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A UNIFIED APPROACH TO QCD PHENOMENOLOGY*

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ABSTRACT

The Fock state wave functions of hadrons defined by quantizing quantum chromodynamics at equal time on the light-cone provides a unifying element in QCD predictions, from low to high momentum transfer. A number of novel QCD effects are reviewed, including heavy quark and higher twist phenomena, initial and final state interactions, direct processes, multiparticle collisions, color transparency, and nuclear target effects. A method for fixing the momentum scale in leading order QCD predictions is also briefly presented.

1. INTRODUCTION

There is now an extraordinary array of experimental observations which support the premise that the basic degrees of freedom of hadrons and their interactions are the confined quark and gluon fields of quantum chromodynamics.[1] The empirical evidence ranges from hadronic spectroscopy (including the heavy quark bound state spectrum, and the emerging evidence for gluonic bound states), the basic phenomena of deep inelastic lepton scattering and massive lepton pair production (consistent with point-like spin 1/2 quarks carrying the electromagnetic and weak currents in hadrons and the QCD-predicted pattern of scale violation), the scaling of $\sigma(e^+e^- \to X)$ (consistent with SU(3) color, asymptotic freedom), and large momentum transfer exclusive processes such as hadronic form factors (consistent with QCD dimensional counting rules and scale-invariant quark-quark interactions at short distances).

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Multiparticle production appears to be consistent with the general patterns of jet development and quantum number flow expected in QCD, including evidence for gluon jets in $e^+e^- \to q\bar{q}g$ events, quark fragmentation in $\ell p \to \ell' HX$, and the dominance of hard scattering QCD mechanisms for high $p_\perp$ direct photon and single hadron production in high energy hadron-hadron collisions. The recent observation of clear jet signals at the SPS collider, reported to the meeting[2] together with the jet data for the ISR and FNAL appear to give a striking confirmation of the quark and gluon hard-scattering processes predicted by QCD. The recent results from PETRA[3] on $\gamma\gamma \to$ jets and the photon structure functions are especially important since they provide immediate and striking verification of the QCD-predicted pointlike coupling of real photons to the quark current at high momentum transfer.

Certainly, at the qualitative level, QCD does provide a viable framework for understanding present hadronic phenomena. The paradox is that despite these successes we are not certain that we are actually testing QCD predictions for hadron dynamics in a truly quantitative way, particularly since many results follow from simple parton ideas or more general principles, independent of the theory.

There are several reasons why quantitative tests of QCD have been so difficult.

i) Even for the simplest processes, many different QCD mechanisms contribute, including initial state corrections (including color correlations) and "higher twist" terms (non-leading terms in $1/Q^2$), some of which involve multiparticle coherent effects and heavy quark phenomena.

ii) It is difficult to gauge the reliability of perturbative QCD predictions in the absence of systematic higher order calculations, and the uncertainties due to the possible influence of non-perturbative dynamics.

iii) Predictions for inclusive hadron production processes are at present based on probabalistic models for quark and gluon jet hadronization, which by necessity contain ad hoc assumptions. It is not clear whether such predictions test QCD or the jet model Monte Carlo.

A prime example of the difficulty in testing QCD quantitatively is the fact that the coupling strength $\alpha_s(Q^2)$, has not been reliably determined to within 50% accuracy at any momentum scale; there is certainly no direct evidence that $\alpha_s(Q^2)$ decreases logarithmically with momentum transfer. In particular, the analysis of CELLO group at PETRA indicates that the value of $\alpha_s(Q^2)$ derived from $e^+e^- \to 3$ jet events is strongly sensitive to the particular model used to simulate quark and gluon jet fragmentation -- the Lund model gives values of α_s from 20% to more than 50% larger than that determined using the conventional Hoyer et al., Feynman-Field type models.[4] Similarly, predictions for energy flow correlations and asymmetries at present accessible energies, $Q^2 \sim 1000\ GeV^2$ are dependent on the model used for the non-perturbative jet hadronization.

The central problem for testing QCD (see Section 2) is that even though virtually all probes of the theory are done within the confines of hadrons or involve hadron production, we have very little knowledge of hadron wave functions or their effect on QCD processes. The traditional solution to this problem has been to develop tools such as the operator product expansion, factorization theories, and evolution equations which can provide tests of the theory independent of the form of the hadron wave functions. The applicability of this program is by necessity limited.

In this talk I will discuss a unified approach to QCD phenomenology in which the hadronic Fock state wave functions are the central denominator.[5,6] By assuming specific parameterizations of these wave functions,[6] one can extend the domain of QCD predictions to exclusive processes,[5] decay matrix elements, certain higher twist subprocesses,[7-13] "direct" hadron reactions (in which the hadrons themselves enter the hard scattering subprocesses is semi-inclusive reactions),[7-9] and possibly soft hadronic reactions and jet fragmentation. Another advantage of the approach is the fact that we can obtain new connections between different processes, in somes cases obtaining results independent of the form of the wave functions. Eventually one can hope to solve the QCD equations of motion [see Eq. (12)] for the hadronic wave function and present a complete unified phenomenology. An automatic method[14] for fixing the momentum scale in leading order QCD productions is briefly discussed in Section 4.

2. COMPLICATIONS IN TESTING QCD

It is in the nature of inclusive hadronic processes that virtually any QCD mechanism which can be drawn as a Feynman graph will contribute to the cross section at some level. Perturbative QCD, the operator product expansion, and the factorization ansatz are important guides to the dominant contributions for large momentum transfer reactions. However, the secondary effects are often not under good theoretical control because of the absence of rigorous bounds or because of parametrization uncertainties. On the other hand, many of these complicating processes constitute novel QCD effects and can be important tests of the theory.

There are a number of reasons why the precise determination of $\alpha_s(Q^2)$ from $e^+e^- \to q\bar{q}g$ jet events is intrinsically difficult. The primary problem is that the separation between 3-jet $q\bar{q}g$ and two-jet qq events requires detailed, certain knowledge of the transverse momentum $k_\perp$ and longitudinal light-cone fraction $z = (k^0 + k^3)/(p_q{}^0 + p_q{}^3)$ dependence of the quark fragmentation distributions $D(z,k_\perp) = dN/dz/d^2k_\perp$. The only theoretical input from perturbative QCD is at large $z \sim 1$ and/or large $k_\perp$ where the hadron wave functions are probed in the far off shell regime. One can show that (modulo logarithmic factors) (see figure 1)

$$D_{M/q}(z) = \frac{dN}{dz}(M/q) \underset{z\to 1}{\sim} A(1-z)^2 + \frac{C_M}{Q^2} \qquad (1)$$

where the C_M/Q^2 term is associated with the longitudinal current.[15] The high twist contribution to jet fragmentation from Fig. 1(b) is[6,13,14]

$$\frac{dN}{dk_\perp^2}(M/q) \sim \frac{dN}{dk_\perp^2}(g/q)\, F_M(k_\perp^2) \propto \frac{\alpha_s}{\pi k_\perp^2} F_M(k_\perp^2) \qquad (2)$$

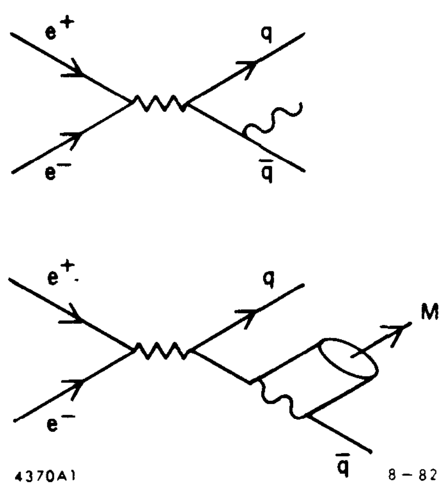

Fig. 1. QCD contributions to the quark fragmentation functions at large $k_\perp$. The direct meson contribution (b) gives a $1/k_\perp^4$ power-law tail normalized to the meson form factor.

for meson production at large $k_\perp$. The scale constant C_M can be computed in terms of the meson form factor [see Eq. (3)].[7] (Calculations also show that the distribution in $k_\perp$ and z does not factorize.) In contrast to the QCD forms, standard jet hadronization parametrizations usually assume that the transverse momentum distribution falls as $\exp(-bk_\perp^2)$, and that $D_{M/q}$ is non-vanishing at $z \to 1$ at large Q^2. The QCD form (2) for $q \to q+M$ predicts hadron production at relatively large transverse momentum, reducing the number of events which should be identified as $q\bar{q}g$; i.e., the $F_M(k_\perp^2)/k_\perp^2$ terms (summed over all meson states) can be a significant background to the $\alpha_s(k_\perp^2)/k_\perp^2$ gluon jet signal. The problem is compounded by heavy quark fragmentation uncertanties (e.g. $e^+e^- \to \bar{c}c \to \bar{c}\Lambda_c X \to \bar{c}pX'$), and the production of gluonium states. The important prediction of Bjorken[16] and Suzuki[17] that charmed hadrons are produced dominately at large z in the c-quark fragmentation region appears to be confirmed by e^+e^- annihilation and deep inelastic lepton scattering data.[18] Charm and beauty quark fragmentation thus could account for a substantial fraction of meson and baryon production at large $k_\perp$ and z, competing with hard gluon bremsstrahlung processes in e^+e^- annihilation.[19] There are also questions of a more fundamental nature which require an understanding of confinement and non-perturbative effects. The analysis of Gupta and Quinn[20] shows that the standard picture of jet hadronization would certainly break down if all quark masses were large compared to the QCD scale Λ. Thus there must be a hidden analytic dependence on the quark mass which controls jet fragmentation, and corrections to the jet cross section beyond that indicated by the perturbative expansion.

In the case of deep inelastic lepton scattering the central test of QCD is the evolution of the structure functions. The background effects include:

i) Higher twist contributions (terms suppressed by powers of $1/Q^2$ at fixed x). Such terms arise from mass insertions, $k_\perp$ effects, and coherent lepton-multiquark scattering contributions. Calculations of the latter contributions require knowledge of the multiquark distribution amplitudes of the target. However, absolutely normalized predictions can be obtained for the meson structure functions at $x \to 1$ since the required valence wave function matrix elements are already determined by the meson electromagnetic form factor. For example, the leading contribution at $x \to 1$ to the meson longitudinal structure function takes the form ($C_F = 4/3$)[7]

$$F_L{}^M(x,Q^2) = \frac{C_F}{2\pi}\,\frac{x^2}{Q^2}\int_{m^2}^{Q^2} dk^2\ \alpha_s(k^2)\ F_M(k^2)\left[1 + O\left(\frac{k^2(1-x)}{Q^2}\right)\right] \quad . \qquad (3)$$

Numerically this gives $F_L{}^M \sim .1x^2/Q^2$ in GeV^2 units. Notice that for $x \to 1$ and fixed Q^2 this contribution will dominate the transverse current meson structure function which is predicted to decrease as $(1-x)^2$ at $x \to 1$, modulo logarithmic factors. The higher twist F_L contribution comes from the direct interaction of the pion with the current; its QCD evolution is analogous to that of the photon structure function.

In the case of the nucleon structure functions, the leading twist contribution to F_2 at $x \to 1$ is predicted to vanish as $(1-x)^3$ or $(1-x)^5$ for quarks with helicity parallel or antiparallel respectively to the nucleon helicity.[15,6] A recent model calculation of the higher twist $[m^2/Q^2(1-x)^2]^{1,2}$ contributions to the nucleon function yields even larger effects than in the meson case. The analyses of Barnett et al.[21] show that present deep ielastic nucleon target data cannot unambiguously separate scale-violating QCD evolution effects in leading twist from the higher twist contributions allowed in QCD. The (remote) possibility that all the observed scale violation is due to higher twist effects or that Λ_{QCD} is very small is not ruled out.

ii) Heavy quark thresholds. As $W^2 = (q+p)^2$ is increased beyond the production threshold for new quark flavors, the structure functions increase at fixed x in a direction opposite to QCD evolution, with an attendant change in the characteristics of the final state. Strong scale-breaking effects asso-ciated with the charm threshold are seen directly in the EMC data[22] for the $c(x,Q^2) + \bar{c}(x,Q^2)$ distribution in the nucleon. As emphasized by D. P. Roy,[23] this effect can simulate an apparent cancellation of QCD evolution effects in deep inelastic structure functions. It is thus essential to accurately parametrize the quark mass effects. One can distinguish two components to the heavy quark distribution functions in the nucleon: (a) the "extrinsic" interaction dependent component generated by standard QCD evolution in association with the deep inelastic scattering, and (b) the "intrinsic" (i.e., initial condition) component generated by the QCD binding potential and equations of state for the proton.[24] The intrinsic component is maximal when all the quarks of the bound state have similar velocities, thus favoring large momentum fractions for the intrinsic c and $\bar{c}$ quarks. This leads to a valence-like distribution for the intrinsic charm quark distribution $c(x) \sim u(x)$ and a possible explanation of the charmed hadron distributions seen at the ISR in pp collisions. The present status of the intrinsic charm contribution is discussed in Section 5 and Ref. 24.

iii) Final state QCD interactions.[25] These effects do not affect the structure functions measured in deep inelastic scattering, but they do lead to changes in the particle distribution in the final state, e.g., $k_\perp$ smearing of the current quark distribution,[26] the production of associated hadrons in the central region, and the attendant degradation of the fast hadron momentum distribution. It is particularly interesting to study these effects in nuclear targets, using the nucleus to perturb the evolution of quark and gluon jets in hadronic matter. A general principle, the "formation zone"[27,28] (which can be derived in perturbative QCD) leads to the prediction that radiation collinear with the current quark cannot be induced during passage through the nucleus at high energies.[29] On the other hand, induced central region hadron production proportional to $A^{1/3}$ is allowed by QCD [see Fig. 2(a)]. (We note that such radiation, if verified, violates the usual assumption made in analyzing hadron-nucleus collisions that the induced central region radiation is always correlated with the number of "wounded" nucleons in the target, and it predicts cascading effects in the central region). [see Fig. 2(b).]

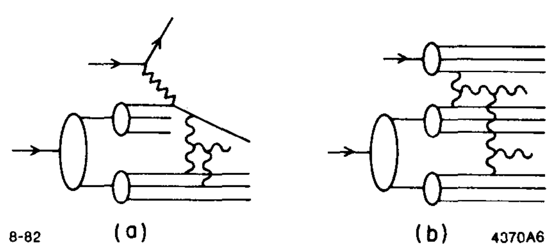

Fig. 2. (a) Production of central rapidity region multiplicity in association with final state interactions in deep inelastic lepton scattering on a nuclear target. (b) Production of central region multiplicity in hadron-nucleon collisions. The cascading interactions lead to a ramp shaped multiplicity distribution rather than a flat plateau in the central rapidity region (see Ref. 30).

iv) Non-additive contributions to the nuclear structure functions. Among the most important effects are shadowing (and possibly antishadowing) at low x and binding and kinematical corrections at $x \sim 1$. An argument that the shadowing due to traditional Glauber processes in suppressed at Q^2 large compared to a scale proportional to $1/x$ is given in Ref. 25. This is in contrast to standard fixed $W^2 = (q+p)^2$ duality arguments[31] which connects shadowing at low x with shadowing of the real photon photoabsorption cross section. In addition, QCD evolution itself may be non-additive in the nuclear number.[32] Non-additivity of the nuclear structure functions can also occur because of meson exchange currents (leading to the anomalous A dependence of the sea quark distributions), because of perturbations of the nucleon Fock states due to nuclear binding, or possibly because of "hidden color" components or six quark correlations[33] in the nuclear state. There is new evidence from the EMC experiment presented at this meeting by Montgomery[34] that the nuclear structure functions are not additive even at large x.

Each of the above QCD complications can lead to significant effects in the cross section for virtually any inclusive process. In the case of Drell-Yan process, initial state interactions[25] of the active quarks with spectators lead to (a) increased smearing of the $Q_\perp$ distribution of the lepton pair beyond that contained in QCD radiative corrections or the hadronic wave functions, (b) target-dependent induced radiation in the central region and the associated degradation of the quark longitudinal momentum distributions, and (c) a modification of the overall normalization of the pair product cross section $d\sigma/dQ^2dx$ due to induced color correlations, (at least at subasymptotic Q^2). The color correlation can have anomalous dependence on the nuclear number A relative to that measured at the corresponding kinematic range in deep inelastic scattering.

The dominance of the longitudinal structure functions in the fixed W limit for mesons [as in Eq. (3)] is an essential prediction of perturbative QCD. Perhaps the most dramatic consequence is in the Drell-Yan process $\pi p \to \ell^+\ell^- X$; one predicts that for fixed pair mass Q, the angular distribution of the ℓ^+ (in the pair rest frame) will change from the conventional $(1 + \cos^2\theta_+)$ distribution to $\sin^2(\theta_+)$ for pairs produced at large $x_L \to 1$. The results of the Chicago-Illinois-Princeton[35] experiment at FNAL appears to confirm the QCD higher twist contribution; but it is not seen in a recent SPS experiment.[36] Striking evidence for such an effect has also been seen in a Gargamelle[37] analysis of the quark fragmentation functions in $\nu p \to \pi^+\mu^- X$. The results yield a quark fragmentation distribution into positive charged hadrons which is consistent with the predicted form: $dN^+/dzdy \sim B(1-z)^2 + (C/Q^2)(1-y)$ where the $(1-y)$ behavior corresponds to a longitudinal structure function. It is also crucial to check that the $e^+e^- \to MX$[38] cross section becomes purely longitudinal ($\sin^2\theta$) at large z at moderate Q^2, and that the observed effects are not kinematical in origin, or due to backgrounds from baryon production.

All of the above QCD complications are compounded in reactions involving large transverse momentum particle or jet production in hadron-hadron collisions. Conventional calculations based on the leading $p_\perp^{-4}$ $(2 \to 2)$ subprocesses suffer severe model-dependent ambiguities due to uncertainties in how to include $k_\perp$ smearing effects, the size of scale breaking effects, and the strong parameter-dependence on the quark and gluon fragmentation processes. In these analyses it is usual to assume on-shell kinematics for the basic parton-parton cross section. However, when $k_\perp$ smearing from the hadron wave functions is introduced, this leads to a divergence at zero momentum transfer in the gluon propagator and thus an anomalous sensitivity to an arbitrary low momentum cutoff. The correct use of off-shell kinematics[39] at high $k_\perp$ removes this divergence and also much of the scale violations associated with $k_\perp$ smearing. In addition, in standard approaches, the $q \to M+q$ fragmentation function is <u>forced</u> to be non-vanishing at $z \to 1$ (e.g., $D_{M/q}(z) \sim (1-z)^2+C$ with $C \neq 0$); otherwise, one predicts more hadronic momentum collinear with the high p_T trigger hadron than that observed by experiment. In fact, as we have emphasized, QCD predicts that the only non-vanishing contributions to the fragmentation function at $z \to 1$ are due to higher twist subprocesses: specifically, direct subprocesses such as $gq \to Mq$ where the high p_T hadron is created in the short distance reaction instead of by fragmentation. In addition, the effects of initial and final state interactions ($k_\perp$ smearing, multiple scattering, color correlation effects, associated central region multiplicites) can severely complicate the model calculations.[40]

There is another serious difficulty with standard QCD phenomenology. If leading twist $2 \to 2$ subprocesses dominate direct photon and hadron production at large transverse momentum then one predicts a ratio $R_{\gamma/\pi} \sim f(x_\perp,\theta_{cm})$, independent of $p_\perp$ at fixed $x_\perp = 2p_\perp/\sqrt{s}$ and θ_{cm}. The ISR data reported to this meeting, however, is consistent with $R_{\gamma/\pi} \sim p_\perp^2$ at fixed $x_\perp$ and θ_{cm} at large $p_\perp$. The simplest interpretation of this result is simply that higher twist $p_\perp^{-6}$-scaling "direct" subprocesses such as $qg \to Mq$ and $q\bar{q} \to Mg$ dominate meson production the ISR kinematic regime whereas direct photons are produced by standard $p_\perp^{-4}$ $q\bar{q} \to \gamma g$ and $qg \to qg$ subprocesses. The cross section for the direct meson processes can be precisely normalized[7,41] in terms of the meson form factor since the same moments of the hadron wave function (distribution amplitude) appear. The direct processes can dominate the leading twist $qq \to qq$ and $qg \to qg$ subprocesses in the ISR regime because quark fragmentation into fast hadrons is not required; the meson M is made directly in the subprocesses. As in the case of direct photon production, the direct meson is unaffected by final state interactions[25] since a point-like component of the meson valence Fock state is involved (see Section 5). The direct processes also lead to significant quantum number correlations with the away-side jet since fermion exchange (rather than gluon exchange) plays an important role in the direct processes. Evidence for such strong correlations has been reported by the BFS and SFM collaborations.[42]

It is clear that a complete formulation of large transverse momentum hadron production which takes into account all the relevant QCD effects, including initial state interactions, $k_\perp$ smearing with off-shell kinematics, realistic fragmentation functions, as well as the higher twist direct subprocesses must be given before further progress can be made. In particular a careful separation of the leading twist and direct hadron production processes is required. It is certainly incorrect to force parametrizations such that all high p_T hadron production are attributed to the simplest quark or gluon scattering process.

An important clue to identifying underlying QCD subprocesses is the pattern of hadronic radiation produced in association with each inclusive reaction. The real difficulty in analyzing multiparticle reactions quantitatively in QCD is that we do not understand color confinement or even how an individual hadron is formed! To understand these problems at a basic level we will have to go beyond a statistical treatment of quark fragmentation and actually analyze hadron production at the amplitude level where coherent effects can be identified. It is clear even from electrodynamics that coherence is crucial for soft photon production; e.g., consider the case where two charged sources are nearly collinear.[43] It is also important to look in detail at exclusive QCD processes such as form factors, high momentum transfer Compton scattering and two photon reactions $\gamma\gamma \to M\bar{M}$ and $\gamma\gamma \to B\bar{B}$ at fixed θ_{cm} which give the simplest and most direct information on the production of individual hadrons.[44,45] It is also conceivable that the basic hadron wave function knowledge required for understanding quark and gluon fragmentation processes can be obtained from non-perturbative QCD calculations, such as lattice gauge theory, or the QCD equation of state on the light-cone. We will discuss this in more detail in Section 3.

3. HADRONIC WAVE FUNCTIONS IN QCD[6,46-48]

Even though quark and gluon perturbative subprocesses are simple in QCD, the complete description of a physical hadronic process requires the consideration of many different coherent and incoherent amplitudes, as well as the effects of non-perturbative phenomena associated with the hadronic wave functions and color confinement. Despite this complexity, it is possible to obtain predictions for many exclusive and inclusive reactions at large momentum transfer provided we make the standard ansatz that the effect of non-perturbative dynamics is negligible in the short-distance and far-off-shell domain. (This assumption appears reasonable since a linear confining potential $V \propto r$ is negligible compared to perturbative $1/r$ contributions.) For many large momentum transfer processes, such as deep inelastic lepton-hadron scattering reactions and meson form factors, one can then isolate the long-distance confinement dynamics from the short-distance quark and gluon dynamics -- at least to leading order in $1/Q^2$. The essential QCD dynamics can thus be computed from (irreducible) quark and gluon subprocesses amplitudes as a perturbative expansion in an asymptotically small coupling constant $\alpha_s(Q^2)$.

For example, the pion form factor at large Q^2 takes the form [see Fig. 3(b)][46]

$$F_\pi(Q^2) = \int_0^1 dx \int_0^1 dy\ \phi^*(x,Q)\ T_H(x,y;Q)\ \phi(y,Q) \qquad (4)$$

where

$$\phi_\pi(x,Q) = \int^Q \frac{d^2k_\perp}{16\pi^3}\ \psi_{q\bar{q}/\pi}{}^Q(x,\vec{k}_\perp) \qquad (5)$$

is the amplitude for finding the q and $\bar{q}$ in the valence state of the pion collinear up to scale Q with light cone longitudinal momentum fractions x and 1-x, and

$$T_H = \frac{16\pi\ C_F\ \alpha_s[Q^2(1-x)(1-y)]}{(1-x)(1-y)Q^2}\left[1 + \frac{O[\alpha_s(Q^2)]}{\pi}\right] \qquad (6)$$

is the probability amplitude for scattering collinear constituents from the initial to the final direction. (The superscript Q in $\psi_{q\bar{q}}{}^Q$ indicates that all internal loop in $\psi_{q\bar{q}}$ are to be cutoff at $k_\perp^2 < Q^2$.) The log Q^2 dependence of the distribution amplitude $\phi(x,Q)$ is determined by the operator product expansion on the light-cone or an evolution equation; its specification at subasymptotic momentum requires the solution to the pion bound state problem. The general form of $F_\pi(Q^2)$ is

$$F_\pi(Q^2) = \left| \sum_{n=0}^{\infty} a_n \log^{-\gamma_n} \frac{Q^2}{\Lambda^2} \right|^2 C_F \frac{\alpha_s(Q^2)}{Q^2}$$
$$\times \left[1 + O\left(\frac{\alpha_s(Q^2)}{\pi}\right) + O\left(\frac{m^2}{Q^2}\right)\right] \qquad (7)$$

where the γ_n are computable anomalous dimensions. Similar calculations determine the baryon form factors, decay amplitudes such as $\Upsilon \to B\bar{B}$[48] and fixed angle scattering processes (see Figs. 3 and 4) such as Compton scattering, photo-production, and hadron-hadron scattering, although the latter calculations are complicated by the presence (and suppression) of pinch singularities.[6,32] It is interesting to note that $\phi_\pi(x,Q^2)$ can be measured directly from the angular θ_{cm} dependence of the $\gamma\gamma \to \pi^+\pi^-$ and $\gamma\gamma \to \pi^0\pi^0$ cross sections at large s.[44] In addition, independent of the form of the mesons wave function we can obtain α_s from the ratio

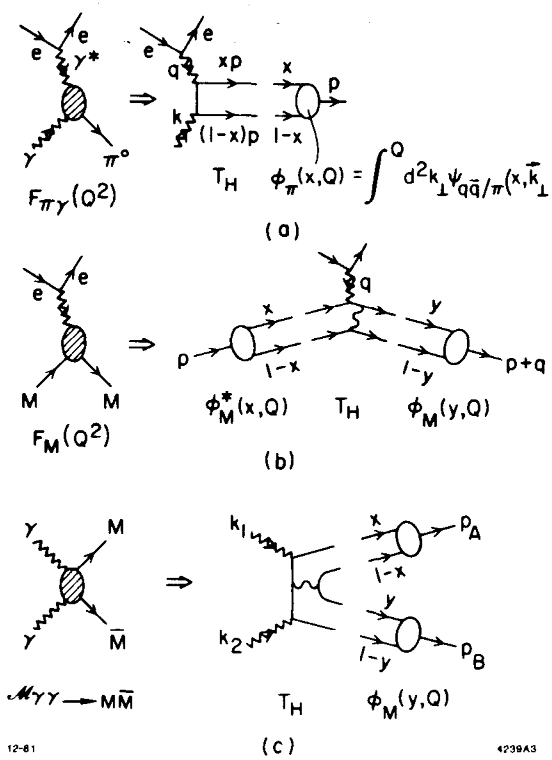

Fig. 3. QCD subprocesses and factorization for high momentum transfer exclusive processes. (a) The $\gamma \to \pi^0$ transition form factor, measureable in $ee \to ee\pi^0$ reactions. (b) The meson form factor in factorized form. (c) Contributions to the $\gamma\gamma \to M\bar{M}$ amplitude. A complete calculation of these processes to leading order in $\alpha_s(Q^2)$ is given in Ref. 44.

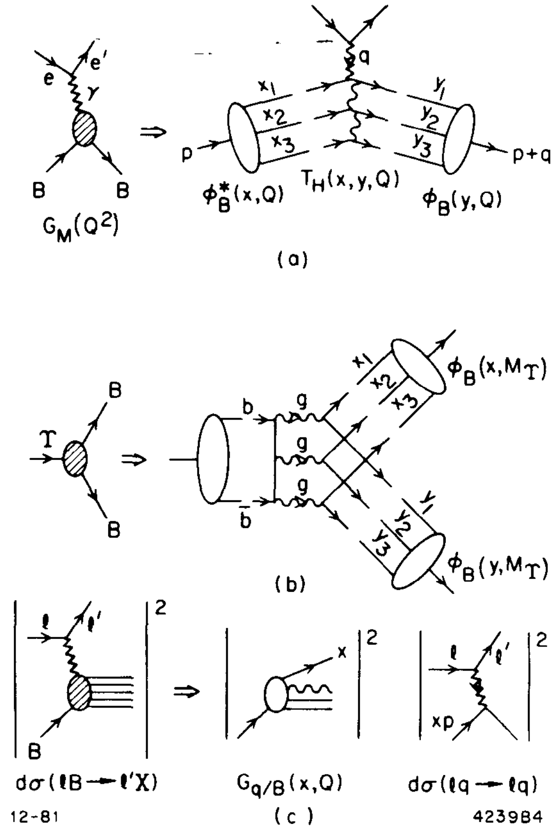

Fig. 4. Constraints on the baryon wave function in QCD. (a) Factorization of the baryon form factor (see Ref. 6). (b) Contribution to quarkonium decay into baryon pair (see Ref. 48). (c) Calculation of the deep inelastic scattering structure functions from light-cone wave functions.

$$\alpha_s(Q^2) = \frac{F_\pi(Q^2)}{4\pi Q^2 \left|F_{\pi\gamma}(Q^2)\right|^2} \left[1 + O\left(\frac{\alpha_s(Q^2)}{\pi}\right) \right] \qquad (8)$$

where the transition form factor $F_{\pi\gamma}(Q^2)$ can be measured in the two photon reaction $\gamma^*\gamma \to \pi^0$ via $ee \to \pi^0 ee$. Equation (8) is in principle one of the cleanest ways to measure α_s. The higher order corrections in α_s are discussed in Ref. 49.

Thus an essential part of the QCD predictions is the hadronic wave functions which determine the probability amplitudes and distributions of the quark and gluons which enter the short distance subprocesses. The hadronic wave functions provide the link between the long distance non-perturbative and short-distance perturbative physics. Eventually, one can hope to compute the wave functions from the theory, e.g., from lattice or bag models, or directly from the QCD equations of motion, as we shall outline below. Knowledge of hadronic wave function will also allow the normalization and specification of the power law (higher twist) corrections to the leading impulse approximation results.

The wave function $\psi_{q\bar{q}}^{\pi}(x,k_\perp)$ which appears in Eq. (5) is related to the Bethe-Salpeter amplitude at equal "time" $\tau = t + z$ on the light-cone in $A^+ = 0$ gauge. The quark has transverse momentum $k_\perp$ relative to the pion direction and fractional "light-cone" momentum $x = (k^0 + k^3)/(p^0 + p^3) = k^+/p^+$. The state is off the light cone $k^- = k^0 - k^3$ energy shell. In general a hadron state can be expanded in terms of a complete set of Fock state at equal τ:

$$|\pi\rangle = |q\bar{q}\rangle\, \psi_{q\bar{q}} + |q\bar{q}g\rangle\, \psi_{q\bar{q}g} + \ldots \tag{9}$$

with

$$\sum_n \int [d^2k_\perp]\, [dx]\, |\psi_n(x_i,\vec{k}_{\perp i})|^2 = 1 \quad .$$

(We suppress helicity labels.) At large Q^2 only the valence state contributes to an exclusive process, since by dimensional counting an amplitude is suppressed by a power of $1/Q^2$ for each constituent required to absorb large momentum transfer. The amplitudes ψ_n are infrared finite for color-singlet bound states. The meson decay amplitude (e.g. $\pi^+ \to \mu^+\nu$) implies a sum rule

$$\frac{a_0}{6} = \frac{1}{2\sqrt{n_c}}\, f_\pi = \int_0^1 dx\, \phi_\pi(x,Q) \quad . \tag{10}$$

This result, combined with the constraint on the wave function from $\pi^0 \to \gamma\gamma$ requires that the probability that the pion is in its valence state is $\lesssim 1/4$.[6,48] Given the $\{\psi_n\}$ for a hadron, virtually any hadronic properties can be computed, including anomalous moments, form factors (at any Q^2), etc.[48]

The $\{\psi_n\}$ also determine the structure functions appearing in deep inelastic scattering at large Q^2 ($a=q,\bar{q},g$)

$$G_{a/p}(x,Q) = \sum_n \int^Q [d^2k_\perp]\, [dx]\, |\psi_n^Q(x_i,k_{\perp i})|^2\, \delta(x-x_i) \tag{11}$$

where one must sum over all Fock states containing the constituent a and integrate over all transverse momentum $d^2k_\perp$ and the light-cone momentum fractions $x_i \neq x_a$ of the (n-2) spectators. The valence state dominates $G_{q/p}(x,Q)$ only at the edge of phase space, $x \rightarrow 1$. All of the multiparticle x and $k_\perp$ momentum distributions needed for multiquark scattering processes can be defined in a similar manner. The evolution equations for the $G_a(x,Q^2)$ can be easily obtained from the high $k_\perp$ dependence of the perturbative contributions to ψ.

There are many advantages obtained by quantizing a renormalizable local $\tau = t+z$. These include the existence of an orthornormal relativistic wave function expansion, a convenient τ-ordered perturbative theory, and diagonal (number-conserving) charge and current operators. The central reason why one can construct a sensible relativistic wave function Fock state expansion on the light cone is the fact that the perturbative vacuum is also an eigenstate of the full Hamiltonian. The equation of state for the $\{\psi_n(x_i,k_{\perp i})\}$ takes the form

$$H_{LC}\,\Psi = M^2\,\Psi \tag{12}$$

where

$$H_{LC} = \sum_{i=1}^{n}\left(\frac{k_\perp^2 + m^2}{x}\right)_i + V_{LC}$$

is derived[5,14] from the QCD Hamiltonian in $A^+ = 0$ gauge quantized at equal τ, and Ψ is a column matrix of the Fock state wave functions. Ultraviolet regularization and invariance under renormalization is discussed in Refs. 5, 6 and 14.

A comparison of the properties of exclusive and inclusive cross sections in QCD is given in Table I. Given the $\{\psi_n\}$ we can also calculate decay amplitudes, e.g. $\Upsilon \rightarrow p\bar{p}$ which can be used to normalize the proton distribution amplitudes [see Fig. 4(b)]. The constraints on hadronic wave functions which result from present experiments are given in Refs. 6, 14 and 48. An approximate connection between the valence wave functions defined at equal τ with the rest frame wave function is also given in Ref. 48, so that one can make predictions from non-perturbative analyses such as bag models, lattice gauge theory, chromostatic approximations, potential models, etc. Other constraints from QCD sum rules are discussed in Ref. 50. Applications to the high momentum transfer behavior of nuclear amplitudes are discussed in Ref. 51.

It is interesting to note that the higher twist amplitudes such as $\gamma q \rightarrow Mq$, $gq \rightarrow Mq$, $q\bar{q} \rightarrow MM$, $q\bar{q} \rightarrow Bq$ which are important for inclusive hadron production reactions at high $x_\perp$ can be absolutely normalized in terms of the distribution amplitudes $\phi_M(x,Q)$, $\phi_B(x_i,Q)$, by using the same analysis as used for the analysis of form factors.[5] In fact "direct" amplitudes

Table I. Comparison of exclusive and inclusive cross sections.

Exclusive Amplitudes	Inclusive Cross Sections
$\mathcal{M} \sim \Pi\, \phi(x_i,Q) \otimes T_H(x_i,Q)$	$d\sigma \sim \Pi\, G(x_a,Q) \otimes d\hat{\sigma}(x_a,Q)$
$\phi(x,Q) = \int^Q \left[d^2k_\perp\right] \psi^Q_{val}(x,k_\perp)$	$G(x,Q) = \sum_n \int^Q \left[d^2k_\perp\right][dx]' \lvert\psi^Q_n(x,k_\perp)\rvert^2$
Measure ϕ in $\gamma\gamma \to M\bar{M}$	Measure G in $\ell p \to \ell X$
$\sum_{i\epsilon H} \lambda_i = \lambda_H$	$\sum_{i\epsilon H} \lambda_i \neq \lambda_H$
	EVOLUTION
$\frac{\partial\phi(x,Q)}{\partial \log Q^2} = \alpha_s \int [dy] V(x,y)\phi(y)$	$\frac{\partial G(x,Q)}{\partial \log Q^2} = \alpha_s \int dy\ P(x/y)G(y)$
$\lim_{Q\to\infty} \phi(x,Q) = \prod_i x_i \cdot C_{flavor}$	$\lim_{Q\to\infty} G(x,Q) = \delta(x)\ C$
	POWER LAW BEHAVIOR
$\frac{d\sigma}{dx}(A+B \to C+D) \cong \frac{1}{s^{n-2}} f(\theta_{CM})$	$\frac{d\sigma}{d^2p/E}(AB \to CX) \cong \sum \frac{(1-x_T)^{2n_s-1}}{(Q^2)^{n_{act}-2}} f(\theta_{CM})$
$n = n_A + n_B + n_C + n_D$	$n_{act} = n_a + n_b + n_c + n_d$
T_H: expansion in $\alpha_s(Q^2)$	$d\hat{\sigma}$: expansion in $\alpha_s(Q^2)$
	COMPLICATIONS
End point singularities Pinch singularities High Fock states	Multiple scales Phase-space limits on evolution Heavy quark thresholds Heavy twist multiparticle processes Initial and final state interactions

such as $\gamma q \to Mq$, $q\bar{q} \to Mg$ and $gq \to Mq$ where the meson acts directly in the subprocess are rigorously related to the meson form factor since the same moment of the distribution amplitude appears in each case.

The light cone Fock state expansion also gives insight into the nature of forward hadron production in soft hadronic collisions. It should be emphasized that the properties of the strong interaction itself may distort the properties of the Fock state wave function in such a way that the quark distribution observed in soft collisions and fragmentation processes may differ significantly from that observed in deep inelastic scattering. For example, in the case of the DTU model,[52] at least one valence quark must be at low x in order to initiate the Pomeron "cylinder" interaction. One can also understand this "held back" feature if one assumes that only these Fock states which are very peripheral (possessing a large impact parameter, low x constituent) can interact strongly. The non-interacting constituents thus have more of the beam momentum on the average than they would have in undistorted wave functions. For example, counting rules[53] for renormalizable theories predict $G_{a/A}(x) \underset{x\to 1}{\sim} (1-x)^{2n_s-1}$ where n_s is the minimum number of spectators require to be stopped. If the hadronic wave function is undistorted in hadronic collisions, then one predicts, for example, $dN(pp \to K^+X)/dx \sim (1-x_L)^5$ (corresponding to 3 spectators). The power-behavior indicated by data however is $\sim(1-x_L)^3$. On the other hand if one quark in the incident hadron is required to be at $x \sim 0$ in order to have a large hadronic cross section, then there is one less spectator to stop. More generally, the held back mechanism reduces[54] the power-law fall off by $(1-x_L)^2$ and systematically gives results in good agreement with data. The above considerations also indicate that the impact space distribution of the Fock state wave function could be strongly correlated with the x distribution; for example, low x sea quarks may be predominantly peripheral in impact space. Such a correlation would be natural if the sea quarks are pictured as constituents of virtual mesons in the nucleon cloud. Other possible QCD mechanisms without the held back mechanism are discussed in Refs. 53 and 55.

4. THE PERTURBATIVE EXPANSION IN QCD

The essential predictive power of perturbative QCD is due to the smallness of the effective coupling $\alpha_s(Q^2)$ and the existence of perturbative expansions for subprocess amplitudes which control larger momentum transfer reactions. There has been significant progress in extending the QCD predictions beyond leading order for a number of inclusive and exclusive reactions.

However, a major ambiguity in the interpretation of these perturbation expansions and assessing the convergence of the series is in the choice of expansion parameter. Given a specific renormalization scheme, one must still specify the argument Q^2 of α_s in order to make a definite prediction from an expansion to finite order. As shown in a recent paper by Lepage, Mackenzie and myself, the scale of momentum Q^* transfer which appears in the leading order term in α_s (in a given renormalization scheme) in the perturbative expansions is not arbitrary, but is in fact, automatically set by the theory. For example, in QED, the running coupling constant $\alpha(Q)$ is defined to take into account the effects of lepton pair vacuum polarization in each photon propagator at $Q^2 \gg 4m_\ell^2$, and the approximate scale Q^* as the QED expansion in $\alpha(Q^*)$ is readily identified. In the case of QCD in the leading non-trivial order, $\alpha_s(Q)$ is only a function of $\beta_0 = 11-2/3\ n_f$. Thus we can use the analytic n_f dependence of the higher order corrections coming from light quark loop contributions to the gluon propagator to identify the correct scale of the α_s expansion in leading order, for almost every QCD processes.[14] The net result is that for almost every reaction except $\Upsilon \to 3g$, the perturbative expansion appears to have good convergence with reasonably small coefficients of α_s/π. An important example is

$$R_{e^+e^-} = 3 \sum_q e_q^2 \left[1 + \frac{\alpha^{\overline{MS}}(Q^*)}{\pi} + 0.0825(\alpha_s/\pi)^2 \right] \qquad (13)$$

where $Q^* = 0.71\ Q$. We can then use this result to define or measure $\alpha_s^R(Q) \equiv \alpha_s^{\overline{MS}}(0.71Q)$, and then write the perturbative expansion for other observables in terms of α_s^R. The QCD predictions for the deep inelastic moments and decay rates for the 0^- and 1^- quarkonium decays in terms of the first two nontrivial orders in α_s^R is given in Ref. 14.

5. NOVEL EFFECTS AND UNEXPECTED EFFECTS IN QCD

One can also test QCD by verifying novel effects which are essentially unique features of the theory. In this section we will list some examples of experimental tests which are specific to gauge theories of the strong interactions:

i) Hadron helicity conservation.

To leading order in $1/Q^2$ an exclusive process at large momentum transfer is dominated by amplitudes which conserve total hadron helicity; independent of photon (or weak boson) polarization. This is a consequence of the vector nature of gluon interactions and the fact that the total quark helicity equals the hadron helicity in the distribution amplitude ($L_z=0$). Many tests and predictions based on this rule are given in Ref. 56. An important prediction is that $\Upsilon \to p\bar{p}$ should have a $(1 + \cos^2\theta_{cm})$ angular distribution.

ii) Direct processes.

A surprising feature of QCD is the existence of inclusive processes such as $\pi_D p \to q\bar{q}X$ where the pion's energy and momentum are completely consumed in the large $p_\perp$ $q\bar{q}$ production; i.e., there is no associated hadron production in the forward fragmentation region [see Fig. 5(c)]. The jet cross section based on the $\pi g \to qq$ (and $\pi q \to gq$) subprocess is absolutely normalized in terms of the pion form factor (see Ref. 57); compared to the leading $qq \to qq$ subprocess one has

$$d\sigma(\pi \to q\bar{q}) \sim F_\pi(p_\perp^2)\, d\sigma(qq \to qq) \qquad (14)$$

independent of α_s and the pion wave function. This process is also identifiable by conservation of the $p^+ = p^0 + p^3$ components between the pion and jet fragments, and the close transverse momentum balance of the q and $\bar{q}$ jets. There is no bremsstrahlung or initial state interactions of the incident pion to leading order in $1/p_\perp^2$ since only the valence qq component of the pion wave function at $b_\perp \sim O(1/p_\perp)$ contributes to this process --

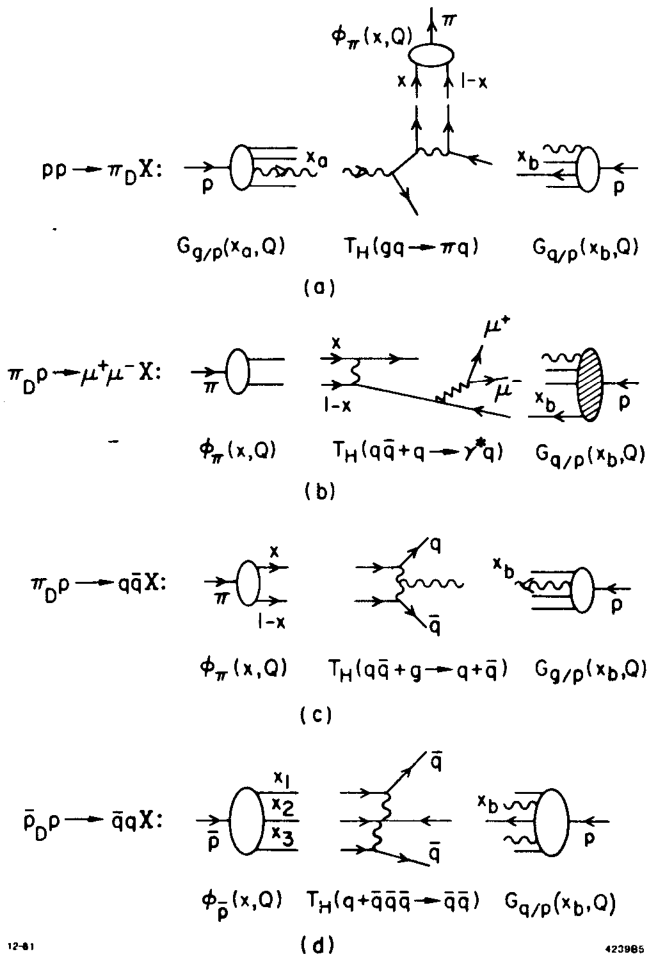

Fig. 5. Direct QCD subprocesses in inclusive high momentum transfer reactions. (a) The $gq \to Mq(p_\perp^{-6})$ contribution for producing high p_T mesons. (b) The $qM \to \ell\bar{\ell}q(Q^{-2})$ contribution to massive lepton pair production. (c) The $gM \to q\bar{q}(p_\perp^{-6})$ contribution to high p_T jet production. No hadrons are produced in the meson beam fragmentation region. (d) The $q\bar{p} \to q\bar{q}(p_\perp^{-8})$ contribution to the $p\bar{p} \to q\bar{q}X$ jet cross section. No hadrons are produced in the $\bar{p}$ fragmentation region.

such a state has negligible color (hadronic) interactions. Similarly, in the case of nuclear targets $\pi_D A \to q\bar{q}X$, the valence pion state at $b_\perp \sim 1/p_\perp$ penetrates throughout the nuclear volume without any nuclear interactions. The corresponding direct baryon induced reaction based on $\bar{p}q \to \bar{q}\bar{q}$ is shown in Fig. 5(d). We also note that the amplitude to find three quarks in the proton at small separation with one accompanying spectator meson can be measured by an analogous process $\bar{p}_D p \to \bar{q}\bar{q}\pi X$. Such amplitudes are of interest for predictions of baryon decay in grand unified theories.

In the case of hadron production at large $p_\perp$, the direct subprocess $gq \to \pi_D q$ produces pions unaccompanied by other hadrons on the trigger side [see Fig. 5(a)]. These $p_\perp^{-6}$ QCD processes are absolutely normalized in terms of the meson form factor and can dominate jet fragmentation processes at large $x_\perp$. More generally, an entire set of hadrons and resonances can be produced by direct subprocesses (e.g., $qq \to Bq$ and $qB \to qB$ for baryon production). Such contributions provide a serious background to any measurement where there is a high $p_\perp$ single particle trigger. Again, the direct process hadrons have no final state interactions or accompanying collinear radiation to leading order in $1/p_T^2$.

iii) Quasi-elastic reactions in nuclei and color transparency.

As we have noted in Section 3, large momentum transfer exclusive reactions are dominated (to leading order in $1/p_\perp^2$) by valence Fock states with small constituent separation. Since such states have negligible hadronic interactions, a large momentum transfer quasi exclusive reaction can take place inside of a nuclear target [e.g. $\pi A \to \pi p(A'-1)$] without any elastic or inelastic initial or final state hadronic interaction.[58] The rate for such "clean" reactions is normalized to A times the nucleon target rate.

iv) Diffractive dissociation and color filtering.

The existence of the light cone Fock state expansion for a meson implies a finite probability for the hadron to exist as a valence state at small relative impact parameters. This component of the state will interact only weakly in nuclear matter, whereas the majority part of the Fock state structure will interact strongly. The nucleus thus acts as a "color filter" absorbing all but the weakly interacting Fock components. The consequence is a computable cross section for the diffractive dissociation by a nucleus of pion into relatively high $k_\perp$ $q\bar{q}$ jets.[59] The rate is normalized to the pion decay constant and $\sigma_{\pi A}^{e\ell}$. The θ_{cm} dependence of the jet in the $q\bar{q}$ rest frame is related to the pion distribution amplitude. Predictions for the $k_\perp$ dependence of the jets are given in Ref. 59. All of the effects (2)-(4) test a basic feature of QCD: the Fock state structure of the color singlet hadronic wave functions and the special features of the valence Fock state.

v) Initial state interactions and color correlations.

A detailed discussion of the expected effects of initial and final state interactions in a non-Abelian theory is given in Refs. 25 and 60. The main novel effects are subasymptotic color correlations (in principle A-dependent), and associated production in the central region. It is also possible that the soft particles produced in an initial state interaction could subsequently interact with other valence quark in the beam hadron to produce low mass lepton pairs ($Q^2 < m\sqrt{s}$), low $p_\perp$ jets, etc. (see Fig. 6).

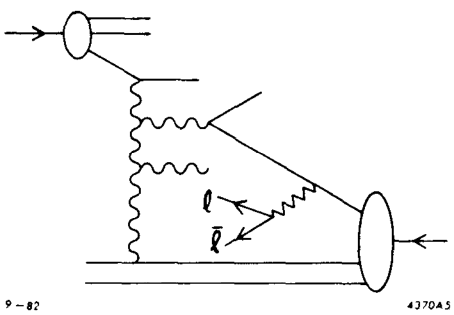

Fig. 6. Intermediate mass lepton pairs produced by the annihilation of a valence quark with a central region $\bar{q}$ produced in association with an initial state collision. Such contributions have anomalous nuclear number and energy dependence.

vi) Hadron multiplicity in e^+e^- collisions.

In a remarkable calculation, Basetto, Ciafaloni, Marchesini, and Mueller[61] have shown from an all orders perturbative analysis that QCD predicts a dip for $dn/dy_{cm} \sim 0$ for particle production in e^+e^- annihilation. This result needs careful experimental confirmation. Although the QCD pre-diction seems special to gauge theory, it should be noted that the statistical model of Ochs[62] based on a simple branching processs also has this feature.

vii) The perturbative coupling strength of gluons to a gluon jet is 9/4 larger than the coupling to a quark jet.[63] The deviation from a factor of 2 is due to color coherence. Perturbative QCD thus evidentally predicts that at some level gluon jets are broader and have a higher multiplcity plateau than quark jets. The fact that gluon jets can be screened by gluonium production may lead to further differences between quark and gluon jets. Polarized gluon jets may also have distinctive features such as oblateness.[63]

Perhaps the most interesting and important experimental observations are the phenomena not readily explainable or predicted by conventional QCD analyses. We will briefly discuss four examples here:

i) High transverse energy, high multiplicity events.

It is possible that this phenomena, first observed by the NA5 collaboration at the SPS, is a collective or fireball effect, or even a signal of a quark-gluon phase at high temperature. Field, Fox and Kelly[64] have attempted to identify these events with the conventional processes, e.g., quark-quark scattering, etc. subprocesses accompanied by multiple hard gluon radiation and subsequent neutralization. However, realistic calculation taking into account energy momentum correlations and interference effects (color correlations, formation zone) between multiple strings seems very difficult. An alternative explanation is that multiple quark/gluon scattering (Glauber) processes[65] are involved at high transverse energies, since at $x_T = E_T/E_{max} \sim 1$ the scattering of all the beam momentum to the transverse direction is required. Using counting rules, the pp → jet production cross section receives contributions of nominal order (R is the hadron transverse size)

$$E \frac{d\sigma}{d^3p} \sim \frac{\alpha_s^2}{p_T^4}(1-x_T)^7 \ , \ \frac{\alpha_s^4}{R^2 p_T^6}(1-x_T)^3 \ , \ \frac{\alpha_s^6}{R^4 p_T^8}(1-x_T)^{-1} \ , \qquad (15)$$

corresponding to the 2 particle, 4 particle, and 6 particle scattering processes shown in Fig. 7. A realistic calculation requires consideration of all the q and g multiparticle subprocesses, account of scale breaking, etc. The multiscattering terms can clearly dominate at x_T large; they are also characterized by high multiplicity, long-range correlations in rapidity and absence of coplanarity.

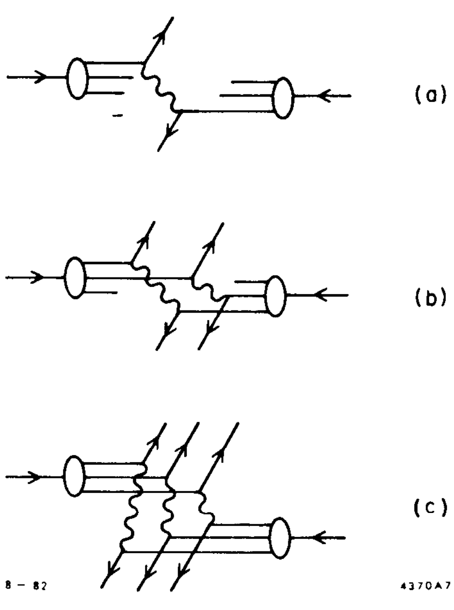

Fig. 7. QCD contributions to high p_T processes. The multiparticle 2+2 and 3+3 reactions produce high transverse energy events with a large fraction of the available energy. More complicated reactions with gluons and sea quarks are not shown.

ii) Heavy flavor production in hadron collsions.

Standard QCD calculations based on $gg \rightarrow Q\bar{Q}$ subprocesses do not account for the x_L dependence, diffractive character, and magnitude of forward charm hadron production observed at the ISR. It is also difficult to reconcile the large cross section ($\sigma_{charm} \sim 1$ mb) observed at the ISR with the much smaller cross sections observed at Fermilab energies.

One possible mechanism[66] for fast charm production in forward pp collisions is that the high momenta of the ud spectators in a nucleon, combined with a centrally produced charmed quark will yield a high momentum $\Lambda_c = |udc\rangle$. This is however contrary to the Bjorken-Suzuki effect, which indicates that the Λ_c should have roughly the same momentum as the charmed quark. This explanation also implies a long-range rapidity separation between the c at $y_{cm} \sim 0$ and a charmed hadron in the fragmentation region; it also does not account for the observed $(1-x_L)^3$ distribution of the $D^+ = |cd\rangle$ meson (unless such mesons always arise from charm baryon decay). Measurements of the x_L distribution of the $\bar{\Lambda}_c$ would be decisive.

The simplest explanation of the observed ISR phenomena is that the proton contains a Fock state $|uudc\bar{c}\rangle$ with the $c\bar{c}$ bound inside the hadron over a relatively long time scale $O(m_c^{-1})$.[24] Such a state (analogous to a $\Lambda_c\bar{D}$ component of the nucleon Fock state) could be diffractively dissociated into charmed hadrons in the forward fragmentation region as well as producing valence-like c and $\bar{c}$ distributions in deep inelastic lepton scattering (see Section 2). If the model is correct, b and t quarks can also be produced at large x_L with substantial cross sections, scaling as $1/M_Q^2$. Such states should contain relatively large transverse momentum--which together with the large quark mass, implies high $k_\perp$ leptons and baryon production; this would also yield a clear signal for t-quark production.

The origin of the intrinsic heavy quark state in the nucleon wave function is the heavy quark loop vacuum polarization (and light-by-light scattering) insertions proportional to $1/M_Q^2$ in the QCD potential. As in the case of the potential which appears in the evolution equation for the distribution amplitude,[5] such contributions do not have a dx/x singularity at small x. The peak of the light-one distribution is at $x_i \sim m_i/\sum_j m$ where the state is minimally off-shell. This corresponds to the fact that in the rest system of the proton, the virtual $Q\bar{Q}$ pair is produced dominantly at threshold, at low velocities. This even to lowest order in $1/M_Q^2$ the heavy quarks are produced with a distribution which peaks at large x_Q. There is thus no question that Fock states containing heavy quarks exist at some level in ordinary hadrons -- the real question is the absolute normalization. The vacuum polarization contributions give $P_{Q\bar{Q}} \sim O\alpha_s^2(\lambda^2)\lambda^2/M_Q^2$ where λ^2 is a typical hadronic scale. The bag model calculation of Donoghue and Golowich[67] gives $P_{c\bar{c}} \sim O$ (1%). One also could hope to relate the normalization of $P_{Q\bar{Q}}$ to the spin-spin splitting in the baryon system.

The origin and normalization of the heavy quark Fock state components and their systematic dependence or the hadron or nuclear state can yield important clues to the nature of hadron wave functions in QCD. Similarly, the production of heavy quark states by diffractive excitation of the intrinsic heavy quark states or a similar mechanism, should also give important insight into the nature of hadronic processes in forward high energy collisions.

iii) The observation of copious baryon production in e^+e^- annihilation is somewhat of a surprise from the standpoint of perturbation QCD ideas. For $z \to 1$ baryon production should be suppressed by at least a power of $(1-z)$ relative to meson production whereas the data indicates a flat baryon/meson ratio out to $z = 0.5$. In the Lund model[68] copious baryon production is accounted for by effective diquark production in a non-perturbative tunneling model. Another possiblity, suggested by T. DeGrand,[69] is that the fast charmed and beauty baryons produced at large z by the Bjorken-Suzuki mechanism leads by decay to a significant fraction of the large x_L baryons.

6. Conclusions

Despite the formal simplicity of its underlying Lagrangian, QCD has turned out to be an extraordinarly complex theory. Definitive tests have turned out to be very difficult, but there is at present no reason to doubt that QCD is at the basis of all hadron and nuclear dynamics accessible at present energies. Fortunately, as experiment has become more definitive, theoretical analyses has also made definite progress. The goal is to make precise predictions, with systematic control of background effects (such as high twist contribution, threshold effects, initial and final state interactions),-as well as to attain a deeper understanding of jet hadronization and the QCD perturbation expansion.

Much of the prevent uncertainty in QCD predictions is due to the absence of detailed information on hadronic matrix elements. It seems likely that direct calculations of hadronic amplitudes will be possible using lattice gauge theory, or as discussed in Section 3, by solving the QCD equation of state for the Fock states of hadrons at equal $\tau = t+z$. This formalism gives a consistent relativistic wave function basis and calculational framework for QCD. One can also use this formalism to obtain many new predictions which are in principle exact, e.g., large momentum transfer exclusive reactions [such as the meson form factors and $\gamma\gamma \to M\bar{M}$], new QCD constraints such as hadron helicity conservation, constraints on wave functions from decay amplitudes, new methods to determine α_s, and methods to calculate the higher twist and direct subprocesses.

We have also seen that QCD diverges in many ways from the expectations of the parton model. Many of the novel phenomena discussed here such as color correlations in initial state interactions, color filtering and transparency, intrinsic heavy quark Fock states, associated production in hard collisions, color coherence, direct processes, the dominant longitudinal component to the meson structure function, and many nuclear effects[69] (reduced form factors, anomalous A dependence) were not anticipated within the parton model framework. In addition, the basic hard scattering mechanism for form factors and the QCD breakdown of the exclusive-inclusive connection is contrary to the mechanisms assumed to be dominant in the parton model framework.

We have also emphasized here the importance of understanding the physics of initial and final state interactions in QCD, particularly the breakdown of factorization at large target length, the physics of the formation zone in QCD, and the interesting effects of color correlations. An important clue toward understanding these phenomena as well as the propagation of quark and gluon jets through hadronic matter will be the careful study of nuclear target effects.

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