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NUCLEAR CHROMODYNAMICS: APPLICATIONS OF QCD TO RELATIVISTIC MULTIQUARK SYSTEMS*

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ABSTRACT

We review the applications of quantum chromodynamics to nuclear multiquark systems. In particular, predictions are given for the deuteron reduced form factor in the high momentum transfer region, hidden color components in nuclear wavefunctions, and the short distance effective force between nucleons. A new antisymmetrization technique is presented which allows a basis for relativistic multiquark wavefunctions and solutions to their evolution to short distances. Areas in which conventional nuclear theory conflicts with QCD are also briefly reviewed.

INTRODUCTION

Nuclear chromodynamics is concerned with the application of quantum chromodynamics to nuclear physics. Its goal is to give a fundamental description of nuclear dynamics and nuclear properties in terms of quark and gluon fields at short distance, and to obtain a synthesis at long distances with the normal nucleon, isobar, and meson degrees of freedom. Nuclear chromodynamics provides an important testing ground for coherent effects in QCD and nuclear effects at the interface between perturbative and non-perturbative dynamics.[♯1]

Among the areas of interest:

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♯1 For additional discussion of applications of QCD to nuclear physics and references, see S. J. Brodsky, published in the proceedings of the conference "New Horizons in Electromagnetic Physics", University of Virginia, April 1982; S. J. Brodsky, T. Huang and G. P. Lepage, SLAC-PUB-2868 (1982) published in Springer Tracts in Modern Physics, Vol. 100, "Quarks and Nuclear Forces", ed. D. Fries and B. Zeitnitz (1982); S. J. Brodsky and G. P. Lepage, in the proceedings of the Eugene Few Body Conference 1980: 247C (Nucl. Phys. A363, 1981) and S. J. Brodsky, in the proceedings of the NATO Pacific Summer Institute Progress in Nuclear Dynamics, Vancouver Island (1982).

Lectures presented by SJB at the International School of Nuclear Physics, Erice, Italy, April 8-20, 1984

1. The representation of the nuclear force in terms of quark and gluon subprocesses.[♮2] The nuclear force between nucleons can in principle be represented at a fundamental level in QCD in terms of quark interchange (equivalent at large distances to pion and other meson exchange) and multiple-gluon exchange. Although calculations from first principles are still too complicated, recent results derived from effective potential, bag, and soliton models suggest that many of the basic features of the nuclear force can be understood from the underlying QCD substructure. At a more basic level we will show directly from QCD that the nucleon-nucleon force must be repulsive at short distances. At high momentum transfer the nucleon-nucleon interactions agree with the scaling laws predicted by the simplest constituent exchange processes.

2. The composition of the nucleon and nuclear state in terms of quark and gluon quanta. The light-cone quantization formalism provides a consistent relativistic Fock state momentum space representation of multiquark and gluon color singlet bound states.

3. The propagation of quarks and gluons through nuclear matter: one is interested in the interplay between multiple scattering,[♮3] induced radiation, the Landau-Pomeranchuk coherence effect,[♮4] shadowing phenomena, and confinement.

4. Factorization theorems for inclusive and exclusive[♮5] reactions: for nuclear reactions one is particularly interested in the origin of the EMC non-additivity effect[♮6] and other nuclear-induced effects in high transverse momentum reactions.

♮2 For a recent discussion of progress in the derivation of nuclear forces from QCD-based models, see K. Maltman and N. Isgur, Phys. Rev. Lett. 50, 1827 (1983), E. L. Lomon, MIT preprint CTP No. 1116 (1983); and references therein. The quark interchange mechanism for $N-N$ scattering is discussed in J. F. Gunion, S. J. Brodsky, and R. Blankenbecler, Phys. Rev. D8, 287 (1973). Qualitative QCD-based arguments for the repulsive $N-N$ potential at short distances are given in C. Detar, HU-TFT-82-6 (1982); M. Harvey, Nucl. Phys. A352, 301 (1981); A352, 326 (1981); R. L. Jaffe, Phys. Rev. Lett. 24, 228 (1983); and G. E. Brown, in Erice 1981, Proceedings, Quarks and the Nucleus. The possibility that the deuteron form factor is dominated at large momentum transfer by hidden color components is discussed in V. A. Matveev and P. Sorba, Nuovo Cimento 45A, 257 (1978); Nuovo Cimento 20, 435 (1977).

♮3 The eikonal scattering of quarks in hard scattering processes such as massive lepton pair production, including $k_\perp$ smearing and induced radiation is discussed by G. Bodwin, S. J. Brodsky, and G. P. Lepage, SLAC-PUB-2927 (1982) and S. J. Brodsky, SLAC-PUB-3263 (1984), to be published in Acta Physica Polonica. Other uses of the Drell-Yan process for studying quark and gluon distributions and shadowing processes in nuclei are discussed by L. L. Frankfurt and M. I. Strikman, Leningrad preprint 838 (1983).

♮4 L. Landau and I. Pomeranchuk, Dok. Akademii Nauk SSSR 92, 535 (1953), and 92, 735 (1953); L. Stodolsky, MPI-PAE/TH 23/75 (1981); I. M. Dremin, Lebedev preprint 250 (1981). The Landau-Pomeranchuk effect, which is an essential component of factorization theorems for inclusive reactions, predicts that there is no induced hard radiation within a nucleus for incident constituents of energy large compared to the nucleon length. This leads to new conditions for the validity of QCD factorization for the Drell-Yan process, etc. Details are found in Ref. 3.

♮5 For a comprehensive review of exclusive processes in QCD and further references, see V. L. Chernyak and A. R. Zhitnisky, Novosibirsk preprints 83-105, 83-105, 83-106, 83-107 and 83-108 (1983).

♮6 R. T. Aubert, et al., Phys. Lett. 105B, 315, 322 (1981); Phys. Lett. 123B, 123, 275 (1983). A. Bodek, et al., Phys. Rev. Lett. 50, 1431; 51, 524 (1983). For recent theoretical discussions and references to the EMC effect see e.g., M. Chemtob and R. Peschanski, Saclay preprint SPh.T/83/116 (1983), and J. Pirner (this volume). Various models are discussed in H. J. Pirner and J. P. Vary, Phys. Rev. Lett. 46, 1376 (1981); R. Jaffe, Phys. Rev. Lett. 50, 228 (1983). L. S. Celenza and C. Shakin, Brooklyn College preprint BCINT-82/111/117 (1982). M. Staszel, J. Roznek, G. Wilk, Warsaw preprint IFT19/83 (1983). F. E. Close, R. B. Robert, G. G. Ross, Rutherford preprint RL-83-051 (1983). O. Nachtmann and J. H. Pirner, Heidelberg preprint HD-THE-8-83-8 (1983).

5. Novel nuclear phenomena in QCD such as (a) color coherence effects in high momentum transfer quasi-elastic reactions in nuclei;[♯7] (b) the nuclear number dependence of strange and charm quarks in the sea; (c) new color singlet multiquark states.

6. The use of reduced nuclear amplitudes[♯8] in order to obtain a consistent and covariant identification of the effects of nucleon compositeness in nuclear reactions.

Because of asymptotic freedom, the effective strength of QCD interactions[♯9] becomes logarithmically weak at short distances and large momentum transfer

$$\alpha_s(Q^2) = \frac{4\pi}{\beta \log(Q^2/\Lambda^2_{\rm QCD})} \qquad (Q^2 \gg \Lambda^2) \; . \tag{1}$$

[Here $\beta = 11 - \frac{2}{3} n_f$ is derived from the gluonic and quark loop corrections to the effective coupling constant; n_f is the number of quark contributions to the vacuum polarizations with $m_f^2 \lesssim Q^2$.] The parameter $\Lambda_{\rm QCD}$ normalizes the value of $\alpha_s(Q_0^2)$ at a given momentum transfer $Q_0^2 \gg \Lambda^2$, given a specific renormalization or cutoff scheme. Recently α_s has been determined fairly unambiguously using the measured branching ratio for upsilon radiative decay $\Upsilon(b\bar{b}) \to \gamma X$:[♯9,♯10]

$$\alpha_s(0.157\; M_\Upsilon) = \alpha_s(1.5\; GeV) = 0.23 \pm 0.13\; . \tag{2}$$

Taking the standard $\overline{MS}$ dimensional regularization scheme, this gives $\Lambda_{\overline{MS}} = 119^{+52}_{-34}$ MeV. In more physical terms, the effective potential between infinitely heavy quarks has the form $[C_F = 4/3$ for $n_c = 3]$,[♯10]

$$\begin{aligned} V(Q^2) &= -C_F \,\frac{4\pi\alpha_V(Q^2)}{Q^2} \\ \alpha_V(Q^2) &= \frac{4\pi}{\beta\log(Q^2/\Lambda_V^2)} \qquad (Q^2 \gg \Lambda_V^2) \end{aligned} \tag{3}$$

where $\Lambda_V = \Lambda_{\overline{MS}}\, e^{5/6} \simeq 270 \pm 100\; MeV$. Thus the effective physical scale of QCD is $\sim 1\; f_m^{-1}$. At momentum transfers beyond this scale, α_s becomes small, QCD perturbation

♯7 Because of color cancellations, QCD predicts no initial or final state corrections to quasi-elastic high momentum transfer reactions such as $\pi A \to \pi N(A-1)$. See A. H. Mueller, to be published in Proceedings of the Moriond Conference (1982). Applications to elastic hadron-nucleus amplitudes are given in S. J. Brodsky and B. T. Chertok, Phys. Rev. Lett. 37, 269 (1976). Color singlet cancellations for valence states interacting inclusively in nuclei are discussed in G. Bertsch, S. J. Brodsky, A. S. Goldhaber and J. F. Gunion, Phys. Rev. Lett. 47, 297 (1981). Further discussion may be found in S. J. Brodsky, SLAC-PUB-2970 (1982), published in the Proceedings of the XIIIth International Symposium on Multiparticle Dynamics, Volendam, The Netherlands (1982). The generalization of y-scaling to such reactions appears to be successful phenomenologically. [S. Gurvitz, private communication.]

♯8 S. J. Brodsky and J. R. Hiller, Phys. Rev. C28, 475 (1983). Fig. 7 is corrected for a phase-space factor $\sqrt{s/(s-m_d^2)}$. S. J. Brodsky and B. T. Chertok, Phys. Rev. Lett. 37, 269 (1976); Phys. Rev. D14, 3003 (1976). S. J. Brodsky, in Proceedings of the International Conference on Few Body Problems in Nuclear and Particle Physics, Laval University, Quebec (1974).

♯9 C. Klopfenstein, et al., CUSB 83-07 (1983).

♯10 S. J. Brodsky, G. P. Lepage, P. B. Mackenzie, Phys. Rev. D28, 228 (1983).

theory becomes applicable, and a microscopic description of short-distance hadronic and nuclear phenomena in terms of quark and gluon subprocesses becomes viable. In these lectures we will particularly emphasize the use of asymptotic freedom and light-cone quantization[♯11] to derive factorization theorems, rigorous boundary conditions, and exact results for nuclear amplitudes at short distances.[♯12] This includes the nucleon form factor[♯12,♯13,♯14] at large momentum transfer, meson photoproduction amplitudes, deuteron photo- and electro-disintegration[♯8] and most important for nuclear physics, exact results for the form of form factors of nuclei at large momentum transfer.[♯15] Eventually it should be possible to construct fully analytic nuclear amplitudes which at low energies fit the standard chiral constraints and low energy theories of traditional nuclear physics while at the same time satisfying the scaling laws[♯16] and anomalous dimension structure predicted by QCD at high momentum transfer.

Conversely, nuclear chromodynamics implies in some cases a breakdown of traditional nuclear physics concepts. For example, we can identify where off-shell effects modify traditional nuclear physics formulas, such as the impulse approximation for elastic nuclear form factors.[♯17] At high momentum transfer nuclear amplitudes are predicted to have a power law fall off in QCD in contrast to the Gaussian or exponential fall off usually assumed in nuclear physics. There are other areas where conventional techniques in nuclear theory, such as the use of local meson nucleon field theories break down. We discuss some of these problems in the last section of this chapter. When we deal with complex nuclear systems we begin to realize how little is understood about the nucleon system: even the size and the shape parameters of the hadron wave functions in terms of quark and gluon degrees of freedom are still unknown. We briefly discuss in this chapter known constraints on the wave functions of mesons and nucleons.

In QCD, the fundamental degrees of freedom of nuclei as well as hadrons are postulated to be the spin-1/2 quark and spin-1 gluon quanta. Nuclear systems are identified as color-singlet composites of quark and gluon fields, beginning with the six-quark Fock component of the deuteron. An immediate consequence is that nuclear states are a mixture of several color representations which *cannot* be described solely in terms of the conventional nucleon, meson, and isobar degrees of freedom: there must also exist "hidden color" multiquark wavefunction components—nuclear states which are not separable at large distances into the usual color singlet nucleon clusters.

The goal of nuclear chromodynamics is thus to understand the fundamental basis of nuclear amplitudes. Solutions to QCD for bound states eventually may be obtained from lattice gauge

♯11 Details of light-cone Fock methods are given in G. P Lepage, S. J. Brodsky, T. Huang, and P. B. Mackenzie CLNS-82/522, published in proceedings of the Banff Summer Institute on Particle Physics, Alberta, Canada; S. J. Brodsky and G. P. Lepage, Phys. Rev. D24, 1808 (1981) and S. J. Brodsky, in proceeding of Quarks and Nuclear Forces, Springer 100, Bad Liebenzell (1981).

♯12 G. P. Lepage and S. J. Brodsky, Phys. Rev. D22, 2157 (1980).

♯13 S. J. Brodsky, G. P. Lepage, S.A.A. Zaidi, Phys. Rev. D22, 1152 (1981); G. R. Farrar, E. Maina, and F. Neri, Rutgers preprint RU-84-13 (1984).

♯14 V. L. Chernyak and A. R. Zhitnitsky, Novosibirsk preprint 83-103. There is an overall sign disagreement between the results of this reference and that of Refs. 12 and 13.

♯15 S. J. Brodsky, C.-R. Ji, G. P. Lepage, Phys. Rev. Lett. 51, 83 (1983).

♯16 S. J. Brodsky and G. R. Farrar, Phys. Rev. Lett. 31, 1153 (1973), and Phys. Rev. D11, 1309 (1975); V. A. Matveev, R. M. Muradyan and A. V. Tavkheldize, Lett. Nuovo Cimento 7, 719 (1973).

♯17 A detailed discussion will be given in S. J. Brodsky and C.-R. Ji (to be published).

theory or the light-cone quantization formalism. Nevertheless, even without explicit solutions, (1) we can use asymptotic freedom to calculate the underlying quark and gluon subprocess amplitudes at short distances, (2) we can derive factorization theorems for both inclusive and exclusive processes which separate the hadronic bound state physics from perturbative dynamics, and (3) we can use the apparatus of light cone quantization (i.e.: equal time $\tau = t + z/c$ wave functions) to represent bound states of composite systems in a consistent covariant manner. In some cases, we can derive exact constraints on the wave functions, or use approximation methods and sum rules to model the wave functions.[♯11,♯12] We can also derive connections with the non-relativistic wavefunctions.[♯8] In the case of multiquark systems we can derive asymptotic constraints such as the form of the deuteron wavefunction. Using these techniques we can analyze the role of hidden color degrees of freedom in ordinary nuclei, and understand the role of QCD relativistic effects. The introduction of reduced nuclear amplitudes then allows the direct phenomenological study of the specific role of QCD in nuclear physics. Finally, we can derive constraints on the hadronic meson nucleon vertices which are required for calculating meson and exchange currents and similar coherent phenomena.

Just as Bohr's correspondence principle played a crucial role in bridging the gap between classical and quantum mechanics, we also need a similar correspondence principle to bridge the gap between nuclear physics at large distances and QCD at short distances. Since QCD has the same natural length scale $\sim 1\ fm$ as nuclear physics it is difficult to argue that nuclear physics can be studied in isolation from QCD. Thus one of the most interesting questions in nuclear physics is the transition between conventional meson-nucleon degrees of freedom to the quark and gluon degrees of freedom of QCD. As one probes distances shorter than $\Lambda_{\rm QCD}^{-1} \sim 1\ fm$ the meson-nucleon degrees of freedom must break down, and we expect new nuclear phenomena, new physics intrinsic to composite nucleons and mesons, and new phenomena outside the range of traditional nuclear physics. One apparent signal for this is the experimental evidence[♯6,♯18] from deep inelastic lepton-nucleus scattering that nuclear structure functions deviate significantly from simple nucleon additivity,[♯6,♯19] much more than would have been expected for lightly bound systems.[♯20] Further, as discussed in later sections, there are many areas where QCD predictions conflict with traditional concepts of nuclear dynamics.

♯18 The experimental situation is reviewed in this volume by E. Gabathler.

♯19 There are many possible origins of non-additive effects in nuclear structure functions including rearrangements of meson, nucleon, and isobar degrees of freedom (see, e.g., J. Szwed, this volume), increasing size of a nucleon inside a nucleus (C. Shakin, Brooklyn College preprints; C. Wilk, this volume), color conductivity and quark flow between nucleons which are related to meson exchange effects (see J. Pirner, this volume). In addition the effect of QCD radiative corrections can be modified by the nucleus at low momentum transfer (see Pirner and Nachtman). [It should be emphasized the nuclear state is not itself changed (in leading twist) by the momentum transfer of the probe.] Strange quarks in the nucleus would be non-additive if the origin of heavy quarks is due to higher dimension operators (S. Brodsky and M. Soldate, unpublished). Theories of the EMC effect also have to account for the indicated absence of deviation from additivity for low x_{bj} in deep inelastic neutrino reactions and the low energy American University/SLAC data. See R. G. Arnold et al., SLAC-PUB-3320 (1984).

♯20 Actually the small value of binding energy per nucleon is not a reliable guide to the magnitude of non-additivity effects, as witnessed by the strikingly large nuclear effects ($O(20\%)$) observed for the g_ℓ magnetic moment of the nucleons in high orbital states. See T. Yamagishi, in Mesons and Nuclei, Vol. II, (North-Holland, Amsterdam), edited by M. Rho and D. Wilkinson (1979), and T. Yamagishi, this volume.

It is apparently true that for distances greater than 1 fm, i.e.: momentum transfer less than 200 MeV, non-relativistic Schroedinger equation and potential theory give an accurate phenomenological description of nuclear matter. Similarly, in the short distance domain (distances less than 0.2 fm or momentum transfer ≥ 1 GeV) the quark-gluon degrees of QCD give a good representation of strong interaction dynamics. The synthesis between nuclear physics and QCD is then the analogue to the correspondence principle. For example, the nuclear potential can now be understood in terms of quark interchange[♯21] and gluon exchange amplitude at the high momentum transfer region.[♯2] At long distances these contributions must merge into the traditional meson and Yukawa force. The nuclear state, which can be primarily represented as meson and nucleon degrees of freedom at large distances; at short distances must give way to a description in terms of quark and gluon degrees of freedom, specifically hidden color components, at very short distances. The electromagnetic and weak interactions of the nucleus, which is traditionally described in terms of nucleon and meson currents, is replaced in QCD by interactions which couple directly to the quark currents at any momentum transfer scale. What we perceive at large distances and refer to as meson and nucleon degrees of freedom are coherent effects of QCD. The form factor in nuclear physics in the non-relativistic domain can be represented as a Fourier transform of a charge distribution. At relativistic energies, this is replaced by an exact QCD calculation of the probability amplitude for the nuclear system to remain intact. Asymptotic results are given in later sections of this chapter.

The joining of nuclear physics at long distances and QCD at short distances also brings a number of new general analytical tools, including (a) light cone quantization, (b) a relativistic Fock state expansion, (c) factorization theorems, (d) evolution equations which give the leading behavior of hadronic amplitudes at short distances, and (e) a system of counting rules for obtaining the leading power behavior[♯12] and leading helicity behavior of nuclear reduced amplitudes.[♯22] In particular, one now has a completely relativistic framework for multi-particle systems applicable to nuclear systems: light cone quantization provides a Hamiltonian formulation for QCD and is an alternative to the Bethe-Salpeter formalism.

Despite its generality, in concept, and often in practice, light-cone quantization is as simple to use as Schroedinger many body theory.[♯11,♯12] Using this formalism one can readily obtain exact results for the form of the nucleon, meson, and nuclear form factors and other exclusive nuclear amplitudes at large momentum transfer, such as the photo-disintegration of the deuteron at large θ_{CM}. One obtains rigorous constraints on the six-quark wave function of the deuteron at small relative distances as well as a value for the percentage of hidden color at short distances in the deuteron wave function. More generally, as we discuss in a later section, one can identify the degrees of freedom of multi-quark system and obtain a completely anti-symmetrized basis Fock state representation for multi-quark states.

The fact that the degrees of freedom and permutation symmetries of the covariant QCD equation of motion for multiquark states on the light-cone are the same as those of the non-relativistic quark model[♯23] can account for the successes of the non-relativistic approach for

♯21 A new systematic treatment of quark interchange effects in the nucleus is given by J. Hiller, Purdue preprint (1984).

♯22 S. J. Brodsky and G. P. Lepage, Phys. Rev. D24, 2848 (1981).

♯23 See the chapter by G. Karl, this volume.

describing the hadronic spectrum despite the dynamical failure of non-relativistic equations for describing wavefunctions and structure functions.

APPLICATIONS OF QCD TO HIGH MOMENTUM TRANSFER HADRONIC AND NUCLEAR PROCESSES

Although we cannot yet compute hadronic wave functions in quantum chromodynamics, it is still possible to make predictions at large momentum transfer directly from the theory. The results are rigorous and can be proved to arbitrary order in perturbation theory.[♯5,♯12,♯24] The processes which are most easily analyzed are those in which all final particles are measured at large invariant masses compared to each other, i.e.: large momentum transfer exclusive reactions. This includes form factors of hadrons and nuclei at large momentum transfer Q and large angle scattering reactions such as photoproduction $\gamma p \to \pi^+ n$, nucleon-nucleon scattering, photodisintegration $\gamma d \to np$ at large angles and energies, etc. A key result is that such amplitudes factorize at large momentum transfer in the form of a convolution of a hard scattering amplitude T_H which can be computed perturbatively from quark-gluon subprocesses multiplied by process-independent "distribution amplitudes" $\phi(x,Q)$ which contain all of the bound-state non-perturbative dynamics of each of the interacting hadrons.[♯12] To leading order in $1/Q$ the scattering amplitude has the form

$$\mathcal{M} = \int_0^1 T_H(x_j, Q) \prod_{H_i} \phi_{H_i}(x_j, Q)[dx] \; . \tag{4}$$

Here T_H is the probability amplitude to scatter quarks with fractional momentum $0 < x_j < 1$ from the incident to final hadron directions, and ϕ_{H_i} is the probability amplitude to find quarks in the wavefunction of hadron H_i collinear up to the scale Q, and

$$[dx] = \prod_{j=1}^{n_i} dx_j \delta\Big(1 - \sum_k^{n_i} x_k\Big) \tag{5}$$

A key to the derivation of this factorization of perturbative and non-perturbative dynamics is the use of a Fock basis $\{\psi_n(x_i, \vec{k}_{\perp i}, \lambda_i)\}$ defined at equal $\tau = t + z/c$ on the light-cone to represent relativistic color singlet bound states. Here λ_i are the helicities; $x_i \equiv (k_i^0 + k_i^z)/(p^0 + p^z)$, $(\sum_{i=1}^n x_i = 1)$, and $\vec{k}_{\perp i}$, $(\sum_{i=1}^n \vec{k}_{\perp i} = 0)$, are the relative momentum coordinates. Thus the proton is represented as a column vector of states ψ_{qqq}, ψ_{qqqg}, $\psi_{qqqq\bar{q}}$ .... In the light-cone gauge, $A^+ = A^0 + A^3 = 0$, only the minimal "valence" Fock state needs to be considered at large momentum transfer since any additional quark or gluon forced to absorb large momentum transfer yields a power-law suppressed contribution to the hadronic amplitude. For example at large Q^2, the baryon form factor takes the form [Fig. 1(a)]

$$F_B(Q^2) = \int_0^1 [dy] \int_0^1 [dx] \; \phi_B^\dagger(y_j, Q) T_H(x_i, y_j, Q) \phi_B(x_i, Q) \quad , \tag{6}$$

where to leading order in $\alpha_s(Q^2)$, T_H is computed from $3q + \gamma^* \to 3q$ tree graph amplitudes:

♯24 M. Peskin, Phys. Lett. 88B, 128 (1979); A. Duncan and A. H. Mueller, Phys. Lett. 90B, 159 (1980); Phys. Rev. D 21, 1636 (1980).

[Fig. 1(b)]

$$T_H = \left[\frac{\alpha_s(Q^2)}{Q^2}\right]^2 f(x_i, y_j) \tag{7}$$

and

$$\phi_B(x_i, Q) = \int [d^2k_\perp]\ \psi_V(x_i, \vec{k}_{\perp i})\theta(k_{\perp i}^2 < Q^2) \tag{8}$$

is the valence three-quark wavefunction [Fig. 1(c)] evaluated at quark impact separation $b_\perp \sim \mathcal{O}(Q^{-1})$. Since ϕ_B only depends logarithmically on Q^2 in QCD, the main dynamical dependence of $F_B(Q^2)$ is the power behavior $(Q^2)^{-2}$ derived from scaling of the elementary propagators in T_H. Thus, modulo logarithmic factors, one obtains a dimensional counting rule for any hadronic or nuclear form factor at large Q^2 ($\lambda = \lambda' = 0$ or $1/2$)

$$F(Q^2) \sim \left(\frac{1}{Q^2}\right)^{n-1}, \tag{9}$$

$$F_1^N \sim \frac{1}{Q^4}, \quad F_\pi \sim \frac{1}{Q^2}, \quad F_d \sim \frac{1}{Q^{10}}, \tag{10}$$

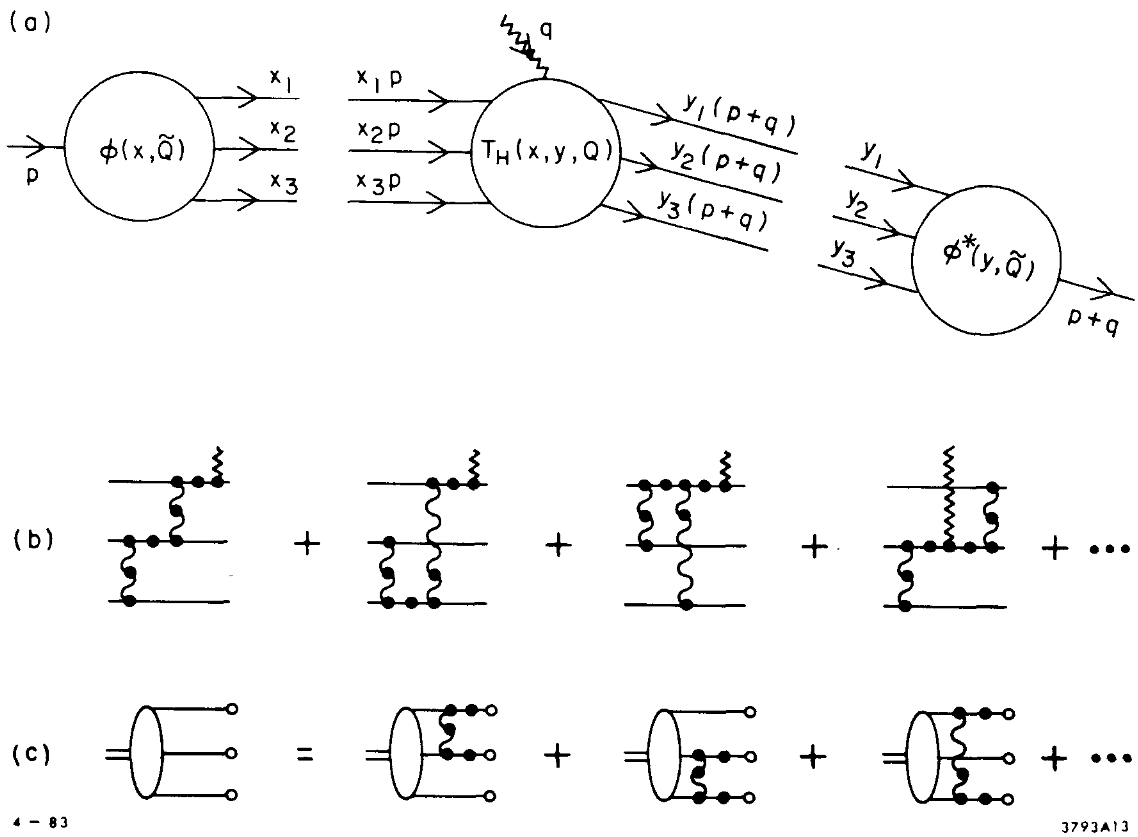


Fig. 1. (a) Factorization of the nucleon form factor at large Q^2 in QCD. The optimal scale $\check{Q}$ for the distribution amplitude $\phi(x, \check{Q})$ is discussed in Ref. 9. (b) The leading order diagrams for the hard scattering amplitude T_H. The dots indicate insertions which enter the renormalization of the coupling constant. (c) The leading order diagrams which determine the Q^2 dependence of $\phi_B(x, Q)$.

where n is the minimum number of fields in the hadron. Since quark helicity is conserved in T_H and $\phi(x_i, Q)$ is the $L_z = 0$ projection of the wavefunction, total hadronic helicity is conserved at large momentum transfer for any QCD exclusive reaction. The dominant nucleon form factor thus corresponds to $F_1(Q^2)$ or $G_M(Q^2)$; the Pauli form factor $F_2(Q^2)$ is suppressed by an extra power of Q^2. In the case of the deuteron, the dominant form factor has helicity $\lambda = \lambda' = 0$, corresponding to $\sqrt{A(Q^2)}$. The general form of the logarithmic dependence of $F(Q^2)$ can be derived from the operator product expansion[♯24] at short distance or by solving an evolution equation[♯12] for the distribution amplitude computed from gluon exchange [Fig. 1(c)], as we discuss in more detail in the next section. The result for the large Q^2 behavior of the baryon form factor in QCD is[♯12,24]

$$F_B(Q^2) = \frac{\alpha_s^2(Q^2)}{Q^4} \sum_{n,m} d_{nm} \left(\ell n \frac{Q^2}{\Lambda^2}\right)^{-\gamma_m - \gamma_n} \tag{11}$$

where the γ_n are computable anomalous dimensions of the baryon three-quark wave function at short distance and the d_{mn} are determined from the value of the distribution amplitude $\phi_B(x, Q_0^2)$ at a given point Q_0^2 and the normalization of T_H. Asymptotically, the dominant term has the minimum anomalous dimension. The dominant part of the form factor comes from the region of the x integration where each quark has a finite fraction of the light cone momentum; the end point region where the struck quark has $x \simeq 1$ and spectator quarks have $x \sim 0$ is asymptotically suppressed by quark (Sudakov) form factor gluon radiative corrections.

As shown in Fig. 2 the power laws predicted by perturbative QCD are consistent with experiment.[♯25] The behavior $Q^4 G_M(Q^2) \sim$ *const* at large Q^2 [see Fig. 3] provides a direct check that the minimal Fock state in the nucleon contains three quarks and that the quark propagator and the $qq \rightarrow qq$ scattering amplitudes are approximately scale-free. More generally, the nominal power law predicted for large momentum transfer exclusive reactions is given by the dimensional counting rule[♯16] $\mathcal{M} \sim Q^{4-n_{TOT}} F(\theta_{cm})$ where n_{TOT} is the total number of elementary fields which scatter in the reaction. The predictions are apparently compatible with experiment. In addition, for some scattering reactions there are contributions from multiple scattering diagrams (Landshoff contributions) which together with Sudakov effects can lead to small power-law corrections, as well as a complicated spin, and amplitude phase phenomenology. Recent measurements of $\gamma\gamma \rightarrow \pi^+\pi^-$, K^+K^- at large invariant pair mass are also consistent with the QCD predictions.[♯26,♯27] In principle it should be possible to use measurements of the scaling and angular dependence of the $\gamma\gamma \rightarrow M\bar{M}$ reactions to measure the shape of the distribution amplitude $\phi_M(x, Q)$.[♯26]

♯25 M. D. Mestayer, SLAC-Report 214 (1978) F. Martin, et al., Phys. Rev. Lett. 38, 1320 (1977); W. P. Schultz, et al., Phys. Rev. Lett. 38, 259 (1977); R. G. Arnold, et al., Phys. Rev. Lett. 40, 1429 (1978); SLAC-PUB-2373 (1979); B. T. Chertok, Phys. Lett. 41, 1155 (1978); D. Day, et al., Phys. Rev. Lett. 43, 1143 (1979). Summaries of the data for nucleon and nuclear form factors at large Q^2 are given in B. T. Chertok, in Progress in Particle and Nuclear Physics, Proceeding of the International School of Nuclear Physics, 5th Course, Erice (1978), and Proceedings of the XVI Rencontre de Moriond, Les Arcs, Savoie, France, 1981.

♯26 S. J. Brodsky and G. P. Lepage, Phys. Rev. D24, 1808 (1981).

♯27 The calculation of $\gamma\gamma \rightarrow B\bar{B}$ is given by G. R. Farrar, E. Maina, and F. Neri, Ref. 13.

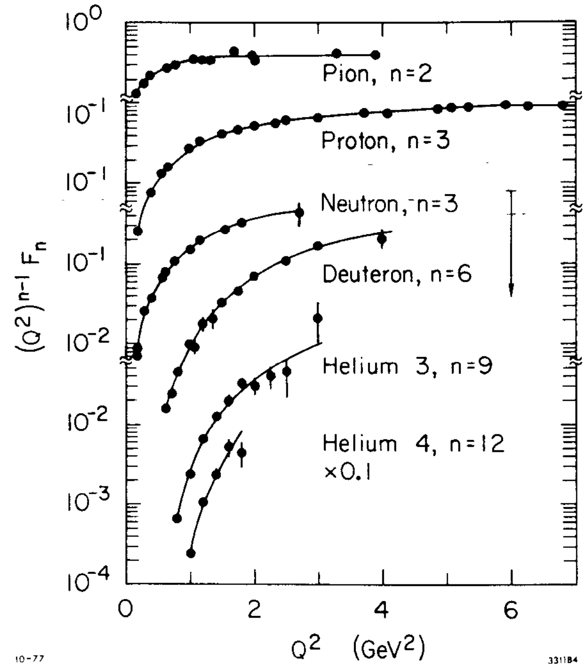
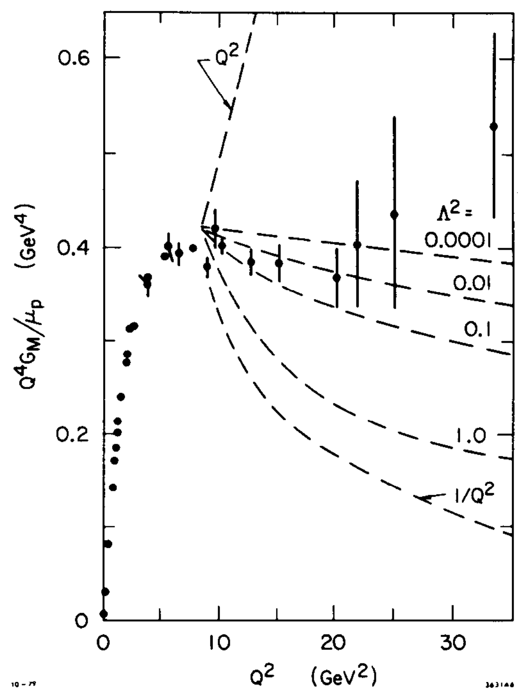


Fig. 2. Comparison of experiment with the QCD dimensional counting rule $(Q^2)^{n-1}F(Q^2) \sim const$ for form factors. The proton data extends beyond 30 GeV^2.

Fig. 3. Prediction for $Q^4G_M^P(Q^2)$ for various QCD scale parameters Λ^2 (in GeV^2). The data are from Ref. 25. The initial wave function is taken as $\phi(x,\lambda) \propto \delta(x_1 - \frac{1}{3})\delta(x_2 - \frac{1}{3})$ at $\lambda^2 = 2$ GeV^2. The factor $(1 + \frac{m_p^2}{Q^2})^{-2}$ is included in the prediction as a representative of mass effects, and the overall normalization is unknown.



An actual calculation of $\phi(x, Q)$ from QCD requires non-perturbative methods such as lattice gauge theory, or more directly, the solution of the light-cone equation of motion[11,12]

$$\Big[M^2 - \sum_{i=1}^{n}\Big(\frac{k_\perp^2 + m^2}{x}\Big)_i\Big]\Psi = V_{\rm QCD}\Psi \tag{12}$$

The explicit form for the matrix representation of V_{QCD} and a discussion of the infrared and ultraviolet regulation required to interpret this result is given in Ref. 12. Thus far experiment has not been sufficiently precise to measure the logarithmic modification of dimensional counting rules predicted by QCD. Checks of the normalization of $(Q^2)^{n-1}F(Q^2)$ require independent determinations of the valence wavefunction. The relatively large normalization of $Q^4G_M^p(Q^2)$ at large Q^2 can be understood if the valence three-quark state has small transverse size, i.e., is large at the origin. The physical radius of the proton measured from $F_1(Q^2)$ at low momentum transfer then reflects the contributions of the higher Fock states $qqqg$, $qqqq\bar{q}$ (or meson cloud), etc. A small size for the proton valence wavefunction (e.g., $R^p_{qqq} \sim 0.2$ to $0.3\ fm$) can also explain[♯11,♯12,♯28] the large magnitude of $\langle k_\perp^2\rangle$ of the intrinsic quark momentum distribution needed to understand hard-scattering inclusive reactions. The necessity for small valence state Fock components can be demonstrated explicitly for the pion wavefunction, since $\psi_{q\bar{q}/\pi}$ is constrained by sum rules derived from $\pi^+ \to \ell^+\nu$, and $\pi^0 \to \gamma\gamma$. One finds[♯12] a valence state radius $R^\pi_{q\bar{q}} \sim 0.4\,fm$, corresponding to a probability $P^\pi_{q\bar{q}} \sim 1/4$.

THE DEUTERON IN QCD

Of the five color-singlet representations of six quarks, only one corresponds to the usual system of two color singlet baryonic clusters.[♯15,♯29,♯30] Notice that the exchange of a virtual gluon in the deuteron at short distance inevitably produces Fock state components where the three-quark clusters correspond to color octet nucleons or isobars. Thus, in general, the deuteron wavefunction will have a complete spectrum of hidden-color wavefunction components, although it is likely that these states are important only at small internucleon separation.

Despite the complexity of the multi-color representations of nuclear wavefunctions, the analysis[♯15] of the deuteron form factor at large momentum transfer can be carried out in parallel with the nucleon case. Only the minimal six-quark Fock state needs to be considered to leading order in $1/Q^2$. The deuteron form factor can then be written as a convolution [see Fig. 4],

$$F_d(Q^2) = \int_0^1 [dx]\ [dy]\ \phi_d^\dagger(y,Q)\ T_H^{6q+\gamma^*\to 6q}(x,y,Q)\ \ \phi_d(x,Q)\ , \tag{13}$$

where the hard scattering amplitude scales as

$$T_H^{6q+\gamma^*\to 6q} = \left[\frac{\alpha_s(Q^2)}{Q^2}\right]^5 t(x,y)\quad [1 + O(\alpha_s(Q^2))] \tag{14}$$

The anomalous dimensions γ_n^d are calculated from the evolution equations for $\phi_d(x_i,Q)$ derived to leading order in QED from pairwise gluon-exchange interactions: $(C_F = 4/3,\ C_d = -C_F/5)$

♯28 S. J. Brodsky, T. Huang, G. P. Lepage, in Particles and Fields 2, Edited by A. Z. Capri, A. N. Kamal, Plenan (1983), and T. Huang, SLAC-PUB-2580 (1980), published in the Proceedings of the XXth International Conference on High Energy Physics, Madison, Wisconsin (1980).

♯29 See, e.g., V. Matveev and P. Sorba, Nuovo Cimento Lett. 20, 435 (1977).

♯30 S. J. Brodsky, C.-R. Ji, and G. P. Lepage (to be published).

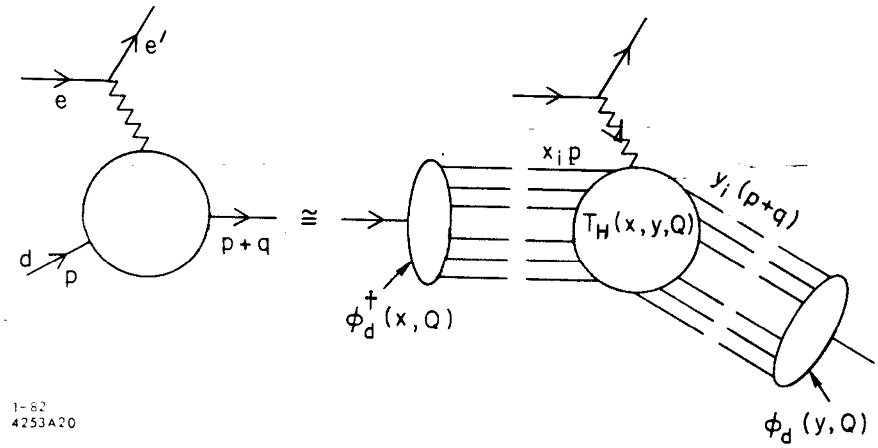


Fig. 4. Factorization of the deuteron form factor at large Q^2.

$$\prod_{k=1}^{6} x_k \left[\frac{\partial}{\partial \xi} + \frac{3C_F}{\beta}\right] \tilde{\Phi}(x_i, Q) = -\frac{C_d}{\beta} \int_0^1 [dy]\, V(x_i, y_i) \tilde{\Phi}(y_i, Q) \ . \tag{15}$$

Here we have defined

$$\Phi(x_i, Q) = \prod_{k=1}^{6} x_k \tilde{\Phi}(x_i, Q), \tag{16}$$

and the evolution is in the variable

$$\xi(Q^2) = \frac{\beta}{4\pi} \int_{Q_0^2}^{Q^2} \frac{dk^2}{k^2} \alpha_s(k^2) \sim ln \left(\frac{ln \frac{Q^2}{\Lambda^2}}{ln \frac{Q_0^2}{\Lambda^2}} \right) \ . \tag{17}$$

The kernel V is computed to leading order in $\alpha_s(Q^2)$ from the sum of gluon interactions between quark pairs. The general matrix representations of γ_n with bases $\left| \prod_{i=1}^{5} x_i^{m_i} \right>$ will be given in Ref. 30. The effective leading anomalous dimension γ_0, corresponding to the eigenfunction $\tilde{\Phi}(x_i) = 1$, is $\gamma_0 = (6/5)(C_F/\beta)$ (see the next section).

In order to make more detailed and experimentally accessible predictions, we will define the "reduced" nuclear form factor in order to remove the effects of nucleon compositeness:[#8]

$$f_d(Q^2) \equiv \frac{F_d(Q^2)}{F_N^2(Q^2/4)} \ . \tag{18}$$

The arguments for each of the nucleon form factors (F_N) is $Q^2/4$ since in the limit of zero binding energy each nucleon must change its momentum from $\sim p/2$ to $(p+q)/2$. Since the leading anomalous dimensions of the nucleon distribution amplitude is $C_F/2\beta$, the QCD prediction for the asymptotic Q^2-behavior of $f_d(Q^2)$ is

$$f_d(Q^2) \sim \frac{\alpha_s(Q^2)}{Q^2}\left(ln\frac{Q^2}{\Lambda^2}\right)^{-\frac{2}{5}\frac{C_F}{\beta}}, \tag{19}$$

where $-(2/5)(C_F/\beta) = -8/145$ for $n_f = 2$.

Although this QCD prediction is for asymptotic momentum transfer, it is interesting to compare it directly with the available high Q^2 data[#25] [see Fig. 5]. In general one would expect corrections from higher twist effects (e.g., mass and $k_\perp$ smearing), higher particle number Fock states, higher order contributions in $\alpha_s(Q^2)$, as well as non-leading anomalous dimensions. However, the agreement of the data with simple $Q^2 f_d(Q^2) \sim const$ behavior for $Q^2 > 1/2$ GeV^2 implies that, unless there is a fortuitous cancellation, all of the scale-breaking effects are small, and the present QCD perturbative calculations are viable and applicable even in the nuclear physics domain. The lack of deviation from the QCD parameterization also suggests that the parameter Λ is small. A comparison with a standard definition such as $\Lambda_{\overline{MS}}$ would require a calculation of next to leading effects. A more definitive check of QCD can be made by calculating the normalization of $f_d(Q^2)$ from T_H and the evolution of the deuteron wave function to short distances. It is also important to confirm experimentally that the helicity $\lambda = \lambda' = 0$ form factor is indeed dominant.

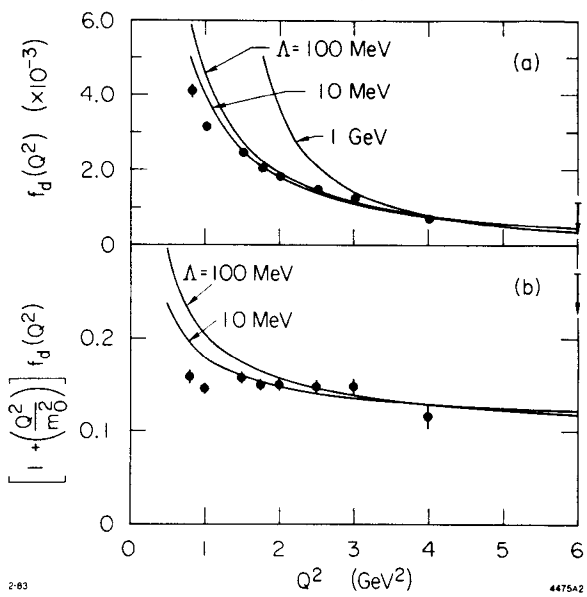


Fig. 5. (a) Comparison of the asymptotic QCD prediction (18) and (19) with experiment using $F_N(Q^2) = [1 + (Q^2/0.71\,\mathrm{GeV}^2)]^{-2}$. The normalization is fit at $Q^2 = 4$ GeV^2. (b) Comparison of the prediction $[1 + (Q^2/m_0^2)]f_d(Q^2) \propto (\ell n\, Q^2)^{-1-(2/5)(C_F/\beta)}$ with data. The value $m_0^2 = 0.28$ GeV^2 is used.

The calculation of the normalization $T_H^{6q+\gamma^*\to 6q}$ to leading order in $\alpha_s(Q^2)$ will require the evaluation of over 300,000 Feynman diagrams involving five exchanged gluons. Fortunately this appears possible using the algebraic computer methods introduced by Farrar and Neri.[#31] The method of setting the appropriate scale $\hat{Q}$ of $\alpha_s^5(\hat{Q}^2)$ in T_H is given in Ref. 10.

We note that the deuteron wave function[#32] which contributes to the asymptotic limit of the form factor is the totally antisymmetric wave function corresponding to the orbital Young

#31 G. R. Farrar and F. Neri, Phys. Lett. 130B, 109 (1983).

#32 For a recent attempt at a phenomenological determination of the deuteron six-quark component, see V. G. Ableev, A. P. Kobushkin, L. N. Strunov, et al., Dubna preprint E1-83-487 (1983), and references therein.

symmetry given by [6] and isospin (T)+ spin (S) Young symmetry given by {33}. The deuteron state with this symmetry is related to the NN, $\Delta\Delta$, and hidden color (CC) physical bases, for both the $(TS) = (01)$ and (10) cases, by the formula[#33]

$$\psi_{[6]\{33\}} = \sqrt{\frac{1}{9}}\,\psi_{NN} + \sqrt{\frac{4}{45}}\,\psi_{\Delta\Delta} + \sqrt{\frac{4}{5}}\,\psi_{CC} \quad (20)$$

Thus the physical deuteron state, which is mostly ψ_{NN} at large distance, must evolve to the $\psi_{[6]\{33\}}$ state when the six quark transverse separations $b^i_\perp \le O(1/Q) \to 0$. Since this state is 80% hidden color, the deuteron wave function cannot be described by the meson-nucleon isobar degrees of freedom in this domain. The fact that the six-quark color singlet state inevitably evolves in QCD to a dominantly hidden-color configuration at small transverse separation also has implications for the form of the nucleon-nucleon ($S_z = 0$) potential, which can be considered as one interaction component in a coupled scattering channel system. As the two nucleons approach each other, the system must do work in order to change the six-quark state to a dominantly hidden color configuration; i.e., QCD requires that the nucleon-nucleon potential must be repulsive at short distances [see Fig. 6].[#11,#34] The evolution equation for the six-quark system suggests that the distance where this change occurs is in the domain where $\alpha_s(Q^2)$ most strongly varies. The general solutions of the evolution equation for multiquark systems is discussed below. Some of the solutions are orthogonal to the usual nuclear configurations which correspond to separated nucleons or isobars at large distances. Such solutions could be connected with the anomalous phenomena observed in heavy ion collisions.

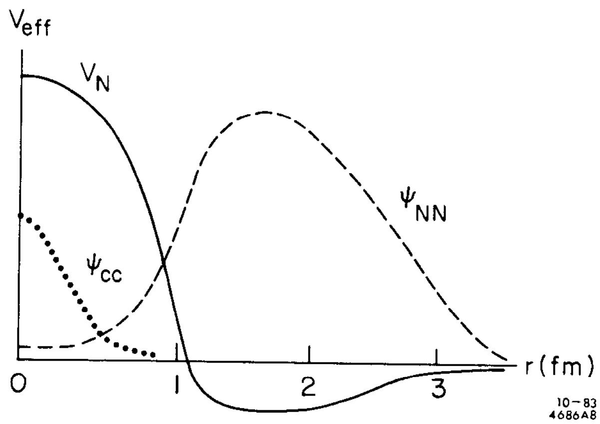


Fig. 6. Schematic representation of the deuteron wave function in QCD indicating the presence of hidden color six-quark components at short distances.

REDUCED NUCLEAR AMPLITUDES

One of the basic problems in the analysis of nuclear scattering amplitudes is how to consistently account for the effects of the underlying quark/gluon component structure of nucleons. Traditional methods based on the use of an effective nucleon/meson local Lagrangian field theory

#33 M. Harvey, Ref. 2.

#34 Similar considerations for nonrelativistic system are given in A Faessler et al., Nucl. Phys. A402, 555 (1983); S. Furui and A. Faessler Nucl. Phys. A397, 413 (1983).

are not really applicable, giving the wrong dynamical dependence in virtually every kinematic variable for composite hadrons. The inclusion of *ad hoc* vertex form factors is unsatisfactory since one must understand the off-shell dependence in each leg while retaining gauge invariance; such methods have little predictive power. On the other hand, the explicit evaluation of the multiquark hard-scattering amplitudes needed to predict the normalization and angular dependence for a nuclear process, even at leading order in α_s requires the consideration of millions of Feynman diagrams. Beyond leading order one must include contributions of non-valence Fock states wavefunctions, and a rapidly expanding number of radiative corrections and loop diagrams.

The reduced amplitude method, although not an exact replacement for a full QCD calculation, provides a simple method for identifying the dynamical effects of nuclear substructure, consistent with covariance, QCD scaling laws and gauge invariance. The basic idea has already been introduced for the reduced deuteron form factor. More generally if we neglect nuclear binding, then the light-cone nuclear wavefunction can be written as a cluster decomposition of collinear nucleons: $\psi_{q/A} = \psi_{N/A} \prod_N \Psi_{q/N}$ where each nucleon has $1/A$ of the nuclear momentum. A large momentum transfer nucleon amplitude then contains as a factor the probability amplitude for each nucleon to remain intact after absorbing $1/A$ of the respective nuclear momentum transfer. We can identify each probability amplitude with the respective nucleon form factor $F\left(\hat{t}_i = \frac{1}{A^2}\, t_A\right)$. Thus for any exclusive nuclear scattering process, we define the reduced nuclear amplitude

$$m = \frac{\mathcal{M}}{\prod_{i=1}^{A} F_N(\hat{t}_i)} \tag{21}$$

The QCD scaling law for the reduced nuclear amplitude m is then identical to that of nuclei with point-like nuclear components: e.g., the reduced nuclear form factors obey

$$f_A(Q^2) \equiv \frac{F_A(Q^2)}{\left[F_N(Q^2/A^2)\right]^A} \sim \left[\frac{1}{Q^2}\right]^{A-1} . \tag{22}$$

Comparisons with experiment and predictions for leading logarithmic corrections to this result are given in Ref. 8. In the case of photo- (or electro-) disintegration of the deuteron one has

$$m_{\gamma d \to np} = \frac{\mathcal{M}_{\gamma d \to np}}{F_n(t_n) F_p(t_p)} \sim \frac{1}{p_T}\, f(\theta_{cm}) \tag{23}$$

i.e., the same elementary scaling behavior as for $\mathcal{M}_{\gamma M \to q\bar{q}}$. Comparison with experiment is encouraging [see Fig. 7], showing that as was the case for $Q^2 f_d(Q^2)$, the perturbative QCD scaling regime begins at $Q^2 \gtrsim 1\ GeV^2$. Detailed comparisons and a model for the angular dependence and the virtual photon-mass dependence of deuteron electrodisintegration are discussed in Ref. 8. Other potentially useful checks of QCD scaling of reduced amplitudes are

$$\begin{aligned} m_{pp \to d\pi^+} &\sim p_T^{-2}\, f(t/s) \\ m_{pd \to H^3\pi^+} &\sim p_T^{-4}\, f(t/s) \\ m_{\pi d \to \pi d} &\sim p_T^{-4}\, f(t/s)\ . \end{aligned} \tag{24}$$

It is also possible to use these QCD scaling laws for the reduced amplitude as a parametrization for the background for detecting possible new dibaryon resonance states.

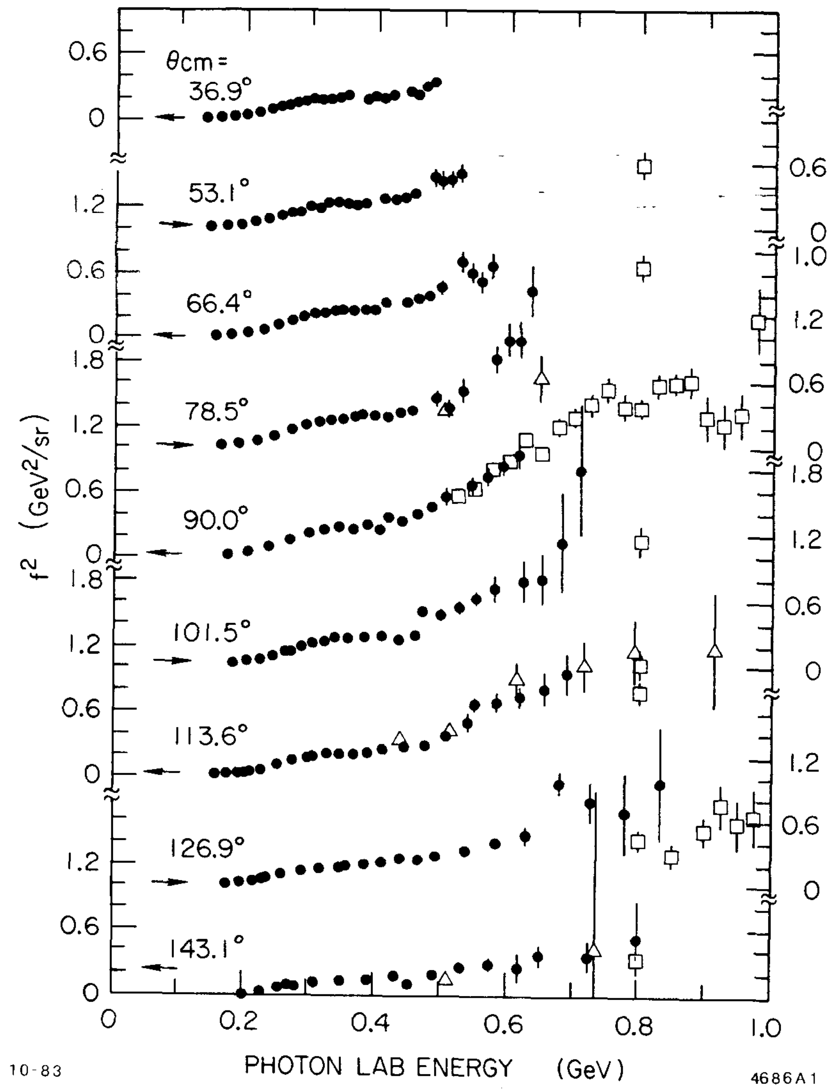


Fig. 7. Comparison of deuteron photo-disintegration data with the scaling prediction which requires $f^2(\theta_{c.m.})$ to be independent of energy at large momentum transfer. The data are from H. Myers et al., Phys. Rev. 121, 630 (1961); R. Ching and C. Schaerf, Phys. Rev. 141, 1320 (1966); P. Dougan et al., Z. Phys. A 276, 55 (1976).

QUANTUM CHROMODYNAMIC EVOLUTION OF MULTIQUARK SYSTEMS

As we have discussed, exclusive processes involving the transfer of large momenta can be systematically analyzed in QCD. A large number of experimentally accessible phenomena including the elastic and inelastic electromagnetic and weak form factors and large-angle elastic scattering processes of hadrons and nuclei can be represented in terms of a simple picture for exclusive processes based on light-cone perturbation theory. For example, as we discussed above, the baryon form factor at large Q^2 is represented by the factorized form (see Fig. 1)[♯12,♯13,♯14,♯24,♯35]

♯35 The leading form factor corresponds to helicity conservation $h = h' = \frac{1}{2}$ or $h = h' = 0$. For nonleading form factors, see C. E. Carlson and F. Gross, preprint of College of William and Mary (1984).

$$F_B(Q^2) = \int_0^1 [dx] \int_0^1 [dy]\, \phi^*(y_i, \check{Q}_y)\, T_H(x_i, y_i, Q)\, \phi(x_i, \check{Q}_x) \left[1 + O\left(\frac{m_i^2}{Q^2}\right)\right]$$

$$= \frac{32\pi^2}{9} \frac{\alpha_s^2(Q^2)}{Q^4} \sum_{n,m} b_{nm} \left(\ell n \frac{Q^2}{\Lambda^2}\right)^{-\gamma_n - \gamma_m} \left[1 + O\left(\alpha_s(Q^2), \frac{m_i^2}{Q^2}\right)\right]$$

$$\rightarrow C \left(\frac{\alpha_s(Q^2)}{Q^2}\right)^2 \left(\ell n \frac{Q^2}{\Lambda^2}\right)^{-2\gamma_0} \quad (\text{as } Q^2 \rightarrow \text{large}) \quad , \tag{25}$$

where x_i is the light-cone longitudinal momentum fraction of i^{th} quark $x_i = (k_i^0 + k_i^3)/(p^0 + p^3)$, $[dx] \equiv dx_1\, dx_2\, dx_3\, \delta\, (1 - \sum_i x_i)$ and $\check{Q}_x \equiv \min_i(x_i Q)$.

The dominant Q^2 dependence $(\alpha_s(Q^2)/Q^2)^2$ is derived from the hard scattering amplitude $T_H(x_i, y_i, Q)$ $(\gamma^* + 3q \rightarrow 3q)$ with only weak (logarithmic) Q^2 dependence coming from quark distribution amplitude $\phi(x_i, Q)$ (γ_0 is the leading anomalous dimension). The essential feature of this result is that a very complicated process can be simply represented by the factorization into product of three amplitudes, and, thus, the main calculation for this process turns out to be the calculation of T_H. The distribution amplitude $\phi(x_i, Q)$ is in principle determined by nonperturbative bound state physics, and it is independent of the process.

The quark distribution amplitude $\phi(x_i, Q)$ is the amplitude for converting the baryon into three valence quarks at impact separation $b_\perp \sim O(1/Q)$. It is related to the equal $\tau = t + z$ hadronic wave function $\psi(x_i, \vec{k}_{\perp i})$:

$$\phi(x_i, Q) \propto \int^Q \prod_{i=1}^{3} d^2\vec{k}_{\perp i}\, \delta^2 \left(\sum_i \vec{k}_{\perp i}\right) \psi(x_i, \vec{k}_{\perp i}) \quad , \tag{26}$$

and contains the essential physics of that part of the hadronic wave function which affects exclusive processes with large momentum transfer. In this section, we present a new technique[♯36] for constructing $\phi(x_i, Q)$ in order to predict the short distance behavior of multiquark systems.

The distribution amplitude for a baryon is determined by an evolution equation which can be derived from the Bethe-Salpeter equation at large transverse momentum projected on the light-cone:[♯12]

$$\left(Q^2 \frac{\partial}{\partial Q^2} + \frac{3C_F}{2\beta}\right) \phi(x_i, Q) = \frac{C_B}{\beta} \int [dy]\, V(x_i, y_i)\, \phi(y_i, Q) \quad , \tag{27}$$

where $C_F = (n_c^2 - 1)/2n_c = 4/3$, $C_B = (n_c + 1)/2n_c = 2/3$, $\beta = 11 - (2/3)n_f$, and $V(x_i, y_i)$ is computed to leading order in α_s from the single-gluon-exchange kernel. The evolution equation automatically sums to leading order in $\alpha_s(Q^2)$ all of the contributions from multiple gluon exchange which determine the tail of the valence wavefunction and thus the Q^2-dependence

♯36 Details will be given in a separate paper.

of the distribution amplitude. The general solution of this equation is

$$\phi(x_i, Q) = x_1 x_2 x_3 \sum_{n=0}^{\infty} a_n \left(\ell n \frac{Q^2}{\Lambda^2} \right)^{-\gamma_n} \tilde{\phi}_n(x_i) \quad , \tag{28}$$

where the anomalous dimensions γ_n and the eigenfunctions $\tilde{\phi}_n(x_i)$ satisfy the characteristic equation:

$$x_1 x_2 x_3 \left(-\gamma_n + \frac{3C_F}{2\beta} \right) \tilde{\phi}_n(x_i) = \frac{C_B}{\beta} \int_0^1 [dy] \, V(x_i, y_i) \, \tilde{\phi}_n(y_i) \quad . \tag{29}$$

In the large Q^2 limit, only the leading anomalous dimension γ_0 contributes.

In the three quark case, the color singlet property of the baryon system guarantees all three quarks have different quantum numbers. Thus, we do not necessarily have to antisymmetrize the system according to Pauli's principle, and $\tilde{\phi}_n(x_i)$ may be derived by expanding $V(x_i, y_i)$ on a polynomial basis $\{x_1^m x_3^n\}_{m,n=0}^{\infty}$. However, if we consider multibaryon (nuclear) systems of $3n$ quarks, then the color singlet requirement does not guarantee that all the quarks of the system have different quantum numbers and antisymmetric representations in the total quantum space are needed. This is the essential point of the new technique.

The antisymmetrization technique will be presented in detail below. If we apply this method to three quark system, then we have consistency with preceding results but additionally we obtain a distinctive classification of nucleon (N) and delta (Δ) wave functions and the corresponding Q^2 dependence which discriminates N and Δ form factors. We derive QCD predictions for the reduced form factor of the deuteron and compare them with the available experimental results. Furthermore, we can combine these results with the fractional parantage technique of Harvey,[♯33] and we can derive constraints on the effective force between two baryons at short distances. This will be explained below.

The Antisymmetrization Technique for Solving Multiquark Evolution Equations

In order to solve the evolution equations, we use the following procedure:

1. Construct the representations in each of the quantum spaces color (C), isospin (T), spin (S), and orbital (O) using Young diagrammatic techniques.[♯37] Each quantum state is constructed by filling up the Young tableaus with corresponding specific quantum numbers. "Orbital" states are classified by polynomials $\prod_i x_i^{n_i}$, with the minimal powers dominant in the high Q^2 region. We use the symmetry of the x_i dependence of $\tilde{\phi}_n(x_i)$ in analogy with the permutation symmetry of orbital dependence of nonrelativistic wave functions. After an orthonormalization procedure, the orbital functions satisfy the condition:

$$\int [dx] \, \omega(x) \, \tilde{\phi}_m^*(x_i) \, \tilde{\phi}_n(x_i) = \delta_{mn} \quad , \tag{30}$$

where $\omega(x) = \prod_i x_i$.

♯37 M. Hamermesh, *Group Theory* (Addison-Wesley, Reading, Massachusetts, 1962).

2. Construct the inner-product of Young diagrams in order to produce completely antisymmetric representations in the CTSO total space. The Clebsch-Gordan coefficients of the permutation group are used. A convenient algebraic method will be given in a separate paper.

3. Calculate the QCD kernel matrix in the basis of completely antisymmetric representations. For example, the one gluon exchange kernel for the three quark system is given by $(i, j, k = 1, 2, 3)$[♯12]

$$
\begin{aligned}
V(x_i, y_i) &= \left(\sum_{a=1}^{8} \frac{\lambda_a}{2} \cdot \frac{\lambda_a}{2}\right) \int_0^1 [dy] \sum_{i \neq j} \theta(y_i - x_i)\, \delta(x_k - y_k) \\
&\quad \times \frac{y_j}{x_j} \left(\frac{\delta_{h_i h_j}}{x_i + x_j} + \frac{\Delta}{y_i - x_i} \right) \quad , \quad k \neq i, j
\end{aligned}
\tag{31}
$$

where the λ_a are $SU(3)_c$ Gell-mann matrices, $\Delta\phi(y_i) = \phi(y_i) - \phi(x_i)$, and $\delta_{h_i h_j} = 0(1)$ when the helicities of constituents are antiparallel (parallel). From this kernel, we find the following QCD evolution properties:

(a) Color singlet states are preserved by the action of V.

(b) Isospin cannot be changed, i.e. N and Δ cannot mix with each other.

(c) Spin states can mix by the spin annihilation term $(\delta_{h_i h_j})$.

(d) Orbital states can also mix, with total $n = \sum_i n_i$ preserved.

As an example, let's consider the leading $n = 1$ amplitude of the s^2p excited nucleon state of $p^*_{(\frac{1}{2},\frac{1}{2})}$ and $p^*_{(\frac{3}{2},\frac{1}{2})}$ (the x dependence of ϕ is given by the orbital Young diagram).

$$
\begin{aligned}
\phi_{\frac{1}{2}} &= \begin{array}{|c|}\hline r \\ \hline y \\ \hline b \\ \hline \end{array} \times \begin{array}{|c|c|}\hline u & u \\ \hline d \\ \cline{1-1} \end{array} \times \begin{array}{|c|c|}\hline \uparrow & \uparrow \\ \hline \downarrow \\ \cline{1-1} \end{array} \times \begin{array}{|c|c|}\hline 0 & 0 \\ \hline 1 \\ \cline{1-1} \end{array} \; , \\
\phi_{\frac{3}{2}} &= \begin{array}{|c|}\hline r \\ \hline y \\ \hline b \\ \hline \end{array} \times \begin{array}{|c|c|}\hline u & u \\ \hline d \\ \cline{1-1} \end{array} \times \begin{array}{|c|c|c|}\hline \uparrow & \uparrow & \downarrow \\ \hline \end{array} \times \begin{array}{|c|c|}\hline 0 & 0 \\ \hline 1 \\ \cline{1-1} \end{array} \; .
\end{aligned}
\tag{32}
$$

If we split the kernel V in terms of a spin annihilation term V_δ and the remainder V_Δ, we find

$$
\begin{aligned}
\int [dy]\, V_\delta(x, y) \begin{bmatrix} \phi_{\frac{1}{2}}(y) \\ \phi_{\frac{3}{2}}(y) \end{bmatrix} &= \begin{bmatrix} \frac{2}{3} & \frac{1}{6} \\ \frac{1}{6} & \frac{2}{3} \end{bmatrix} \begin{bmatrix} \phi_{\frac{1}{2}}(x) \\ \phi_{\frac{3}{2}}(x) \end{bmatrix} \\
\int [dy]\, V_\Delta(x, y) \begin{bmatrix} \phi_{\frac{1}{2}}(y) \\ \phi_{\frac{3}{2}}(y) \end{bmatrix} &= \begin{bmatrix} -\frac{3}{2} & 0 \\ 0 & -\frac{3}{2} \end{bmatrix} \begin{bmatrix} \phi_{\frac{1}{2}}(x) \\ \phi_{\frac{3}{2}}(x) \end{bmatrix} \quad .
\end{aligned}
\tag{33}
$$

4. Diagonalize the kernel matrix $(V = V_\delta + V_\Delta)$ to determine the eigenvalues and eigenso-

lutions. From the above example, we find the following results:

$$b = \begin{cases} \frac{2}{3} \\ 1 \end{cases} \quad \text{for} \quad \check{\phi} = \begin{cases} \frac{1}{\sqrt{2}}\phi_{\frac{1}{3}} + \frac{1}{\sqrt{2}}\phi_{\frac{2}{3}} \\ -\frac{1}{\sqrt{2}}\phi_{\frac{1}{3}} + \frac{1}{\sqrt{2}}\phi_{\frac{2}{3}} \end{cases} , \tag{34}$$

where $\gamma = (2bC_B + (3/2)C_F)/\beta$.

Following the above procedures (1 through 4), we can find the anomalous dimensions and construct the corresponding eigenfunctions for arbitrary multiquark systems. In the three quark case, we find that the results coincide with preceding calculations, but here we can unambiguously resolve the N and Δ state wave functions and discriminate their form factors.

We emphasize that the methods used here systematically determine all of the relativistic ($L_z = 0$) bound states which can be constructed from the multiquark degrees of freedom. Although the evolution equation is not an eigenvalue equation for the mass spectrum its eigenstates form a relativistic wavefunction basis which is diagonalized by the one-gluon exchange kernel. The diagonalization of the kernel on the polynomial basis gives a construction of eigensolutions of the generalization of angular momenta to states on the light-cone.

The Effective Force Between Baryons: SU(2) Color Examples

In the preceeding section, we have discussed how we can solve QCD evolution equations in order to predict the short distance behavior of multiquark systems using Young diagrammatic methods. Since the eigensolutions obtained in this way have definite permutation symmetry, we can apply the fractional parantage technique for the multibaryon system in order to relate the eigensolutions to cluster representations which have physical baryon, or alternatively, "hidden-color" degrees of freedom.

For example, if we apply this technique[♯38] to the simple case of the four quark system under $SU(2)_c$ then we find the transition matrix given by Table 1 (T = S = 0 case) which relates the symmetry basis represented by four-quark eigensolutions and the physical basis represented by "toy"-dibaryon and hidden-color degrees of freedom. From this table we can expand the distribution amplitudes of the physical basis in terms of eigensolutions:

$$\begin{aligned} \phi_{NN}(x_i,Q) &= 0.07\phi_1(x_i)\left(\ell n\frac{Q^2}{\Lambda^2}\right)^{0.13C_F/\beta} - 0.64\phi_2(x_i)\left(\ell n\frac{Q^2}{\Lambda^2}\right)^{-0.02C_F/\beta} + \dots \\ \phi_{\Delta\Delta}(x_i,Q) &= -0.07\phi_1(x_i)\left(\ell n\frac{Q^2}{\Lambda^2}\right)^{0.13C_F/\beta} - 0.59\phi_2(x_i)\left(\ell n\frac{Q^2}{\Lambda^2}\right)^{-0.02C_F/\beta} + \dots \\ \phi_{CC}(x_i,Q) &= -0.70\phi_1(x_i)\left(\ell n\frac{Q^2}{\Lambda^2}\right)^{0.13C_F/\beta} - 0.35\phi_2(x_i)\left(\ell n\frac{Q^2}{\Lambda^2}\right)^{-0.02C_F/\beta} + \dots \end{aligned} \tag{35}$$

where $C_F = 3/4$ in this case.

♯38 The generalization to $SU(3)$ color will be presented in a later paper.

TABLE 1. The relationship between four-quark antisymmetric SU(2) color representations and effective two-cluster representations (T = S = 0 case). Isospin singlet and triplet states both with color singlet are denoted N and Δ, while color triplet state is represented by C. The square and curly brackets represent orbital (O) and spin-isospin (TS) symmetries separately.

	[4] {22}	[22] {22}	[22] {4}
NN	$-\frac{1}{2}$	$-\frac{1}{2}$	$\frac{1}{\sqrt{2}}$
$\Delta\Delta$	$\frac{1}{2}$	$\frac{1}{2}$	$\frac{1}{\sqrt{2}}$
CC	$\frac{1}{\sqrt{2}}$	$-\frac{1}{\sqrt{2}}$	0

Thus, we find that the NN, $\Delta\Delta$ and CC states have completely different Q^2 evolution. As Q^2 goes to infinity, the NN and $\Delta\Delta$ components are negligible but the CC components are large. In other words, the dominant degrees of freedom at the origin of the dibaryon system at zero impact separation are hidden-color states rather than physical baryon states. This indicates that the physical dibaryons have a repulsive core at the origin[♯39] while the colorful hidden-color clusters behave as in an attractive well [see Fig. 6]. In this way, we derive constraints on the effective force between two baryons. Analogous results hold for the six-quark states in $SU(3)_c$.

In summary, using the above techniques based on completely antisymmetrized representations, we can analyze multi-quark distribution amplitudes $\phi(x_i, Q)$ in QCD in order to predict the short distance behavior of multiquark systems. Since the new technique is based on permutation symmetry, we can readily classify the multiquark systems. In the 3-quark case, we can resolve the N and Δ form factors. In the multibaryon system, this technique is essential since it cannot be guaranteed that all quarks have different quantum numbers.

As we have discussed in the previous section, the QCD predictions for the Q^2 dependence of the deuteron reduced form factor in the high Q^2 regime above 1 GeV^2 agree well with the available experimental data. We have also decomposed the multiquark systems into multibaryon physical components and hidden color components, and expanded each component in terms of the QCD eigensolutions. Through the evolution of each component we can derive constraints on the effective force between the clusters. Using the toy-$SU(2)_c$-dibaryon analysis, we find that colorless clusters tend to be repulsive but colorful clusters are attractive at short distances.

LIMITATIONS OF TRADITIONAL NUCLEAR PHYSICS[♯40]

The fact that the QCD prediction for the reduced form factor $Q^2 f_d(Q^2) \sim const$ appears to be an excellent agreement with experiment for $Q^2 > 1\ GeV^2$ provides an excellent check on

♯39 See Ref. 2.

♯40 A more detailed discussion of the material of this section is given in S. J. Brodsky, to be published in the proceedings of the NATO Pacific Summer Institute "Progress in Nuclear Dynamics", Vancouver Island (1982).

the six-quark description of the deuteron at short-distance as well as the scale-invariance of the $qq \to qq$ scattering amplitude. On the other hand, the impulse approximation form[♯41] used in standard nuclear physics calculations

$$F_d(Q^2) = F_N(Q^2) \times F_{\rm Body}(Q^2) \tag{36}$$

is invalid in QCD at large Q^2 since off-shell nucleon form factors enter [see Fig. 8(a)]. The usual treatment of nuclear form factors also overestimates the contribution of meson exchange currents [Fig. 8(b)] and $N\bar{N}$ contributions [Fig. 8(c)] since they are strongly suppressed by vertex form factors as we shall show in this section.

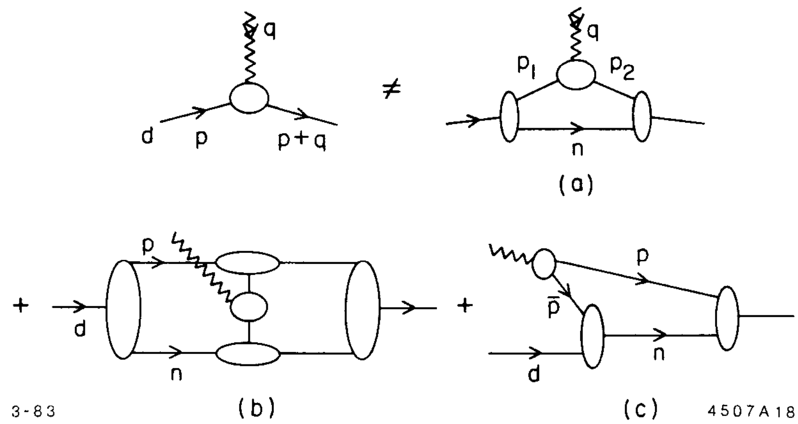


Fig. 8. Critique of the standard nuclear physics approach to the deuteron form factor at large Q^2. (a) The effective nucleon form factor has one or both legs off-shell: $|p_1^2 - p_2^2| \sim q^2/2$. (b) Meson exchange currents are suppressed in QCD because of off-shell form factors. (c) The nucleon pair contribution is suppressed because of nucleon compositeness. Contact terms appear only at the quark level.

At long distances and small, non-relativistic momenta, the traditional description of nuclear forces and nuclear dynamics based on nucleon, isobar, and meson degrees of freedom appears to give a viable phenomenology of nuclear reactions and spectroscopy. It is natural to try to extend the predictions of these models to the relativistic domain, e.g., by utilizing local meson-nucleon field theories to represent the basic nuclear dynamics, and to use an effective Dirac equation to describe the propagation of nucleons in nuclear matter. An interesting question is whether such approaches can be derived as a "correspondence" limit of QCD, at least in the low momentum transfer ($Q^2R_p^2 \ll 1$) and low excitation energy domain ($M\nu \ll M'^2 - M^2$).

The existence of hidden-color Fock state components in the nuclear state in principle precludes an exact treatment of nuclear properties based on meson-nucleon-isobar degrees of freedom since these hadronic degrees of freedom do not form a complete basis on QCD. Since the deuteron form factor is dominated by hidden color states at large momentum transfer, it cannot be described by np, $\Delta\Delta$ wavefunction components on meson exchange currents alone.

♯41 See, e.g., S. A. Gurvitz, Phys. Rev. C22, 1650 (1980). The derivations in this paper require that hadronic interactions have Gaussian fall off.

It is likely that the hidden color states give less than a few percent correction to the global properties of nuclei; nevertheless, since extra degrees of freedom lower the energy of a system it is even conceivable that the deuteron would be unbound were it not for its hidden color components!

Independent of hidden color effects, we can still ask whether it is possible—in principle—to represent composite systems such as mesons and baryons as local fields in a Lagrangian field theory, at least for sufficiently long wavelengths such that internal structure of the hadrons cannot be discerned. Here we will outline a method to construct an effective Lagrangian of this sort. First, consider the ultraviolet-regulated QCD Lagrangian density $\mathcal{L}^{\kappa}_{\rm QCD}$ defined such that all internal loops in the perturbative expansion are cut off below a given momentum scale κ. Normally κ is chosen to be much larger than all relevant physical scale. Because QCD is renormalizable, $\mathcal{L}^{\kappa}_{\rm QCD}$ is form-invariant under changes of κ provided that the coupling constant $\alpha_s(\kappa^2)$ and quark mass parameter $m(\kappa^2)$ are appropriately defined. However, if we insist on choosing the cutoff κ to be as small as hadronic scales then extra ("higher twist") contributions will be generated in the effective Lagrangian density:

$$\begin{aligned}\mathcal{L}^{\kappa} = \mathcal{L}^{\kappa}_0 + \frac{em(\kappa)}{\kappa^2}\,\bar{\psi}_N\sigma_{\mu\nu}\partial^{\mu}\psi_N A^{\nu}_{\rm em} + e\,\frac{f_\pi^2}{\kappa^2}\,\phi^{\dagger}_{\pi}\partial_{\mu}\phi_{\pi}A^{\mu}_{\rm em} \\ + e\,\frac{f_p^2}{\kappa^2}\,\bar{\psi}_N\gamma_{\nu}\psi_N A^{\nu}_{\rm em} + \frac{f_p^2 f_\pi}{\kappa^6}\,\partial_{\nu}\bar{\psi}_N\gamma_5\gamma^{\nu}\psi_N\phi_{\pi} + \cdots\end{aligned} \tag{37}$$

where $\mathcal{L}^{\kappa}_0$ is the standard Lagrangian and the "higher twist" terms of order κ^{-2}, κ^{-4}, ... are schematic representations of the quark Pauli form factor, the pion and nucleon Dirac form factors, and the pion nucleon-antinucleon coupling. The pion and nucleon fields ϕ_π and ψ_N represent composite operators constructed and normalized from the valence Fock amplitudes and the leading interpolating quark operators. One can use the above equation to estimate the effective asymptotic power law behaviors of the couplings, e.g., $F^{\rm quark}_{\rm Pauli} \sim 1/Q^2$, $F_\pi \sim f_\pi^2/Q^2$, $G_M \sim f_p^2/Q^4$ and the effective $\pi\bar{N}\gamma_5 N F_{\pi N\bar{N}}$ coupling: $F_{\pi N\bar{N}}(Q^2) \sim M_N f_p^2 f_\pi/Q^6$. The net pion exchange amplitude for $NN - NN$ scatterings thus falls off very rapidly at large momentum transfer $M^{\pi}_{NN\to NN} \sim (Q^2)^{-7}$ much faster than the leading quark interchange amplitude $M^{qq}_{NN\to NN} \sim (Q^2)^{-4}$. Similarly, the vector exchange contributions give contributions $M^{\rho}_{NN\to NN} \sim (Q^2)^{-6}$. Thus meson exchange amplitudes and currents, even summed over their excited spectra, do not contribute to the leading asymptotic behavior of the nucleon-nucleon scattering amplitudes or deuteron form factors once proper account is taken of the off-shell form factors which control the meson-nucleon-nucleon vertices.

Aside from such estimates, an effective Lagrangian only has utility as a rough tree graph approximation; in higher order the hadronic field terms give loop integrals highly sensitive to the ultraviolet cutoff because of their non-renormalizable character. Thus an effective meson-nucleon Lagrangian serves to organize and catalog low energy constraints and effective couplings, but it is not very predictive for obtaining the actual dynamical and off-shell behavior of hadronic amplitudes due to the internal quark and gluon structure.

Local Lagrangian field theories for systems which are intrinsically composite are however misleading in another respect. Consider the low-energy theorem for the forward Compton amplitude on a (spin-average) nucleon target[♯42]

$$\lim_{\nu \to 0} \mathcal{M}_{\gamma p \to \gamma' p}(\nu, t=0) = -2\hat{\epsilon} \cdot \hat{\epsilon}' \, \frac{e^2}{M_p} \; . \tag{38}$$

One can directly derive this result from the underlying quark currents as indicated in Fig. 9(b). However, if one assumes the nucleon is a local field, then the entire contribution to the Compton amplitude at $\nu = 0$ would arise from the nucleon pair z-graph amplitude, as indicated in Fig. 9(a). Since each calculation is Lorentz and gauge invariant, both give the desired result. However, in actuality, the nucleon is composite and the $N\bar{N}$ pair term is strongly suppressed: each $\gamma p\bar{p}$ vertex is proportional to

$$\langle 0 | \, J^\mu(0) \, | p\bar{p} \rangle \;\; \propto \;\; F_p(Q^2 = 4M_p^2) \;\; ; \tag{39}$$

i.e., the timelike form factor as determined from $e^+e^- \to p\bar{p}$ near threshold. Thus, as would be expected physically, the $N\bar{N}$ pair contribution is highly suppressed for a composite system (even for real photons). Clearly a Lagrangian based on local nucleon fields gives an inaccurate description of the actual dynamics and cannot be trusted away from the forward scattering, low energy limit.

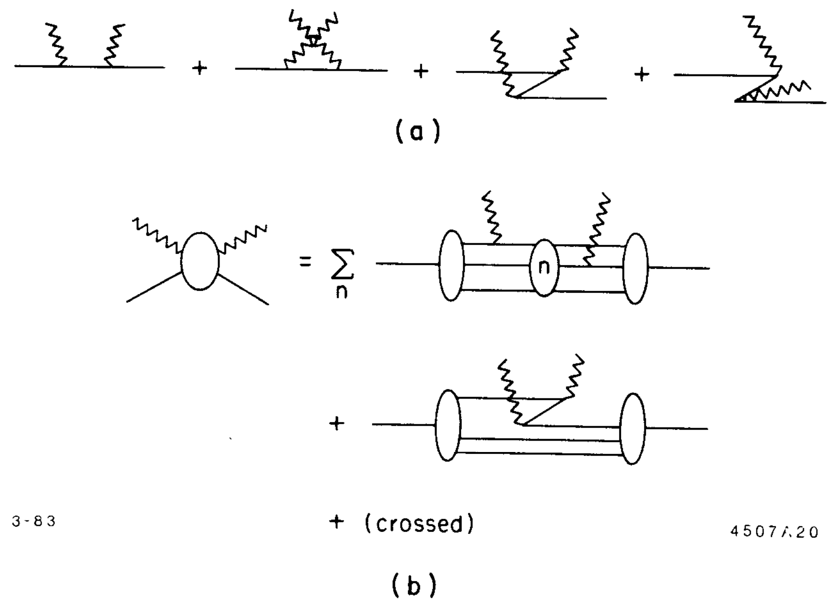


Fig. 9. Time-ordered contributions to (a) The Compton amplitudes in a local Lagrangian theory such as QED. Only the Z-graphs contribute in the forward low energy limit. (b) Calculation of the Compton amplitude for composite systems.

We can see from the above discussion that a necessary condition for utilizing a local Lagrangian field theory as a dynamical approximation to a given composite system H is that its timelike

♯42 A general analysis of Compton scattering on composite systems, including the essential effects of Lorentz boosts, is given by S. J. Brodsky and J. R. Primack, Ann. Phys. 52, 315 (1969). See also D. Drechsel, this volume.

form factor at the Compton scale must be close to 1:

$$F_H(Q^2 \simeq 4M^2) \simeq 1 . \qquad (40)$$

For example, even if it turns out that the electron is a composite system at very short distances, the QED Lagrangian will still be a highly accurate tool. The above condition on the timelike form factor of threshold fails for all hadrons, save the pion. This result does suggest that effective chiral field theories which couple point-like pions to quarks could be a viable approximation to QCD.

More generally, one should be critical of any use of point-like couplings for nucleon-antinucleon pair production, e.g., in calculations of deuteron form factors, photo- and electro-disintegration since such contributions are always suppressed by the timelike nucleon form factor. Notice that $\gamma N\bar{N}$ point-like couplings are not needed for gauge invariance, once all quark current contributions including pointlike $q\bar{q}$ pair terms are taken into account.

We also note that a relativistic composite fermionic system, whether it is a nucleon or nucleus, does not obey the usual Dirac equation—with a momentum-independent potential—beyond first Born approximation.[♯43] Again, the difficulty concerns intermediate states containing $N\bar{N}$ pair terms: the validity of the Dirac equation requires that $\langle p|V_{\rm ext}|p'\rangle$ and $\langle 0|V_{\rm ext}|p'\bar{p}\rangle$ be related by simple crossing, as for leptons in QED. For composite systems the pair production terms are again suppressed by the timelike form factor. It is however possible that one can write an effective, approximate relativistic equation for a nucleon in an external potential of the form

$$(\vec{\alpha}\cdot\vec{p} + \beta m_N + \Lambda_+ V_{\rm eff}\Lambda_+)\Psi_N = E\Psi_N \qquad (41)$$

where the projection operator Λ_+ removes the $N\bar{N}$ pair terms, and $V^{\rm eff}$ includes the local (seagull) contributions from $q\bar{q}$-pair intermediate states, as well as contributions from nucleon excitation.

An essential property of a predictive theory is its renormalizability, the fact that physics at a very high momentum scale $k^2 > \kappa^2$ has no effect on the dynamics other than to define the effective coupling constant $\alpha(\kappa^2)$ and mass terms $m(\kappa^2)$. Renormalizability also implies that fixed angle unitarity is satisfied at the tree-graph (no-loop) level. In addition, it has recently been shown that the tree graph amplitude for photon emission for any renormalizable gauge theory has the same amplitude zero structure as classical electrodynamics. Specifically, the tree graph amplitude for photon emission caused by the scattering of charged particles *vanishes* (independent of spin) in the kinematic region where the ratios $Q_i/p_i\cdot k$ for all the external charged lines are identical. This "null zone"[♯44] of zero radiation is not restricted to soft photon momentum, although it is identical to the kinematic domain for the complete destructive interference of the radiation associated with classical electromagnetic currents of the external charged particles. Thus the tree graph structure of gauge theories, in which each

♯43 S. J. Brodsky, in New Horizons in Electromagnetic Physics, Charlottesville, Virginia (1982). Related difficulties with Dirac-equation applications have also been noted by D. A. Adams and M. Bleszynski, Los Alamos preprint LA-UR-83-2749 (1983).

♯44 S. J. Brodsky, R. W. Brown and K. L. Kowalski, Phys. Rev. D28, 624 (1983), and references therein.

elementary charged field has zero anomalous moment ($g = 2$) is properly consistent with the classical ($\hbar = 0$) limit. On the other hand, local field theories which couple particles with non-zero anomalous moments violate fixed angle unitarity and the above classical correspondence limit at the tree graph level. The anomalous moment of the nucleon is clearly a property of its internal quantum structure; by itself, this precludes the representation of the nucleon as a local field.

The essential conflict between quark and meson-nucleon field theory is thus at a very basic level: because of Lorentz invariance a conserved charge must be carried by a local (point-like) current; there is no consistent relativistic theory where fundamental constituent nucleon fields have an extended charge structure.

WHEN IS PERTURBATIVE QCD APPLICABLE?

An important phenomenological question for the application of QCD to nuclear physics is the momentum transfer scale μ where perturbative predictions become reliable. Ignoring heavy quark thresholds, the natural scale parameters of QCD are $\Lambda_{\overline{MS}}$ ($\sim 100 \pm 50\ MeV$), the mass scale of the light hadrons ($\lesssim 1$ GeV), and the constituent transverse momenta $\langle k_\perp^2 \rangle^{1/2} \sim 300\ MeV$. Thus *a priori* we expect the nominal power-law behavior predicted by QCD hard subprocesses to be reliable for $Q^2 \gg \mu^2$; i.e., Q^2 beyond a few GeV^2.

In the case of deep inelastic lepton-nucleon scattering, Bjorken scaling, which reflects the scale-invariant behavior of incoherent lepton quark scattering becomes evident for $Q^2 > 1\ GeV^2$, $W > 1.8\ GeV$. Coherent contributions, which occur when, e.g., two different struck quarks interfere become relevant for $Q^2 \lesssim \mathcal{O}\langle k_\perp^2 \rangle$ such that there is significant overlap in the final state. In the case of exclusive processes, the leading QCD power law dominates when the nucleon valence $|qqq\rangle$ or meson valence $|q\bar{q}\rangle$ Fock state contributions overtake the faster falling contributions from higher Fock states. Phenomenologically, the onset of the leading power law occurs at $Q^2 \sim$ few GeV^2. In the case of the deuteron form factor, $(Q^2)^5 F_d(Q^2)$ cannot be expected to approach constant behavior until considerably larger Q^2 since the virtual photon's momentum in the underlying hard subprocess is divided six ways: i.e. one requires $(Q/6)^2 \gg \langle k_\perp^2 \rangle$. (Detailed numerical estimates are given in Ref. 8.) On the other hand, reduced nuclear amplitudes such as the reduced deuteron form factor can be expected to approach scaling law $Q^2 f_d(Q^2) \to const$ quickly since the relevant hard propagators route to large momentum transfer $\sim 5/6\,Q$. All other sources of scale breaking by definition divide out via the nucleon form factors.

The scaling behavior of the form factors shown in Fig. 2 bear out these expectations. Other QCD predictions for the leading power behavior, including $pp \to pp$, $\pi p \to \pi p$, $\gamma p \to \pi^+ n$, etc. are consistent with the predicted nominal scaling law at momentum transfer $p_T^2 \gtrsim 3\ GeV^2$. Recent measurements of the basic QCD process $\gamma\gamma \to \pi^+\pi^-$ at large angles agree with the QCD predicted scaling behavior for invariant mass $W > 1.8\ GeV$.

An essential question is whether QCD also correctly accounts for the normalization as well as the scaling behavior of high momentum transfer form factors and other exclusive process data. The $\gamma\gamma \to \pi^+\pi^-$, K^+K^-[♮45] agree very well with the absolutely normalized QCD predictions.

♮45 J. R. Smith et al., SLAC-PUB-3205 (1983).

Meson form factor predictions for $Q^2 F_M(Q^2)$ are within a factor or two of the QCD prediction using the most naive form of the meson distribution amplitude, i.e.: $\phi(x) = Cx(1-x)$. Since this distribution is relevant only at asymptotic momentum scales (where the leading anomalous dimension dominates), there is no conflict with existing data.

In the case of the nucleon form factor, the normalization of the QCD prediction $Q^4 G_M(Q^2) \sim$ const is sensitive to three effects.

1. The average interquark separation d_v of the valence wavefunction: it should be emphasized that $G_M(Q^2)$ is proportional to the fourth inverse power of d_v.
2. The shape of the nucleon distribution amplitude $\phi_N(x_i)$ especially near $x_i \sim 1$.
3. The use of the running coupling constant $\alpha_s(Q^2)$ at the correct scale in the $3q + \gamma^* \to 3q$ hard scattering amplitudes. This removes an "accidental" cancellation in $G_M^P(Q^2)$ when the asymptotic distribution amplitude $\phi_N(x) = Cx_1x_2x_3$ is assumed.

If one uses fixed coupling constant, together with a naive non-relativistic wavefunction with symmetric quark distributions and the standard rms radius $\sim 0.8\ fm$, then the predicted normalization of $Q^4 G_M^P(Q^2)$ is two orders of magnitude below experiment.[♯46] However, by taking into account the above three effects it is straightforward to fit the experimental normalization of G_M as well as νW_2^P at $x \to 1$ and the decay rate for $\psi \to p\bar{p}$.

In our work we have noted that the valance wavefunction of the nucleon is likely to be much more compact than indicated by the physical proton radius as derived from an average over all Fock states. More precisely, we define the QCD Fock state expansion for the proton wavefunction at equal $\tau = t + z$ on the light-cone and $A^+ = 0$ gauge, at a given renormalization scale κ.

$$|p\rangle = \begin{pmatrix} \psi^\kappa_{qqq}(x_i, k_{\perp i})\,|qqq\rangle \\ \psi^\kappa_{qqqg}(x_i, k_{\perp i})\,|qqqg\rangle \\ \cdot \\ \cdot \\ \cdot \end{pmatrix}$$

At small Q^2 where the proton rms radius is determined (from $6\,\frac{d}{dQ^2}\,F_1(Q^2)|_{Q^2=0}$) all Fock states contribute. At large Q^2 where $Q^4 G_M^p(Q^2)$ becomes nearly constant, the valence Fock state $\psi_{qqq}(x_i, k_{\perp i})$ is dominant. Since the higher Fock states are analogous to states containing a meson cloud, it is reasonable that the valence state radius $\langle k_\perp^2 \rangle_{3q}^{-1/2}$ is smaller than the total radius. In fact in the case of the pion, this statement can be demonstrated explicitly. Using normalization constraints from the decay amplitudes $\pi^- \to \mu^- \nu$ and $\pi^0 \to \gamma\gamma$ we can determine the valence state probability and radius ($P_{q\bar{q}/\pi} \sim \frac{1}{4}$, $\langle r^2 \rangle^{1/2}_{q\bar{q}/\pi} \sim 0.42\ f_m$).[♯12] This suggests that the nucleon form factors be parametrized with at least two components, one soft, falling at least as fast as $(Q^2)^{-3}$ and the other with a large $> 1\ GeV^2$ mass scale falling as $(Q^2)^{-2}$ asymptotically. It is in fact easy to find parametrizations of this type which fit the standard dipole form.

♯46 This problem with the normalization was discussed in Ref. 11, 28, and in a different context by B. L. Ioffe and A. V. Smilga, Phys. Lett. 114B, 353 (1982).

More recently, the possibility of significant asymmetry in the x-distribution of the valence wavefunction of the nucleon has been investigated by Chernyak and Zhitnitsky[♯14] using the constraints of the ITEP sum rules. The result is the prediction that the u-quark with spin parallel to the proton carries $\sim 2/3$ of the nucleon momentum, leaving 1/6 each for the other two quarks. The predicted normalization for $Q^4 G_M^p$ and $\psi \to p\bar{p}$ is again in good agreement with the data, even using normal a valence state radius. The actual solution for the nucleon form factor will presumably involve a combination of effects. A final resolution to the problem will require more phenomenological input, and actual solutions to the nucleon bound state wavefunction in QCD. Certainly at this point there is no evidence of any difficulty or conflict[♯47] with the predictions of perturbative QCD for either the scaling behavior or the normalization of exclusive processes in the few GeV^2 momentum transfer region.

FUTURE DIRECTIONS

QCD can be regarded as the underlying theory of nuclear phenomena in the same sense that QED is the basis for atomic and molecular physics. At this point we are only at the beginning of quantitative calculations, and further progress will require the development of new theoretical techniques for solving strong coupling theories, boundary condition models, etc. We will also need new experimental input, especially in the transition region between coherent and incoherent quark processes. At this point, theoretical progress is being made in the following areas:

1. One can now find approximate analytic solutions to the light-cone equation of motion in the valence quark sector, thus allowing model computations of hadronic and nuclear wavefunctions, distribution amplitudes, and structure functions. Recently, together with M. Sawicki, we have developed fixed particle number equations analogous to the non-retardation approximation in atomic physics. Because this approach consistently restricts the equation to fixed particle number, non-analytic "cusps" of the type derived by Karmonov[♯48] are avoided.

2. The normalization of the deuteron form factor at large Q^2 is a formidable though feasible task. The calculation of $T_H(6q + \gamma^* \to 6q)$ should be feasible using the methods of Farrar, Maina, and Neri.[♯13] The normalization of $\phi_d(x_i, Q)$ requires a careful study of the six-quark evolution equation and matching to the non-relativistic regime.

3. The above study also allows estimates of corrections to the reduced amplitude formalism and the effects of hidden color states. Recently we have examined the transition from reduced form factors to the low momentum transfer regime ($Q^2 \lesssim 2M\epsilon_{BE}$) where the impulse approximation form becomes valid. More generally the reduced amplitude formalism can be used to redefine the nuclear potential in such a way that nucleon structure is consistently removed.

4. A systematic analysis of all sources of non-additivity in inelastic lepton-nucleus scattering, lepton pair production, and their effects on specific final states is needed before we have a clear understanding of the EMC effect. For example, one possible origin of

♯47 We thus differ with the conclusions of N. Isgur and C. H. Llewellyn Smith, Phys. Rev. Lett. 52, 1080 (1984). See also N. Isgur, this volume.

♯48 V. A. Karmonov, Nucl. Phys. A362, 331 (1981), B166, 378 (1980).

anomalous A-dependence is non-additivity of strange and other sea quarks in the nucleus, which will be apparent in the A-dependence of K^- electroproduction.[♯19] The decrease of the mean value of x with increasing nucleon number implies[♯49] by rotational symmetry a corresponding decrease of $\langle k_\perp^2 \rangle^{1/2}$. The decrease of the intrinsic transverse momentum should then lead to a contribution to the average transverse momentum of lepton pairs produced in hadron-nucleus collisions decreasing with A.

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♯49 G. Bodwin, S. J. Brodsky, and G. P. Lepage, to be published.