SLAC-PUB-2595
August 1980
(T/E)

THE SYNTHESIS OF QUANTUM CHROMODYNAMICS AND NUCLEAR PHYSICS

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Abstract: The asymptotic freedom behavior of quantum chromodynamics allows the rigorous calculation of hadronic and nuclear amplitudes at short distances using perturbative methods. The implications of QCD for large momentum transfer nuclear form factors and scattering processes, as well as for the structure of nuclear wavefunctions and nuclear interactions at short distances are discussed. The necessity for color-polarized internal nuclear states is also discussed.

1. Introduction

In quantum chromodynamics the fundamental degrees of freedom of hadrons and their interactions are the quanta of quark and gluon fields which obey an exact internal SU(3) "color" symmetry. It now seems possible that quantum chromodynamics[1]) is the theory of the strong interactions in the same sense that quantum electrodynamics accounts for electromagnetic interactions. It is well known that the general structure of QCD meshes remarkably well with the facts of the hadronic world, especially quark-based spectroscopy, current algebra, the approximate point-like structure of large momentum transfer lepton-hadron reactions, and the logarithmic violation of scale-invariance in deep-inelastic reactions. The theory is particularly successful in predicting the features of electron-positron annihilation into hadrons: the magnitude and scaling of the total cross section, the production of hadronic jets with a pattern conforming to elementary quark and gluon processes, and heavy quark phenomena. The empirical results are consistent with the basic postulates of QCD, that the charge and weak currents within hadrons are carried by the quarks, and that the strength of the quark-gluon couplings become weak at short distances (asymptotic freedom).

It is clear that if QCD is the correct theory of the strong interactions it must account for the features and interactions of nuclei as well as mesons and baryons. Because of asymptotic freedom[1]), we can in fact make detailed predictions for nuclear form factors and scattering processes at large momentum transfer, as well as predict the asymptotic short distance features of the nucleon-nucleon interaction and nuclear wavefunctions. We will also discuss here a particularly novel feature of QCD: the necessity for color-polarized (or "hidden-color") nuclear states.

In terms of their Fock state description, the hadrons, including nuclei, are (color singlet) composites of quark and gluon quanta; e.g.,

$$|\pi^+\rangle = a^{\pi}_{(2)}|u\bar{d}\rangle + a^{\pi}_{(3)}|u\bar{d}g\rangle + \ldots$$

$$|p\rangle = a^{p}_{(3)}|uud\rangle + a^{p}_{(4)}|uudg\rangle + \ldots$$

$$|D\rangle = a^{D}_{(6)}|uud\,ddu\rangle + \ldots$$

DISCLAIMER

This is in exact analogy to the Fock state expansion

$$|\text{positronium}\rangle = a_{(2)}|e^+e^-\rangle + a_{(3)}|e^+e^-\gamma\rangle + \ldots$$

* Work supported in part by the Department of Energy under contract number DE-AC03-76SF00515 and by the National Science Foundation.
(Invited talk presented at the 9th International Conference on the Few Body Problem, Eugene, Oregon, August 17-23, 1980. To be published in Nucl. Phys. A.)

in QED. For definiteness we shall specify the states at equal "time" $\tau = t+z$ on the light-cone[2]). The total 3-momenta, $\vec{k}_\perp$ and $k^+ \equiv k^0 + k^z$, are then conserved $\left(\sum_i \vec{k}_{\perp i} = 0,\ \sum_i k_i^+ = p^+\right)$, and the momentum-space wavefunction for each n-particle Fock state component of a hadron with total momentum p^μ is a function (spin-labels are suppressed) $\Psi' = \Psi(x_i, \vec{k}_{\perp i})$, where $x_i = k_i^+/p^+$, $\sum_i \vec{k}_{\perp i} = 0$, and $\sum_i x_i = 1$. The states are off the "energy" shell ($k^- \equiv k^0 - k^z$)

$$p^- - \sum_{i=1}^{n} k_i^- = \frac{1}{p^+}\left[M^2 - \sum_{i=1}^{n}\left(\frac{\vec{k}_\perp^2 + m^2}{x}\right)_i\right] < 0 \quad . \qquad (1)$$

This "light-front" formalism greatly simplifies relativistic bound state calculations since it is very similar to the ordinary non-relativistic many-particle theory; essentially all theorem proved in the Schroedinger theory hold for the relativistic theory at equal time on the light-cone. The "relativistic kinetic energy" $(\vec{k}_\perp^2 + m^2)/x$ plays the role of the non-relativistic kinetic energy $\vec{k}^2/2m$. Although the results are independent of the choice of Lorentz frame, the variable $x = k^+/p^+$ can be conveniently interpreted as the longitudinal momentum fraction k^z/p^z in an "infinite momentum" frame where $p^z \to \infty$. A summary of the calculational rules in light-cone perturbation theory, and examples of its use are given in ref. 3 and references therein.

QCD has a renormalizable perturbation theory which is an elegant generalization of QED. In addition to the $g_s \bar{\psi}\gamma_\mu \psi A^\mu$ Dirac coupling of the color octet gluon to the color triplet and anti-triplet q and $\bar{q}$, there are also non-Abelian 3-point ggg and 4-point gggg couplings of the gluons. The crucial distinction is in the net sign of the vacuum polarization: in QED, the effective Coulomb interaction $\alpha(Q^2)$ increases with the momentum transfer; in QCD the added self-gluon couplings causes the effective strength of the quark and gluon interactions to decrease logarithmically at large momentum transfer, as summarized by the asymptotic expression ($Q^2 \gg \Lambda^2_{QCD}$)

$$\alpha_s(Q^2) = \frac{4\pi}{\left(11 - \frac{2}{3} n_f\right)\log\left(Q^2/\Lambda^2_{QCD}\right)} \quad . \qquad (2)$$

The parameter Λ_{QCD} sets the basic mass scale for QCD, and n_f is the number of types or "flavors" of quarks with effective mass much smaller than Q. The fact that QCD interactions tend to zero at large momentum transfer $Q^2 \gg \Lambda^2_{QCD}$ ("asymptotic freedom") allows the rigorous perturbative calculation of hadronic interactions at short distances. The fact that Λ_{QCD} is empirically not very large-current estimates give

$$\Lambda_{QCD} \lesssim 100 - 300 \text{ MeV} \quad ,$$

indicates that perturbative QCD calculations may start to become relevant at hadronic and nuclear distances of a fermi or less. In fact, as we shall see in the following discussions, QCD dynamics must be taken into account in any nuclear process where nucleon structure is relevant.

2. Short distance processes in QCD

In a series of recent papers[3,4]) we have shown that the predictions of perturbative QCD can be extended to the whole domain of large momentum transfer exclusive reactions, including the form factors $F(Q^2)$ of hadrons as measured in large momentum transfer electron scattering reactions $eA \to eB$ ($Q^2 = -t$), and fixed angle scattering processes $d\sigma/dt$ ($A+B \to C+D$) for $s \gg M^2$ with $z = \cos\theta_{c.m.}$ fixed. In general A, B, C, and D can be mesons, baryons, photons, and nuclei! The results for meson form factors have also been derived using different methods by Efremov and Radyushkin[5]) and by Duncan and Mueller[6]).

As an example, let us briefly consider the calculation of the nucleon form factor[3,4]. Only the "valence" |qqq⟩ Fock state needs to be considered to leading order in 1/Q since (in a physical gauge such as $A^+ = 0$) any additional quark or gluon forced to absorb large momentum transfer (proportional to Q) yields a power-law suppression M/Q to the form factor. Further, because of the spin-1, helicity-conserving couplings of the gauge gluon, overall hadronic helicity is conserved, $h_I = h_F$, again to leading order in 1/Q. Thus QCD predicts the suppression of the Pauli form factor: $F_2(Q^2)/F_1(Q^2) \sim \mathcal{O}(M^2/Q^2)$.

The calculation of the nuclear form factor thus reduces to a 3-body problem. The helicity-conserving form factor $G_M(Q^2)$ can be written to leading order in 1/Q in the factorized form $[\tilde{Q}_x = (\min_i x_i)Q]$:

$$G_M(Q^2) = \int_0^1 [dx] \int_0^1 [dy]\, \phi^\dagger(y_i,\tilde{Q}_y)\, T_H(x,y,Q)\, \phi(x_i,\tilde{Q}_x) \quad , \qquad (3)$$

with $[dx] = dx_1 dx_2 dx_3 \delta(1-x_1-x_2-x_3)$. Here $T_H(x,y,Q)$ is the probability amplitude for scattering three quarks collinear with p to the final direction p + q, as illustrated in fig. 1. To leading order in $\alpha_s(Q^2)$,

$$T_H(x,y,Q) = \left[\frac{\alpha_s(Q^2)}{Q^2}\right]^2 t(x,y) \qquad (4)$$

where t(x,y) is a rational function of the light-cone longitudinal momentum fractions x_i and y_i. The function $\phi(x_i,Q)$ gives the probability amplitude for finding the valence quarks in the nucleon with light-cone fractions x_i at small relative distances $b_\perp \sim \mathcal{O}(1/Q)$:

$$\phi(x_i,Q) \propto \int^{k_{\perp i}<Q} [d^2k_\perp]\, \psi(x_i,\vec{k}_{\perp i}) \quad . \qquad (5)$$

The "distribution amplitude" $\phi(x_i,Q)$ is the fundamental wavefunction which controls high-momentum transfer exclusive reactions in QCD; it is the analogue of the wavefunction at the origin in the non-relativistic theory. The (logarithmic) Q-dependence of ϕ can be completely determined by the operator product expansion at short distances[7] and the renormalization group, or by "evolution equations" computed from perturbative quark-quark scattering kernels at large momentum transfer[3,4]. The high momentum tail of the wavefunction for each hadron is thus controlled by QCD perturbation theory.

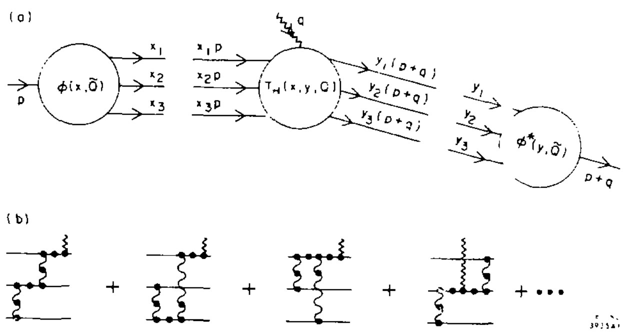

Fig. 1. (a) The general factorized structure of the nucleon form factor at large Q^2 in QCD. (b) Leading contributions to the hard-scattering amplitude T_H.

The QCD prediction for nucleon form factors at large Q^2 to leading order in $\alpha_s(Q^2)$ is[7,3])

$$G_M^N(Q^2) = \frac{\alpha_s^2(Q^2)}{Q^4} \sum_{n,m=0}^{\infty} b_{nm} \left[\ln \frac{Q^2}{\Lambda^2}\right]^{-\gamma_n-\gamma_m} \quad . \tag{6}$$

The γ_n^N anomalous dimensions are known positive constants determined by the evolution equation for the nucleon distribution amplitude. Since $\alpha_s(Q^2)$ is slowly varying at large Q^2, the most important dynamical behavior is the Q^{-4} power-law dependence of $G_M(Q^2)$ which reflects the basic scale invariance of quark and gluon interactions and the fact that the minimal Fock state of the nucleon contains three quarks -- both non-trivial features of QCD. The prediction (6) for $G_M^p(Q^2)$ agrees with data for $3 \lesssim Q^2 \lesssim 25$ GeV^2 provided $\Lambda_{QCD} \lesssim 300$ MeV, as shown in fig. 2. A detailed discussion is given in ref. 3. The QCD predictions for weak and electromagnetic elastic and transition baryon form factors and the n/p ratio are given in ref. 8.

The power-law behavior of the QCD predictions for exclusive processes at large momentum transfer can be summarized by simple counting and helicity rules. To leading order in 1/Q: (i) Total hadron helicity is conserved[3]). In particular this implies that weak and electromagnetic form factors are helicity conserving, and independent of total spin. The dominant form factor corresponds to $h_I = h_F = 0$ or $h_I = h_F = \pm\frac{1}{2}$. (ii) Dimensional counting[9]) predicts the power-law scaling of fixed-angle scattering processes:

$$\frac{d\sigma}{dt}\,(AB \to CD) \sim \frac{1}{s^{n-2}}\, f(\theta_{c.m.}) \quad , \tag{7}$$

where n = total number of constituent fields in A,B,C, and D, and the power-law fall-off of helicity conserving form factors:

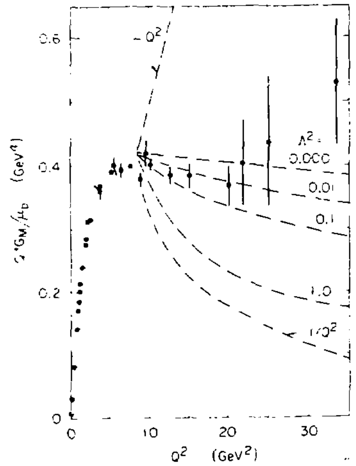

Fig. 2. QCD prediction for $Q^4 G_M^p(Q^2)$ for various scale parameters Λ^2 (in GeV^2). The data are from M. D. Mestayer, SLAC Report 214 (1978) and references therein.

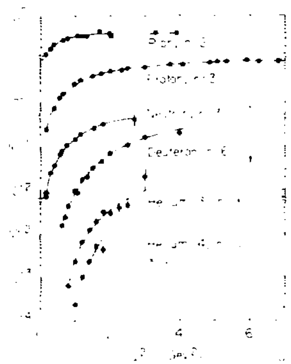

Fig. 3. Comparison of the dimensional counting rule $(Q^2)^{n-1}F(Q^2) \to$ const. $(Q^2 \gg M^2)$ with data. See ref. 21 and references therein.

$$F_H(Q^2) \sim \frac{1}{(Q^2)^{n_H - 1}} \quad , \tag{8}$$

where n_H is the number of constituent fields in H.

In particular, for the deuteron QCD predicts[illegible]

$$F_D(Q^2) \sim \left(\frac{1}{Q^2}\right)^5 \tag{9}$$

for the helicity zero → helicity zero elastic form factor. The other helicity form factors are suppressed by powers. All of these results are modified by calculable logarithm corrections, as in eq. (6). The derivations utilize the fact that Sudakov factors suppress possible anomalous contributions from end-point integration regions and pinch singularities[1,6,11]. The predictions for the power-law behavior reflect the scale-invariance of renormalizable interactions and appear to be in accord with large momentum transfer experiments. A recent comparison with data is given in fig. 3.

3. Nuclear applications of quantum chromodynamics

The deuteron's Fock state structure is much richer in QCD than it would be in a theory in which the only degrees of freedom are hadrons. Restricting ourselves to the six-quark valence state, we can readily generate states like

$$|D\rangle_6 = a\,|(uud)_{1C}(ddu)_{1C}\rangle + b\,|(uud)_{8C}(ddu)_{8C}\rangle$$
$$+\ c\,|(uuu)_{1C}(ddd)_{1C}\rangle + d\,|(uuu)_{8C}(ddd)_{8C}\rangle \quad . \tag{10}$$

The first component corresponds to the usual n-p structure of the ducteron. The second component corresponds to "hidden color" or "color polarized" configurations where the three-quark clusters are in color-octets, but the overall state is a color-singlet. The last two components are the corresponding isobar configurations. If we suppose that at low relative momentum the deuteron is dominated by the n-p configuration, then quark-quark scattering via single gluon exchange generates the color polarized states (b) and (d) at high $k_\perp$; i.e., there must be mixing with color-polarized states in the deuteron wavefunction at short distances.

It is interesting to speculate on whether the existence of these new configurations in normal nuclei could be related to the repulsive core of the nucleon-nucleon potential[12], and the enhancement[13] of parity-violating effects in nuclear capture reactions. One may also expect that there are resonance states with nuclear quantum numbers which are dominantly color-polarized. The mass of these states is not known; if in the unlikely case they are nearly degenerate with ordinary nucleons then they could be long-lived and play havoc with detailed balance experiments. It has also been speculated[14,15] that such long-lived states could have an anomonously large interaction cross section, and thus account for the Judek[16] anomaly in cosmic ray and heavy ion experiments[17]. Independent of these (wild) speculations, it is clearly important that detailed high-resolution searches for these states be conducted, particularly in inelastic electron scattering and tagged photon nuclear target experiments.

In analogy with the nuclear form factor calculation, the QCD prediction for the leading helicity zero deuteron form factor has the form

$$F_D(Q^2) \sim \left[\frac{\alpha_s(Q^2)}{Q^2}\right]^5 \sum_{n,m=0}^{\infty} d_{nm} \left[\ln \frac{Q^2}{\Lambda^2}\right]^{-\gamma_n^D - \gamma_m^D} \tag{11}$$

where the first factor is computed from the sum of hard-scattering $6q + \gamma^* \to 6q$ diagrams. [See fig. 4.] The anomalous dimensions can be calculated from a system of evolution equations for the coupled six-quark components of the deuteron form factor at short distances[illegible].

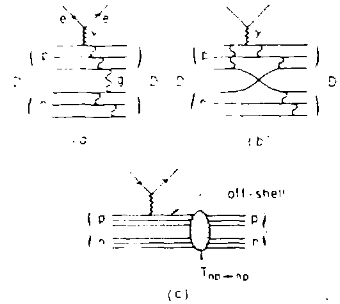

Fig. 4. Hard-scattering contributions to the deuteron form factor. The contribution of diagram (a) requires an internal color-polarized state. Diagram (b) shows the relationship of the nuclear form factor to the N-N off-shell scattering amplitude.

For a general nucleus, the asymptotic power behavior is[18,19] $F_A(Q^2) \sim (Q^2)^{1-3A}$, reflecting the fact that one must pay a penalty of $\alpha_s(Q^2)/Q^2$ to move each quark constituent from p to p+q. The fact that the momentum transfer must be partitioned among the constituents implies that the truly asymptotic regime increases with the nucleon number A. Nevertheless, the QCD perturbative structure is still relevant even in the subasymptotic domain where the nucleus can still be regarded to first approximation as a bound state of nucleons.

In order to make quantitative predictions, let us consider elastic electron-deuteron scattering in a general Lorentz frame. The deuteron form factor $F_D(Q^2)$ is the probability amplitude for the nucleus to stay intact after absorbing momentum transfer Q. Clearly $F_D(Q^2)$ must fail at least as fast as $F_p(Q^2/4)F_n(Q^2/4)$ since each nucleon must change momentum from $\sim \frac{1}{2}p$ to $\sim \frac{1}{2}(p+q)$ *and stay intact*. Thus we should consider a "reduced form factor" $f_D(Q^2)$ defined via[10,19,20]

$$F_D(Q^2) \equiv F_N^2(Q^2/4)\, f_D(Q^2) \quad . \tag{12}$$

Note that $f_D(Q^2)$ must decrease at large Q^2 since it can be identified as the probability amplitude for the final n-p system to remain a ground state deuteron. In fact, the QCD counting rule eq. (8) predicts $f_D(Q^2) \sim 1/Q^2$ for a scale invariant theory. It is easy to see that a T_H diagram such as fig. 4b, where a gluon immediately transfers momentum $\frac{1}{2}q^{L}$ to the other nucleo., gives the form

$$F_D(Q^2) \sim F_N^2(Q^2/4)\, \frac{\alpha_s(Q^2)}{1+Q^2/m^2} \quad . \tag{13}$$

The mass parameter m can be estimated[19]) from the parameters in the meson and nucleon *form factors and is expected to be small* ($m^2 \sim 0.3$ GeV^2). *The comparison* of data[21]) for $f_D(Q^2)$ with the prediction[19]) $(Q^2+0.3\ \text{GeV}^2)f_D(Q^2) \to$ const. is given *in fig. 5. Remarkably, the prediction* seems to be accurate from Q^2 below 1 GeV^2 *out to the limits of the* experimental data.

In general, we can define reduced nuclear form factors[10])

$$f_A(Q^2) \equiv \frac{F_A(Q^2)}{[F_N(Q^2/A^2)]^A} \quad . \tag{14}$$

QCD then predicts the power-behavior $f_D(Q^2) \sim (Q^2)^{1-A}$ (as if the nucleons were elementary). Comparisons with the data for D, H_e^3 and H_e^4 are given in ref. 10. The fact that $f_D^{expt}(Q^2) \sim (Q^2)^{-1}$ is a remarkable success for QCD. We note that the usual *nuclear physics formula*

$$F_A(Q^2) = F_N(Q^2)F_{body}(Q^2) \tag{15}$$

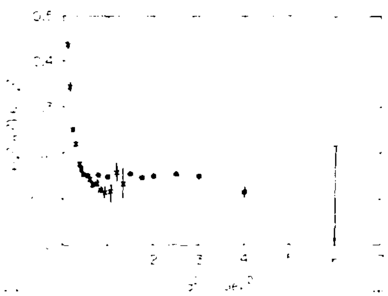

Fig. 5. Comparison of deuteron form factor data with the QCD prediction $(1+Q^2/m^2) \times f_D(Q^2) \to$ const. at large Q^2. The data are *from ref. 21.*

which is supposed to remove the effects of the struck nucleon's structure is <u>invalid</u> in QCD. [A nucleon with momentum $\frac{1}{2}p^{\mu}$ which absorbs momentum transfer q^{μ} as in fig. 4c becomes far off-shell and spacelike $(\frac{1}{2}p+Q)^2 \sim \frac{1}{2}q^2$, so that the total $\gamma^{*}+N \to N^{*}$ off-shell Q^2-dependence is not given by $F_N(Q^2)$.] The same dynamics which controls the nucleon form factor also controls the nuclear physics mechanism which transfers momentum to the other constituents in the nucleus. The definition of the reduced form factor $f_A(Q^2)$ takes into account the correct partitioning of the nuclear momenta, and thus, to first approximation, represents the nuclear form factor in the limit of point-like nuclear constituents. It may be of interest to see whether a consistent parameterization of nuclear amplitudes can be obtained if in each nuclear scattering process, reduced "point" amplitudes are defined by dividing out the nuclear form factors at the correct partitioned momentum).

4. QCD and the nucleon -- nucleon interaction

The asymptotic-freedom property of QCD implies that the nuclear force at short distances can be computed directly in terms of perturbative QCD hard-scattering diagrams. The basic prediction for the nucleon-nucleon amplitude is (modulo logarithmic factors))

$$T_{NN \to NN} \sim \left(\frac{1}{Q^2}\right)^4 f(\theta_{c.m.}) \qquad (16)$$

where $Q^2 = -t$ is the square of the momentum transfer. The predicted fixed angle scaling behavior),

$$\frac{d\sigma}{dt}(pp \to pp) \sim \frac{1}{s^{10}} f^2(\theta_{c.m.}) \qquad (17)$$

is consistent with the high momentum transfer data. [See fig. 6.] To actually compute the angular distribution of the N-N amplitude is a formidable task since even at the Born level there are[4]) of the order of $3 \cdot 10^6$ connected Feynman diagrams in which five gluons interact with six quarks; in addition, Sudakov

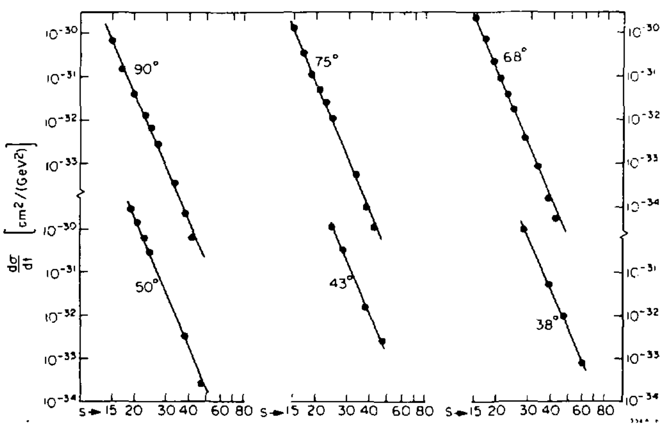

Fig. 6. Differential cross sections for $pp \to pp$ scattering at wide center of mass angles. The straight lines correspond to the predicted power-law fall-off of $1/s^{10}$. The data compilation is from P. V. Landshoff and J. C. Polkinghorne, Phys. Lett. 44B, 293 (1973).

suppressed pinch singularities must be circumvented[4,?]). Considerable phenomenological progress has, however, been made simply by assuming that the dominant diagrams involve quark interchange[?]). This ansatz seems to yield a good approximation to the observed large-angle baryon and meson angular distributions and charge correlations, as well as the observed crossing behavior between amplitudes such as $pp \to pp$ and $p\bar{p} \to p\bar{p}$. Applications to spin correlations are discussed in ref. 2[illegible]. The relation of the scaling behavior of the N-N amplitude to the Q^2-dependence of nucleon form factors (as in fig. 4e) is discussed in detail in ref. 10.

The perturbative structure of QCD at short distances can also be used to determine the far off-shell behavior of hadronic and nuclear wavefunctions and their momentum distributions. For example, the x near 1 behavior of particle distributions in the bound state (at the kinematical end of the Fermi distribution where one constituent has nearly all of the available longitudinal momentum) requires the far off-shell dependence of the wavefunction. Modulo logarithms, the power behavior of perturbative QCD contributions to inclusive distributions is given by the "spectator rule"[?]) $[x_a = (k_a^0 + k_a^z)/(p^0 + p^z)]$

$$\left.\frac{dN_{a/A}}{dx_a}\right|_{x_a \to 1} = C_{a/A}(1-x)^{2n_s-1} \qquad (18)$$

where n_s is the number of spectator constituents in the bound state forced to carry small light-cone momentum fractions. The rule holds for the case where the helicities of a and A are identical; otherwise there is additional power-law suppression[?]). Examples of the spectator rule are $dN/dx \sim (1-x)^3$ for q/p, $(1-x)^{15}$ for q/H_e^3 and $(1-x)^{11}$ for p/H_e^3. These rules can be tested in forward inclusive reactions for particles produced with large longitudinal momenta, and in deep inelastic lepton scattering on hadron and nuclei[?]). In general, the impulse approximation implies[?,?])

$$\frac{d\sigma}{dQ^2 dx}(eA \to eX) = \sum_a \left(\frac{d\sigma}{dQ^2}\right)_{ea \to ea} \frac{dN_{a/A}}{dx} \quad , \qquad (19)$$

representing the sum of incoherent contributions each of which correspond to scattering on one quark or clusters of quarks in the nuclear or hadronic target. Further discussions, applications, and tests can be found in the refs. 10, 19, 20, 21 and 2[illegible]. The transverse momentum distributions $dN_{a/A}/dk_\perp^2$ can also be predicted from the perturbative QCD processes which control the high momentum tail of the bound-state wavefunction[4]).

5. The continuity of nuclear physics and quantum chromodynamics

The synthesis of nuclear dynamics with the quark and gluon processes of quantum chromodynamics is clearly a fascinating fundamental problem in hadron physics. The short distance behavior of the nucleon-nucleon interaction, which is rigorously determined by QCD, must join smoothly and analytically with the large distance constraints of nuclear physics. As we have emphasized here, the fundamental mass scale of QCD is comparable with the inverse nuclear radius; it is thus difficult to argue that nuclear physics at distances below ~1 fm can be studied in isolation from QCD.

The constraints of asymptotic QCD behavior -- especially its power-law scaling and helicity selection rules -- have only begun to be exploited. For example, dispersion relations and superconvergence relations for the hadronic helicity amplitudes should yield sum rules and constraints on hadronic couplings and their spectra[16]). The imposition of duality between the q-$\bar{q}$-g and meson-nucleon degrees of freedom implies even stronger constraints. However, it should be noted that the existence of hidden-color states implies that duality in terms of ordinary hadrons cannot be a true identity. Proof of the existence of the color-polarized configurations -- whether mixed with ordinary nuclear states or appearing as resonance excitations -- clearly would have dramatic implications for QCD and nuclear physics. Speculations on possible phenomena associated with hidden color states have been discussed in section 3.

The most important outstanding problem in the synthesis program is the actual derivation of the nucleon-nucleon interaction from first principles in QCD. Interesting calculational attempts using the MIT bag model are given in refs. 27 and 28. The qualitative similarities between quark interchange amplitudes and meson exchange processes is also evident [ ].

In this talk we have tried to emphasize the continuity between nuclear and elementary particle dynamics. There are many other interesting aspects of the interactions of quarks and gluons within the nuclear environment which are reviewed in ref. 20. There are also important areas in hadron physics which require solutions of few-body problems: e.g., the derivation of the baryon spectrum for three quarks interacting via a QCD confining potential[ ]. It would also be useful to have explicit solutions for positronium-positronium scattering (including pinch singularities) in analogy with meson-meson scattering processes in QCD. We also note that the calculation[4]) of the meson and baryon form factors at large momentum transfer in QCD represents a non-trivial solution of a relativistic few body problem.

Acknowledgements

Part of the work is based on collaborations with R. Blankenbecler, C. Carlson, R. Chertok, G. Farrar, T. Huang, J. Gunion and H. Lipkin. This work was supported in part by the Department of Energy under contract number DE-AC03-76SF00515 and by the National Science Foundation.

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