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ANTIPROTON ANNIHILATION IN QUANTUM CHROMODYNAMICS*

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1. INTRODUCTION

Quantum chromodynamics^[1] has been extensively tested in high-momentum transfer inclusive reactions, where factorization theorems, asymptotic freedom and jet algorithms provide semiquantitative perturbative predictions. Tests of the confining nonperturbative aspects of the theory are either quite qualitative or at best indirect. In fact QCD is a theory of relatively low mass scales ($\Lambda_{MS} \sim 200 \pm 100$ MeV, $(k_{\perp}^2)^{1/2} \sim 300$ MeV, etc.) and its most critical tests of the theory as a viable theory of strong and nuclear interactions must involve relatively low energies and momentum transfer.

Anti-proton annihilation has a number of important advantages as a probe of QCD in the low energy domain. Exclusive reaction in which *complete* annihilation of the valance quarks occur ($\bar{p} \rightarrow \ell\bar{\ell}, \gamma\gamma, \phi\phi$, etc.) necessarily involve impact distances $b_{\perp}$ smaller than $1/M_p = 5 \text{ fm}^{-1}$ since baryon number is exchanged in the t -channel. There are a number of exclusive and inclusive $\bar{p}$ reactions in the intermediate momentum transfer domain which provide useful constraints on hadron wavefunctions or test novel features of QCD involving both perturbative and nonperturbative dynamics. In several cases ($\bar{p} \rightarrow \ell\bar{\ell}, \bar{p} \rightarrow J/\psi, \bar{p} \rightarrow \gamma\gamma$), complete leading twist (leading power law) predictions are available. These reactions not only probe the subprocesses $\bar{q}q\bar{q}qq \rightarrow \gamma\gamma$, etc., but they also are sensitive to the normalization and shape of the proton distribution amplitude $\phi_p(x_1, x_2, x_3; Q)$, the basic measure of the proton's three-quark valance wavefunction. Additionally, one can explore such processes in terms of quasielastic reactions inside of nuclear targets, e.g., $\bar{p}A \rightarrow (J/\psi) (A-1)$, and study an extraordinary feature of QCD: "color transparency." There is another class of exclusive reactions in QCD involving light nuclei, such as $\bar{p}d \rightarrow \gamma n$ and $\bar{p}d \rightarrow \pi^- p$ which can probe quark and gluon degrees of freedom of the nucleus at surprisingly low energy. These will be discussed in sec. 11.

Inclusive reactions involving antiprotons have the advantage that the parton distributions are well understood. In these lectures, I will particularly focus on lepton pair production $\bar{p}A \rightarrow \ell\bar{\ell}X$ as a means to understand specific nuclear features in QCD, including collision broadening, breakdown of the QCD "target length condition." Thus studies of low to moderate energy antiproton reactions with laboratory energies under 10 GeV could give further insights into the full structure of QCD.

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2. QCD TESTS AND HADRON LIGHT-CONE WAVEFUNCTIONS

QCD has two essential properties which make calculations of processes at short distance or high-momentum transfer tractable and systematic. The critical feature is asymptotic freedom: the effective coupling constant $\alpha_s(Q^2)$ which controls the interactions of quarks and gluons at momentum transfer Q^2 vanishes logarithmically at large Q^2 . Complementary to asymptotic freedom is the existence of factorization theorems for both exclusive and inclusive processes at large momentum transfer. In the case of exclusive processes (in which the kinematics of all the final state hadrons are fixed at large invariant mass), the hadronic amplitude can be represented as the product of a hard-scattering amplitude for the constituent quarks convoluted with a distribution amplitude for each incoming or outgoing hadron.^[12-14] The distribution amplitude contains all of the bound-state dynamics and specifies the momentum distribution of the quarks in the hadron.^[21] The hard-scattering amplitude can be calculated perturbatively as a function of $\alpha_s(Q^2)$. The analysis can be applied to form factors, exclusive photon-photon reactions, photoproduction, fixed-angle scattering, etc. In the case of the simplest processes, $\gamma\gamma \rightarrow M\bar{M}$ and the meson form factors, rigorous all-order proofs can be given. As we shall see, many of these predictions are directly applicable to antiproton-initiated reactions.

The predictions of perturbative QCD have been strikingly confirmed in inclusive e^+e^- and $\gamma\gamma$ collisions, deep inelastic lepton reactions, massive lepton pair production, and the whole array of large p_T jet and photon reactions. Measurements of exclusive processes at high-momentum transfer, especially form factors and two-body photon-photon reactions have led to detailed checks on the scaling behavior of the theory. Recent results^[17] for $\gamma\gamma \rightarrow M\bar{M}$ are shown in fig. 1. In general, the experimental results on the scaling behavior of exclusive and inclusive reactions appear consistent with short-distance subprocesses based on the elementary scattering of spin 1/2 quarks and spin 1 gluons, the fundamental degrees of freedom of QCD.

The key to understanding hadronization and hadron matrix elements is the hadron wavefunction itself. A convenient description of hadron wavefunctions is given by the set of n-body momentum space amplitudes, $\psi_n(x_i, k_{\perp i}, \lambda_i)$, $i = 1, 2, \dots, n$, defined on the free quark and gluon Fock basis at equal "light-cone time" $\tau = t + z/c$ in the physical "light-cone" gauge $A^+ \equiv A^0 + A^3 = 0$. (Here $x_i = k_i^+/p^+$, $\sum x_i = 1$, is the light-cone momentum fraction of quark or gluon i in the n — particle Fock state; $k_{\perp i}$, with $\sum k_{\perp i} = 0$, is its transverse momentum relative to the total momentum p^μ ; and λ_i is its helicity.) The quark and gluon structure functions $G_{q/H}(x, Q)$ and $G_{g/H}(x, Q)$ which control hard inclusive reactions and the hadron distribution amplitudes $\phi_H(x, Q)$ which control hard exclusive reactions are simply related to these wavefunctions:

$$G_{q/H}(x, Q) = \sum_n \int \prod d^2 k_{\perp i} \int \prod dx_i |\psi_n(x_i, k_{\perp i})|^2 \delta(x_q - x) \quad ,$$

and

$$\phi_H(x_i, Q) = \int \prod d^2 k_{\perp i} \psi_{\text{valence}}(x_i, k_{\perp i}) \quad .$$

Thus an important tool is the use of light-cone quantization to construct a consistent relativistic Fock state basis for the hadrons and their observables in terms of quark and gluon quanta. The distribution amplitudes and the structure functions are defined directly in terms of these light-cone wavefunctions.^[18] The form factor of a hadron can be computed exactly in terms of a convolution of initial and final light-cone Fock state wavefunctions.^[19] In the case of inclusive reactions all of the

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Fig. 1. Comparison of $\gamma\gamma \rightarrow \pi^+\pi^-$ and $\gamma\gamma \rightarrow K^+K^-$ meson pair production data with the parameter free perturbative QCD prediction of ref 2. The data are from ref 7.

hadron Fock states generally participate; the necessity for higher-particle Fock states in the proton is apparent from its large gluon momentum fraction and the recent results from the EMC collaboration^[9] suggesting that, on the average, little of the proton's helicity is carried by the light quarks.^[10] In the case of high-momentum transfer Q exclusive reactions perturbative QCD predicts that only the lowest particle number (valence) Fock state contributes to leading order in $1/Q$. The essential gauge-invariant input is the distribution amplitude^[11] $\phi_H(x, Q)$. Its dependence in $\log Q$ is controlled by evolution equations derivable from perturbation theory^[12] or the operator product expansion.^[13] A more detailed discussion of the light-cone Fock state wavefunctions and their relation to observables is given in ref. 13.

The phenomenology of hadron wavefunctions in QCD is now just beginning. Constraints on the baryon and meson distribution amplitudes have been recently obtained using QCD sum rules and lattice gauge theory. The results are expressed in terms of gauge-invariant moments $\langle x_j^m \rangle = \int \Pi dx_i x_j^m \phi(x_i, \mu)$ of the hadron's distribution amplitude.

A particularly important challenge relevant to antiproton exclusive processes is the construction of baryon distribution amplitudes. A three-dimensional "snapshot" of the proton's uud wavefunction at equal light-cone time as deduced from QCD sum rules at $\mu \sim 1$ GeV by Chernyak et al.^[14] is shown in fig. 2. The moments

of the proton distribution amplitude computed by Chernyak et al., have now been confirmed in an independent analysis by Sachrajda and King.^[11]

The moments of distribution amplitudes can also be computed using lattice gauge theory.^[12] In the case of the pion distribution amplitudes, there is good agreement of the lattice gauge theory computations of Martinelli and Sachrajda^[13] with the QCD sum rule results. This check has strengthened confidence in the reliability of the QCD sum rule method, although the shape of the meson distribution amplitudes are unexpectedly structured: the pion distribution amplitude is broad and has a dip at $x = 1/2$. In the case of the proton, the QCD sum rule prediction suggests that the u quark with helicity parallel to the proton helicity carries nearly 2/3 of the momentum in the three-quark valence Fock state. In fact, the QCD sum rule distributions, combined with the perturbative QCD factorization predictions, account well for the scaling, normalization of the pion form factor and also the branching ratio for $J/\psi \rightarrow p\bar{p}$. In addition, Maina has found that the data for large angle Compton scattering $\gamma p \rightarrow \gamma p$ are also well described.^[14] However, a very recent lattice calculation of the lowest two moments by Martinelli and Sachrajda^[15] does not show skewing of the average fraction of momentum of the valence quarks in the proton.

These initial results are interesting — suggesting a highly structured oscillating momentum-space valence wavefunctions. The results from both the lattice calculations and QCD sum rules demonstrate that the light quarks are highly relativistic, at least in mesons. This gives further indication that while nonrelativistic potential models are useful for enumerating the spectrum of hadrons (because they express the relevant degrees of freedom), they may not be reliable in predicting wave function structure.

Fig. 2. The proton distribution amplitude $\phi_p(x_i, \mu)$ determined at the scale $\mu \sim 1 \text{ GeV}$ from QCD sum rules by Chernyak, Ogloblin and Zhitnitski.

The sum rule model form for the nucleon distribution amplitude together with the QCD factorization formulae, predicts the correct sign and magnitude as well as scaling behavior of the proton and neutron form factors.^[11] (See fig. 3.)

Fig. 3. Comparison of the scaling behavior of the proton magnetic form factor with the theoretical predictions of refs. 2 and 4. The CZ predictions^[4] are normalized in sign and magnitude. The data are from ref. 19.

A new nonperturbative method “discretized light-cone quantization,” (DLCQ)^[30] has been developed which has the potential for providing detailed information on all the hadron’s Fock light-cone components. The basic idea is to diagonalize the QCD Hamiltonian on the light-cone Fock states, using a computationally-convenient discrete momentum space basis. The eigenvalues M^2 of H_{LC} provide the spectrum of the theory; the eigenvectors yield the Fock state wavefunctions $\psi_n(x_i, k_{\perp i}, \lambda_i)$. So far the method has been successfully applied to gauge theories and Yukawa theory (scalar gluons) in one-space and one-time dimension. New results for the spectrum and wavefunctions for QCD[1+1] are presented in sec. 15.

3. THE TIME-LIKE PROTON FORM FACTOR

As an introduction to the application of QCD to exclusive antiproton reactions, I will first review the leading-twist perturbative QCD predictions for $\bar{p}p \rightarrow \bar{\ell}\ell$, i.e., the time-like proton form factors $F_1(s)$ and $F_2(s)$. According to the QCD factorization analysis for exclusive processes, the main dynamics is contained [to leading order in $1/s$, $s = (\bar{p} + p)^2$] in the $\bar{q}\bar{q}\bar{q} qqq \rightarrow \gamma^* \rightarrow \bar{\ell}\ell$ amplitude. To leading order in $\alpha_s(Q^2)$ the latter can be computed from minimally connected PQCD tree graphs containing two off-shell gluons and two off-shell quark lines. (See fig. 4.) To leading order in $1/s$, it is sufficient to compute the subprocess amplitude taking the incident antiquarks and quarks collinear with their respective incident hadron direction, $p_i^\mu = x_i p^\mu$. One easily finds $T_H \sim \alpha_s^2(Q^2)/(Q^2)^2 f(x_i, y_j)$, where f is a rational function of the momentum fractions. By definition, higher-loop corrections to T_H are computed such that in intermediate states all quark propagators are noncollinear, and thus off-shell by order Q . This generates corrections of higher order in $\alpha_s(Q^2)$, with no additional logarithms. The contributions where partons in the intermediate state are collinear yields the evolution of the proton and antiproton distribution amplitudes.

After computing the hard-scattering amplitude T_H to the desired order in $\alpha_s(Q^2)$, one then convolutes the amplitude with the distribution amplitudes of the $\bar{p}$ and p . The distribution amplitude $\phi_p(x_i, s)$, ($\sum_{i=1,2,3} x_i = 1$) is a nonperturbative wavefunction, but its logarithm dependence on s is computable from first principles (e.g.,

Fig. 4. Calculation of electromagnetic $\bar{p}p$ annihilation in PQCD.

renormalization group plus operator product expansion) since this is controlled by the large momentum tail of the wavefunction:

$$\frac{\partial}{\partial \ln s} \phi(x_i, s) = \int_0^1 dy_i V(y_i, x_i, \alpha_s(s)) \phi(y_i, s) .$$

This equation is of the form of an evolution equation where $\ln s$ plays the role of the "time." The solution has the form

$$\phi(x_i, s) = \sum_n C_n(x_i) \left(\ln \frac{s}{\Lambda^2} \right)^{\gamma_n} ,$$

where the $C_n(x_i)$ are a known set of orthonormal polynomials. The γ_n are computable fractional powers, anomalous dimensions characteristic of the interpolating local operators for three quarks with the proton quantum numbers.

Thus the $\bar{p}p \rightarrow \ell\bar{\ell}$ time-like Fock Dirac factor (helicity conserving) has the form

$$F_1(Q^2) = \left[\frac{\alpha_s(Q^2)}{Q^2} \right]^2 \sum_{n,m} b_{nm} \left(\frac{\ln s}{\Lambda^2} \right)^{-\gamma_n - \gamma_m} .$$

The coefficients $b_{n,m}$ reflect the nonperturbative input and are determined by the initial data for $\phi(x_i, s_0)$ in the evolution equation. The power law scaling, $F_1(s) \sim 1/s^2$, and $d\sigma/dt (\bar{p}p \rightarrow \ell\bar{\ell}) \sim 1/s^6$ are consistent with the quark counting rules $F(Q^2) \sim (1/Q^2)^{n-1}$, $(d\sigma/dt)_{AB \rightarrow CD} \sim f(\theta_{cm})/s^{N_{tot}-2}$ where n is the number of fields in the bound state, and N_{tot} is the total number of incident and outgoing fields in $A+B \rightarrow C+D$; e.g., ($n=3$, $N_{tot}=8$ for $\bar{p}p \rightarrow \ell\bar{\ell}$).

It is important to note that the leading power-law behavior originates in the minimum three-particle Fock state of the $\bar{p}$ and p , at least in physical gauge, such as $A^+ = 0$. Higher Fock states give contributions higher order in $1/s$. For $\bar{p}p \rightarrow \ell\bar{\ell}$

this means that initial-state interaction such as one gluon exchange are dynamically suppressed. (See fig. 5.) Soft-gluon exchange is suppressed since the incident p or $\bar{p}$ color neutral wavefunction in the three-parton state with impact operation $b_{\perp} \sim 0(1/\sqrt{s})$. Hard-gluon exchange is suppressed by powers of $[\alpha s(s)]$.

Fig. 5. Analysis of initial-state interactions in PQCD.

The absence of a soft initial-state interaction in $\bar{p}p \rightarrow \bar{\ell}\ell$ is a remarkable consequence of gauge theory and is quite contrary to normal treatments of initial interactions based on Glauber theory. This novel effect has experimental consequences: it can be studied in quasielastic $\bar{p}A \rightarrow \bar{\ell}\ell (A-1)$ reaction. Here we have in mind reactions in which there are no extra hadrons produced and the produced leptons are coplanar with the beam. (The nucleus $(A-1)$ can be left excited). Since PQCD predicts the absence of initial-state elastic and inelastic interactions, the number of such events could be strictly additive in the number Z of protons in the nucleus, every proton in the nucleus is equally available for short-distance annihilation. In traditional Glauber theory only the surface protons can participate because of the strong absorption of the $\bar{p}$ as it traverses the nucleus.

The above description is the ideal result for large s . QCD predicts that additivity is approached monotonically with increasing energy, corresponding to two effects: a) the effective transverse size of the $\bar{p}$ wavefunction is $b_{\perp} \sim 1/\sqrt{s}$, and b) the formation time for the $\bar{p}$ is sufficiently long, such that the Fock state stays small during transit of the nucleus.

The above example is an important example of the PQCD effect called "color transparency;" similar behavior is expected for all hard annihilation processes $\bar{p}p \rightarrow \gamma M$, $\bar{p}p \rightarrow J/\psi$, etc. In the case of exclusive high P_T processes in which hadrons are produced in the final state, each the final state hadron is produced with a small

color singlet wavefunction; thus one predicts negligible attenuation of these hadrons in the quasielastic nucleus reaction.

A test of color transparency in quasielastic $pp \rightarrow pp$ scattering has recently been performed at BNL. This is discussed in more detail in secs. 4 and 9. In the next section I will discuss the general problem of hadronization in the nuclear environment for both exclusive and inclusive reactions. More detailed discussions of exclusive reactions in QCD are then given in secs. 5-13.

4. QCD HADRONIZATION IN NUCLEI

The least-understood process in QCD is *hadronization* — the mechanism which converts quark and gluon quanta to color-singlet integrally-charged hadrons. One way to study hadronization is to perturb the environment by introducing a nuclear medium surrounding the hard-scattering short distance reaction. This is obviously impractical in the theoretically simplest processes — e^+e^- or $\gamma\gamma$ annihilation. However, for large momentum transfer reactions occurring in a nuclear target, such as deep inelastic lepton scattering or massive lepton pair production, the nuclear medium provides a nontrivial perturbation to jet evolution through the influence of initial- and/or final-state interactions. In the case of large momentum transfer quasiexclusive reactions, one can use a nuclear target to filter and influence the evolution and structure of the hadron wavefunctions themselves. The physics of such nuclear reactions is surprisingly interesting and subtle — involving concepts and novel effects quite orthogonal to usual expectations.

In the case of inclusive reactions, the essential test of QCD involving p reactions is the Drell-Yan process $\bar{p}A \rightarrow \ell^+\ell^-X$ and $\bar{p}A \rightarrow \gamma\gamma X$. One of the remarkable consequences of QCD factorization for inclusive reactions at large p_T is the absence of inelastic initial- or final-state interactions of the high-energy particles in a nuclear target. Since structure functions measured in deep inelastic lepton scattering are essentially additive (up to the EMC deviations), factorization implies that the $q\bar{q} \rightarrow \mu^+\mu^-$ subprocesses in Drell-Yan reactions occurs with equal probability on each nucleon throughout the nucleus. At first sight this seems surprising since one expects energy loss from inelastic initial-state interactions.

Fig. 6. Induced radiation from the propagation of an antiquark through a nuclear target in massive lepton production. Such inelastic interactions are coherently suppressed at parton energies large compared to a scale proportional to the length of the target.

In fact, inelastic reactions such as hard gluon bremsstrahlung induced in the nucleus which could potentially decrease the incident parton energy (illustrated in fig. 6) are suppressed by coherence if the quark energy (in the laboratory frame) is large compared to the target length: $E_q > \mu^2 L_A$. Here μ^2 is the difference of mass squared between the incident quark and the quark-gluon pair produced in the initial or final state collision. This phenomenon has its origin in studies of QED

processes by Landau and Pomeranchuk. The QCD analysis is given by Bodwin, Lepage and myself.^[24] The result can be derived by showing that the hard inelastic radiation emitted from differing scattering centers destructively interferes provided the target length condition is maintained. The destructive interference occurs when the momentum transfer μ^2/E_q due to the induced radiation is smaller than the inverse of the separation between two scattering centers in the nucleus. Soft radiation and elastic collisions, however, are still allowed, so one predicts collision broadening of the initial parton transverse momentum. Recent measurements of the Drell-Yan process $\pi A \rightarrow \mu^+ \mu^- X$ by the NA-10 group^[25] at the CERN-SPS confirm that the cross section for muon pairs at large transverse momentum is increased in a tungsten target relative to a deuteron target. (See fig. 7). Since the total cross section for lepton-pair production scales linearly with A (aside from relatively small EMC-effect corrections), there must be a corresponding decrease of the ratio of the differential cross section at low values of the di-lepton transverse momentum. This is also apparent in the data. Further measurements of low-energy $\bar{p}$ Drell-Yan reactions are needed to understand the limits of validity of QCD factorization and to explore the re-emergence of traditional Glauber inelastic scattering at low-antiquark energies. We discuss this further in sec. 14.

Fig. 7. The ratio $\sigma(\pi^- W \rightarrow \mu^+ \mu^- X)/\sigma(\pi^- D \rightarrow \mu^+ \mu^- X)$ as a function of the pair transverse momentum (from ref. 25).

The nucleus thus plays two complimentary roles in quantum chromodynamics:

1. A nuclear target can be used as a control medium or background field to modify or probe quark and gluon subprocesses. Some novel examples are *color transparency*, the predicted transparency of the nucleus to hadrons participating in high-momentum transfer exclusive reactions, and *formation zone phenomena*, the absence of hard, collinear, target-induced radiation by a quark or gluon interacting in a high-momentum transfer inclusive reaction if its energy is large compared to a scale proportional to the length of the target. (Soft radiation and elastic initial-state interactions in the nucleus still occur.) *Coalescence* with co-moving spectators is discussed as a mechanism which can lead to increased open charm hadroproduction, but which also suppresses forward charmonium production (relative to lepton pairs) in heavy ion collisions. There are also novel

features of nuclear diffractive amplitudes — high energy hadronic or electromagnetic reactions which leave the entire nucleus intact and give nonadditive contributions to the nuclear structure function at low x_{Bj} .

2. Conversely, the nucleus can be studied as a QCD structure. At short distances nuclear wavefunctions and nuclear interactions necessarily involve *hidden color*, degrees of freedom orthogonal to the channels described by the usual nucleon or isobar degrees of freedom. At asymptotic momentum transfer, the deuteron form factor and distribution amplitude are rigorously calculable. One can also derive new types of testable scaling laws for exclusive nuclear amplitudes in terms of the reduced amplitude formalism. We discuss this topic in detail in sec. 11.

5. PERTURBATIVE QCD ANALYSIS OF EXCLUSIVE REACTIONS

Perturbative QCD predictions for exclusive processes such as $p\bar{p}$ annihilation into two photons at high-momentum transfer and high invariant pair mass can provide severe tests of the theory.^[17] The simplest, but still very important example,^[18] of the QCD analysis of an exclusive reaction is the calculation of the Q^2 -dependence of the process $\gamma^*\gamma \rightarrow M$ where M is a pseudoscalar meson such as the η . The invariant amplitude contains only one form factor: $M_{\mu\nu} = \epsilon_{\mu\nu\sigma\tau} p_q^\sigma q^\tau F_{\gamma\eta}(Q^2)$.

It is easy to see from power counting at large Q^2 that the dominant amplitude (in light-cone gauge) gives $F_{\gamma\eta}(Q^2) \sim 1/Q^2$ and arises from diagrams (see fig. 8) which have the minimum path carrying Q^2 ; i.e., diagrams in which there is only a single quark propagator between the two photons. The coefficient of $1/Q^2$ involves only the two-particle $q\bar{q}$ component of the meson wavefunction. More precisely the wavefunction is the distribution amplitude $\phi(x, Q)$, defined below, which evolves logarithmically on Q . Higher particle number Fock states give higher power-law falloff contributions to the exclusive amplitude.

Fig. 8. Calculation of the $\gamma - \eta$ transition form factor in QCD from the valence $q\bar{q}$ and $q\bar{q}g$ Fock states.

The TPC/ $\gamma\gamma$ data^[20] shown in fig. 9 are in striking agreement with the predicted QCD power: a fit to the data gives $F_{\gamma\eta}(Q^2) \sim (1/Q^2)^n$ with $n = 1.05 \pm 0.15$. Data for the η' from Pluto and the TPC/ $\gamma\gamma$ experiments give similar results, consistent with scale-free behavior of the QCD quark propagator and the point coupling to the quark current for both the real and virtual photons. In the case of deep inelastic lepton scattering, the observation of Bjorken scaling tests these properties when both photons are virtual.

The QCD power law prediction, $F_{\gamma\eta}(Q^2) \sim 1/Q^2$, is consistent with dimensional counting^[20] and also emerges from current algebra arguments (when both photons are very virtual).^[18] On the other hand, the $1/Q^2$ falloff is also expected in vector meson dominance models. The QCD and VDM predictions can be readily discriminated by

Fig. 9. Comparison of TPC/ $\gamma\gamma$ data^[22] for the $\gamma - \eta$ and $\gamma - \eta'$ transition form factors with the QCD leading twist prediction of ref. 27. The vector meson dominance predictions are also shown.

studying $\gamma^*\gamma^* \rightarrow \eta$. In VDM one expects a product of form factors; in QCD the falloff of the amplitude is still $1/Q^2$ where Q^2 is a linear combination of Q_1^2 and Q_2^2 . It is clearly very important to test this important feature of QCD.

The analysis of $\gamma^*\gamma \rightarrow \eta$ given here is the prototype of the general QCD analysis of exclusive amplitudes at high-momentum transfer:^[22] at large p_T the power behavior of the amplitude is controlled by the minimum tree diagram connecting the valence quarks in the initial and final state — this is the hard scattering amplitude T_H which shrinks to a local operator at asymptotic momentum transfer — effectively the quarks interact when they are all at relative impact separation $b_T \sim 1/p_T$. One then convolutes T_H with the distribution amplitudes $\phi(x_i, Q)$ of the hadrons — analogs of the “wavefunction at the origin” in nonrelativistic quantum mechanics — to construct the hadronic amplitude. (See also sec. 3) This convolution is the basis of the factorization theorem for QCD exclusive reactions: to leading order in $1/p_T$, the nonperturbative dynamics associated with the hadronic bound states is isolated in universal, process-independent distribution amplitudes.^[23] In cases such as $\gamma\gamma$ annihilation into meson pairs and meson form factors, the analysis is completely rigorous in the sense that it can be carried out systematically to all orders in perturbation theory.

A striking feature of the QCD description of exclusive processes is “color transparency.”^[22] The only part of the hadronic wavefunction that scatters at large momentum transfer is its valence Fock state where the quarks are at small relative impact separation. Such a fluctuation has a small color-dipole moment and thus has negli-

gible interactions with other hadrons. Since such a state stays small over a distance proportional to its energy, this implies that quasielastic hadron-nucleon scattering at large momentum transfer as illustrated in fig. 10 can occur additively on all of the nucleons in a nucleus with minimal attenuation due to elastic or inelastic final state interactions in the nucleus, i.e., the nucleus becomes "transparent." By contrast, in conventional Glauber scattering, one predicts strong, nearly energy-independent initial- and final-state attenuation.

Fig. 10. Quasielastic pp scattering inside a nuclear target. Normally one expects such processes to be attenuated by elastic and inelastic interactions of the incident proton and the final-state interaction of the scattered proton. Perturbative QCD predicts minimal attenuation; i.e., "color transparency," at large momentum transfer.

A recent experiment^[34] at BNL measuring quasielastic $pp \rightarrow pp$ scattering at $\theta_{cm} = 90^\circ$ in various nuclei appears to confirm the color transparency prediction — at least for p_{lab} up to 10 GeV/c. (See fig. 11.) Descriptions of elastic scattering which involve soft hadronic wavefunctions cannot account for the data. However, at higher energies, $p_{lab} \sim 12$ GeV/c, normal attenuation is observed in the BNL experiment. This is the same kinematical region $E_{cm} \sim 5$ GeV where the large spin correlation in A_{NN} are observed.^[35] Both features may be signaling new s -channel physics associated with the onset of charmed hadron production^[36] or interference with Landshoff pinch singularity diagrams.^[37] Much more testing of the color transparency phenomena is required, particularly in quasielastic lepton-proton scattering, Compton scattering, antiproton-proton scattering, etc.

As I have discussed in the introduction, the essential nonperturbative input for exclusive reactions at high-momentum transfer is the hadron "distribution amplitude" $\phi(x, Q)$ which describe the longitudinal momentum distribution of the quarks in the valence, lowest particle-number Fock state.^[38] Hadron wavefunctions can be conveniently defined as coefficients on a Fock basis at fixed $\tau = t + z/c$ in the light-cone gauge. Then

$$\phi(x, Q) = \int d^2 k_\perp \theta(Q^2 - k_\perp^2) \psi_V(x, k_\perp) ;$$

i.e., $\phi(x, Q)$ is the probability amplitude to find the quark and antiquark in the meson (or three quarks in a baryon) collinear up to the transverse momentum scale Q . Here $x = (k^0 + k^*)/(p^0 + p^*)$. More generally, the distribution amplitude can be defined as a gauge-invariant matrix-element product of quark fields evaluated between the QCD vacuum and the hadron state. At large Q^2 one can use an operator product expansion or an evolution equation to determine $\phi(x, Q)$ from an initial value $\phi(x, Q_0)$ determined by nonperturbative input. The distribution amplitude contains all of the bound-state dynamics and specifies the momentum distribution of the quarks in the

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Fig. 11. Measurements of the transparency ratio $T = \frac{Z_{eff}}{Z} = \frac{d\sigma}{dt}[pA \rightarrow p(A-1)] / \frac{d\sigma}{dt}[pA \rightarrow pp]$ near 90° on Aluminum (from ref. 84). Conventional theory predicts that T should be small and roughly constant in energy. Perturbative QCD^[33] predicts a monotonic rise to $T = 1$.

hadron. The hard-scattering amplitude can be calculated perturbatively as a function of $\alpha_s(Q^2)$. The analysis can be applied to form factors, exclusive photon-photon reactions, photoproduction, fixed-angle scattering, etc.

Exclusive two-body processes $\gamma\gamma \rightarrow H\bar{H}$ at large $s = W_{\gamma\gamma}^2 = (q_1 + q_2)^2$ and fixed $\theta_{cm}^{\gamma\gamma}$ provide a particularly important laboratory for testing QCD, since the large momentum-transfer behavior, helicity structure, and often even the absolute normalization can be rigorously predicted.^[37,38] The angular dependence of some of the $\gamma\gamma \rightarrow H\bar{H}$ cross sections reflects the shape of the hadron distribution amplitudes $\phi_H(x_i, Q)$. The $\gamma_\lambda \gamma_{\lambda'} \rightarrow H\bar{H}$ amplitude can be written as a factorized form

$$M_{\lambda\lambda'}(W_{\gamma\gamma}, \theta_{cm}) = \int_0^1 [dy_i] \phi_H^*(x_i, Q) \phi_{\bar{H}}^*(y_i, Q) T_{\lambda\lambda'}(x, y; W_{\gamma\gamma}, \theta_{cm}) \quad ,$$

where $T_{\lambda\lambda'}$ is the hard scattering helicity amplitude. To leading order $T \propto \alpha(\alpha_s/W_{\gamma\gamma}^2)^{1,2}$ and $d\sigma/dt \sim W_{\gamma\gamma}^{-4,-6} f(\theta_{cm})$ for meson and baryon pairs, respectively.

Lowest order predictions for pseudoscalar and vector-meson pairs for each helicity amplitude are given in ref. 27. In each case the helicities of the hadron pairs are equal and opposite to leading order in $1/W^2$. The normalization and angular dependence of the leading order predictions for $\gamma\gamma$ annihilation into charged meson pairs are almost model independent; i.e., they are insensitive to the precise form of the meson distribution amplitude. If the meson distribution amplitudes is symmetric in x and $(1-x)$, then the same quantity $\int_0^1 dx [\phi_\pi(x, Q)/(1-x)]$ controls the x -integration for both $F_\pi(Q^2)$ and to high accuracy $M(\gamma\gamma \rightarrow \pi^+\pi^-)$. Thus for charged pion pairs Lepage and I found the relation:

$$\frac{\frac{d\sigma}{dt}(\gamma\gamma \rightarrow \pi^+\pi^-)}{\frac{d\sigma}{dt}(\gamma\gamma \rightarrow \mu^+\mu^-)} \cong \frac{4|F_\pi(s)|^2}{1 - \cos^4\theta_{cm}}$$

Note that in the case of charged kaon pairs, the asymmetry of the distribution amplitude may give a small correction to this relation.

The scaling behavior, angular behavior and normalization of the $\gamma\gamma$ exclusive pair production reactions are nontrivial predictions of QCD. Recent Mark II meson pair data and PEP4/PEP9 data for separated $\pi^+\pi^-$ and K^+K^- production in the range $1.6 < W_{\gamma\gamma} < 3.2$ GeV near 90° are in satisfactory agreement with the normalization and energy dependence predicted by QCD. (See fig. 1.) In the case of $\pi^0\pi^0$ production, the $\cos\theta_{cm}$ dependence of the cross section can be inverted to determine the x -dependence of the pion distribution amplitude. The one-loop corrections to the hard-scattering amplitude for meson pairs have been calculated by Nizic.^[100] The QCD predictions for mesons containing admixtures of the $|gg\rangle$ Fock state is given by Atkinson, Sucher and Tsokos.^[101]

The perturbative QCD analysis has been extended to baryon-pair production in comprehensive analyses by Farrar et al.^[102] and by Gunion et al.^[103] Predictions are given for the "sideways" Compton process $\gamma\gamma \rightarrow p\bar{p}$, $\Delta\bar{\Delta}$ pair production, and the entire decuplet set of baryon pair states. The arduous calculation of 280 $\gamma\gamma \rightarrow qq\bar{q}\bar{q}\bar{q}\bar{q}$ diagrams in T_H required for calculating $\gamma\gamma \rightarrow B\bar{B}$ is greatly simplified by using two-component spinor techniques. The doubly charged Δ pair is predicted to have a fairly small normalization. Experimentally, such resonance pairs may be difficult to identify under the continuum background.

The normalization and angular distribution of the QCD predictions for proton-antiproton production shown in fig. 12 depend in detail on the form of the nucleon distribution amplitude, and thus provide severe tests of the model form derived by Chernyak, Ogloblin and Zhitnitsky from QCD sum rules.^[100]

Fig. 12. Perturbative QCD predictions by Farrar and Zhang for the $\cos(\theta_{cm})$ dependence of the $\gamma\gamma \rightarrow p\bar{p}$ cross section assuming the King-Sachrajda (KS), Chernyak, Ogloblin and Zhitnitsky (COZ), and original Chernyak and Zhitnitsky (CZ) forms for the proton distribution amplitude, $\phi_p(x_i, Q)$.

An important check of the QCD predictions can be obtained by combining data from $\gamma\gamma \rightarrow p\bar{p}$ and the annihilation reaction, $p\bar{p} \rightarrow \gamma\gamma$, with large angle Compton scattering $\gamma p \rightarrow \gamma p$.^[101]

This comparison checks in detail the angular dependence and crossing behavior expected from the theory. Furthermore, in $p\bar{p}$ collisions one can study time-like photon production into e^+e^- and examine the virtual photon mass dependence of the Compton amplitude. Predictions for the q^2 dependence of the $p\bar{p} \rightarrow \gamma\gamma^*$ amplitude can be obtained by crossing the results of Gunion and Millers.^[30]

The region of applicability of the leading power-law predictions for $\gamma\gamma \rightarrow p\bar{p}$ requires that one be beyond resonance or threshold effects. It presumably is set by the scale where $Q^4 G_M(Q^2)$ is roughly constant, i.e., $Q^2 > 3 \text{ GeV}^2$. Present measurements may thus be too close to threshold for meaningful tests.^[32] It should be noted that unlike the case for charged meson-pair production, the QCD predictions for baryons are sensitive to the form of the running coupling constant and the endpoint behavior of the wavefunctions.

The QCD predictions for $\gamma\gamma \rightarrow H\bar{H}$ can be extended to the case of one or two virtual photons, for measurements in which one or both electrons are tagged. Because of the direct coupling of the photons to the quarks, the Q_1^2 and Q_2^2 dependence of the $\gamma\gamma \rightarrow H\bar{H}$ amplitude for transversely polarized photons is minimal at W^2 large and fixed θ_{cm} , since the off-shell quark and gluon propagators in T_H already transfer hard momenta; i.e., the 2γ coupling is effectively local for $Q_1^2, Q_2^2 \ll p_T^2$. The $\gamma^*\gamma^* \rightarrow \bar{B}B$ and $\bar{M}M$ amplitudes for off-shell photons have been calculated by Millers and Gunion.^[33] New results on charged $\pi\rho$ pair production were also presented to this meeting by Kessler and Tamazoust. In each case, the predictions show strong sensitivity to the form of the respective baryon and meson distribution amplitudes.

We also note that photon-photon collisions provide a way to measure the running coupling constant in an exclusive channel, independent of the form of hadronic distribution amplitudes. The photon-meson transition form factors $F_{\gamma \rightarrow M}(Q^2)$, $M = \pi^0, \eta^0, f$, etc., are measurable in tagged $e\gamma \rightarrow e'M$ reactions. QCD predicts

$$\alpha_s(Q^2) = \frac{1}{4\pi} \frac{F_\pi(Q^2)}{Q^2 |F_{\pi\gamma}(Q^2)|^2} \quad ,$$

where to leading order the pion distribution amplitude enters both numerator and denominator in the same manner.

6. APPLICABILITY OF PERTURBATIVE QCD TO EXCLUSIVE PROCESSES

Isgur and Llewellyn Smith^[41] have recently raised a number of questions concerning the application of perturbative QCD to exclusive reactions in the momentum transfer range presently accessible to experiment. The issues involved are very important for understanding the basis of virtually all perturbative QCD predictions. In the following I will review and discuss the main points at issue:

(1) Isgur and Llewellyn Smith, and also Radyshkin,^[44] argue that the normalization of the PQCD amplitude is of order $(\alpha_s/\pi)^n (\lambda^2/Q^2)^n$ where λ is a typical hadronic scale. If this were the correct estimate, the perturbative contributions would be too small to compete with the rapidly-falling "soft" nonperturbative contributions until very large momentum transfers Q .

In fact, the PQCD prediction for the pion form factor at large Q^2 is nominally of order $16\pi\alpha_s f_\pi^2$, a factor of order $16\pi^2$ times larger than the above estimate. The actual coefficient of the leading twist, leading power law term depends on the integral $\int_0^1 dx \frac{\phi_\pi(x/Q)}{(1-x)}$, and is thus only moderately sensitive to the shape of the meson distribution amplitude in the endpoint region.

The normalization and sign of the leading power law terms predicted by PQCD are in agreement with the measurements of the meson and baryon form factors as well as large invariant mass exclusive photon-photon meson pair production cross sections if one uses the hadron distribution amplitudes predicted by Chernyak et al.^[40] and Sachrajda and King^[44] from QCD sum rules. As discussed in sec. 2, the recent lattice gauge theory analysis of the moments of the meson distribution amplitude by Martinelli and Sachrajda^[47] give results consistent with those of Chernyak and Zhitnitsky.

It might also be noted that in QED, the "soft" contributions to the positronium form factor from Coulomb photon exchange are the same order in α as the "hard" contributions from transverse photon exchange. There are no extra powers of α in the hard amplitude! Once the electrons are relativistic, i.e., for $Q^2 \sim M^2$, the hard, perturbative contribution dominates.^[48]

(2) Isgur and Llewellyn Smith argue that the momentum transfer flowing through the gluon propagator in the hard scattering amplitude for an exclusive reaction is typically too small to trust the perturbative expansion. This seems to be of particular concern for the skewed, highly relativistic distribution amplitudes obtained from the QCD sum rule analysis of Chernyak et al. since the integration region where x is large tends to be emphasized. In the case of the hard-scattering T_H amplitude for the pion form factor (illustrated in fig. 13), the struck quark is off-shell at order $(1-x)Q^2$ whereas the momentum transfer of the exchanged gluon is of order $(1-x)(1-y)Q^2$, which can be considerably smaller.

Fig. 13. Leading twist contribution to the meson form factor in QCD.

In fact, as shown by Lepage and myself,^[49] the momentum transfer scale where one can analyze amplitudes perturbatively in QCD is controlled by the virtuality of the quark propagator not that of the exchanged gluon. [The range of the gluon virtuality is of course important in setting the scale of the effective coupling constant $\alpha_s(Q^2)$.] If the struck quark is sufficiently off-shell, $|k_q^2| > \Lambda_{\text{QCD}}^2$, one can easily show that multiple soft gluon exchange contributions are suppressed by powers of Q^2 relative to one-gluon exchange. The same considerations apply to the analysis of the evolution of deep inelastic structure functions: the critical scale is the off-shell virtuality of the quark propagators — not the minimum virtuality of the gluons. Even though the radiated gluons have low virtuality, one can compute the form of QCD evolution with elementary vector gluon couplings provided that the struck quark is sufficiently off-shell. Similarly, in computations of quark jet evolution, the perturbative gluon coupling dominates even though the gluon can be radiated near its mass shell. Requiring the gluon to have a minimum virtual mass corresponds to multiple jet production.

How can one reconcile the PQCD analysis with the concept that at low momentum transfer the interaction between quarks is nonperturbative? The concept of a nonperturbative potential (and estimates of scales involving the gluonium mass) can only be applicable to situations in which quarks are close to their mass shells and scatter at low relative velocity so that there is sufficient time to interact strongly. However, in the high momentum transfer form factor and deep inelastic scattering reactions, the struck quark is relatively far off its mass shell and interacts at high momentum relative to the spectator quarks. Thus its interactions may be computed perturbatively.

The above observations form the basis of the application of renormalization group equations and the operator product expansion to these reactions, and allows one to calculate the leading power behavior and the QCD logarithmic evolution of exclusive amplitudes for the pion form factor and $\gamma\gamma$ annihilation into meson pairs to all orders in perturbation theory.

The predictions^[123] for the leading twist term in exclusive QCD hadronic amplitudes are thus unambiguous. Higher twist corrections to the quark and gluon propagator due to mass terms and intrinsic transverse momenta of a few hundred MeV give nominal corrections of higher order in $1/Q^2$. These finite mass corrections combine with the leading twist results to give a smooth approach to small Q^2 . The PQCD scaling laws thus become valid at relatively low momentum transfer, the few GeV scale, consistent with what is observed in experiment, as in the results shown in figs. 1 and 9.^[144]

(3) *Independent of the underlying theory, the form factor of a hadron can be computed from the overlap of light-cone wavefunctions, summed over Fock states, as shown by Drell and Yan.^[147] This is the starting point for all relativistic calculations including the PQCD analysis. Isgur and Llewellyn Smith, and also Radyshkin, argue that one can obtain reasonable agreement with the form factor data by parameterizing the three-point vertex amplitude using various models for the bound state wavefunctions.*

However, phenomenological agreement with a parameterization of the vertex amplitude is not in contradiction with the PQCD analysis unless one can show that the QCD wavefunction with gluon exchange can be excluded in favor of purely nonperturbative forms. The analyses^[148] of Dziembowski and Mankiewicz (which is consistent with QCD sum rules), Carlson and Gross, and Jacob and Kisslinger show that strictly soft wavefunctions, consistent with rotational invariance in the rest frame, and normalized correctly, cannot account for the pion or proton form factors in the power-law scaling regime.

Perhaps the most compelling evidence for the validity of the PQCD approach to exclusive processes is the observation^[144] of color transparency in pp quasielastic scattering in nuclei, as discussed in sec. 4. The BNL data exclude models in which the scattering is dominated by soft wavefunctions.

The perturbative QCD predictions for the leading twist power-law contributions are generally consistent with data for exclusive processes when the momentum transfer exceeds several GeV.^[149] It is difficult to understand the claim that these data are explained by higher twist or soft nonperturbative contributions since such effects necessarily fall at least one power of Q^2 faster than the dimensional counting prediction.

7. SUMMARY OF QCD PREDICTIONS FOR $\bar{p}p$ EXCLUSIVE PROCESSES

Dimensional counting rules^[150] give a direct connection between the degree of hadron compositeness and the power-law fall of exclusive scattering amplitudes at

fixed center of mass angle: $M \sim Q^{4-n} F(\theta_{cm})$ where n is the minimum number of initial and final state quanta. This rule gives the QCD prediction for the nominal power law scaling, modulo corrections from the logarithmic behavior of α_s , the distribution amplitude and small power-law corrections from Sudakov-suppressed Landshoff multiple scattering contributions. For $\bar{p}p$ one predicts

$$\frac{d\sigma}{d\Omega} (\bar{p}p \rightarrow \gamma\gamma) \simeq \frac{\alpha^2}{(p_T^2)^5} f^{\gamma\gamma}(\cos\theta, \ln p_T)$$

$$\frac{d\sigma}{d\Omega} (\bar{p}p \rightarrow \gamma M) \simeq \frac{\alpha^2}{(p_T^2)^6} f^{\gamma M}(\cos\theta, \ln p_T)$$

$$\frac{d\sigma}{d\Omega} (p\bar{p} \rightarrow M\bar{M}) \simeq \frac{1}{(p_T^2)^7} f^{M\bar{M}}(\cos\theta, \ln p_T)$$

$$\frac{d\sigma}{d\Omega} (p\bar{p} \rightarrow B\bar{B}) \simeq \frac{1}{(p_T^2)^8} f^{B\bar{B}}(\cos\theta, \ln p_T) .$$

The angular dependence reflects the structure of the hard-scattering perturbative T_H amplitude, which in turn follows from the flavor pattern of the contributing duality diagrams. For example, a minimally-connected diagram such as that illustrated in fig. 14 is approximately characterized^[51] as

$$T_H \sim \frac{1}{t^2} \frac{1}{s} \frac{1}{u} .$$

Fig. 14. A perturbative contribution to the hard-scattering amplitude in nucleon-nucleon collisions.

We emphasize that comparisons between channels related by crossing of the Mandelstam variables places a severe constraint on the angular dependence and analytic form of the underlying QCD exclusive amplitude. For example, it is possible to measure and compare

$$\bar{p}p \rightarrow \gamma\gamma : \gamma p \rightarrow \gamma p : \gamma\gamma \rightarrow \bar{p}p$$

$$\bar{p}p \rightarrow \gamma\pi^0 : \gamma p \rightarrow \pi^0 p : \pi^0 p \rightarrow \gamma p .$$

SLAC measurements^[52] of the $\gamma p \rightarrow \pi^+ n$ cross section at $\theta_{CM} = \pi/2$ are consistent with the normalization and scaling (see fig. 15) $d\sigma/dt (\gamma p \rightarrow \pi^+ n) \simeq [1 \text{ nb}/(s/10 \text{ GeV})^7] f(t/s)$. We thus expect similar normalization and scaling for $d\sigma/dt (\bar{p}p \rightarrow \gamma\pi^0)$; all angle measurements up to $s \lesssim 15 \text{ GeV}^2$ appear possible given a high luminosity $\bar{p}$ beam.

Extensive measurements of the $pp \rightarrow pp$ cross section have been made at ANL, BNL and other laboratories.^[51] The fixed-angle data on a log-log plot (see fig. 16) appears consistent with the nominal $s^{-10} f(\theta_{CM})$ dimensional counting production.

Fig. 15. Comparison of photoproduction data with the dimensional counting power-law prediction. The data are summarized in ref. 52.

Fig. 16. Test of fixed θ_{CM} scaling for elastic pp scattering. The best fit gives the power $N = 9.7 \pm 0.5$ compared to the dimensional counting prediction $N=10$. Small deviations are not readily apparent on this log-log plot. The compilation is from Landshoff and Polkinghorne.

However, as emphasized by Hendry,^[53] the $s^{10}d\sigma/dt$ cross section exhibits oscillatory behavior with p_T . Even more serious is the fact that polarization measurements^[54]

show significant spin-spin correlations (A_{NN}), and the single spin asymmetry (A_N) is not consistent with predictions based on hadron helicity conservation (see sec. 6) which is expected to be valid for the leading power behavior.^[66] Recent discussions of these effects have been given by Farrar^[66] and Lipkin.^[67] I discuss a new explanation of all of these effects in sec. 9. Clearly, $\bar{p}p \rightarrow \gamma\gamma$ data in the large-angle, large-energy regime will also be helpful in clarifying these fundamental issues.

The simplest exclusive channels accessible to a $\bar{p}p$ facility are $\bar{p}p \rightarrow e^+e^-$, $\mu^+\mu^-$, $\tau^+\tau^-$ which to leading order in α provides a direct measurement of the Dirac and Pauli time-like proton form factor. The θ_{CM} angular dependence can be used to separate F_2 and F_1 and check the basic prediction,^[2] $F_2(s)/F_1(s) \sim M^2/s$.

As discussed in sec. 3, perturbative QCD predicts asymptotic scaling of the form^[2]

$s^2 F_1(s) \sim f(\ln s)$. A high-luminosity $\bar{p}$ facility could push time-like measurements of both form factors well beyond those available from e^+e^- storage rings. Since the normalization is similar to that of $p\bar{p} \rightarrow \gamma\gamma$, one should be able to measure the proton form factors out to center-of-mass energy squared as large as $s \sim 10 \text{ GeV}^2$.

An important example of an exclusive process in QCD is the process $p\bar{p} \rightarrow \gamma\gamma$ as illustrated in fig. 17. To leading order in $1/p_T^2$,

$$M_{p\bar{p} \rightarrow \gamma\gamma}(p_T^2, \theta_{CM}) = \int_0^1 [dx] \int_0^1 [dy] \phi_p(x, p_T) T_H(qqq + \bar{q}\bar{q}\bar{q} \rightarrow \gamma\gamma) \phi_{\bar{p}}(y, p_T) ,$$

where $\phi_p(x, p_T)$ is the antiproton distribution amplitude and $T_H \sim \alpha_s^2(p_T^2)/(p_T^2)$ gives the scaling behavior of the minimally connected tree-graph amplitude for the two-photon annihilation of three quarks and three antiquarks collinear with the initial hadron directions. (See fig. 18.) QCD thus predicts

$$\frac{d\sigma}{d\Omega_{CM}}(p\bar{p} \rightarrow \gamma\gamma) \simeq \frac{\alpha_s^4(p_T^2)}{(p_T^2)^5} f(p_T, \theta_{CM}, \ln p_T^2) .$$

Fig. 17. Application of QCD factorization to $\bar{p}p$ annihilation into photons.

Fig. 18. Example of a lowest-order perturbative contribution to T_H for the process $\bar{p}p \rightarrow \gamma\gamma$.

Fig. 19. Application of QCD to two-photon production of meson pairs.^[58]

Fig. 20. Next-to-leading perturbative contribution to T_H for the process $\gamma\gamma \rightarrow M\bar{M}$. The calculation has been done by Nizic.^[59]

Fig. 21. Leading diagrams for $\gamma + \gamma \rightarrow \bar{p} + p$ calculated in refs. 59 and 60.

The complete calculations of the tree-graph structure (see figs. 19-21) of both $\gamma\gamma \rightarrow M\bar{M}$ and $\gamma\gamma \rightarrow B\bar{B}$ amplitudes has now been completed. One can use

crossing to compute T_H for $p\bar{p} \rightarrow \gamma\gamma$ to leading order in $\alpha_s(p_T^2)$ from the calculations reported by Farrar, Maina and Neri^[60] and Gunion and Millers.^[60] Examples of the predicted angular distributions are shown in figs. 22 and 23. The region of applicability of the leading power-law results is presumed to be set by the scale where $Q^4 G_M(Q^2)$ is roughly constant, i.e., $Q^2 > 3 \text{ GeV}^2$. (See fig. 3.) Present two-photon collision measurements^[61] are at energies too close to the $p\bar{p}$ threshold to meaningfully test the predictions.

Fig. 22. QCD prediction for the scaling and angular distribution for $\gamma + \gamma \rightarrow \bar{p} + p$ calculated by Farrar et al.^[60] The dashed-dot curve corresponds to $4\Lambda^2/s = 0.0016$ and a maximum running coupling constant $\alpha_s^{\text{max}} = 0.8$. The solid curve corresponds to $4\Lambda^2/s = 0.016$ and a maximum running coupling constant $\alpha_s^{\text{max}} = 0.5$. The dashed curve corresponds to a fixed $\alpha_s = 0.3$. The results are very sensitive to the endpoint behavior of the proton distribution amplitude. The CZ form is assumed.

As discussed in sec. 2, a model form for the proton distribution amplitude has been proposed by Chernyak and Zhitnitskii^[62] based on QCD sum rules which leads to normalization and sign consistent with the measured proton form factor. (See fig. 2.) The CZ sum rule analysis has been recently corrected and modified by King and Sachrajda^[63] but the final results are not known at this time. The CZ proton distribution amplitude yields predictions for $\gamma\gamma \rightarrow p\bar{p}$ in rough agreement with the experimental normalization, although the production energy is too low for a clear test. It should be noted that unlike meson pair production^[64] the QCD predictions

Fig. 23. QCD prediction for the scaling and angular distribution for $\gamma + \gamma \rightarrow \bar{p} + p$ calculated by Gunion, Sparks and Millers.^[60] CZ distribution amplitudes^[61] are assumed. The solid and running curves are for real photon annihilation. The dashed and dot-dashed curves correspond to one photon space-like, with $Q^2/s = 0.1$.

for baryons are highly sensitive to the form of the running coupling constant and the endpoint behavior of the wavefunctions.

The $\gamma^*\gamma^* \rightarrow \bar{B}B$ and $M\bar{M}$ amplitudes for off-shell photons have also been calculated by Millers and Gunion.^[61] The results show important sensitivity to the form of the respective baryon and meson distribution amplitudes. The consequences of $[gg]$ mixing in singlet mesons in $\gamma\gamma$ processes is discussed in ref. 64. It is possible that data from $p\bar{p}$ collisions at energies up to 10 GeV could greatly clarify the question of whether the perturbative QCD predictions are reliable at moderate momentum transfer.^[65,66] As emphasized in sec. 4, an important check of the QCD predictions can be obtained by combining data from $p\bar{p} \rightarrow \gamma\gamma$, $\gamma\gamma \rightarrow p\bar{p}$ with large angle Compton scattering $\gamma p \rightarrow \gamma p$. This comparison checks in detail the angular dependence and crossing behavior expected from the theory. Furthermore, in $p\bar{p}$ collisions one can even study time-like photon production into e^+e^- and examine the virtual photon mass dependence of the Compton amplitude. Predictions for the q^2 dependence of the $p\bar{p} \rightarrow \gamma\gamma^*$ amplitude can be obtained by crossing the results of Gunion and Millers.^[62]

6. HELICITY SELECTION RULE AND EXCLUSIVE CHARMONIUM DECAYS

One of the simplest predictions of perturbative QCD for exclusive processes including $p\bar{p}$ reactions is hadron helicity conservation: to leading order in $1/Q$, the total helicity of hadrons in the initial state must equal the total helicity of hadrons in the final state. This selection rule is independent of any photon or lepton spin appearing in the process. The result follows from (a) neglecting quark mass terms, (b) the vector coupling of gauge particles, and (c) the dominance of valence Fock states with zero angular momentum projection.^[67]

Hadron helicity conservation may be relevant to an interesting puzzle concerning the exclusive decays J/ψ and $\psi' \rightarrow \rho\pi, K^*\bar{K}$ and possibly other Vector-Pseudoscalar

(VP) combinations. One expects $J/\psi(\psi')$ to decay to hadrons via three gluons or, occasionally, via a single direct photon. In either case the decay proceeds via $|\Psi(0)|^2$, where $\Psi(0)$ is the wave function at the origin in the nonrelativistic quark model for $c\bar{c}$. Thus it is reasonable to expect on the basis of perturbative QCD, that for any final hadronic state h :

$$Q_h \equiv \frac{B(\psi' \rightarrow h)}{B(J/\psi \rightarrow h)} \cong \frac{B(\psi' \rightarrow e^+e^-)}{B(J/\psi \rightarrow e^+e^-)} = 0.135 \pm 0.023$$

Usually this is true, as is well documented in ref. 68 for $p\bar{p}\pi^0$, $2\pi^+2\pi^-\pi^0$, $\pi^+\pi^-\omega$, and $3\pi^+3\pi^-\pi^0$, hadronic channels. The startling exceptions occur for $\rho\pi$ and $K^*\bar{K}$ where the present experimental limits⁽⁶⁹⁾ are $Q_{\rho\pi} < 0.0063$ and $Q_{K^*\bar{K}} < 0.0027$.

Recently San Fu Tuan, Peter Lepage and I⁽⁷⁰⁾ have proposed an explanation of the puzzle by assuming (a) the general validity of the perturbative QCD theorem that total hadron helicity is conserved in high-momentum transfer exclusive processes, but supplemented by (b) violation of the QCD theorem when the J/ψ decay to hadrons via three hard gluons is modulated by the gluons forming an intermediate gluonium state O before transition to hadrons. In essence the model of Hou and Soni⁽⁷⁰⁾ takes over in this latter stage.

Since the vector state O has to be produced with helicity $\lambda = \pm 1$, the VP decays should be suppressed by a factor $1/s$ in the rate. The ψ' seems to respect this rule. The J/ψ does not, and that is the mystery. Put in more quantitative terms, we expect on the basis of perturbative QCD⁽⁷¹⁾ $Q_{\rho\pi} \equiv [B(\psi' \rightarrow \rho\pi)/B(J/\psi \rightarrow \rho\pi)] \sim [M_{J/\psi}/M_{\psi'}]^6$ assuming quark helicity is conserved in strong interactions. This includes a form factor suppression proportional to $[M_{J/\psi}/M_{\psi'}]^4$. This suppression is not nearly large enough, though, to account for the data.

One can question the validity of the QCD helicity conservation theorem at the charmonium mass scale. Helicity conservation has received important confirmation in $J/\psi \rightarrow p\bar{p}$ where the angular distribution is known experimentally to follow $|1 + \cos^2 \theta|$ rather than $\sin^2 \theta$ for helicity flip. The ψ' decays clearly respect hadron helicity conservation. It is difficult to understand how the J/ψ could violate this rule since the J/ψ and ψ' masses are so close. Corrections from quark mass terms, soft gluon corrections and finite energy corrections would not be expected to lead to large J/ψ differences. It is hard to imagine anything other than a resonant or interference effect that could account for such dramatic energy dependence.

A relevant violation of the QCD theorem which does have significance to this problem, is the recognition that the theorem is built on the underlying assumption of short-range "point-like" interactions among the constituents throughout. For instance, $J/\psi(c\bar{c}) \rightarrow 3g$ has a short-range $\cong 1/m_c$ associated with the short time scale of interaction. If, however, subsequently the three gluons were to resonate forming a gluonium state O which has large transverse size $\cong 1/M_H$ covering an extended (long) time period, then the theorem is invalid. Note that even if the gluonium state O has large mass, close to $M_{J/\psi}$, its size could still be the standard hadronic scale of 1 fm, just as the case for the D -meson and B -mesons.

We have thus proposed, following Hou and Soni, that the enhancement of $J/\psi \rightarrow K^*\bar{K}$ and $J/\psi \rightarrow \rho\pi$ decay modes is caused by a quantum mechanical mixing of the J/ψ with a $J^{PC} = 1^{--}$ vector gluonium state O which causes the breakdown of the QCD helicity theorem. The decay width for $J/\psi \rightarrow \rho\pi(K^*\bar{K})$ via the sequence $J/\psi \rightarrow O \rightarrow \rho\pi(K^*\bar{K})$ must be substantially larger than the decay width for the (non-pole) continuum process $J/\psi \rightarrow 3 \text{ gluons} \rightarrow \rho\pi(K^*\bar{K})$. In the other channels (such as $p\bar{p}, p\bar{p}\pi^0, 2\pi^+2\pi^-\pi^0$, etc.), the branching ratios of the O must be so small that the continuum contribution governed by the QCD theorem dominates over that

of the O pole. For the case of the ψ' the contribution of the O pole must always be inappreciable in comparison with the continuum process where the QCD theorem holds. The experimental limits on $Q_{\rho\pi}$ and $Q_{K^*\bar{K}}$ are now substantially more stringent than when Hou and Soni made their estimates of M_O , $\Gamma_{O \rightarrow \rho\pi}$ and $\Gamma_{O \rightarrow K^*\bar{K}}$ in 1982.

It is interesting, indeed, that the existence of such a gluonium state O was first postulated by Freund and Nambu^[7] based on OZI dynamics soon after the discovery of the J/ψ and ψ' mesons. In fact, Freund and Nambu predicted that the O would decay copiously precisely into $\rho\pi$ and $K^*\bar{K}$ with severe suppression of decays into other modes like e^+e^- as required for the solution of the puzzle.

Final states h which can proceed only through the intermediate gluonium state satisfy the ratio:

$$Q_h = \frac{B(\psi' \rightarrow e^+e^-) (M_{J/\psi} - M_O)^2 + \frac{1}{4} \Gamma_O^2}{B(J/\psi \rightarrow e^+e^-) (M_{\psi'} - M_O)^2 + \frac{1}{4} \Gamma_{\psi'}^2}.$$

We have assumed that the coupling of the J/ψ and ψ' to the gluonium state scales as the e^+e^- coupling. The value of Q_h is small if the O is close in mass to the J/ψ . Thus we require $(M_{J/\psi} - M_O)^2 + \frac{1}{4} \Gamma_O^2 \lesssim 2.6 Q_h \text{ GeV}^2$. The experimental limit for $Q_{K^*\bar{K}}$ then implies $[(M_{J/\psi} - M_O)^2 + \frac{1}{4} \Gamma_O^2]^{1/2} \lesssim 80 \text{ MeV}$. This implies $|M_{J/\psi} - M_O| < 80 \text{ MeV}$ and $\Gamma_O < 160 \text{ MeV}$. Typical allowed values are $M_O = 3.0 \text{ GeV}$, $\Gamma_O = 140 \text{ MeV}$ or $M_O = 3.15 \text{ GeV}$, $\Gamma_O = 140 \text{ MeV}$. Notice that the gluonium state could be either lighter or heavier than the J/ψ . The branching ratio of the O into a given channel must exceed that of the J/ψ .

It is not necessarily obvious that a $J^{PC} = 1^{--}$ gluonium state with these parameters would necessarily have been found in experiments to date. One must remember that though $O \rightarrow \rho\pi$ and $O \rightarrow K^*\bar{K}$ are important modes of decay, at a mass of order 3.1 GeV many other modes (all be it less important) are available. Hence, a total width $\Gamma_O \cong 100$ to 150 MeV is quite conceivable. Because of the proximity of M_O to $M_{J/\psi}$, the most important signatures for an O search via exclusive modes $J/\psi \rightarrow K^*\bar{K}h$, $J/\psi \rightarrow \rho\pi h$; $h = \pi\pi, \eta, \eta'$, are no longer available by phase-space considerations. However, the search could still be carried out using $\psi' \rightarrow K^*\bar{K}h$, $\psi' \rightarrow \rho\pi h$; with $h = \pi\pi$, and η . Another way to search for O in particular, and the three-gluon bound states in general, is via the inclusive reaction $\psi' \rightarrow (\pi\pi) + X$, where the $\pi\pi$ pair is an isosinglet. The three-gluon bound states such as O should show up as peaks in the missing mass (i.e., mass of X) distribution.

Perhaps the most direct way to search for the O is to scan $p\bar{p}$ or e^+e^- annihilation at $\sqrt{s}$ within $\sim 100 \text{ MeV}$ of the J/ψ , triggering on vector/pseudoscalar decays such as $\pi\rho$ or $\bar{K}K^*$.

The fact that the $\rho\pi$ and $K^*\bar{K}$ channels are strongly suppressed in ψ' decays but not in J/ψ decays clearly implies dynamics beyond the standard charmonium analysis. As we have shown, the hypothesis of a three-gluon state O with mass within $\cong 100 \text{ MeV}$ of the J/ψ mass provides a natural, perhaps even compelling, explanation of this anomaly. If this description is correct, then the ψ' and J/ψ hadronic decays are not only confirming hadron helicity conservation (at the ψ' momentum scale) but are also providing a signal for bound gluonic matter in QCD. An alternative model by Chaichian and Tornquist, based on nonperturbative exponential vertex functions, has also been proposed.

The production of heavy quark resonances $p\bar{p} \rightarrow \psi, \chi, \eta_c$, etc. can be analyzed in a systematic way in QCD using the exclusive amplitude formalism of ref. 2. Since quark helicity is conserved in the basic subprocesses to leading order, and the distribution amplitude is the azimuthal angle symmetric $L_z = 0$ projection of the valence hadron

Fock wavefunction, total hadron helicity is conserved for $A+B \rightarrow C+D$:^[66] $\lambda_A + \lambda_B = \lambda_C + \lambda_D$. The result is predicted to hold to all orders in $\alpha_s(Q^2)$. Thus an essential feature of the perturbative QCD is the prediction of hadron helicity conservation up to kinematical and dynamical corrections of order m/Q and $\langle\psi\bar{\psi}\rangle^{1/3}/Q$ where Q is the momentum transfer or heavy mass scale, m is the light quark mass, and $\langle\psi\bar{\psi}\rangle$ is a measure of nonperturbative effects due to chiral symmetry breaking of the QCD vacuum. Applying this prediction to $p\bar{p}$ annihilation, one predicts $\lambda_p + \lambda_{\bar{p}} = 0$, i.e., $S_z = J_z = \pm 1$ is the leading amplitude for heavy resonance production.^[66] Thus the ψ is expected to be produced with $J_z = \pm 1$, whereas the χ and η_c cross sections should be suppressed, at least to leading power in the heavy quark mass. The analogous tests in e^+e^- annihilation appear to be verified for ψ' decays but not the ψ . Hou and Soni^[72] have suggested this effect may be due to the ψ mixing with $J = 1$ gluonium states. Antiproton-proton production of narrow resonances should be able to clarify these basic QCD issues.

9. SPIN CORRELATIONS, QCD COLOR TRANSPARENCY, AND HEAVY QUARK THRESHOLDS IN pp SCATTERING

One of the most serious challenges to quantum chromodynamics is the behavior of the spin-spin correlation asymmetry $A_{NN} = \frac{d\sigma(11) - d\sigma(\bar{1}\bar{1})}{d\sigma(11) + d\sigma(\bar{1}\bar{1})}$ measured in large momentum transfer pp elastic scattering. (See fig. 24.) At $p_{lab} = 11.75$ GeV/c and $\theta_{cm} = \pi/2$, A_{NN} rises to $\simeq 60\%$, corresponding to four times more probability for protons to scatter with their incident spins both normal to the scattering plane and parallel, rather than normal and opposite.

The polarized cross section shows a striking energy and angular dependence not expected from the slowly-changing perturbative QCD predictions.^[74] However, the unpolarized data is in first approximation consistent with the fixed angle scaling law $s^{10}d\sigma/dt(pp \rightarrow pp) = f(\theta_{CM})$ expected from the perturbative analysis. (See fig. 16.) The onset of new structure^[74] at $s \simeq 23$ GeV² is a sign of new degrees of freedom in the two-baryon system. In this section, I will discuss an explanation by Guy de Teramond and myself³⁶ for (1) the observed spin correlations, (2) the deviations from fixed-angle scaling laws, and (3) the anomalous energy dependence of absorptive corrections to quasielastic pp scattering in nuclear targets, in terms of a simple model based on two $J = L = S = 1$ broad resonances (or threshold enhancements) interfering with a perturbative QCD quark-interchange background amplitude. The structures in the $pp \rightarrow pp$ amplitude may be associated with the onset of strange and charmed thresholds. If this view is correct, large angle pp elastic scattering would have been virtually featureless for $p_{lab} \geq 5$ GeV/c, had it not been for the onset of heavy flavor production. As a further illustration of the threshold effect, we also show the effect in A_{NN} due to a narrow 3F_3 pp resonance at $\sqrt{s} = 2.17$ GeV ($p_{lab} = 1.26$ GeV/c) associated with the $p\Delta$ threshold.

The perturbative QCD analysis^[74] of exclusive amplitudes assumes that large momentum transfer exclusive scattering reactions are controlled by short distance quark-gluon subprocesses, and that corrections from quark masses and intrinsic transverse momenta can be ignored. The main predictions are fixed-angle scaling laws^[77] (with small corrections due to evolution of the distribution amplitudes, the running coupling constant, and pinch singularities), hadron helicity conservation,^[78] and the novel phenomenon, "color transparency."

As discussed in secs. 2 and 5, the power-law scaling quark-counting predictions for form factors, two-body elastic hadron-hadron scattering,^[79] and exclusive two-photon reactions are generally consistent with experiment at transverse momenta

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Fig. 24. The spin-spin correlation A_{NN} for elastic pp scattering with beam and target protons polarized normal to the scattering plane.^[7a] (From ref. 73.) $A_{NN} = 60\%$ implies that it is four times more probable for the protons to scatter with spins parallel rather than antiparallel.

beyond a few GeV. (See figs. 3, 9, 15 and 16). In leading order in $1/p_T$, only the lowest particle-number "valence" Fock state wavefunction with all the quarks within an impact distance $b_\perp \leq 1/p_T$ contributes to the high momentum transfer scattering amplitude in QCD. Such a Fock state component has a small color dipole moment and thus interacts only weakly with hadronic or nuclear matter.^[2a] This minimally interacting proton configuration can retain its small size as it propagates in the nucleus over a distance which grows with energy. Thus, unlike traditional Glauber theory, QCD predicts that large momentum transfer quasielastic reactions occurring in a nucleus suffer minimal initial and final state attenuation; i.e., one expects a volume rather than surface dependence in the nuclear number. This is the QCD "color transparency" prediction.

As discussed in the introduction, a test of color transparency in large momentum transfer quasielastic pp scattering at $\theta_{cm} \simeq \pi/2$ has recently been carried out at BNL using several nuclear targets (C, Al, Pb).^[24] The attenuation at $p_{lab} = 10$ GeV/c in the various nuclear targets was observed to be in fact much less than that

predicted by traditional Glauber theory. (See fig. 11.) This appears to support the color transparency prediction. However at $p_{lab} = 12$ GeV/c, normal attenuation was observed, in contradiction to the expectation from perturbative QCD that the transparency effect should become even more apparent! Our observation is that one can explain this surprising result if the scattering at $p_{lab} = 12$ GeV/c ($\sqrt{s} = 4.93$ GeV), is dominated by an s -channel $B=2$ resonance (or resonance-like structure) with mass near 5 GeV, since unlike a hard-scattering reaction, a resonance couples to the fully-interacting large-scale structure of the proton. If the resonance has spin $S = 1$, this can also explain the large spin correlation A_{NN} measured nearly at the same momentum, $p_{lab} = 11.75$ GeV/c. Conversely, in the momentum range $p_{lab} = 5$ to 10 GeV/c we predict that the perturbative hard-scattering amplitude is dominant at large angles. The experimental observation of diminished attenuation at $p_{lab} = 10$ GeV/c thus provides support for the QCD description of exclusive reactions and color transparency.

What could cause a resonance at $\sqrt{s} = 5$ GeV, more than 3 GeV beyond the pp threshold? We can think of several possibilities: (a) a multigluonic excitation such as $|qqqqqqggg\rangle$, (b) a "hidden color" color singlet $|qqqqqq\rangle$ excitation,^[80] or (c) a "hidden flavor" $|qqqqqqQ\bar{Q}\rangle$ excitation, which is the most interesting possibility, since it is so predictive. As in QED, where final state interactions give large enhancement factors for attractive channels in which $Z\alpha/v_{rel}$ is large, one expects resonances or threshold enhancements in QCD in color-singlet channels at heavy quark production thresholds since all the produced quarks have similar velocities.^[81] One thus can expect resonant behavior at $M^* = 2.55$ GeV and $M^* = 5.08$ GeV, corresponding to the threshold values for open strangeness: $pp \rightarrow AK^+p$, and open charm: $pp \rightarrow \Lambda_c D^0 p$, respectively. In any case, the structure at 5 GeV is highly inelastic: we find that its branching ratio to the proton-proton channel is $B^{pp} \simeq 1.5\%$.

We now proceed to a description of the model. We have purposely attempted not to overcomplicate the phenomenology; in particular, we have used the simplest Breit-Wigner parameterization of the resonances, and we have not attempted to optimize the parameters of the model to obtain a best fit. It is possible that what we identify as a single resonance is actually a cluster of resonances.

The background component of the model is the perturbative QCD amplitude. Although complete calculations are not yet available, many features of the QCD predictions are understood, including the approximate s^{-4} scaling of the $pp \rightarrow pp$ amplitude at fixed θ_{cm} and the dominance of those amplitudes that conserve hadron helicity.^[78] Furthermore, recent data comparing different exclusive two-body scattering channels from BNL^[79] show that quark interchange amplitudes^[82] dominate quark annihilation or gluon exchange contributions. Assuming the usual symmetries, there are five independent pp helicity amplitudes: $\phi_1 = M(++++)$, $\phi_2 = M(---+)$, $\phi_3 = M(+--+)$, $\phi_4 = M(-+-)$, $\phi_5 = M(+++-)$. The helicity amplitudes for quark interchange have a definite relationship.^[74] For definiteness, we will assume the following form

$$\begin{aligned}\phi_1(\text{PQCD}) &= 2\phi_3(\text{PQCD}) = -2\phi_4(\text{PQCD}) \\ &= 4\pi C F(t) F(u) \left[\frac{t - m_d^2}{u - m_d^2} + (u \leftrightarrow t) \right] e^{i\delta}.\end{aligned}$$

The hadron helicity nonconserving amplitudes, $\phi_2(\text{PQCD})$ and $\phi_5(\text{PQCD})$ are zero. This form is consistent with the nominal power-law dependence predicted by perturbative QCD^[74] and also gives a good representation of the angular distribution over a broad range of energies.^[82] Here $F(t)$ is the helicity conserving proton form factor, which for simplicity, we take as the standard dipole form, $F(t) = (1 - t/m_d^2)^{-2}$, with

$m_d^2 = 0.71 \text{ GeV}^2$. As shown in ref. 74, the PQCD-quark-interchange structure alone predicts $A_{NN} \simeq 1/3$, nearly independent of energy and angle.

Because of the rapid fixed-angle s^{-4} falloff of the perturbative QCD amplitude, even a very weakly-coupled resonance can have a sizeable effect at large momentum transfer. The large empirical values for A_{NN} suggest a resonant $pp \rightarrow pp$ amplitude with $J = L = S = 1$ since this gives $A_{NN} = 1$ (in absence of background) and a smooth angular distribution. Because of the Pauli principle, an $S = 1$ di-proton resonances must have odd parity and thus odd orbital angular momentum. We parameterize the two non-zero helicity amplitudes for a $J = L = S = 1$ resonance in Breit-Wigner form:

$$\phi_3(\text{resonance}) = 12\pi \frac{\sqrt{s}}{p_{cm}} d_{1,1}^1(\theta_{cm}) \frac{\frac{1}{2} \Gamma^{pp}(s)}{M^* - E_{cm} - \frac{i}{2} \Gamma},$$

$$\phi_4(\text{resonance}) = -12\pi \frac{\sqrt{s}}{p_{cm}} d_{1,1}^1(\theta_{cm}) \frac{\frac{1}{2} \Gamma^{pp}(s)}{M^* - E_{cm} - \frac{i}{2} \Gamma}.$$

(The 3F_3 resonance amplitudes have the same form with $d_{\pm 1,1}^3$ replacing $d_{\pm 1,1}^1$.) Since we are far from threshold, threshold factors in the pp channel can be treated as constants. As in the case of a narrow resonance like the Z^0 , we expect that the partial width into nucleon pairs is proportional to the square of the time-like proton form factor: $\Gamma^{pp}(s)/\Gamma = B^{pp}|F(s)|^2/|F(M^{*2})|^2$, corresponding to the formation of two protons at this invariant energy. The resonant amplitudes then die away by one inverse power of $(E_{cm} - M^*)$ relative to the dominant PQCD amplitudes. (In this sense, they are higher twist contributions relative to the leading twist perturbative QCD amplitudes.) The model is thus very simple: each pp helicity amplitude ϕ_i is the coherent sum of PQCD plus resonance components: $\phi = \phi(\text{PQCD}) + \Sigma \phi(\text{resonance})$. Because of pinch singularities and higher-order corrections, the hard QCD amplitudes are expected to have a nontrivial phase;^[84] we have thus allowed for a constant phase δ in $\phi(\text{PQCD})$. Because of the absence of the ϕ_5 helicity-flip amplitude, the model predicts zero single spin asymmetry A_N . This is consistent with the large angle data at $p_{lab} = 11.75 \text{ GeV}/c$.^[83]

At low transverse momentum, $p_T \leq 1.5 \text{ GeV}$, the power-law fall-off of $\phi(\text{PQCD})$ in s disagrees with the more slowly falling large-angle data, and we have little guidance from basic theory. Our interest in this low-energy region is to illustrate the effects of resonances and threshold effects on A_{NN} . In order to keep the model tractable, we have simply extended the background quark interchange and the resonance amplitudes at low energies using the same forms as above but replacing the dipole form factor by a phenomenological form $F(t) \propto e^{-1/2\beta\sqrt{|t|}}$. We have also included a kinematic factor of $\sqrt{s}/2p_{cm}$ in the background amplitude. The value $\beta = 0.85 \text{ GeV}^{-1}$ then gives a good fit to $d\sigma/dt$ at $\theta_{cm} = \pi/2$ for $p_{lab} \leq 5.5 \text{ GeV}/c$.^[84] The normalizations are chosen to maintain continuity of the amplitudes.

The predictions of the model and comparison with experiment are shown in figs. 25-30. The following parameters are chosen: $C = 2.9 \times 10^3$, $\delta = -1$ for the normalization and phase of $\phi(\text{PQCD})$. The mass, width and pp branching ratio for the three resonances are $M_d^* = 2.17 \text{ GeV}$, $\Gamma_d = 0.04 \text{ GeV}$, $B_d^{pp} = 1$; $M_s^* = 2.55 \text{ GeV}$, $\Gamma_s = 1.6 \text{ GeV}$, $B_s^{pp} = 0.65$; and $M_c^* = 5.08 \text{ GeV}$, $\Gamma_c = 1.0 \text{ GeV}$, $B_c^{pp} = 0.0155$, respectively. As shown in figs. 25 and 26, the deviations from the simple scaling predicted by the PQCD amplitudes are readily accounted for by the resonance structures. The cusp which appears in fig. 26 marks the change in regime below $p_{lab} = 5.5 \text{ GeV}/c$ where PQCD becomes inapplicable. It is interesting to note that in this energy region normal attenuation of quasielastic pp scattering is observed.^[84] The angular distribution (normalized to the data at $\theta_{cm} = \pi/2$) is predicted to broaden relative to the steeper perturbative QCD form, when the resonance dominates. As

shown in fig. 27 this is consistent with experiment, comparing data at $p_{lab} = 7.1$ and 12.1 GeV/c.

Fig. 25. Prediction (solid curve) for $d\sigma/dt(pp \rightarrow pp)$ at $\theta_{cm} = \pi/2$ compared with the data of Akerlof et al.^[14] The dotted line is the background PQCD prediction.

Fig. 26. Ratio of $d\sigma/dt(pp \rightarrow pp)$ at $\theta_{cm} = \pi/2$ to the PQCD prediction. The data^[14] are from Akerlof et al. (open triangles), Allaby et al. (solid dots) and Cocconi et al. (open square). The cusp at $p_{lab} = 5.5$ GeV/c indicates the change of regime from PQCD.

The most striking test of the model is its prediction for the spin correlation A_{NN} shown in fig. 28. The rise of A_{NN} to $\approx 60\%$ at $p_{lab} = 11.75$ GeV/c is correctly reproduced by the high energy $J=1$ resonance interfering with ϕ (PQCD). The narrow peak which appears in the data of fig. 28 corresponds to the onset of the $pp \rightarrow p\Delta(1232)$ channel which can be interpreted as a $uuuuddq\bar{q}$ resonant state. Because of spin-color statistics one expects in this case a higher orbital momentum state, such as a pp^3F_3 resonance. The model is also consistent with the recent high-energy data point for A_{NN} at $p_{lab} = 18.5$ GeV/c and $p_T^2 = 4.7$ GeV² (see fig. 29). The data show a dramatic decrease of A_{NN} to zero or negative values. This is explained in our model by the destructive interference effects above the resonance region. The same effect accounts for the depression of A_{NN} for $p_{lab} \approx 6$ GeV/c shown in fig. 28. The comparison of the angular dependence of A_{NN} with data at $p_{lab} = 11.75$ GeV/c is shown in fig. 30. The agreement with the data^[17] for the longitudinal spin correlation A_{LL} at the same p_{lab} is somewhat worse.

Fig. 27. The $pp \rightarrow pp$ angular distribution normalized at $\theta_{cm} = \pi/2$. The data are from the compilation given in Sivers et al., ref. 79. The solid and dotted lines are predictions for $p_{lab} = 12.1$ and 7.1 GeV/c, respectively, showing the broadening near resonance.

Fig. 28. A_{NN} as a function of p_{lab} at $\theta_{cm} = \pi/2$. The data^[1] are from Croasbie et al. (solid dots), Lin et al. (open squares) and Bhatia et al. (open triangles). The peak at $p_{lab} = 1.26$ GeV/c corresponds to the $p\Delta$ threshold. The data are well reproduced by the interference of the broad resonant structures at the strange ($p_{lab} = 2.35$ GeV/c) and charm ($p_{lab} = 12.8$ GeV/c) thresholds, interfering with a PQCD background. The value of A_{NN} from PQCD alone is $1/3$.

Thus far we have not attempted a global fit to all the pp elastic scattering data, but rather have tried to show that many features can be naturally explained with only a few ingredients: a perturbative QCD background plus resonant amplitudes associated with rapid changes of the inelastic pp cross section. The model provides a good description of the s and t dependence of the differential cross section, including its "oscillatory" dependence^[81] in s at fixed θ_{cm} , and the broadening of the angular distribution near the resonances. Most important, it gives a consistent explanation for the striking behavior of both the spin-spin correlations and the anomalous energy dependence of the attenuation of quasielastic pp scattering in nuclei. We predict that color transparency should reappear at higher energies ($p_{lab} \geq 16$ GeV/c), and also at smaller angles ($\theta_{cm} \approx 60^\circ$) at $p_{lab} = 12$ GeV/c where the perturbative QCD amplitude dominates. If the $J=1$ resonance structures in A_{NN} are indeed associated with heavy quark degrees of freedom, then the model predicts inelastic pp cross

Fig. 29. A_{NN} at fixed $p_T^2 = (4.7 \text{ GeV}/c)^2$. The data point^[11] at $p_{lab} = 18.5 \text{ GeV}/c$ is from Court et al.

Fig. 30. A_{NN} as a function of transverse momentum. The data^[11] are from Crabb et al. (open circles) and O'Fallon et al. (open squares). Diffractive contributions should be included for $p_T^2 \leq 3 \text{ GeV}^2$.

sections of the order of 1 mb and $1\mu\text{b}$ for the production of strange and charmed hadrons near their respective thresholds.^[88] Thus a crucial test of the heavy quark hypothesis for explaining A_{NN} , rather than hidden color or gluonic excitations, is the observation of significant charm hadron production at $p_{lab} \geq 12 \text{ GeV}/c$.

10. HEAVY HADRON PAIR PRODUCTION IN $p\bar{p}$ EXCLUSIVE REACTIONS

One of the few areas of high-energy phenomenology which is apparently not well-understood in perturbative QCD is the production of charmed hadrons. The simple fusion subprocesses $q\bar{q} \rightarrow Q\bar{Q}$ and $gg \rightarrow Q\bar{Q}$ are expected to dominate heavy quark inclusive reactions at least for very large M_Q .^[90] However, in the case of charm production cross sections, the predictions for the energy and x_L dependence appear to contradict experiment.^[91] It is possible that there are significant nonperturbative contributions to charm production such as the intrinsic heavy quark contributions^[92] associated with loop interactions in the hadronic wavefunction, strong binding effects

at low relative velocity,^[10] and other nonperturbative or higher twist effects. A review of some of these issues is given in ref. 91.

Here we want to address the equally provocative question of heavy flavor production in exclusive $p\bar{p}$ reactions, e.g., $p\bar{p} \rightarrow \bar{\Lambda}_Q \Lambda_Q$ where $Q = s, c, b$. The following arguments are heuristic, but they may give a guide to the expected scaling laws and features of these reactions.

Consider the diagram of fig. 31. for the $p\bar{p} \rightarrow \bar{\Lambda}_Q \Lambda_Q$ matrix element. If the Λ 's are produced in the forward direction with $p_T^2 \lesssim \mu^2 \sim (300 \text{ MeV})^2$ then there is maximal kinematic overlap for the light quarks between the initial and final light wavefunctions. The hard subprocess cross section $\bar{u}u \rightarrow c\bar{c}$ would normally give cross sections of order

$$\frac{d\sigma}{d\Omega} \sim \frac{\alpha_s^2(s)f(\Omega)}{s} \sim \frac{\alpha_s^2(4m_Q^2)}{4M_Q^2} f(\Omega) ,$$

but the alignment restriction $p_T^2 < \mu^2$ gives an extra $\mu^2/4m_Q^2$ suppression in the angular integral. Therefore we expect the scaling law

$$\sigma(p\bar{p} \rightarrow \bar{\Lambda}_Q \Lambda_Q) \sim \mu^2 \frac{\alpha_s^2(4M_Q^2)}{m_Q^4} F\left(\frac{4m_Q^2}{s}\right) ,$$

i.e., $\bar{\Lambda}_s \Lambda_s : \bar{\Lambda}_c \Lambda_c : \bar{\Lambda}_b \Lambda_b = 1 : 10^{-2} \text{ to } 10^{-3} : 10^{-4} \text{ to } 10^{-4}$ for $s \gg 4m_Q^2$ similar rates have been recently estimated in the "diquark" model of Kroll and collaborators. Thus it may not be hopeless to actually measure exclusive pairs of heavy charmed baryons in $p\bar{p}$ collisions. The above analysis can be readily extended to other heavy flavor baryon and meson pair exclusive cross sections. The issues are important for clarifying the OZI rule in QCD and the connection between exclusive and inclusive production mechanisms.

Fig. 31. Perturbative diagram for exclusive production of heavy baryon pairs in $p\bar{p}$ annihilation.

Mass corrections to QCD hard scattering amplitudes for a number of heavy quark production amplitudes have been computed. Exclusive pair production of heavy hadrons $|Q_1 \bar{Q}_2\rangle$, $|Q_1 Q_2 Q_3\rangle$ consisting of higher generation quarks ($Q_i = t, b, c$ and possibly s) can be reliably predicted^[10] within the framework of perturbative QCD, since the required wave function input is essentially determined from nonrelativistic considerations. The results can be applied to e^+e^- annihilation, $\gamma\gamma$ annihilation, and W and Z decay into higher generation pairs. The normalization, angular dependence and helicity structure can be predicted away from threshold, allowing a detailed study of the basic elements of heavy quark hadronization. A particularly striking feature of the QCD predictions is the existence of a zero in the form factor and e^+e^- annihilation cross section for zero-helicity hadron pair production close to a specific time-like value $q^2/4M_H^2 = m_h/2m_l$ where m_h and m_l are the heavier and lighter quark masses, respectively. This zero reflects the destructive interference between the spin-dependent and spin-independent (Coulomb exchange) couplings of

the gluon in QCD. In fact, all pseudoscalar meson form factors are predicted in QCD to reverse sign from space-like to time-like asymptotic momentum transfer because of their essentially monopole form. For $m_h > 2m_l$ the form factor zero occurs in the physical region. An interesting question is whether this type of numerator zero structure applies to the gluonic diagram amplitudes appropriate to $\bar{p}p$ reactions

11. EXCLUSIVE NUCLEAR REACTIONS — REDUCED AMPLITUDES

An ultimate goal of QCD phenomenology is to describe the nuclear force and the structure of nuclei in terms of quark and gluon degrees of freedom. Explicit signals of QCD in nuclei have been elusive, in part because of the fact that an effective Lagrangian containing meson and nucleon degrees of freedom must be in some sense equivalent to QCD if one is limited to low-energy probes. On the other hand, an effective local field theory of nucleon and meson fields cannot correctly describe the observed off-shell falloff of form factors, vertex amplitudes, Z-graph diagrams, etc. because hadron compositeness is not taken into account.

The distinction between the QCD and other treatments of nuclear amplitudes is particularly clear in the reaction $\gamma d \rightarrow np$; i.e., photodisintegration of the deuteron at fixed center of mass angle. Using dimensional counting, the leading power-law prediction from QCD is simply $\frac{d\sigma}{d\Omega}(\gamma d \rightarrow np) \sim \frac{1}{s^{1/2}} F(\theta_{cm})$. Of course in the nuclear amplitude, the virtual momenta are partitioned among many quarks and gluons, so that finite mass corrections will be significant at low to medium energies. Nevertheless, there is an elegant way to test the basic QCD dynamics in these reactions taking into account much of the finite-mass, higher-twist corrections by using the "reduced amplitude" formalism.^[10] The basic observation is that for vanishing nuclear binding energy $\epsilon_d \rightarrow 0$, the deuteron can be regarded as two nucleons sharing the deuteron four-momentum (see fig. 32). The momentum ℓ is limited by the binding and can thus be neglected. Thus the photodisintegration amplitude contains the probability amplitude (i.e., nucleon form factors) for the proton and neutron to each remain intact after absorbing momentum transfers $p_p = 1/2 p_d$ and $p_n = 1/2 p_d$, respectively. After the form factors are removed, the remaining "reduced" amplitude should scale as $F(\theta_{cm}/p_T)$. The single inverse power of transverse momentum p_T is the slowest conceivable in any theory, but it is the unique power predicted by PQCD.

Fig. 32. Construction of the reduced nuclear amplitude for two-body inelastic deuteron reactions.^[10]

A test of the prediction that $f(\theta_{cm})$ is energy dependent at high-momentum transfer is compared with experiment in fig. 33. It is particularly striking to see the QCD prediction verified at incident photon lab energies as low as 1 GeV. A comparison with a standard nuclear physics model with exchange currents is also shown for comparison in fig. 33(a). The fact that this prediction falls less fast than

Fig. 39. Comparison of deuteron photodisintegration data with the scaling prediction which requires $f^2(\theta_{cm})$ to be at most logarithmically dependent on energy at large momentum transfer. The data in (a) are from the recent experiment of ref. 46. The nuclear physics prediction shown in (a) is from ref. 96. The data in (b) are from ref. 97.

the data suggests that meson and nucleon compositeness are not taken into account correctly.

There are a number of related tests of QCD and reduced amplitudes which require $\bar{p}$ beams,¹⁰⁴ including $\bar{p}d \rightarrow \gamma n$ and $\bar{p}d \rightarrow \pi^- p$ in the fixed θ_{cm} region. These reactions are particularly interesting tests of QCD in nuclei. Dimensional counting rules predict the asymptotic behavior $\frac{d\sigma}{dt}(\bar{p}d \rightarrow \pi^- p) \sim \frac{1}{(s_T)^{1/2}} f(\theta_{cm})$ since there are 14 initial and final quanta involved. One cannot expect the onset of such scaling laws until p_T is well into the multi-GeV regime since each hard propagator must carry significant momentum transfer. However, again one notes that the $\bar{p}d \rightarrow \pi^- p$ amplitude contains a factor representing the probability amplitude (i.e., form factor) for the proton to remain intact after absorbing momentum transfer squared $\hat{t} = (p - 1/2pd)^2$ and the $\bar{N}N$ time-like form factor at $\hat{s} = (\bar{p} + 1/2pd)^2$. Thus $M_{\bar{p}d \rightarrow \pi^- p} \sim F_{1N}(\hat{t}) F_{1N}(\hat{s}) M_r$, where M_r has the same QCD scaling properties as quark meson scattering. We thus predict

$$\frac{\frac{d\sigma}{dt}(\bar{p}d \rightarrow \pi^- p)}{F_{1N}^2(\hat{t}) F_{1N}^2(\hat{s})} \sim \frac{f(\theta)}{p_T^2}.$$

The analogous analysis of the deuteron form factor as defined in

$$\frac{d\sigma}{dt}(\ell d \rightarrow \ell d) = \frac{d\sigma}{dt} \Big|_{point} (F_d(Q^2))^2$$

yields a scaling law for the reduced form factor (see fig. 34):

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

i.e., the same scaling law as a meson form factor. As shown in fig. 35, this scaling is consistent with experiment for $Q = p_T \gtrsim 1$ GeV. We have also seen the evidence⁴⁶

for reduced amplitude scaling for $\gamma d \rightarrow pn$ at large angles and $p_T \gtrsim 1$ GeV. (see fig. 33). We thus expect similar precocious scaling behavior to hold for $\bar{p}d \rightarrow \pi^+ p$ and other $\bar{p}d$ exclusive reduced amplitudes.

Fig. 34. Application of the reduced amplitude formalism to the deuteron form factor at large momentum transfer.

Fig. 35. Scaling of the deuteron reduced form factor. The data are summarized in ref. 95.

12. QUASI-EXCLUSIVE NUCLEAR PROCESSES

As discussed in sec. 2, one of the most novel features of QCD is "color transparency" which predicts a small absorption cross section for hadrons in specific kinematic configurations.^[100] This concept can be tested in quasiexclusive antiproton-nuclear reactions. For large p_T one predicts

$$\frac{d\sigma}{dt dy} (\bar{p}A \rightarrow \pi^+ \pi^- + (A-1)) \simeq \sum_{p \in A} G_{p/A}(y) \frac{d\sigma}{dt} (\bar{p}p \rightarrow \pi^+ \pi^-) ,$$

where $G_{p/A}(y)$ is the probability distribution to find the proton in the nucleus with light-cone momentum fraction $y = (p^0 + p^x)/(p_A^0 + p_A^x)$, and

$$\frac{d\sigma}{dt}(\bar{p}p \rightarrow \pi^+\pi^-) \simeq \left(\frac{1}{p_T^2}\right)^8 f(\cos \theta_{cm}) .$$

The distribution $G_{p/A}(y)$ can be measured in $eA \rightarrow ep(A-1)$ quasiexclusive reactions. A remarkable feature of the above prediction is that there are no corrections required from initial-state absorption of the $\bar{p}$ as it traverses the nucleus, nor final-state interactions of the outgoing pions. The point is that the only part of hadron wavefunctions which is involved in the large p_T reaction is $\psi_H(b_\perp \sim O(1/p_T))$, i.e., the amplitude where all the valence quarks are at small relative impact parameter. These configurations correspond to small color singlet states which, because of color cancellations, have negligible hadronic interactions in the target. Measurements of these reactions thus test a fundamental feature of the Fock state description of large p_T exclusive reactions.

Another interesting feature which can be probed in such reactions is the behavior of $G_{p/A}(y)$ for y well away from the Fermi distribution peak at $y \sim m_N/M_A$. For $y \rightarrow 1$ spectator counting rules^[101] predict $G_{p/A}(y) \sim (1-y)^{2N_s-1} = (1-y)^{6A-7}$ where $N_s = 3(A-1)$ is the number of quark spectators required to "stop" ($y_i \rightarrow 0$) as $y \rightarrow 1$. This simple formula has been quite successful in accounting for distributions measured in the forward fragmentation of nuclei at the BEVALAC.^[102]

13. COLOR TRANSPARENCY AND QUASIEXCLUSIVE J/ψ PRODUCTION IN $\bar{p}A$ COLLISIONS

Novel features of QCD, including color transparency, can be studied by measuring quasiexclusive J/ψ production by anti-protons in a nuclear target. We are particularly interested in the quasiexclusive annihilation process $\bar{p}A \rightarrow J/\psi(A-1)$ where the nucleus is left in a ground or excited state, but extra hadrons are not created. (See fig. 36.) The cross section involves a convolution of the $\bar{p}p \rightarrow J/\psi$ subprocess cross section with the distribution $G_{p/A}(y)$ where $y = (p^0 + p^3)/(p_A^0 + p_A^3)$ is the boost-invariant light-cone fraction for protons in the nucleus. This distribution can be determined from quasiexclusive lepton-nucleon scattering $\ell A \rightarrow \ell p(A-1)$.

Fig. 36. Schematic representation of quasielastic charmonium production in $\bar{p}A$ reactions.

In first approximation $\bar{p}p \rightarrow J/\psi$ involves $qqq + \bar{q}\bar{q}\bar{q}$ annihilation into three charmed quarks. The transverse momentum integrations are controlled by the charm mass scale and thus only the Fock state of the incident antiproton which contains three antiquarks at small impact separation can annihilate. Since this state has a

relatively small color dipole moment it should have a longer than usual mean-free path in nuclear matter; i.e., "color transparency." Thus unlike traditional expectations, QCD predicts that the $\bar{p}p$ annihilation into charmonium is not restricted to the front surface of the nucleus. The exact nuclear dependence also depends on the formation time for the physical $\bar{p}$ to couple to the small $\bar{q}q\bar{q}$ configuration, $\tau_F \propto E_p$. It may be possible to study the effect of finite formation time by varying the beam energy, E_p , and using the Fermi-motion of the nucleon to stay at the J/ψ resonance.

Since the J/ψ is produced at nonrelativistic velocities in this low energy experiment, it is formed inside the nucleus. The A -dependence of the quasiexclusive reaction can thus be used to determine the J/ψ -nucleon cross section at low energies. For a normal hadronic reaction $\bar{p}A \rightarrow HX$, we expect $A_{eff} \sim A^{1/3}$, corresponding to absorption in the initial and final state. In the case of $\bar{p}A \rightarrow J/\psi X$ we expect A_{eff} much closer to A^1 if color transparency is fully effective and $\sigma(J/\psi N)$ is small.

14. INCLUSIVE $\bar{p}$ REACTIONS AND THE QCD CRITICAL LENGTH

The factorization structure of QCD implies that the structure functions and distribution amplitudes that control high momentum transfer reactions are process-independent. The proofs are highly nontrivial. In the case of inclusive massive lepton-pair production (The Drell-Yan process), the $\bar{p}p \rightarrow \ell\bar{\ell}X$ cross section to leading order in $1/Q^2$ takes the form (see fig. 37):

$$\frac{d\sigma}{dx_a dx_b d\Omega} = \frac{1}{3} \sum_{q\bar{q}} G_{q/p}(x_a, Q) G_{\bar{q}/p}(x_b, Q) \frac{d\sigma}{d\Omega}(q\bar{q} \rightarrow \mu^+\mu^-).$$

This factorization separates the long-distance (nonperturbative) dynamics contained in the universal-process independent structure functions $G_{q/p} = G_{\bar{q}/p}$ from the short-distance perturbative physics contained in the subprocess $q\bar{q} \rightarrow \mu^+\mu^-$ cross section. Antiproton tests of this classic QCD prediction are crucial since the beam and target structure functions for the valence quark and antiquarks are measured directly in deep inelastic lepton scattering.

Fig. 37. Schematic representation of factorization of the Drell-Yan cross section in QCD.

Despite the simple form of the inclusive cross section, all-orders factorization for the Drell-Yan process has only just recently been analyzed to all orders in perturbation theory (by G. Bodwin^[109] and J. Collins, D. Soper and G. Sterman.^[104]) The most serious complications are due to the elastic and inelastic initial-state hadronic interactions which potentially could affect the color correlations, and momentum distribution of the annihilating q and $\bar{q}$. (See fig. 38.) Clearly such effects ruin factorization in a macroscopic target. In fact, as shown in ref. 26 a necessary condition to eliminate the initial state effects is that the incident parton energy must be large compared to a scale proportional to the length of the target. This translates into a necessary condition for factorization: $Q^2 = M_N L \mu^2/x_b$. For a uranium target this implies that factorization can only be valid if the lepton pair mass is greater than

Fig. 38. Induced radiation from the interaction of the active antiquark with target spectators in the Drell-Yan process. The inelastic interactions are suppressed at parton energies which are large compared to a scale set by the length of the target.^[26]

a few GeV at large x_b . It is clearly interesting to study this phenomena experimentally, since it involves the transition between perturbative and soft dynamics and the propagation of antiquarks in nuclear matter. This important area of physics could be studied systematically using a low to medium energy $\bar{p}$ beam.

These results have striking implications for the interaction of the recoil quark jet in deep inelastic electron-nucleus scattering, and other inclusive reactions involving recoil jets. For the quark (and gluons) satisfying the length condition, there should be no extra radiation induced as the parton traverses the nucleus. However, low-energy gluons, emitted in the deep inelastic electron-quark collision, can suffer radiative losses, leading to cascading of soft particles in the nucleus. It is clearly very important to study this phenomena as a function of recoil quark energy and nuclear size. It should be emphasized that the absence of inelastic initial- or final-state collisions for high-energy partons does not preclude collision broadening due to elastic initial- or final-state interactions. The elastic corrections are unitary to leading order in $1/Q$ and do not effect the normalization of the deep inelastic cross section. Thus one predicts that the mean square transverse momentum of the recoil quark and its leading particles will increase as $A^{1/3}$.

The transverse momentum of the recoil quark reflects the intrinsic transverse momentum of the nucleon wave function. The EMC effect^[100] implies that quarks in a nucleus have smaller average longitudinal momentum than in a nucleon.^[101] Independent of the specific physical mechanism underlying the EMC effect,^[102] the quarks in a nucleus would also be expected to have smaller transverse momentum. This effect can counteract to a certain extent the collision broadening of the outgoing jet.

Unlike the struck quark the remnant of the target system does not evolve with the probe momentum Q . However, the quantum numbers of the spectator system is $\bar{3}$ in color, so nonperturbative hadronization must occur. Since the transverse momentum of the leading particles in the spectator jet is not affected by the QCD radiative corrections, it more closely reflects the intrinsic transverse momentum of the hadron state.

It is also interesting to study the behavior of the transverse momentum of the quark and spectator jets as a function of x_{Bj} . For $x_{Bj} \sim 1$, the three-quark Fock

state dominates the reaction. If the valence state has a smaller transverse size^[103] than that of the nucleon, averaged over all of its Fock components, then one expects an increase of $\langle k_{\perp}^2 \rangle$ in that regime. Evidence for a significant increase of $\langle k_{\perp}^2 \rangle$ in the projectile fragmentation region at large quark momentum fractions has been reported by the SFM group^[107] at the ISR for $pp \rightarrow \text{di-jet} + X$ reactions.

It should be noted that the factorization proofs have not yet been extended to reactions such as $p\bar{p} \rightarrow \text{Jet} + \text{Jet} + X$ where the subprocess channel is not in a color singlet. In addition, due to infrared noncancellations, factorization is known to break down beyond leading twist.^[108]

15. DISCRETIZED LIGHT-CONE QUANTIZATION

A central goal of QCD analysis is not only to obtain a complete description of the hadronic spectrum but also to evaluate their current matrix elements. Thus a key problem in the application of QCD to hadron and nuclear physics is how to determine the wavefunction of a relativistic multiparticle composite system. This is obviously a formidable task. Although composite systems in QCD can be represented formally in terms of the covariant Bethe-Salpeter formalism, calculations beyond ladder approximation appear intractable, and the ladder approximation itself is usually inadequate. For example, in order to derive the Dirac equation for the electron in a static Coulomb field from the Bethe-Salpeter equation for muonium with $m_{\mu}/m_e \rightarrow \infty$, one requires an infinite number of irreducible crossed-graph kernel contributions to the QED potential. Similarly, the matrix elements of currents and the wave function normalization also require, at least formally, the consideration of an infinite sum of irreducible kernels. The relative-time dependence of the Bethe Salpeter amplitudes for states with three or more constituent fields adds even more complexities.

A more intuitive procedure would be to extend the Schrödinger wavefunction description of bound states to the relativistic domain by developing a relativistic many-body Fock expansion for the hadronic state. Formally this can be done by quantizing QCD at equal time, and calculating matrix elements from the time-ordered expansion of the S -matrix. However, the calculation of each covariant Feynman diagram with n -vertices requires the calculation of $n!$ frame-dependent time-ordered amplitudes. Even worse, the calculation of the normalization of a bound state wave function (or the matrix element of a charge or current operator) requires the computation of contributions from all amplitudes involving particle production from the vacuum. (Note that even after normal ordering, the interaction Hamiltonian density for QED, $H_I = e : \bar{\psi} \gamma_{\mu} \psi A^{\mu} :$, contains contributions $b^{\dagger} d^{\dagger} a^{\dagger}$ which create particles from the perturbative vacuum.) For this reason, it is not possible to represent a relativistic field-theoretic bound system limited to a fixed number of constituents at a given time in a standard Hamiltonian framework since the interactions create new quanta from the vacuum.^[109] Lorentz invariance is also difficult to incorporate in an equal time formalism.

Fortunately, there is a natural and consistent covariant framework, originally due to Dirac,^[110] (quantization on the "light front") for describing bound states in gauge theory analogous to the Fock state in nonrelativistic physics. This framework is the light-cone quantization formalism in which

$$|\pi\rangle = |q\bar{q}\rangle \psi_{q\bar{q}}^{\pi} + |q\bar{q}g\rangle \psi_{q\bar{q}g}^{\pi} + \dots$$

$$|p\rangle = |qqq\rangle \psi_{qqq}^p + |qqqg\rangle \psi_{qqqg}^p + \dots$$

Each wavefunction component ψ_n describes a state of fixed number of quark and gluon quanta evaluated in the interaction picture at equal light-cone "time" $\tau = t + z/c$. As discussed in the introduction, given the $\{\psi_n\}$, virtually any hadronic property can

be computed, including anomalous moments, form factors, structure functions for inclusive processes, distribution amplitudes for exclusive processes, etc. As shown by Drell and Yan, space-like form factors are given by a simple overlap of the light-cone wavefunctions, summed over Fock states.^[111] At high momentum transfer only the valence Fock-state enters, to leading order in $1/Q$.

As noted above, in an equal time formalism one must allow for fluctuations in which three or four particles appear with zero total three-momentum. In the light-cone formalism such fluctuations cannot appear since the total k^+ is conserved and each particle has to have positive k^+ . Accordingly, the perturbative vacuum is an eigenstate of the total Hamiltonian on the light cone. Light-cone quantization and equal τ wave functions, rather than equal t wavefunctions, thus provides a sensible Fock state expansion. It also turns out to be convenient to use τ -ordered light-cone perturbation theory in place of covariant perturbation theory to analyze light cone dominated processes such as deep inelastic scattering and large momentum transfer exclusive reactions. Light-cone quantization and perturbation theory are developed in detail in ref. 23.

H. C. Pauli and I^[20] have proposed a direct approach to solving QCD by attempting to diagonalize the light-cone Hamiltonian on a free particle discretized momentum Fock state basis. Since H_{LC} , P^+ , $\vec{P}_\perp$, and the conserved charges all commute, H_{LC} is block diagonal. By choosing periodic (or antiperiodic) boundary conditions for the basis states along the negative light-cone $\psi(x^- = +L) = \pm\psi(x^- = -L)$, the Fock basis becomes restricted to finite dimensional representations. The eigenvalue problem thus reduces to the diagonalization of a finite Hermitian matrix. To see this, note that periodicity in x^- requires $P^+ = \frac{2\pi}{L}K$, $k_i^+ = \frac{2\pi}{L}n_i$, $\sum_{i=1}^n n_i = K$. The dimension of the representation corresponds to the number of partitions of the integer K as a sum of positive integers n . For a finite resolution K , the wavefunction is sampled at the discrete points

$$x_i = \frac{k_i^+}{P^+} = \frac{n_i}{K} = \left\{ \frac{1}{K}, \frac{2}{K}, \dots, \frac{K-1}{K} \right\}.$$

The continuum limit is clearly $K \rightarrow \infty$.

One can easily show that P^- scales as L . We thus define $P^- \equiv \frac{L}{2\pi}H$. The eigenstates with $P^2 = M^2$ at fixed P^+ and $\vec{P}_\perp = 0$ thus satisfy $H_{LC}|\Psi\rangle = KH|\Psi\rangle = M^2|\Psi\rangle$, independent of L (which corresponds to a Lorentz boost factor).

The basis of the DLCQ method is thus conceptually simple: one quantizes the independent fields at equal light-cone time τ and requires them to be periodic or antiperiodic in light-cone space with period $2L$. The commuting operators, the light-cone momentum $P^+ = \frac{2\pi}{L}K$ and the light cone energy $P^- = \frac{L}{2\pi}H$ are constructed explicitly in a Fock space representation and diagonalized simultaneously. The eigenvalues give the physical spectrum: the invariant mass squared $M^2 = P^+P^-$. The eigenfunctions give the wavefunctions at equal τ and allow one to compute the current matrix elements, structure functions, and distribution amplitudes required for physical processes. All of these quantities are manifestly independent of L , since $M^2 = P^+P^- = HK$. Lorentz-invariance is violated by periodicity, but re-established at the end of the calculation by going to the continuum limit: $L \rightarrow \infty$, $K \rightarrow \infty$ with P^+ finite. In the case of gauge theory, the use of the light-cone gauge $A^+ = 0$ eliminates negative metric states in both Abelian and non-Abelian theories.

Since continuum as well as single hadron color singlet hadronic wavefunctions are obtained by the diagonalization of H_{LC} , one can also calculate scattering amplitudes as well as decay rates from overlap matrix elements of the interaction Hamiltonian for the weak or electromagnetic interactions. An important point is that all higher Fock amplitudes including spectator gluons are kept in the light-cone quantization

approach; such contributions cannot generally be neglected in decay amplitudes involving light quarks.

Eller, Pauli and I^[112] have used DLCQ to obtain detailed results for the bound state and continuum spectrum and wavefunctions for QED in one-space and one-time dimension for arbitrary mass and coupling constant. I will give here only a brief discussion of the method. The commuting operators K , Q and $H = H_0 + V$ have the form

$$K = \sum n(b_n^\dagger b_n + d_n^\dagger d_n) + n(a_n^\dagger a_n)$$

$$Q = \sum (b_n^\dagger b_n - d_n^\dagger d_n)$$

$$H_0 = \sum \frac{m^2}{n} (b_n^\dagger b_n + d_n^\dagger d_n) + \frac{k_\perp^2}{n} a_n^\dagger a_n$$

$$V = \frac{g^2}{\pi} \sum_{n \neq m, k \neq \ell} b_k^\dagger b_\ell d_m^\dagger d_n \frac{\delta_{n+k, m+\ell}}{(n-m)^2} + \dots$$

Only the one fermion antifermion (Abelian) interaction, corresponding to "instantaneous" gluon exchange, is displayed. The $Q = 0$ Fock state basis states are of the form $b_n^\dagger d_m^\dagger a_\ell^\dagger |0\rangle = |n; m; \ell\rangle$, ($n + m + \ell = K$) where $|0\rangle$ is the perturbative vacuum. (Spin, color and transverse momentum for any number of dimensions are represented as extra internal variables.) We then solve $HK|\Psi\rangle = M^2|\Psi\rangle$ on the free particle basis $|\Psi\rangle = \sum_i C_i |i\rangle$. Note that the eigenvalues of H_{LC} give not only the bound state spectrum, but also all of the multiparticle scattering states with the same quantum numbers.

In the case of gauge theory in 3+1 dimensions, one also takes the $k_\perp^i = (2\pi/L_\perp)n_\perp^i$ as discrete variables on a finite cartesian basis. The theory is covariantly regulated if one restricts states by the condition

$$\sum_i \frac{k_{\perp i}^2 + m_i^2}{x_i} \leq \Lambda^2, \quad$$

where Λ is the ultraviolet cutoff. In effect, states with total light-cone kinetic energy beyond Λ^2 are cut off. In a renormalizable theory physical quantities are independent of physics beyond the ultraviolet regulator; the only dependence on Λ appears in the coupling constant and mass parameters of the Hamiltonian, consistent with the renormalization group.^[113] The resolution parameters need to be taken sufficiently large such that the theory is controlled by the continuum regulator Λ , rather than the discrete scales of the momentum space basis.

The simplest application of DLCQ to local gauge theory is QED in one-space and one-time dimensions. Since $A^+ = 0$ is a physical gauge there are no photon degrees of freedom. Explicit forms for the matrix representation of H_{QED} are given in ref. 112. The basic interactions which occur in $H_{LC}(\text{QCD})$ are illustrated in fig. 39.

For the general case $m^2 \neq 0$, $(\text{QED})_{1+1}$ can be solved by numerical diagonalization. The complete charge-zero spectrum (normalized to the ground state mass) for $K = 16$ is shown as a function of coupling constant in fig. 40. Since the physics can only depend on the ratio m/g , it is convenient to introduce the parametrization $\lambda = \sqrt{\frac{1}{1+\pi(m/g)^2}}$ which maps the entire range of m and g onto the finite interval $0 \leq \lambda \leq 1$.

In the zero coupling limit the spectrum is that of the free theory. In the infinite coupling limit $\lambda = 1$ the theory is essentially equivalent to the limit of zero fermion

Fig. 39. Diagrams which appear in the interaction Hamiltonian for QCD on the light cone. The propagators with horizontal bars represent instantaneous gluon and quark exchange which arise from reduction of the dependent fields in $A^+ = 0$ gauge.

mass. Schwinger has shown that massless $(QED)_{1+1}$ is equivalent to a free boson theory. In the light-cone formalism one can solve the $m = 0$ theory explicitly. One defines^[114] bilinear operators in the fermion fields a_n and $a_n^\dagger$ which have normal boson commutation rules. Then for $Q = 0$

$$H = m^2 \sum_{n=1}^{\infty} \frac{1}{n} (b_n^\dagger b_n + d_n^\dagger d_n) + \frac{g^2}{\pi} \sum_{n=1}^{\infty} \frac{1}{n} a_n^\dagger a_n .$$

Thus for $m^2 = 0$ (or $g^2/\pi \gg m^2$), H_{QED} is equivalent to free boson theory with $m_b^2 = g^2/\pi$. The distinction between the theories in the limit of zero fermion mass is discussed by McCartor.^[115]

Fig. 40. Spectrum of QED in one-space and one-time dimension for harmonic resolution $K = 16$. The ratios M_i/M_1 are plotted as a function of the scaled coupling constant λ . The Schwinger limit is $\lambda = 1$. (From ref. 112.)

Figure 41 shows the structure function for the ground state of $(QED)_{1+1}$ as a function of λ . In the weak binding limit $g \rightarrow 0$ or $(m \rightarrow \infty)$, the structure function becomes a delta function at equal partition of the constituent momentum, as expected. In the strong coupling limit $g \rightarrow \infty$ ($m \rightarrow 0$) the structure function becomes flat. This is consistent with the interpretation of the Schwinger boson as a point-like composite of a fermion and antifermion. The contribution to higher Fock states to the lowest mass structure function is strikingly small; the probability of nonvalence states is less than 1% for any value of λ .

Fig. 41. The structure function of the lowest mass bound state for QED in $1+1$ space-time dimensions, as calculated in the $DLCQ$ formalism.^[112]

It is interesting that there is analytic agreement between the $DLCQ$ results and the exact solutions of the Schwinger model for finite K , as well as in the continuum limit. This can be traced to the fact that the structure function of the Schwinger boson is flat and thus needs minimal resolution. In the case of the massive Schwinger model (QED_2), we established the existence of the continuum limit numerically; for sufficiently large resolution K the results become independent of K . The essential criteria for convergence is that the intrinsic dynamical structure of the wavefunctions is sufficiently resolved at the rational values $x = n/K$, $n = 1, 2, \dots, K-1$ accessible at a given K .

In the large K limit, the eigenvalues agree quantitatively with the results of Bergknoff^[114] and with those of a lattice gauge calculation by Crewther and Hamer.^[116] This result is important in establishing the equivalence of different complementary nonperturbative methods. We also verified numerically that different Fock space representations yield the same physical results. In particular, we solved the QED_2 spectrum in the space corresponding to the solutions of the free, massive Dirac equation $(i\gamma^\mu \partial_\mu + m_F)\psi = 0$ as well as of the massless equation $i\gamma^\mu \partial_\mu \psi = 0$. Convergence is slow in $1/K$ only at very large coupling λ near 1.

Even for moderately large values of the resolution, $DLCQ$ provides one with a qualitatively correct picture of the whole spectrum of eigenfunctions. This aspect becomes important for the development of scattering theory within the $DLCQ$ approach. For example, we have found the rather surprising result that the lowest eigenfunction has very small probability (less than 1%) for $|2f; 2\bar{f}\rangle$ and higher particle Fock states (i.e., no 'sea quarks'). We have also obtained the spectrum of the Yukawa theory with spin-zero bosons, a theory with a more complicated Fock structure. Also Harindranath and Vary^[117] have recently used a $DLCQ$ approach to analyze ϕ^4 theory, a model with a nontrivial vacuum structure.

Recently Hornbostel^[118] has used DLCQ to obtain the complete color-singlet spectrum of QCD in one space and one time dimension for $N_C = 2, 3, 4$. The hadronic spectra are obtained as a function of quark mass and QCD coupling constant. (See fig. 42.) Where they are available, the spectra agree with results obtained earlier; in particular, the lowest meson mass in SU(2) agrees within errors with lattice Hamiltonian results.^[120] The meson mass at $N_C = 4$ is close to the value obtained in the large N_C limit. The method also provides the first results for the baryon spectrum in a non-Abelian gauge theory. The lowest baryon mass is shown in fig. 42 as a function of coupling constant. The ratio of meson to baryon mass as a function of N_C also agrees at strong coupling with results obtained by Frishman and Sonnenschein.^[119] Precise values for the mass eigenvalue can be obtained by extrapolation to large K since the functional dependence in $1/K$ is understood.

Fig. 42. The baryon and meson spectrum in QCD [1+1] computed in DLCQ for $N_C = 2, 3, 4$ as a function of quark mass and coupling constant. (From ref. 118.)

As emphasized above, when the light-cone Hamiltonian is diagonalized for a finite resolution K , one gets a complete set of eigenvalues corresponding to the total dimension of the Fock state basis. A representative example of the spectrum is shown in fig. 43 for baryon states ($B = 1$) as a function of the dimensionless variable $\lambda = 1/(1 + \pi m^2/g^2)$. Antiperiodic boundary conditions are used. Note that spectrum automatically includes continuum states with $B = 1$.

The structure functions for the lowest meson and baryon states in SU(3) at two different coupling strengths $m/g = 1.6$ and $m/g = 0.1$ are shown in figs. 44 and 45. Higher Fock states have a very small probability; representative contributions to the baryon structure functions are shown in figs. 46 and 47. For comparison, the valence wavefunction of a higher mass state which can be identified as a composite of meson

Fig. 43. Representative baryon spectrum for QCD in one-space and one-time dimension. (From ref. 118.)

Fig. 44. The meson quark momentum distribution in QCD[1+1] computed using DLCQ. (From ref. 118.)

Fig. 45. The baryon quark momentum distribution in QCD[1+1] computed using DLCQ. (From ref. 118.)

pairs (analogous to a nucleus) is shown in fig. 48. The interactions of the quarks in the pair state produce Fermi motion beyond $x = 0.5$.

There are a number of important advantages of the DLCQ method which have emerged from this study of two-dimensional field theories. They are as follows:

Fig. 46. Contribution to the baryon quark momentum distribution from $qqqq\bar{q}$ states for $QCD[1+1]$. (From ref. 118.)

Fig. 47. Contribution to the baryon quark momentum distribution from $qqqqq\bar{q}$ states for $QCD[1+1]$. (From ref. 118.)

Fig. 48. Comparison of the meson quark distributions in the $qqq\bar{q}$ Fock state with that of a continuum meson pair state. The structure in the former may be due to the fact that these four-particle wavefunctions are orthogonal. The analysis is for $N_C = 2$ in $1+1$ dimensions. (From ref. 118.)

1. The Fock space is denumerable and finite in particle number for any fixed resolution K . In the case of gauge theory in $3+1$ dimensions, one expects

that photon or gluon quanta with zero four-momentum decouple from neutral or color-singlet bound states, and thus need not be included in the Fock basis. The transverse momenta are additive and can be introduced on a cartesian grid. Hornbostel^[110] has developed methods to implement the color degrees of freedom for the non-Abelian theories. Tang^[121] is currently studying QED[3+1] in DLCQ as a function of the QED coupling constant.

2. Because we are using discrete momentum space representation, rather than a space-time lattice, there are no special difficulties with fermions: e.g., no fermion doubling, fermion determinants, or necessity for a quenched approximation. Furthermore, the discretized theory has basically the same ultraviolet structure as the continuum theory. It should be emphasized that unlike lattice calculations, there is no constraint or relationship between the physical size of the bound state and the length scale L .
3. The DLCQ method has the remarkable feature of generating the complete spectrum of the theory; bound states and continuum states alike. These can be separated by tracing their minimum Fock state content down to small coupling constant since the continuum states have higher particle number content. In lattice gauge theory it appears intractable to obtain information on excited or scattering states or their correlations. The wavefunctions generated at equal light cone time have the immediate form required for relativistic scattering problems.
4. DLCQ is basically a relativistic many-body theory, including particle number creation and destruction, and is thus a basis for relativistic nuclear and atomic problems. In the nonrelativistic limit the theory is equivalent to the many-body Schrödinger theory.

The immediate goal is gauge theory in 3+1 dimensions. Already Klabucar and Pauli^[122] have studied QCD[3+1] in the $q\bar{q}$ sector for strong coupling. In the Abelian case it will be interesting to analyze QED and the positronium spectrum in the large α limit. Whether the non-Abelian theory can be solved using DLCQ — considering its greater number of degrees of freedom and its complex equal-time vacuum is an open question. The studies for Abelian and non-Abelian gauge theory carried out so far in 1+1 dimensions give grounds for optimism.

16. SUMMARY

With the advent of new methods to attack nonperturbative QCD, such as sum rule constraints, implementation of effective Lagrangians such as the Skyrme model, extensions of lattice gauge theory, and promising methods to solve the light-cone Hamiltonian for its spectrum and Fock state solutions, a renaissance of interest is developing for understanding hadron and nuclear dynamics from first principles.

As I have outlined in these lectures, An experimental program with antiprotons of energies under 10 GeV can serve as an important test of QCD dynamics, especially for exclusive channels. Already there are extensive calculations available for $\bar{p}p \rightarrow \gamma\gamma$ for both real and virtual channels. Fixed angle scattering, meson-pair and lepton-pair final states also give sensitive tests of the theory. We have emphasized the possibility that the production of charmed hadrons in exclusive $\bar{p}p$ channels may have a non-negligible cross section. All of these channels bear on the question at what momentum scale perturbative factorization methods apply.

Inclusive measurements are usually studied at much higher energies than those potentially available at a SUPER-LEAR or AMPLÉ facility. Nevertheless, as discussed in these lectures, there are interesting novel effects involving the interface between perturbative and nonperturbative dynamics and quark propagation in hadronic

matter — all of which can be explored at $\bar{p}$ energies below 10 GeV. In particular, $\bar{p}$ — nuclear collisions can play an important role in clarifying fundamental QCD issues such as color transparency, critical length phenomena, and the validity of the reduced nuclear amplitude phenomenology.

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