminate initial and final state interactions of hadrons interacting directly in hard scattering inclusive reactions. For example, the "direct pion"[22] has no initial state interactions in $\pi_D g \to q\bar{q}$ (in $\pi p \to q\bar{q}X$), and no final state interactions in $(pp \to \pi X)$. There are thus no accompanying spectator hadrons accompanying along the meson in such processes. Similarly the higher twist $p_T^{-8} f(x_T, \theta_{cm})$ subprocess $\bar{p}_D\, q \to \bar{q}\,\bar{q}$ leads to the production of two jets at large p_T in $\bar{p}\,p \to$ Jet + Jet + X without beam spectators. We also note that the higher twist $F_L \sim 1/Q^2$ contribution to the meson structure function[23] is unaffected by initial and final state interactions. On the other hand, although they are power law suppressed at large momentum transfer, initial and final state interactions are expected to play an important role at moderate kinematic values, possibility leading to non-trivial helicity and interference effects.[24] Part of the difference between time-like and space-like form factors, e.g., $e^+e^- \to \pi^+\pi^-$ and $e^-\pi^+ \to e^-\pi^+$ is attributable to final state interactions, although the difference is suppressed by $\sim 1/Q^2$.

In contrast, virtually every large momentum transfer inclusive process in QCD is affected by initial and/or final state interactions. It is important to study the phenomenology of these interactions since they bear on the dynamics of quarks and gluons in hadronic matter and are evidentally related to the confinement mechanisms and the space-time "inside-outside" development of QCD jets.[25] Analysis of the role played by nuclear targets is clearly crucial in this study. Although the structure function measured in deep inelastic lepton scattering are unaffected by initial and final state interactions, the development of the final state jet distribution is modified by multiple scattering in the target. The transverse momentum of the struck quark relative to the current direction will obviously be broadened and multiplicity in the central region will be increased, thus affecting the fragmentation distribution of quarks into hadrons $D_{H/Q}(x, k_\perp)$. These effects should increase with the number of collisions in a nuclear target:

$$\delta\langle k_\perp^2 \rangle = A^{1/3}\,, \quad \delta\langle n_{\rm central}\rangle \alpha\, A^{1/3}\,. \tag{3.12}$$

In addition, for long targets, energy-momentum conservation implies a correlated degradation of the leading particle distribution at large z. For low quark energies, collinear radiation can be induced in the target and can drastically alter the longitudinal momentum fraction distributions.

The development of hadronic multiplicity in deep inelastic lepton scattering in the nucleus is particularly interesting since one is studying the influence of hadronic matter of quark jet propagation. As we have emphasized, formation of the leading particle in the jet occurs outside the nuclear volume at high energies. The inelastic final state interactions amount to cascading in the nucleus and demonstrate that, contrary to the usual assumptions made for the analysis of hadron-nucleus collisions, particle production in the target and central rapidity region cannot be correlated with the number of nucleons "wounded" by the beam. a model for the shape of the rapidity distribution based on "color cascading" is given in Ref. 26.

More generally any hard scattering inclusive process is accompanied by soft hadrons in the central rapidity region, which are the result of the initial state or final state interaction of the quark and gluon constituents. We emphasize that, even though the hard scattering cross section can be computed as if a single interactions occurs, the associated multiplicity reflects the full scope of the actual QCD dynamics.

In the case of hadron production at large transverse momentum in a nucleon or nuclear target collisions the inclusive cross section is increased by the $k_\perp$ smearing effects of the initial and final state interactions. The multiple scattering series in a nucleus[27] leads to terms roughly of order A^1, $A^{4/3}/p_\perp^2$, $A^{5/3}/p_\perp^4$, etc. A coefficient of the A^α terms with $\alpha > 1$ can be quite large, since one is smearing a cross section that falls very rapidly with $p_\perp$. Thus, strongly suppressed cross sections such as $pA \to \bar{p}X$ and $pA \to K^-X$ obtain a much larger nuclear enhancement from quark and gluon scattering effects than channels such as $pA \to \pi^+X$ or $pA \to K^+X$. In the case of direct γ production, the photon has no final state interactions, so only initial state interactions of the active q and g constituents are important. Similarly, at large x_T where direct subprocesses such as $gq \to \pi_D q$ or $qq \to \pi_D g$ are expected to dominate $pA \to \pi X$ production, only initial state interactions are important. Thus one can use direct photon reactions, photoproduction, Compton scattering, and direct hadron interactions, especially the A-dependence of the cross sections, to eliminate and effectively isolate the effect of initial and final state interactions.

Nuclear initial and final state effects are, of course, enhanced in processes such as $A_1A_2 \to HX$. Nuclear targets also enhance the effects of multiple scattering processes that lead to multiple jets in the final state.[28] On the other hand, if the valence state of a hadron consists of constituents at small transverse separation, then the hadron can pass through the target with no color or hadronic interactions. An application of this idea to diffractive dissociation processes in nuclei is discussed by G. Bertsch et al.[29]

Processes such as $pp \to pp\mu\bar{\mu}$,[30] which occur via $\gamma\gamma \to \mu\bar{\mu}$ subprocesses, are also sensitive to the nature of initial state interactions. Unlike the corresponding lepton-induced reaction $ee \to ee\mu\bar{\mu}$, the initial state interactions of the two nucleons smear the transverse momentum distribution of the $\mu\bar{\mu}$ pair and can eliminate the strong peaking at $Q_\perp = 0$ associated with the γ poles. However, the cross section integrated over all $Q_\perp$ is unchanged.

4. Other Novel QCD Effects

In this section we will briefly review several other novel QCD effects; further details may be found in the referenced papers.

Higher Twist Anomalies. As described in Ref. 31, there are now a large number of higher twist photon and direct hadron subprocesses which can be absolutely normalized using the analysis of Ref. 20. In particular, the longitudinal structure function of the pion can be absolutely normalized in terms of the pion form factor.[31]

$$F_{2\pi}(x,Q^2) = A(1-x)^2 + \frac{C}{Q^2}\,x^2$$

where

$$C = \sum e_q^2 \frac{C_F}{2\pi} \int\limits_{\mu^2/(1-x)}^{Q^2/x} d\ell^2 \alpha_s(\ell^2) F_\pi(\ell^2) \simeq 0.1\ GeV^2\ .$$

This basic QCD prediction can be tested for the dominance of $(C/Q^2)\sin^2\theta$ dependence of the Drell-Yan $\pi p \to \ell\bar{\ell}X^-$ cross section at $x_1 \sim 1$. In addition one expects contributions of order $(C/Q)(1-x_a)\sin 2\theta\cos\phi$ (θ and ϕ are the angles of the ℓ_+ in the γ^* rest frame) from longitudinal-scalar interference; as emphasized by Pire and Ralston,[24] these phase-sensitive contributions can have an interesting interference pattern due to Sudakov form factor effects.

The higher twist contributions to the nucleon structure function at $x \sim 1$ have now been systematically computed to lowest order in α_s by Blankenbecler, Gunion, and Nason.[32] The result computed from the set of two-gluon exchange diagrams has the form

$$F_{2N}(x,Q^2) = A(1-x)^3 + \frac{B(1-x)}{Q^2} + \frac{C}{Q^2}\,(1-x)^2$$

where $B \simeq -6\mu^2$, and $C \simeq 800\ \mu^2$. Here μ^2 is a typical hadronic scale, estimated as $\mu^2 = 0.01\ GeV$. The astonishingly large size of the $C(1-x)^2/Q^2$ term implies large power-law contributions to the scale-breaking of deep inelastic lepton-nucleon structure functions, consistent with those parametrized by Barnett et al.[32]

Nuclear Chromodynamics.[33] One of the most interesting areas of application of QCD is to nuclear dynamics. These include corrections to nuclear additivity of nuclear structure functions (the EMC effect), calculations of nuclear amplitudes at large momentum transfer (e.g., the deuteron form factor); the application of "reduced" nuclear amplitudes which are defined to remove the effects of nuclear compositions in a covariant fashion; evolution equations for nuclear wavefunctions — e.g the deuteron 6-quark wavefunction evolves to a state which is 80% hidden color at small internucleon separation.[34] Many traditional concepts of standard nuclear physics phenomenology (e.g. the impulse approximation to nuclear form factors, point-like nuclear pair and meson-exchange current contributions to electromagnetic nuclear amplitudes, local meson-nucleon field theory, and simple Dirac equations for relativistic nucleons), require substantial modification. Details, discussion, and further applications may be found in Ref. 33.

Intrinsic Charm. The dynamical origin of heavy quark states in hadron collisions such as charm production is still not satisfactorily understood. The data[35] suggests large contributions of a diffractive nature leading to fast forward charm production $pp \to \Lambda_c X$ at large x_L at ISR energies; the nuclear A-dependence of the charm production cross section at large x_L appears similar to that of total cross sections. It has been suggested[36] that the magnitude and forward behavior of the charm production cross section at high energies can be understood if the wave function of the proton contains charmed quarks at the 1% probability level with a valence-like $G_{c/p} \sim (1-x)^3$ behavior ("intrinsic" charm) in contrast to the soft dependence $\sim (1-x)^7$ usually associated with sea quarks, $g \to c\bar{c}$, or QCD evolution. One can motivate the valence-like distribution of heavy virtual quarks in a light hadron using an atomic physics analogy: Consider the $\mu^+\mu^-e^+e^-$ contribution to the positronium Fock state wavefunction generated by $\mu^+\mu^-$ vacuum polarization, the same contributions which yield the Serber-Uehling potential. The μ^+ and μ^- are produced dominantly at low velocities since this minimizes the off-shell energy of the virtual state. If one now views the atom from a moving relativistic frame, then the fact that the leptons all have nearly the same velocity implies that the muons carry a large momentum fraction since momentum increases with mass. Similarly, a Fock state of an ordinary hadrons containing heavy quarks will be dominated by configurations in which heavy quarks have low velocity relative to the valence quarks, since this increases the QCD binding as well as decreases the off-shell energy. Again this implies that the intrinsic heavy quark components associated with the hadron wavefunction are produced dominantly at large x. If the atomic physics analogy is correct, then the probability for such Fock states scales as $1/m_Q^2$ and is controlled by $\alpha_s(\mu^2)$ where μ is the ordinary hadronic scale. The diffractive excitation of such states at very high energies where $t_{\min}$ effects are negligible then leads to large diffractive cross sections proportional to $1/m_Q^2$ with A^α $(\alpha < 1)$ dependence. A complete understanding of the charm production cross section requires an understanding of both central region $gg \to c\bar{c}$ (with A^1 dependence) and the diffractive intrinsic heavy quark components.

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