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NUCLEAR CHROMODYNAMICS: APPLICATIONS OF QUANTUM CHROMODYNAMICS TO FEW NUCLEON SYSTEMS*

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ABSTRACT

The implications of quantum chromodynamics to nuclear systems is discussed, with particular emphasis on applications to short-distance, high-momentum-transfer processes. Exact QCD predictions for the deuteron form factor at large momentum transfer and new phenomena outside the range of traditional nuclear physics are discussed.

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1. INTRODUCTION

One of the most interesting problems in hadron physics is the application of quantum chromodynamics[1] (QCD) to multiquark systems, i.e. nuclei. If QCD is correct, then it must provide a fundamental description of nuclear forces and dynamics.[2] As we will discuss in this article, QCD provides new, in some cases dramatic, perspectives to nuclear physics, especially in the high momentum transfer domain ($Q > 1$ GeV) where quark and gluon degrees of freedom and "hidden color" wavefunction components are essential. These applications include corrections to nucleon additivity of nuclear structure functions (the EMC effect), calculations of nuclear amplitudes at large momentum transfer (e.g., the deuteron form factor); the application of "reduced" nuclear amplitudes which are defined to remove the effects of nuclear compositeness in a covariant fashion; evolution equations for nuclear wavefunctions—e.g., the deuteron 6-quark waavefunction evolves to a state which is 80% hidden color at small internucleon separation. Many traditional concepts of standard nuclear physics phenomenology (e.g. the impulse approximation to nuclear form factors, point-like nucleon pair and meson-exchange current contributions to electromagnetic nuclear amplitudes, local meson-nucleon field theory, and simple Dirac equations for relativistic nucleons), require substantial modification.

Conversely, the nucleus provides an important tool[2] for studying central problems of particle physics, such as the evolution of quark and gluon jets in nuclear matter, "color transparency" phenomena which predicts that hadrons which interact at large momentum transfer have negligible rescattering in a nucleus, and "formation zone" phenomena which predicts the absence of inelastic hadronic interactions for hadrons with energy large compared to a scale proportional to the target length.

QCD is a renormalizable non-Abelian gauge theory of color-triplet quark and color-octet gluon fields invariant under color-SU(3) transformations. The theory provides natural explanations for the basic features of hadronic physics: the meson and baryon spectra, quark statistics, the structure of the weak and electromagnetic currents of hadrons, the scale-invariance of hadronic interactions at short distances, and evidently, color (i.e., quark and gluon) confinement at large distances. Many different and diverse tests have confirmed the basic predictions of QCD; however, since tests of quark and gluon interactions must be done within the confines of hadrons there have been few truly quantitative checks. Nevertheless, it appears likely that QCD is the fundamental theory of hadronic and nuclear interactions in the same sense that QED gives a precise description of electrodynamic interactions.

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" multi-quark wavefunction components—nuclear states which are not separable at large distances into the usual color singlet nucleon clusters. There are a number of immediate consequences for nuclear dynamics:

1. The electromagnetic and weak currents within a nucleus are carried solely by the quark fields at any momentum transfer scale $Q^2 = -q_\mu^2$. In the

deep inelastic, large momentum and energy transfer domain, the lepton scatters essentially incoherently off of the individual quark constituents of the nucleus, giving point-like cross sections characteristics of Bjorken scaling,modified by logarithmic corrections to scale-invariance due to QCD radiative corrections. At low momentum transfer the quark currents become coherent, giving cross sections characteristics of multi-quark, nucleonic, or mesonic currents.

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.[3] Although calculations from first principles are still too complicated, recent results derived from effective potential, bag, and soliton models[4] suggests that many of the basic features of the nuclear force can be understood from the underlying QCD substructure. At a more basic level one can give a direct proof[5] from perturbative QCD that the nucleon-nucleon force must be repulsive at short distances (see Section 3).

3. Because of asymptotic freedom, the effective strength of QCD interactions becomes logarithmically weak at short distances and large momentum transfer

$$\alpha_s(Q^2) = \frac{4\pi}{\beta_0 \log(Q^2/\Lambda^2_{\rm QCD})} \qquad (Q^2 \gg \Lambda^2) \, . \tag{1.1}$$

[Here $\beta_0 = 11 - \frac{2}{3}\, n_f$ 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 $\Upsilon(b\bar{b}) \to \gamma X$:[6,7]

$$\alpha_s(0.157\ M_\Upsilon) = \alpha_s(1.5\ GeV) = 0.23 \pm 0.13\ . \tag{1.2}$$

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 $n_c = 3]$

$$\begin{aligned} V(Q^2) &= -C_F\ \frac{4\pi\alpha_V(Q^2)}{Q^2} \\ \alpha_V(Q^2) &= \frac{4\pi}{\beta_0 \log(Q^2/\Lambda_V^2)} \qquad (Q^2 \gg \Lambda_V^2) \end{aligned} \tag{1.3}$$

where[7] $\Lambda_V = \Lambda_{\overline{MS}}\, e^{5/6} \simeq 270 \pm 100\ MeV$. 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 theory becomes applicable, and a microscopic description of short-distance hadronic and nuclear phenomena in terms of quark and gluon subprocesses becomes viable. In this lecture we will particularly emphasize the use of asymptotic freedom and light-cone quantization to derive factorization theorems,[8–10] rigorous boundary conditions, and exact results for nuclear amplitudes at short distances.[5,11,12] This includes the nucleon form factor at large momentum transfer,[10] meson photoproduction amplitudes, deuteron photo- and electro-disintegration,[12] and most important for nuclear physics, exact results for the form of the form factor of nuclei at large momentum transfer.[5,11] 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 and anomalous dimension structure predicted by QCD at high momentum transfer.

4. 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[13] from deep inelastic lepton-nucleus scattering that nuclear structure functions deviate significantly from simple nucleon additivity, much more than would have been expected for lightly bound systems. Further, as we discuss in Section 5, there are many areas where QCD predictions conflict with traditional concepts of nuclear dynamics.

2. EXCLUSIVE PROCESSES IN QCD

One area of important progress in hadron physics in the past few years has been the extension of QCD predictions to the domain of large momentum transfer hadronic and nuclear amplitudes including nuclear form factors, deuteron photodisintegration, etc.[8] 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. To leading order in $1/Q$ the scattering amplitude has the form [see Fig. 1(a)]

$$\mathcal{M} = \int_0^1 T_H(x_j, Q) \prod_{H_i} \phi_{H_i}(x_j, Q)[dx_j] \; . \tag{2.1}$$

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

$$[dx_j] = \prod_{j=1}^{n_i} dx_j \delta\Big(1 - \sum_k^{n_i} x_k\Big) \tag{2.2}$$

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.[9] Here λ_i is the helicity; $x_i \equiv (k^0 + k^z)/(p^0 + p^z)$, $(\sum_{i=1}^n x_i = 1)$, 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}, $\psi_{qqq\bar{q}q} \ldots$. 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[10] [Fig. 1(a)]

$$F_B(Q^2) = \int_0^1 [dy] \int_0^1 [dx] \; \phi_B^\dagger(y_j, Q) T_H(x_i, y_j, Q) \phi_B(x_i, Q) \quad , \tag{2.3}$$

where to leading order in $\alpha_s(Q^2)$, T_H is computed from $3q + \gamma^* \to 3q$ tree graph amplitudes: [Fig. 1(b)]

$$T_H = \left[\frac{\alpha_s(Q^2)}{Q^2}\right]^2 f(x_i, y_j) \tag{2.4}$$

and

$$\phi_B(x_i, Q) = \int [d^2k_\perp]\ \psi_V(x_i, \vec{k}_\perp)\theta(k_\perp^2 < Q^2) \tag{2.5}$$

is the valence 3-quark wavefunction evaluated at quark impact separation $b_\perp \sim O(Q^{-1})$. 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 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 $1/2$) [14]

$$F(Q^2) \sim \left(\frac{1}{Q^2}\right)^{n-1}, \tag{2.6}$$

$$F_1^N \sim \frac{1}{Q^4}, \quad F_\pi \sim \frac{1}{Q^2}, \quad F_d \sim \frac{1}{Q^{10}}, \tag{2.7}$$

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[15] 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 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 at short distance or by solving an evolution equation for the distribution amplitude computed from gluon exchange [Fig. 1(c)], as we discuss in Section 3 for the deuteron. The result for the large Q^2 behavior of the baryon form factor in QCD is[8,10]

$$F_B(Q^2) = \frac{\alpha_s^2(Q^2)}{Q^4} \sum_{n,m} d_{nm} \left(\ell n \frac{Q^2}{\Lambda^2}\right)^{-\gamma_m - \gamma_n} \tag{2.8}$$

where the γ_n are computable anomalous dimensions of the baryon 3-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. Asymptotically the dominant term has the minimum anomalous dimension. The predicted sign of $G_M^p(Q^2)$ at large Q^2 is the same as $G_M^p(0)$. 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 (2.6, 2.7) predicted by perturbative QCD are consistent with experiment.[16] The behavior $Q^4 G_M(Q^2) \sim$ *const* at large Q^2 provides a direct check that the minimal Fock state in the nucleon contains 3 quarks and that the quark propagator and the $qq \to qq$ 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 $\mathcal{M} \sim Q^{4-n_{TOT}} F(\theta_{cm})$ 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[17] of

$\gamma\gamma \to \pi^+\pi^-$, K^+K^- at large invariant pair mass appear to confirm the QCD predictions.[18]

In principle it should be possible to use measurements of the scaling and angular dependence of the $\gamma\gamma \to M\bar{M}$ reactions to measure the shape of the distribution amplitude $\phi_M(x,Q)$. 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

$$\left[M^2 - \sum_{i=1}^{n}\left(\frac{k_\perp^2 + m^2}{x}\right)_i\right]\Psi = V_{\rm QCD}\Psi \tag{2.9}$$

The explicit form for the matrix representation of $V_{\rm QCD}$ and a discussion of the infrared and ultraviolet regulation required to interpret (2.9) is given in Ref. 9. Thus far experiments has not been sufficiently precise to measure the logarithmic variation from 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 3 quark state has small transverse size, i.e., is large at the origin.[9,19] 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 $qqqg$, $qqq\bar{q}q$ (or meson cloud), etc. A small size for the proton valence wavefunction (e.g. $R_{qqq}^p \sim 0.3\,fm$) can also explain the large magnitude of $\langle k_\perp^2\rangle$ of the intrinsic quark momentum distribution needed to understand in 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 $\pi^- \to \gamma\gamma$. One finds a valence state radius $R_{q\bar{q}}^\pi \sim 0.2\,fm$, corresponding to a probability $P_{q\bar{q}}^\pi \sim 1/4$. A detailed discussion is given in Ref. 19.

3. 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. (Explicit representations are given in Ref. 20). Notice that the exchange of a virtual gluon in the deuteron at short distance inevitably produces Fock state components where the 3-quark clusters correspond to color octet nucleons or isobars. Thus, in general, the deuteron wavefunction should have a complete spectrum of hidden-color wavefunction components, although it is likely that these states are important only at small internucleon separation.[21]

Despite the complexity of the multi-color representations of nuclear wavefunctions, the analysis[5] of the deuteron form factor at large momentum transfer can be carried out in parallel with the nucleon case outlined in Section 2. Only the minimal 6-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. 3),

$$F_d(Q^2) = \int_0^1 [dx]\,[dy]\;\phi_d^\dagger(y,Q)\;T_H^{6q+\gamma^*\to 6q}(x,y,Q)\;\;\phi_d(x,Q)\,, \qquad (3.1)$$

where the hard scattering amplitude scales as

$$T_H^{6q+\gamma^*\to 6q} = \left[\frac{\alpha_s(Q^2)}{Q^2}\right]^5 t(x,y)\;\;[1+\mathcal{O}(\alpha_s(Q^2))] \qquad (3.2)$$

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: ($C_F = 4/3$, $C_d = -C_F/5$)

$$\prod_{k=1}^{6} x_k\left[\frac{\partial}{\partial\xi}+\frac{3C_F}{\beta}\right]\tilde{\Phi}(x_i,Q) = -\frac{C_d}{\beta}\int_0^1 [dy]\;V(x_i,y_i)\,\tilde{\Phi}(y_i,Q)\,. \qquad (3.3)$$

Here we have defined

$$\Phi(x_i, Q) = \prod_{k=1}^{6} x_k \, \tilde{\Phi}(x_i, Q), \tag{3.4}$$

and the evolution is in the variable

$$\xi(Q^2) = \frac{\beta_0}{4\pi} \int_{Q_0^2}^{Q^2} \frac{dk^2}{k^2} \alpha_s(k^2) \sim ln\left(\frac{ln \frac{Q^2}{\Lambda^2}}{ln \frac{Q_0^2}{\Lambda^2}}\right) \quad . \tag{3.5}$$

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 $\prod_{i=1}^{5} x_i^{m_i} >$ are given in Ref. 20. The leading anomalous dimension γ_0, corresponding to the eigenfunction $\tilde{\Phi}(x_i) = 1$, is $\gamma_0 = (6/5)(C_F/\beta_0)$.

In order to make more detailed and experimentally accessible predictions, we will define the "reduced" nuclear form factor[11,12] in order to remove the effects of nucleon compositeness (see Section 4):

$$f_d(Q^2) \equiv \frac{F_d(Q^2)}{F_N^2(Q^2/4)} \quad . \tag{3.6}$$

The arguments for the nucleon form factors (F_N) are $Q^2/4$ 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[5]

$$f_d(Q^2) \sim \frac{\alpha_s(Q^2)}{Q^2}\left(ln \frac{Q^2}{\Lambda^2}\right)^{-\frac{2}{5}\frac{C_F}{\beta}} \quad , \tag{3.7}$$

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 (3.7) directly with the available high Q^2 data[16] (see

Fig. 4). In general one would expect corrections from higher twist effects (e.g., mass and $k_\perp$ smearing), 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 in (3.7) 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.

The calculation of the normalization $T_H^{6q+\gamma^* \to 6q}$ to leading order in $\alpha_s(Q^2)$ will require the evaluation of $\sim$ 300,000 Feynman diagrams involving five exchanged gluons. Fortunately this appears possible using the algebraic computer methods introduced in Ref. 22. The method of setting the appropriate scale $\hat{Q}$ of $\alpha_s^5(\hat{Q}^2)$ in T_H is given in Ref. 7.

We note that the deuteron wave function which contributes to the asymptotic limit of the form factor is the totally antisymmetric wave function corresponding to the orbital Young 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[23]

$$\psi_{[6]\{33\}} = \sqrt{\frac{1}{9}}\ \psi_{NN} + \sqrt{\frac{4}{45}}\ \psi_{\Delta\Delta} + \sqrt{\frac{4}{5}}\ \psi_{CC} \quad . \tag{3.8}$$

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 $b^i_\perp \leq O(1/Q) \to 0$. Since this state is 80-percent 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 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[3,5] (see Fig. 5). 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.

4. 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 are not really applicable (see Section 5), 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 contribution of non-valence Fock states wavefunctions, and a rapidly expanding number of radiative corrections and loop diagrams.

The reduced amplitude method,[11,12] 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 in Section 3 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: $\psi_{q/A} = \psi_{N/A} \Pi_N \Psi_{q/N}$ 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 $\sqrt{-t}/A$. 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

$$m = \frac{\mathcal{M}}{\prod_{i=1}^{A} F_N(\hat{t}_i)} \tag{4.1}$$

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

$$f_A(Q^2) \equiv \frac{F_A(Q^2)}{\left[F_N(Q^2/A^2)\right]^A} \sim \left[\frac{1}{Q^2}\right]^{A-1} . \tag{4.2}$$

Comparisons with experiment and predictions for leading logarithmic corrections to this result are given in Refs. 5 and 12. In the case of photo- (or electro-) disintegration of the deuteron one has

$$m_{\gamma d \to np} = \frac{\mathcal{M}_{\gamma d \to np}}{F_n(t_n)F_p(t_p)} \sim \frac{1}{p_T}\, f(\theta_{cm}) \tag{4.3}$$

i.e., the same elementary scaling behavior as for $\mathcal{M}_{\gamma M \to q\bar{q}}$. Comparison with experiment[26] is encouraging (see Fig. 6), showing that as was the case for $Q^2 f_d(Q^2)$, the perturbative QCD scaling regime begins at $Q^2 \geq 1\ GeV^2$. Detailed comparisons and a model for the angular dependence and the virtual photon-mass dependence of deuteron electrodisintegration are discussed in Ref. 12. Other potentially useful checks of QCD scaling of reduced amplitudes are

$$\begin{aligned} m_{pp \to d\pi^+} &\sim p_T^{-2}\, f(t/s) \\ m_{pd \to H^3\pi^+} &\sim p_T^{-4}\, f(t/s) \\ m_{\pi d \to \pi d} &\sim p_T^{-4}\, f(t/s) . \end{aligned} \tag{4.4}$$

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.

5. LIMITATIONS OF TRADITIONAL NUCLEAR PHYSICS[27]

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 the six-quark description of the deuteron at short-distance as well as the scale-invariance of the $qq \to qq$ scattering amplitude. It should also be emphasized that the impulse approximation form used in standard nucleon physics calculations

$$F_d(Q^2) = F_N(Q^2) \times F_{\text{Body}}(Q^2) \tag{5.1}$$

is invalid in QCD at large Q^2 since off-shell nucleon form factors enter [see Fig. 7(a)]. The region of validity[25] of (5.1) is restricted to $Q^2 < \lambda_H^2$ where λ_H^2 is a hadronic scale. The traditional treatment of nuclear form factors also overestimates the contribution of meson exchange currents [Fig. 7(b)] and $N\bar{N}$ contributions [Fig. 7(c)] since they are strongly suppressed by vertex form factors as we shall show in this section.

At long distances and low, 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.[26] 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 ($M\nu \ll M'^2 - M^2$).

As we have discussed in Sections 2 and 4, the existence of hidden-color Fock state components in the nuclear state 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. 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 is it possible—in principle—to represent composite systems such as meson 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:[9]

$$\begin{aligned}\mathcal{L}^{\kappa} = \mathcal{L}^{\kappa}_0 + \frac{em(\kappa)}{\kappa^2}\ \bar{\psi}_N \sigma_{\mu\nu}\partial^{\mu}\psi_N A^{\nu}_{\rm em} + e\ \frac{f_\pi^2}{\kappa^2}\ \phi^{\dagger}_{\pi}\partial_{\mu}\phi_{\pi} A^{\mu}_{\rm em} \\ + e\ \frac{f_p^2}{\kappa^2}\ \bar{\psi}_N\, \gamma_{\nu}\psi_N A^{\nu}_{\rm em} + \frac{f_p^2 f_\pi}{\kappa^6}\ \partial_{\nu}\bar{\psi}_N\, \gamma_5\gamma^{\nu}\psi_N\phi_{\pi} + \cdots\end{aligned} \tag{5.2}$$

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 $\pi - N - N$ 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 Eq. (5.2) 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 $M^{q\bar{q}}_{NN\to NN} \sim (Q^2)^{-4}$. 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 $N - N$ 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, the effective Lagrangian (5.2) 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 nonrenormalizable 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 Lagrangians 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

$$\lim_{\nu \to 0} M_{\gamma p \to \gamma' p}(\nu, t=0) = -2\,\hat{\epsilon} \cdot \hat{\epsilon}'\, \frac{e^2}{M_p}\ . \qquad (5.3)$$

One can directly derive this result from the underlying quark currents as indicated in Fig. 8(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. 8(a). Since each calculation is Lorentz and gauge invariant, both give the desired result (5.3). 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

$$\langle 0|J^{\mu}(0)|p\,\bar{p}\rangle \propto F_p(Q^2 = 4M_p^2) \quad ; \tag{5.4}$$

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.

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

$$F_H(Q^2 \simeq 4M^2) \simeq 1 \; . \tag{5.5}$$

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. Equation (5.5) fails for all hadrons, save the pion, suggesting 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. Note $\gamma N\bar{N}$ 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. Again, the difficulty concerns intermediate states containing $N\bar{N}$ pair terms: the identity 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 suppressed by the timelike form factor (5.4). It is however possible that one can write an effective, approximate relativistic equation for a nucleon in an external potential of the form

$$(\vec{\alpha}\cdot\vec{p}+\beta m_N+\Lambda_+V_{\rm eff}\Lambda_+)\Psi_N=E\Psi_N \tag{5.6}$$

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.[28] This "null zone" 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 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.

6. CONCLUSIONS

The synthesis of nuclear dynamics with the quark and gluon processes of quantum chromodynamics is clearly a basic fundamental problem in hadron physics. The short distance behavior of the nucleon-nucleon interaction as determined by QCD must join smoothly and analytically with the large distance constraints of nuclear physics. As we have emphasized, 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: meson and nucleon degrees of freedom of traditional nuclear physics models must become inadequate at momentum transfer scales $\gtrsim 200$ MeV where nucleon substructure becomes evident.

Thus the essential question for nuclear as well as particle physics is to understand the transition between the meson-nucleon and quark-gluon degrees of freedom. There should be no illusion that this is a simple task; one is dealing with all the complexities and fascinations of QCD such as the effects of confinement and non-perturbative effects intrinsic to the non-Abelian theory. Such considerations also enter the physics associated with the propagation of quarks and gluons in nuclear matter and the phenomenology of hadron and nuclear wavefunctions.[29]

Despite the difficulty of the non-perturbative domain, there is reason for optimism that "nuclear chromodynamics" is a viable endeavor. For example, as we have shown in Section 4 we can use QCD to make predictions for the short distance behavior of the deuteron wavefunction and the deuteron form factor at large momentum transfer. The predictions give a remarkably accurate description of the scaling behavior of the available deuteron form factor data for Q^2 as low as 1 GeV^2. The QCD approach also allows the definition of "reduced" nuclear amplitudes which can be used to consistently and covariantly remove the effect of nucleon compositeness from nuclear amplitudes. An important feature of such predictions is that they provide rigorous constraints on exclusive nuclear amplitudes which have the correct analytic, gauge-invariant, and scaling properties predicted by QCD at short distances. This suggests the construction of boundary condition model amplitudes which simultaneously satisfy low energy and chiral theorems at low momentum transfer as well as the rigorous QCD constraints at high momentum transfers.[30] In addition, by using the light cone formalism, one can obtain a consistent relativistic Fock state wave function description of hadrons and nuclei which ties on to the Schroedinger theory in the non-relativistic regime. One can also be encouraged by progress in non-perturbative methods in QCD such as lattice gauge theories or chromostatics;[31] eventually these approaches should be able to deal with multi-quark source problems.

It is essential to have direct experiment guidance in how to proceed as one develops nuclear chromodynamics. A high duty factor electron accelerator[32] with laboratory energy beyond 4 GeV is an important tool because of the simplicity of the probe and the fact that we understand the coupling of the electron and quark current in QCD. It is also clear that

1. One must have sufficient energy to extend electron scattering measurements from low momentum transfer to the high momentum transfer region with sufficient production energy such that Bjorken scaling can be observed. One certainly does not want to stop at an intermediate momentum transfer domain—a regime of maximal complexity from the standpoint of both QCD and nuclear physics. The recent EMC and SLAC data[13] showing breakdown of simple nucleon additivity in the nuclear structure functions also demonstrates that there is non-trivial nuclear physics even in the high momentum transfer domain.

2. One must have sufficient electron energy to separate the longitudinal and

transverse currents. The σ_L/σ_T separation is essential for resolving individual dynamical mechanisms; e.g. single quark and multiple quark (meson current) contributions.

3. One wishes to study each exclusive channel in detail in order to verify and understand the emergence of QCD scaling laws and to understand how the various channels combine together to yield effective Bjorken scaling. Helicity information is also very valuable. For example QCD predicts that at large momentum transfer, the helicity-0 to helicity-0 deuteron form factor is dominant and that for any large momentum transfer reaction, total hadronic helicity is conserved.[15]

4. One wishes to make a viable search for dibaryon states which are dominantly of hidden color. The argument that such resonances exist in QCD is compelling just from counting of degrees of freedom. The calculation of the mass and width of such resonances is clearly very difficult, since the detailed dynamics is dependent on the degree of mixing with ordinary states, the availability of decay channels, etc. Since hidden color states have suppressed overlap with the usual hadronic amplitudes it may be quite difficult to find such states in ordinary hadronic collisions. On the other hand, the virtual photon probe gives a hard momentum transfer to a single struck quark, and it is thus more likely to be sensitive to the short-distance hidden color components in the target wave function. Adequate electron energy is essential not only to produce dibaryon resonances but also to allow sufficient momentum transfer Q^2 to decrease backgrounds and to provide σ_L/σ_T separation.

5. One wishes to probe and parametrize the high momentum transfer depen-

dence of the deuteron $n-p$ and $\Delta-\Delta$ components, as a clue toward a complete description of the nuclear wavefunction.

6. One wishes to measure the neutron, pion, and kaon form factors.

7. The region well beyond $x=1$ for deep inelastic electron-nucleus scattering is important QCD physics since the virtual quark and gluon configurations in the nuclear wave function are required to be far off shell. Understanding the detailed mechanisms which underlie this dynamics will require coincident measurements and the broadest kinematic region available for σ_L/σ_T separation. The y-variable approach which attributes the electron scattering to nucleon currents is likely to break down even at moderate Q^2. Coincidence measurements which can examine the importance of the nucleon component are well worth study.

8. One wishes to study the emergence of strangeness in the nuclear state.

The fact that QCD is a viable theory for hadronic interactions implies that a fundamental description of the nuclear force is now possible. Although detailed work on the synthesis of QCD and nuclear physics is just beginning, it is clear from the structure of QCD that several traditional concepts of nuclear physics will have to be modified. These include conventional treatments of meson and baryon-pair contributions to the electromagnetic current and analyses of the nuclear form factor in terms of factorized on-shell nucleon form factors. On the other hand, the reduced nuclear form factor and scattering matrix elements discussed in Section 4 give a viable prescription for the extrapolation of nuclear amplitudes to zero nucleon radius. There is thus the possibility that even the low momentum transfer phenomenology of nuclear parameters will be significantly modified.

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FIGURE CAPTIONS

1. (a) Factorization of the nucleon form factor at large Q^2 in QCD. The optimal scale $\tilde{Q}$ for the distribution amplitude $\phi(x, \tilde{Q})$ is discussed in Ref. 8.

   (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 $\phi_B(x, Q)$.

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 (see Ref. 16).

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

4. (a) Comparison of the asymptotic QCD prediction (3.7) with experiment (16) using $F_N(Q^2) = (1 + Q^2/0.71\,GeV^2)^{-2}$. 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.

5. Schematic representation of the deuteron wavefunction in QCD indicating the presence of hidden color 6-quark components at short distances.

6. Comparison of deuteron photodisintegration data[24] with the scaling prediction (4.3) which requires $f^2(\theta_{cm})$ to be independent of energy at large momentum transfer.

7. Critique of the standard nuclear physics approach to the deuteron form factor at large Q^2.

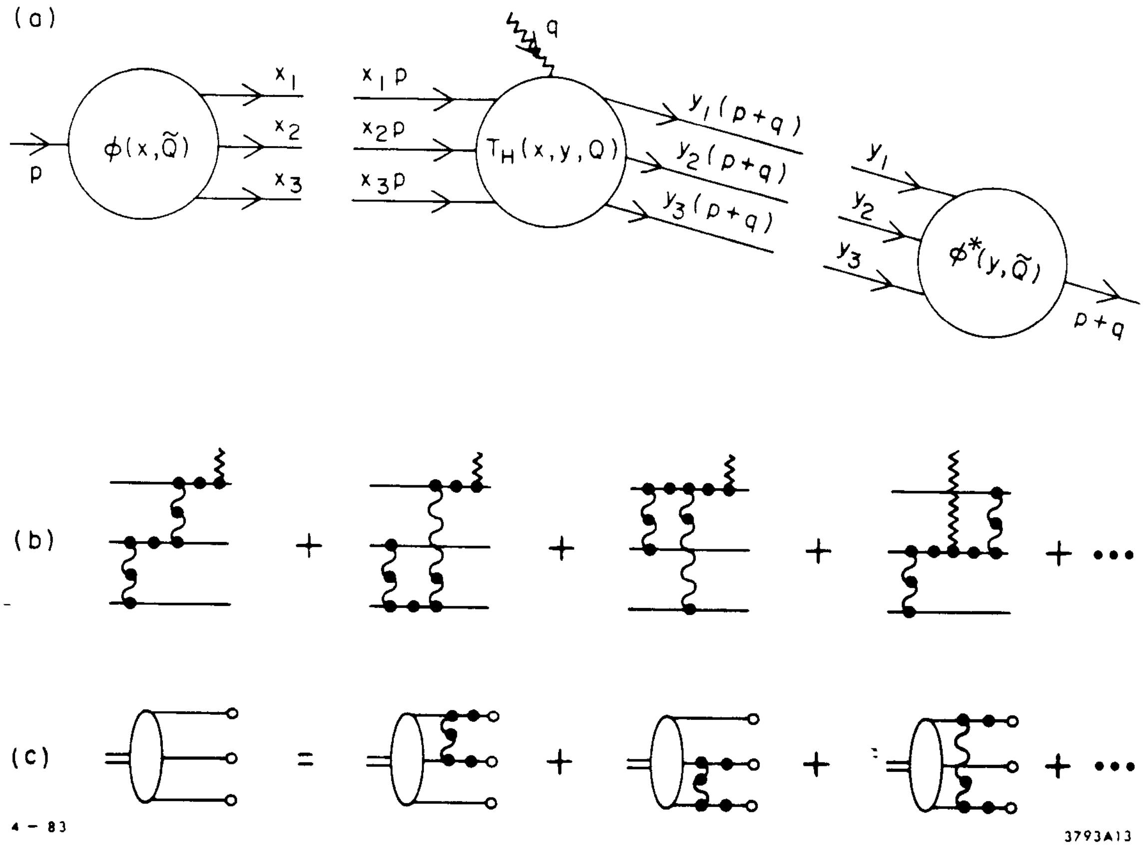
(a) The effective nucleon form factor has one or both legs off shell $|p_1^2 - p_2^2| \sim q^2/2$.

(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.

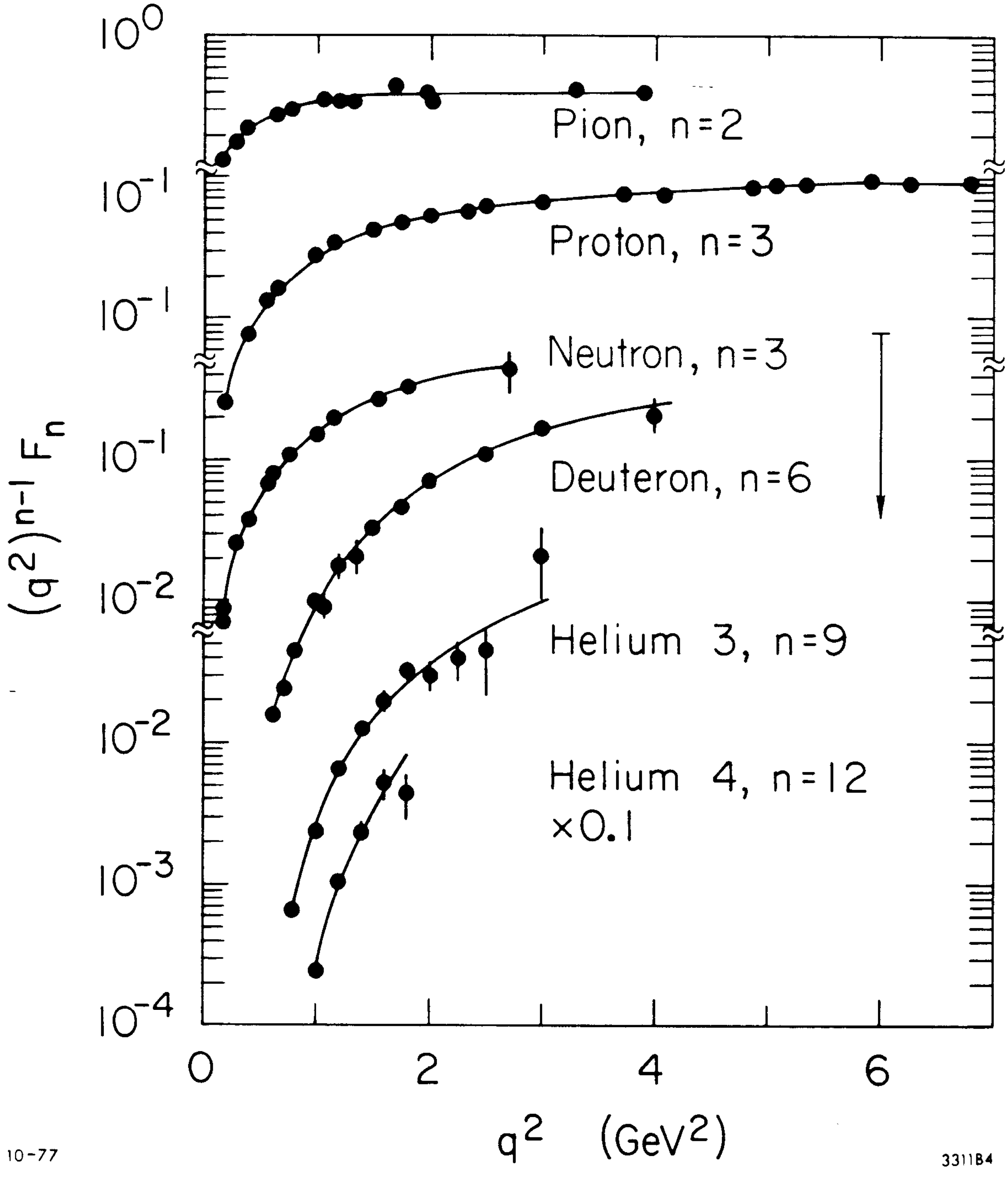
8. Time-ordered contributions to (a) the Compton amplitude in a local Lagrangian theory such as QED. Only the z-graphs contribute in the forward low energy limit.

   (b) Calculation of the Compton amplitude for composite systems.



(a)
x_1
x_2
x_3
p
$\phi(x,\tilde{Q})$
q
$x_1 p$
$x_2 p$
$x_3 p$
$T_H(x,y,Q)$
$y_1(p+q)$
$y_2(p+q)$
$y_3(p+q)$
y_1
y_2
y_3
$\phi^*(y,\tilde{Q})$
p+q
(b)
+ + + + •••
(c)
= + + + •••
4 − 83
3793A13

Fig. 1



10^0
10^{-1}
10^{-1}
10^{-1}
10^{-2}
10^{-2}
10^{-3}
10^{-4}
$(q^2)^{n-1} F_n$
Pion, n=2
Proton, n=3
Neutron, n=3
Deuteron, n=6
Helium 3, n=9
Helium 4, n=12
×0.1
0
2
4
6
q^2 (GeV^2)
10-77
3311B4

Fig. 2

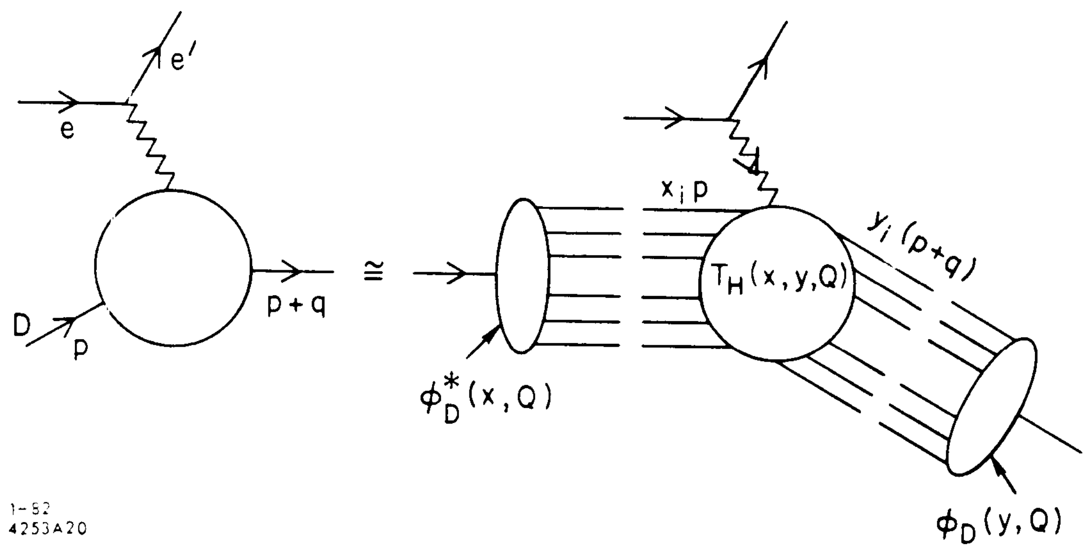


Fig. 3

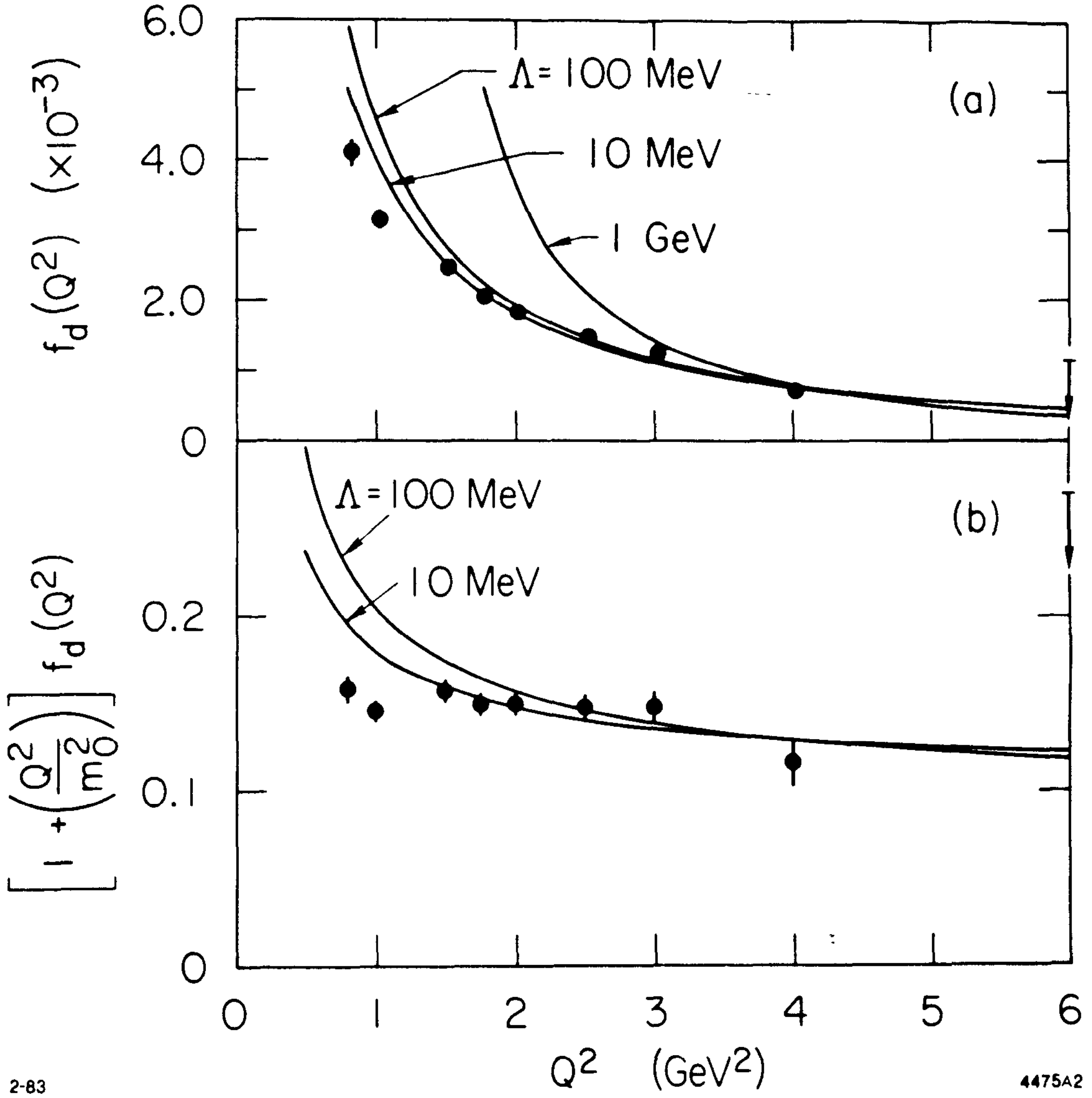

6.0
4.0
2.0
0
$f_d(Q^2)$ $(\times 10^{-3})$
Λ= 100 MeV
10 MeV
1 GeV
(a)
Λ=100 MeV
10 MeV
(b)
0.2
0.1
0
$\left[1+\left(\frac{Q^2}{m_0^2}\right)\right] f_d(Q^2)$
0
1
2
3
4
5
6
Q^2 (GeV^2)
2-83
4475A2

Fig. 4

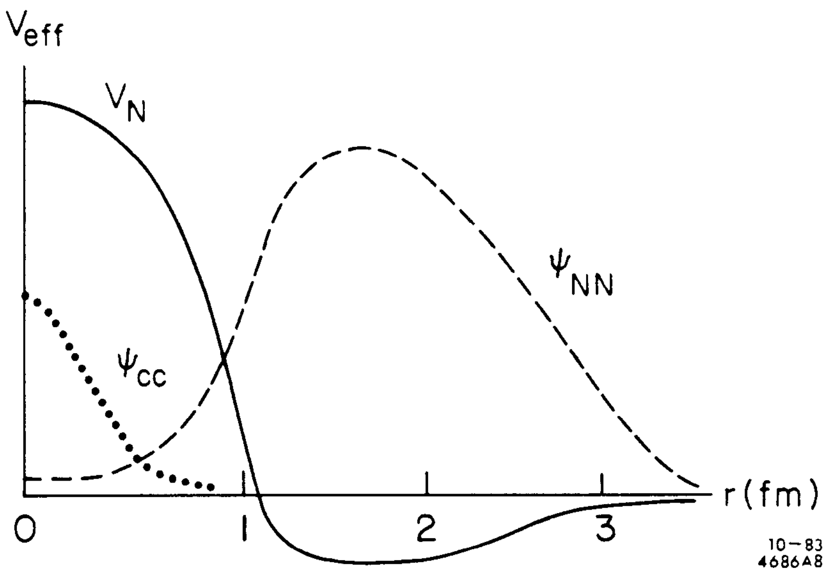


Fig. 5

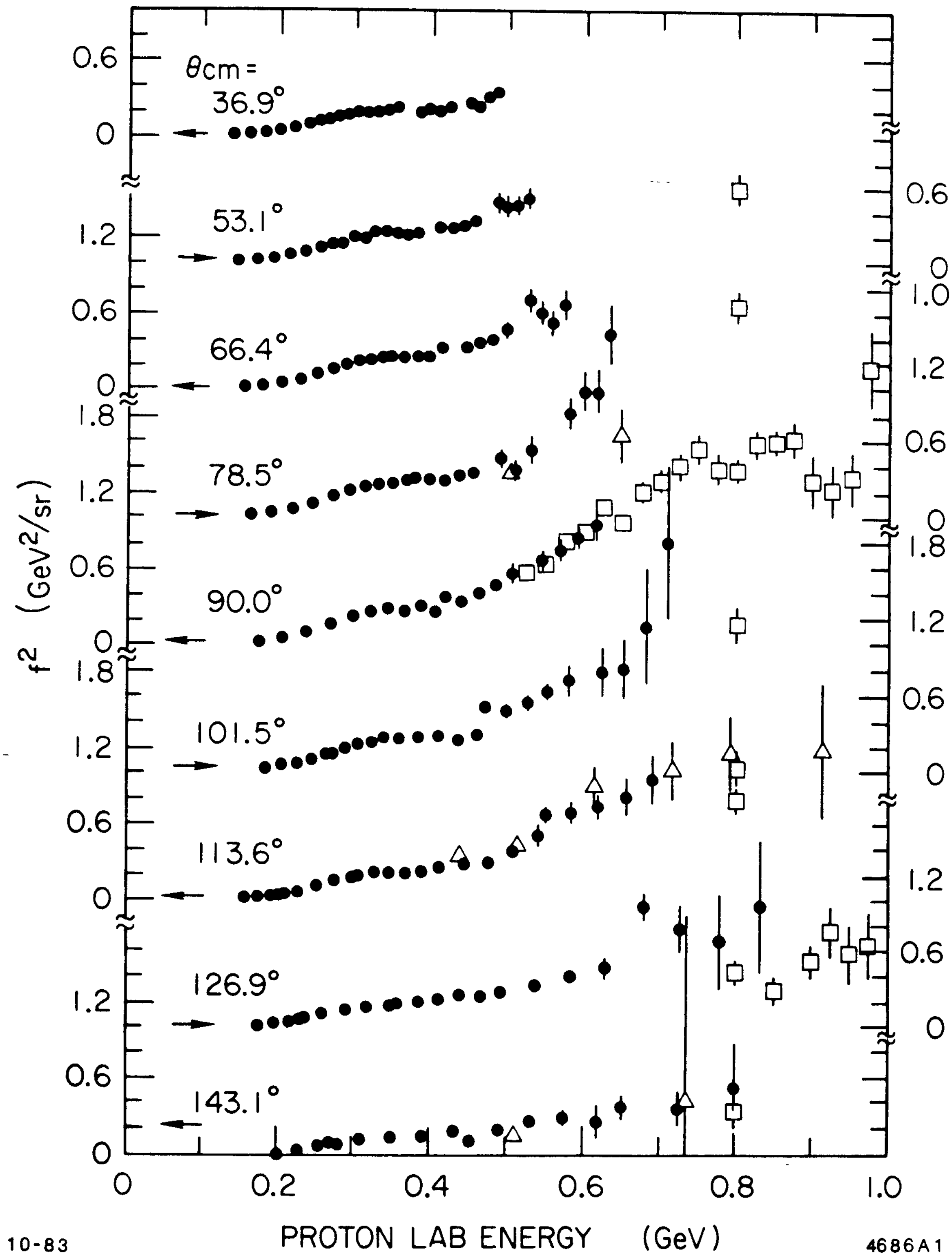

θcm =
36.9°
53.1°
66.4°
78.5°
90.0°
101.5°
113.6°
126.9°
143.1°
f² (GeV²/sr)
PROTON LAB ENERGY (GeV)
10-83
4686A1

Fig. 6

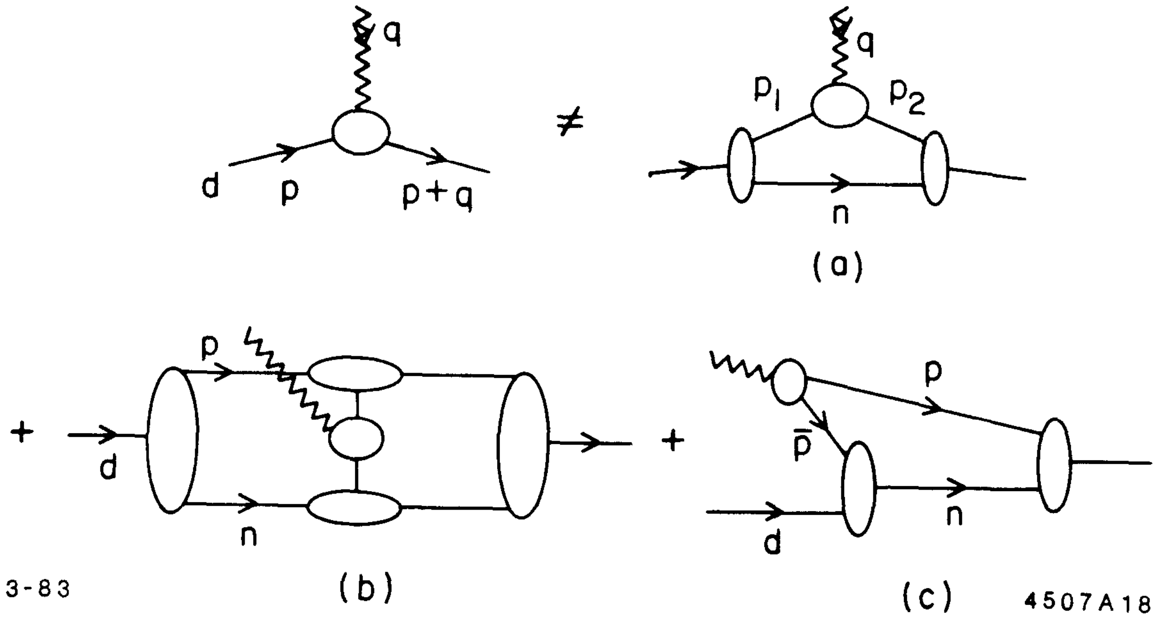


Fig. 7

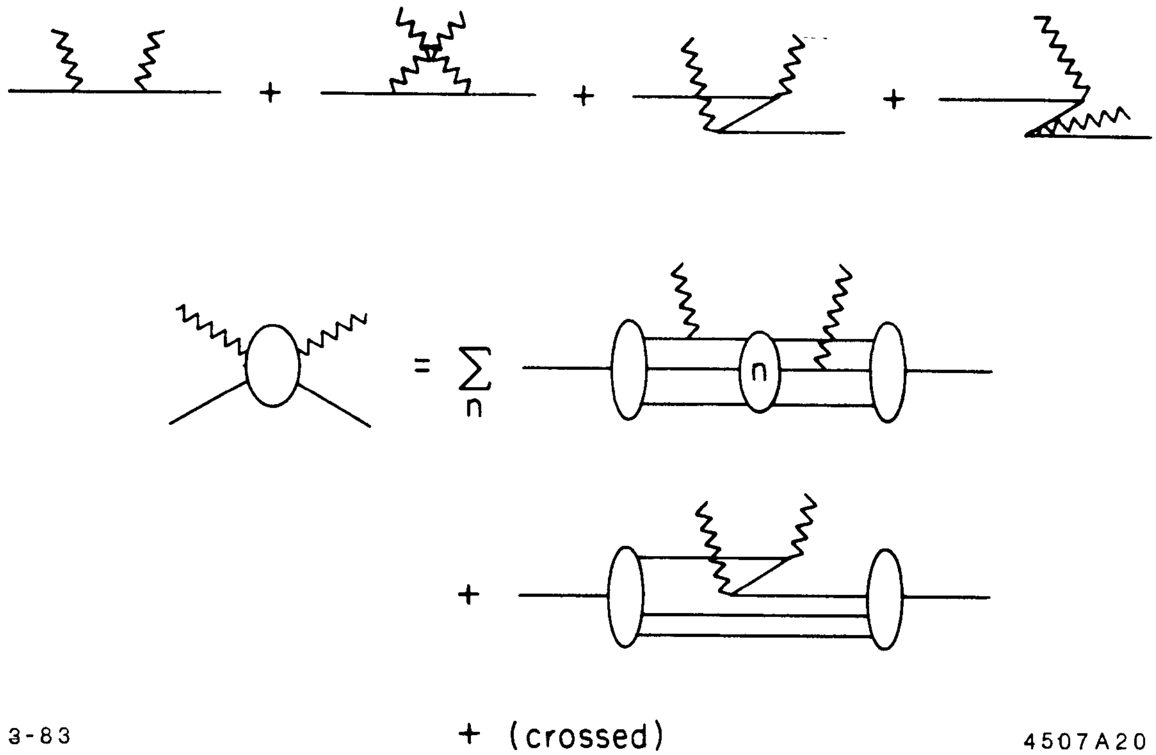


Fig. 8