Quantum chromodynamics and the derivation of a microscopic theory of the nucleus

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The progress which has already been made in the construction of a microscopic theory of the nucleus on the basis of quantum chromodynamics, the problems remaining, and the outlook for future progress are analyzed. The problem of nuclear forces, the role played by a high-momentum component in the nuclear wave function, and the role played by relativistic effects in various hard nuclear processes are discussed.

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1. INTRODUCTION

Although the nucleus has been the subject of research for more than 50 years now, we still lack answers to several fundamental questions, e.g., the origin of nuclear forces and the extent to which nuclei can be described as systems of nucleons. To a large extent the difficulties stem from the need to deal with the internal structure of nucleons, since the distances between nucleons in a nucleus, ~ 1.7 fm, are comparable to the radius of a nucleon, ~ 0.8 fm (Ref. 1). This problem is apparently unsolvable within the framework of the meson theory of nuclear forces, because of the so-called charge-zero problem.[2] These questions are of practical importance because, on the one hand, heavy nuclei are close to a state of collapse[3,4] and, on the other, these questions are pertinent to astrophysics in a realistic theory of superdense nuclear matter. Progress in quantum chromodynamics (QCD) in describing the properties and structure of hadrons[5,6] raises the hope that QCD will lead to answers to basic questions in nuclear physics. However, since it is not possible to avoid the as yet unresolved problem of the confinement of quarks and gluons, the only way at present to analyze phenomena in nuclear physics theoretically is to use the general properties of QCD. This is not, however, a trivial amount of help; quantum chromodynamics will tell us the answers to some of the theoretical problems which arise in the derivation of a theory of the nucleus and of nuclear forces. In particular, the spontaneous breaking of chiral symmetry in quantum chromodynamics (the π-meson-goldstone) is responsible for the dominance of baryon degrees of freedom in the wave function of a nucleus. Quantum chromodynamics makes it easy to understand how to deal with the relativistic motion of nucleons in nuclei, recoil and retardation in the interaction of fast particles with nuclei, etc. The quark-gluon origin of nuclear forces raises some new questions: In which nuclear phenomena do we see the quark and gluon degrees of freedom? What is the spatial scale of fluctuations of the color charge? What roles are played by color screening and by asymptotic freedom?

The nonrelativistic nuclear theory gives values for the basic characteristic of nuclei which are in reasonable agreement with experiment. Quantitative progress in the cases of intermediate and heavy nuclei has resulted from the use of phenomenological approaches where the inadequately studied dynamics of the interaction between nucleons, particularly the short-range nucleon correlations, and a possible admixture of isobars, etc., are "hidden" in the parameters of the effective nucleon-nucleon potential.[4,7,8] On the other hand, attempts to derive a microscopic theory of the nucleus on the basis of the vacuum nucleon-nucleon interaction have run into difficulties. For example, calculations using relativistic nucleon-nucleon potentials underestimate the binding energy of the extremely light nuclei ^{3}He and ^{4}He, and calculations of the elastic and inelastic electromagnetic form factors of ^{3}He and ^{4}He disagree with experiment.[9] One of the most puzzling effects is the existence of a strong repulsion in nuclear forces at internucleon distances $r \lesssim 0.5$ fm.

Until recently it seemed impossible to subject to direct experimental test the basic approximations used in deriving

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a microscopic theory of the nucleus. In particular, it has not been possible to clarify experimentally a possible role of short-range nucleon correlations, three-body internucleon forces, nonnucleon degrees of freedom, etc.—all topics which have been discussed widely by theoreticians. There is a good reason for this situation: In processes which occur at low energies, $\lesssim 100$ MeV (with a small momentum transfer to the nucleons of the nucleus), it is difficult to distinguish effects of short-range nucleon correlations from the background of the more probable processes involving intermediate internucleon distances. As the initial energy is increased, however, processes involving the transfer of momenta ~ 1 GeV/c to individual nucleons of a nucleus become dominant in the inelastic scattering of leptons (or hadrons) by nuclei. It is this circumstance which makes it possible to use high-energy processes to study nuclear structure.[10] The analogy with the history of the detection of quarks (or gluons) in a hadron is somewhat pertinent here. Although a quark structure of hadrons has been suspected on the basis of an analysis of the static characteristics of hadrons, it was only the move to higher energies, to processes with a large momentum transfer, which made it possible to "resolve" the quarks (or gluons) in a hadron. Our purposes in the present review are to outline what has been achieved in research on the nucleus by the theoretical and experimental methods of high-energy physics, to point out what else can be achieved by this approach, and to point out what we can expect in the way of applications to the traditional branches of nuclear physics.

The most direct information about nuclear structure is extracted from an analysis of the total cross sections for deep inelastic scattering of leptons by nuclei, $l + A \to l' + X$, where X is a system of hadrons not detected experimentally. Analysis of the perturbation-theory diagrams in QCD shows that these cross sections can be expressed directly in terms of the momentum distribution of quarks in the ground state of the nucleus (Ref. 11, for example). The European Muon Collaboration (EMC) recently discovered a significant difference between the momentum distributions of quarks in iron and in the deuteron[12] (Fig. 1).[1)] This "EMC effect" has been confirmed[13,14] in electron scattering at $x > 0.2$. It is also consistent with neutrino data.[15,16] At $x < 0.2$, the experimental situation with regard to the EMC effect is less definite.

At $x > 0.2$, the longitudinal distances which are important in the interaction of a γ^* with a nucleus are[17,18] $z \sim 1/m_N x$: much smaller than the typical internucleon distances. At these values of x, therefore, the cross section $\sigma_{\gamma^* A}$ should be equal to the sum of the cross sections for the scattering of γ^* by the individual nucleons of the nucleus (see the discussion in Subsection 2.3). Nevertheless, the observed quark distribution is qualitatively at odds with the predictions of the nonrelativistic theory of the nucleus and thus constitutes *unambiguous* evidence for the presence of significant nonnucleon components in the wave function of the nucleus. It is reasonable to divide the various hypotheses regarding the nature of the EMC effect at $x > 0.2$ into three groups, as follows.

1. In QCD it is expected that the structure function of a nucleon, $F_{2N}(x,Q^2)$, will be dominated in the limit $x \to 1$ by a component of the nucleon wave function $N - |3q\rangle$ which

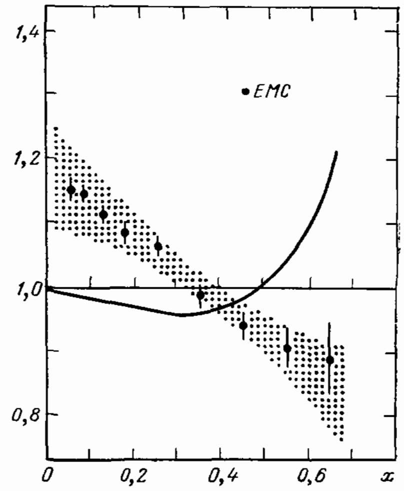

FIG. 1. Ratio of the structure functions of iron and the deuteron.[12] $R = (2/A)F_{2Fe}(x,Q^2)/F_{2D}(x,Q^2)$ Circles—Systematic uncertainties in the dependence of R on x; solid line—expectations of the Fermi-motion models.

does not contain a meson cloud and is accordingly of small radius. Analysis of the elastic and inelastic form factors of the nucleon indicates that this component is dominant in $F_{2N}(x,Q^2)$ at $x \geqslant 0.5$ (Section 2). Since the interaction of the quark-gluon components of small radius in a nucleon with other nucleons is suppressed, the nuclear medium polarizes the nucleon, suppressing the probability for these configurations in a bound nucleon at an intermediate internucleon distance.[19,20] This interpretation has the EMC effect corresponding to a $\lesssim 5\%$ admixture of resonances ($N^*, \Delta, N\pi, \ldots$) in the wave function of the nucleus and, possibly, a 1–3% increase in the radius of the bound nucleon.[2)]

2. There is a probability $\sim$20–30% that a nucleus will contain nonnucleon degrees of freedom: six-, nine-, twelve-quark bags,[23–38] etc. Jaffe has offered the qualitative suggestion[23] that a quark in a bag of this sort is distributed over a volume greater than that in a nucleon. As a result there is a decrease in its momentum, which leads to a value $R_A = (2/A)F_{2A}(x,Q^2)/F_{2D}(x,Q^2) < 1$ at $x \geqslant 0.3$ and to a value $R_A > 1$ at $x \leqslant 0.2$.

3. Close *et al.* and Nachtmann and Pirner[29] have suggested that the confinement radius in the nucleus is 10–20% greater than that in the free nucleon. In this case the reason for the EMC effect at all values of x is that there is a more intense emission of gluons in the scattering of leptons by nuclei than in scattering by the free nucleon.

On the whole, according to essentially all the interpretations which have been offered for the EMC effect, at $x \geqslant 0.3$ this effect may be thought of as a definite experimental indication of the possibility of a phase transition in a superdense nuclear medium.

In Section 2 it is shown that the dependence of the EMC effect on the atomic number at $x > 0.3$ which has been observed experimentally[14] is in reasonable agreement both with the assumption that a suppression of small-radius configurations is dominant in the bound nucleon and with assumption that the EMC effect is determined by a nontrivial quark structure of short-range binary nucleon correlations.

As x decreases, the important longitudinal distances z become comparable (at $x \sim 0.1$) to the typical internucleon distances in a nucleus, ~ 2 fm. At $x < 0.1$, the cross section $\sigma_{\gamma^* A}$ therefore does not reduce to the sum of the cross sections for scattering by the individual nucleons of the nucleus. As x is reduced, the density of partons increases. Incorporating fluctuations of the parton density leads to a contribution to R_A which is comparable to the observed increase in R_A at $x \lesssim 0.1$. This correction falls off rapidly with increasing x, becoming negligible at $x > 0.2$.

The naive application of QCD perturbation theory indicates that the coalescence of partons belonging to different nucleons of a nucleus[30] leads primarily to an enhancement of the distribution of valence quarks at $x \sim 0.1$ (Ref. 20). However, the balance equations cannot be solved reliably because perturbation-theory methods are inapplicable.

The increase in the meson field in a nucleus in comparison with that for a free nucleon (models of pion condensation in nuclei, etc.) leads to an increase in the quark-antiquark sea in a nucleus).[23,31–33] Experiments on ν or $\bar{\nu}$ scattering by the proton, deuteron, and nuclei,[15] however, have not yet revealed any significant increase in the sea in a nucleus. These data may, as theoretical calculations become more reliable and as the role of nonscaling effects is clarified by experiments, prove to be a serious problem for the hypothesis of meson nature of nuclear forces.

Since essentially all the hypotheses which have been advanced to explain the deficiency of quarks with $0.3 < x < 0.7$ in nuclei assume a significant high-momentum component in the wave function of the nucleus, we need to study the magnitude and structure of this component in order to explain the EMC effect. Several leptonic processes which have been discovered would be difficult (impossible?) to interpret in any way other than to a manifestation of a significant high-momentum component of the wave function of a nucleus. This implication is made in the case of the deuteron by an analysis[34–36] of data on elastic eD scattering at $Q^2 \leqslant 4$ GeV^2 (Ref. 37). A significantly greater high-momentum component has been expected for nuclei,[10,34] as has been confirmed by an analysis[38] of near-threshold $e + A \to e' + X$ reactions at $Q^2 \leqslant 6$ GeV^2 (Refs. 39 and 40) and deep inelastic $\mu + A$ scattering at $x > 1$ is a region kinematically forbidden to μN scattering.[41] The first correlation experiments carried out in the deep inelastic reaction $\bar{\nu} + \mathrm{Ne} \to \mu^+ +$ fast proton backward + X (Ref. 42) indicate that the high-momentum component of the Ne nucleus is dominated by short-range binary and ternary nucleon correlations.[38]

Historically the first regular indications of a significant high-momentum component in the wave function of a nucleus came from limiting fragmentation reactions

$$a + A \to p\,(\pi) + X, \tag{1.1}$$

where a is a hadron, a γ ray, a light nucleus, a ν, or a $\bar{\nu}$, with the p or π being detected in the fragmentation region of the nucleus, which is kinematically forbidden to scattering by hydrogen. A reaction of the type in (1.1) in which protons were detected was first observed in Ref. 43, and a reaction in which pions were detected was first observed in Ref. 44 (these experiments are reviewed in Refs. 45–50). These particles are called "cumulative" after Baldin.[51] The complexity of "soft" hadronic processes resulted in the appearance of several hypotheses regarding the mechanism responsible for reaction (1.1): few-nucleon (binary, ternary) correlations,[52–54] multiquark bags,[34,55] a fluctuon model,[51,45,56–59] and a mean field model.[60–63] Reaction (1.1) was interpreted simultaneously in several studies as the result of a final-state interaction (Refs. 64–67, for example).

Comparison of the characteristic properties of reaction (1.1) with those of the leptonic processes listed above reveals that these two types of reactions are related.[38] [A crucial role is played in the interpretation of hadronic reactions (1.1) by the theoretical observation that the Glauber screenings which substantially reduce the total cross sections for hadronic processes at nuclei ($\sigma_{hA} \sim A^{2/3}$) cancel out in the inclusive spectrum.[10]] We thus have a significant set of different high-energy leptonic and hadronic processes which indicate that in nuclei with $A \geqslant 12$ only 70–80% of the nucleons are below the Fermi surface[3)] and that the high-momentum component of the nuclear wave function is determined primarily by short-range correlations, which are apparently dominated by nucleon degrees of freedom.[38,68]

Our interest is thus attracted by recent attempts to derive a theoretical description of short-range nuclear forces in the spirit of QCD by the methods which have proved successful in describing the statistical properties of hadrons: various versions of the bag model[36,69–73] and the potential quark models (Refs. 73–75, for example). A distinctive feature of nuclear physics is the need to take into account the "energy advantage" to the multiquark system of the formation of "white" subsystems.[77,72]

This review is organized as follows: Section 2 deals with the analysis of the EMC effect. Section 3 discusses new approaches to the description of nuclear forces and the possibilities of a search for exotic degrees of freedom in nuclei. Section 4 reviews the results on short-range nucleon correlations in nuclei in high-energy leptonic and hadronic processes. The Appendix explains the distinctive features of high-energy processes which make it necessary to use wave functions on the light cone in order to describe the nucleus. These wave functions can be expressed directly in terms of the wave functions of nonrelativistic nuclear theory over a broad kinematic region.[38]

2. THE EMC EFFECT AND NUCLEAR STRUCTURE

The information which has been acquired in recent years on the microscopic structure of the nucleus by the methods of high-energy physics cannot be considered in isolation, without reference to our previous experience in nuclear physics. Before we take up the high-energy processes, we will therefore briefly summarize the results which have been obtained on nuclear structure in nonrelativistic nuclear theory.

a) Resolved and unresolved problems in nonrelativistic nuclear theory. The success of nonrelativistic nuclear theory in describing the basic properties of nuclei shows that the wave functions of nuclei are dominated by nucleon degrees of freedom.

1) The nonrelativistic theory gives a reasonable description of the basic characteristics of the deuteron: its magnetic moment (within $\sim 1\%$), its electromagnetic form factors up to $Q^2 \sim 1$ GeV^2 (Ref. 78), etc. (We wish to stress that in momentum space the realistic wave functions of the deuteron—the Reid, Paris-potential, Hamada-Johnston, etc., wave functions—differ seriously only at $k \gtrsim 0.6$–0.8 GeV/c, where the nonrelativistic approximation breaks down.)

2) The independent-particle model and the shell model work reasonably well in processes in which momentum transfers $\lesssim 0.2$ GeV/c are important (Ref. 79, for example).

3) In the case of heavy nuclei, approaches based on a phenomenological binary interaction between nucleons, also incorporating some many-particle correlations, have proved useful (see the reviews in Refs. 4, 7, and 8). Comparison of the seemingly different effective potentials used for these purposes reveals that at the momenta which are important, $\lesssim 200$ MeV/c, they are approximately the same in terms of their basic parameters.[80]

4) The data on elastic proton-nucleus scattering at $T_p \lesssim 1$ GeV agree within $\sim 2\%$ with the standard Glauber model with the vacuum value of the pN scattering cross section.[81]

5) Analysis of the quasielastic (e,e′) reactions at ^{12}C(Fe) shows that only $\sim 75\%$ (60%) of the nucleons are below the Fermi surface.[82]

6) Analysis of experiments on the photodisintegration of nuclei (Ref. 83, for example) and on the absorption of slow pions by nuclei (Ref. 84, for example) suggests that at $A \geqslant 12$ a significant number of nucleons are above the Fermi surface,

$$\int n_A(k)\, \theta(k - k_F)\, d^3k \sim 20\text{—}40\ \%,$$

and basically belong to binary nucleon correlations.

7) Calculations based on realistic nucleon-nucleon potentials underestimate the binding energy of the very light nuclei ^{3}He and ^{4}He (the discrepancy in the case of the Reid potential is 20%; Ref. 9).

8) It has not been possible to obtain correct values of the binding energy per nucleon for ^{16}O or of the density of the nuclear medium in the case of infinite nuclear matter.[85] The latter difficulty may stem from the neglect of three-body forces.[85]

9) We have no explanation for the origin of nuclear forces, in particular, the nature of the nuclear core.

b) The EMC effect. The region $x > 0.3$. The most direct way to study nuclear structure is the deep inelastic scattering of leptons by nuclei in the Bjorken limit, where $q^2 \equiv -Q^2$ [the square of the momentum transferred to the hadronic system (the square of the "mass" of the virtual photon)] and the mass (W^2) of the hadronic system which is produced are large, and

$$x = \frac{Q^2}{2m_N q_0} \tag{2.1}$$

is fixed. (All the notation corresponds to Fig. 2). It can be shown in QCD that a fixed value of x in the limit $Q^2 \to \infty$ the

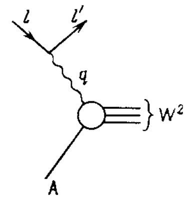

FIG. 2.

scattering by the individual quarks of the nucleus becomes dominant (Ref, 11, for example). In the deep inelastic scattering of leptons by nuclei, we would thus be directly measuring the momentum distribution of the quarks (antiquarks) in the nucleus.

The European Muon Collaboration recently measured the total cross section for the deep inelastic scattering of muons by Fe and D,

$$\mu + \mathrm{Fe}\ (\mathrm{D}) \to \mu + X, \tag{2.2}$$

in the range 9 GeV$^2 < Q^2 < 170$ GeV^2 at a fixed value of the Bjorken variable x [see (2.1)], and discovered a significant difference[4)] in the x distributions of the quarks in Fe and D (Fig. 1).[12] The existence of the effect at $x \gtrsim 0.3$ was soon confirmed in experiments on electron scattering.[13,14] Neutrino data are also consistent with an effect at $x > 0.3$ (Refs. 15 and 16). The wide-spread interest in this discovery is due to the qualitative discrepancy between the observations and the expectations of nonrelativistic nuclear theory, which has nuclei consisting exclusively of nucleons [cf. the discussion following Eq. (2.7)].

Reaction (2.2) is usually described by expanding the amplitude for the process $\gamma^* A \to \gamma^* A$ in inelastic form factors:

$$\mathrm{Im}\, A_{\mu\nu}^{\gamma^*+A \to \gamma^*+A} = \frac{F_{2A}(x, Q^2)}{(p_A q)} \left(p_{A\mu} - q_\mu \frac{(p_A q)}{q^2}\right) \left(p_{A\nu} - q_\nu \frac{(p_A q)}{q^2}\right) - F_{1A}(x, Q^2) \left(g_{\mu\nu} - \frac{q_\mu q_\nu}{q^2}\right). \tag{2.3}$$

The cross section for the inclusive process $\mu + A \to \mu' + X$ (i.e., a process in which no particles other than the lepton are detected) is expressed directly in terms of these form factors:

$$\frac{d\sigma}{d\Omega\, dE'} = \frac{\alpha^2}{4E^2 \sin^4(\theta/2)} \left(\frac{F_{2A}(x, Q^2)}{(p_A q)} \cos^2 \frac{\theta}{2} + \frac{2F_{1A}(x, Q^2)}{M_A} \sin^2 \frac{\theta}{2}\right). \tag{2.4}$$

Here E is the energy of the initial lepton, and θ is the scattering angle in the rest frame of the nucleus. The most detailed data are available on $F_{2A}(x, Q^2)$.

We consider first the kinematic region $x > 0.2$, where effects of large longitudinal distances are small in the interaction of γ^* with the nucleus. In this region, if the nucleus consists exclusively of nucleons, $\sigma_{\gamma^* A}$ reduces to the sum of the cross sections for the scattering by the individual nucleons of the nucleus (cf. the discussion in Section 2c):

$$F_{2A}(x,\ Q^2)=\int \rho_A^N(\alpha,\ k_t)\, F_{2N}\left(\frac{x}{\alpha},\ Q^2\right)\frac{d\alpha}{\alpha}\, d^2k_t; \qquad (2.5)$$

here $\rho_A^N(\alpha,k_t)$ is the density matrix of the nucleons in the nucleus on the light cone. In the nonrelativistic approximation, α should be replaced by[38]

$$\alpha=1+\frac{k_3}{m_N}, \qquad (2.6)$$

and $(1/Am_N)\rho_A^N(\alpha,k_t)$ should be replaced by $n_A^N(k)$—the nonrelativistic density matrix of the nucleons with momentum k in the nucleus. The need to use the wave functions on the light cone to describe high-energy processes is explained in the Appendix, as is the relationship between these wave functions and nonrelativistic nuclear theory. Here $F_{2N}(x,Q^2)=(1/2)(F_{2p}(x,Q^2)+F_{2n}(x,Q^2)$ is the structure function of the nucleon. At $x\ll 1$, it is useful to expand $F_{2N}(x/\alpha)$ in a power series in $(1-\alpha)$. The sum rules for the total angular momentum of the nucleus[38] cause the term linear in $(1-\alpha)$ to cancel out in the summation of the contributions from scattering by all the nucleons of the nucleus. Accordingly, for $F_{2N}(x)=c(1-x)^3$ (a reasonable approximation of the experimental data at large values of W^2) the structure function of the nucleus at $x\ll 1$ becomes[5)]

$$\frac{1}{A}F_{2A}(x,\ Q^2)$$
$$=F_{2N}(x,\ Q^2)\left[1+\frac{2x}{(1-x)^2}\left(x-\frac{1}{2}\right)\langle k_A^2\rangle\, m_N^{-2}\right]; \qquad (2.7)$$

here

$$\langle k_A^2\rangle=\int k^2 n_A^N(k)\, d^3k.$$

Since $\langle k_A^2\rangle > \langle k_D^2\rangle$, we would have $R_A=(2/A)F_{2A}F_{2D}<1$ at $x<0.5$ and $R_A>1$ at $x>0.5$, in contradiction of experiment (Fig. 1),[6)] regardless of the form of the nuclear wave function. The EMC effect at $x>0.3$ is thus unambiguous evidence for the presence of a nonnucleon component in the wave function of the nucleus.[7)] We wish to stress that the equations of the renormalization group can be used to reconstruct the quark distribution in a nucleus at low values of Q^2, comparable to the typical values of hard nuclear processes, from the measured quark distribution at large values of Q^2.

Essentially all the hypotheses which have been offered as explanations for the decrease in $R_A(x,Q^2)$ at $x\sim 0.5$ assume that there is, with a probability $\sim$20–30%, a high-energy component in the iron nucleus due to either (a) a high-momentum $(k>k_F)$ component with a small nonnucleon admixture[19,20] or (b) exotic states: multiquark bags, isobars, etc.[8)] (Refs. 23–28 and 89).

Jaffe[23] was the first to offer a hypothesis to explain the EMC effect; his hypothesis was followed immediately by many others (Refs. 24–28, for example). Jaffe suggested that when nucleons approach to within small distances of each other a common 6q,9q,..., bag forms. (The same hypothesis had been advanced earlier to explain the observed production of cumulative particles and the elastic form factor of D; Refs. 34–36 and 55–59.) As a result, the quarks move in a volume larger than that in a nucleon, so that (by virtue of the uncertainty principle) the quark distribution in a 6q bag is softer than that in a nucleon. As an illustration Jaffe suggested the following expression for the structure function of a six-quark bag of radius of R_{6q}:

$$F_{26q}(x)=2F_{2N}\left(\frac{R_{6q}}{R_N}x\right). \qquad (2.8)$$

Equation (2.8) artificially cuts off the region $x>0.8$ and contradicts the data of Refs. 13, 14, and 41. The functional dependence of $R_{26q}(x)$ on x is thus frequently described[24] by an estimate from QCD perturbation theory[90] or the equations of a dimensionality calculation[91]: $F_{26q}(x)\sim[1-(x/2)]^{10}$ but $F_{2N}(x)\sim(1-x)^3$. We would thus have

$$R_A(x)-1\approx c(A)\left(1-\frac{x}{2}\right)^{10}(1-x)^{-3}. \qquad (2.8')$$

This expression can be generalized to the cases of nine- and twelve-quark systems (Ref. 24, for example). Since the derivation of the quark continuing formulas leans heavily on the boundedness of the phase volume in the limit $x\to 1(x\to 2)$, we do not know of any theoretical basis for the validity of (2.8'); even the sign of $R_A(0.5–0.6)-1$ is not obvious.[92] It is easy to show that (2.8') has a minimum at $x\sim 0.5$. At $x\geqslant 0.85$, $R_A(x)$ increases rapidly with x. A fit of the data of Ref. 12 leads to probability $P_{6q}\sim 20\%$ for a 6q component in iron.

These models are particular cases of the hypothesis that the MEC effect results from the quark structure of short-range nucleon correlations. We can thus estimate the dependence of the EMC on the atomic number. On the basis of the standard nuclear theory we would expect[93,94] that the probability for correlations of j nucleons would be proportional to

$$c_j(A)\sim\int \rho_A^j(r)\, d^3r\sim A^{0,15},\ A^{0,22},\ A^{0,27}$$
$$\text{for}\quad j=2,\ 3,\ 4\quad \text{at}\quad A\geqslant 12, \qquad (2.9)$$

where $\rho_A(r)$ is the density of nucleons in the nucleus.[38] The A dependence of $c_j(A)$ at $A<12$ can be determined from near-threshold (e,e') reactions, from the production of cumulative particles (Section 4). Figure 4 shows the expected dependence of R_A-1 on A for the cases of binary and ternary correlations; this dependence is consistent with the data of Ref. 14. In contrast, the suggestion[26] that the EMC effect is dominated by a scattering by α clusters in the nuclei contradicts experiment, since it predicts that $1-R_A$ will reach a maximum in scattering by ^{4}He, ^{12}C, and ^{16}O (Ref. 26).

The assumption—typical of these models—that nonnucleon degrees of freedom are dominant in the wave functions of short-range correlations corresponds to a significant exotic component in the wave function of the nucleus (per nucleon)[9)]:

$$P_{ex}\geqslant[1-R_A(x)]\left(1-\frac{F_{2ex}(x)}{F_{2N}(x)}\right)^{-1}\sim 20-30\%|_{Fe},$$
$$30-40\%|_{Pb},$$

where $F_{2ex}(x)$ is the structure function of the exotic component. It would seem difficult to reconcile the introduction of such a large exotic component in a nuclear wave function with the succesful phenomenology of nonrelativistic nuclear physics or with the physics of hard nuclear processes (Sections 3 and 4).

An alternative hypothesis[19,20] reduces to the assumption that the EMC effect is caused by the suppression of a quark component [rare but dominant in $F_{2N}(x,Q^2)$ at $x>0.5$] in the wave function of a bound nucleon at typical internucleon distances in a nucleus. The scale of the suppression expected on this basis can be estimated by noting that the magnitude of the correction to the description of a nucleus as a system of nucleons with fixed structure can be characterized by the ratio of the typical frequencies in the nucleus and in the nucleon:

$$\varkappa = \left|\frac{U}{\Delta E}\right| \sim \frac{1}{10}; \tag{2.10}$$

here U is the average value of the potential for a nucleon in a nucleus ($U \approx 40$ MeV), and $\Delta E \approx M_{N^*,\Delta} - M_N \sim 0.3$–$0.5$ GeV is a typical excitation energy of a nucleon.

A manifestation of QCD is the existence in hadrons of quark-gluon components of different spatial sizes, as a consequence of a spontaneously broken chiral symmetry in QCD. The pion cloud around a nucleon has the largest radius (which becomes infinite in the limit of zero quark masses). As a result, we find a significant difference between the electromagnetic and axial radii of the nucleon, $r^2_{ax}/r^2_{em} \approx 0.5$, since the axial current does not interact with the π (because the latter has a zero spin). Furthermore, the idea of weakly interacting compressed configurations in ordinary hadrons has been raised in QCD perturbation theory.[95] It has been found that the cross section for the interaction of "white" particles falls off with decreasing radius (r) of the spatial region occupied by the color quarks in the hadron, in proportion to r^2 in the limit $r\to 0$. For a quantitative description of the EMC effect it is sufficient that the radius of the configurations in the nucleon which contain a quark with $x>0.5$ be significantly smaller than the average radius of a nucleon:

$$\frac{r^2_N(x_q>0{,}5)}{r^2_N} \leqslant \frac{1}{3}. \tag{2.11}$$

A fit of the elastic and inelastic form factors of the nucleon[96] indicates that the $|3q\rangle$ component in the nonperturbative wave function of the nucleon is present in the nucleon with a probability $P\sim 0.1$–0.2 and is dominant in $F_{2N}(x,a^2)$ at[10] $x \gtrsim 0.5$. The radius of this component turns out to be small: $r^2_{3q}/r^2_N \sim 0.1$. The presence of compressed configurations in hadrons should also be manifested in the low-energy interactions of hadrons. For example, if nuclear forces are the result of the exchange of mesons M (M $= \pi,\omega,\ldots$), then arguments like those ordinarily used to calculate hadron form factors[90,91] lead us to conclude[20] that in the case of a "compressed" nucleon ($r_N \to 0$) the MNN vertex would fall of as[11] r_N^2. In fact, as the nucleon is compressed ($r_N\to 0$), the characteristic momentum of the quarks in the nucleon, $k\sim r_N^{-1}$, increases, and the additional gluon propagator (in comparison with that in the case of a point meson) experiences asymptotic conditions (cf. Fig. 3). The internucleon interaction potential should therefore fall off in the limit $r_N\to 0$ in accordance with [12]

$$V_{NN}(r_N,\ R) \sim r_N^2 \tilde{V}(R); \tag{2.12}$$

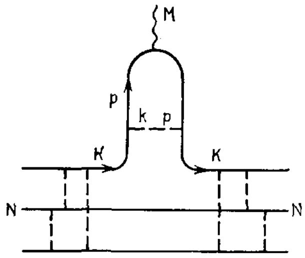

FIG. 3.

here R is the distance between nucleons.

Since the interaction between nucleons takes different forms for different quark-gluon configurations in the wave function of the nucleon, the presence of the nuclear medium deforms the wave function of a nucleon bound in a nucleus. In the case of compressed configurations, the ratio of the probabilities for compressed configurations in the bound and free nucleons can be expressed in terms of $U(R_{ij})$—the ordinary internucleon potential—and $\Delta E \sim 0.3$–0.5 GeV—the typical energy of the compressed configurations[19]—as follows:

$$\delta = 1 + 2\sum_{j(j\neq i)} U(R_{ij})(\Delta E)^{-1}. \tag{2.13}$$

The suppression factor in the case of a nuclear wave function in momentum space is given in the self-consistent-field approximation by[20]

$$\delta(k) = 1 - 4\left(\frac{k^2}{2m_N} + \varepsilon_A\right)(\Delta E)^{-1} \quad \text{for} \quad k < 0.4 - 0.5\ \text{GeV}/c, \tag{2.14}$$

where ε_A is the binding energy of the nucleon. In the approximation of binary nucleon correlations, $\delta(k)$ would be

$$\delta(k) = 1 - \frac{4k^2}{2m_N\Delta\tilde{E}}, \tag{2.15}$$

where $\Delta\tilde{E}$ is a typical excitation energy for the binary nucleon correlations. Since the binary correlations are dominated by correlations with the quantum numbers of the deuteron, for which transitions "D"$\to$NΔ are forbidden, $\Delta\tilde{E}$ should be significantly larger than ΔE ($\Delta\tilde{E}\sim 1.5$–$2\,\Delta E$). As a result, (2.14) effectively applies to binary correlations also. Because of the way in which it was derived, expression (2.14) holds only at $|\delta(k)-1| \ll 1$, i.e., outside the nuclear core.

The nature of the deformation of the wave function of a bound nucleon in the nuclear core can be predicted on the basis of a variational principle. Since an increase in the probability for pointlike configurations leads to an increase in the binding energy in this region, they are intensified; i.e., we have $\delta > 1$.

Accordingly, at $x \gtrsim 0.5$, where the effects of the Fermi motion are not yet large, and where scattering by compressed configurations is dominant, we find the estimate

$$R_A(x) \equiv \frac{2}{A} F_{2A}(x,\ Q^2)\, F_{2D}^{-1}(x,\ Q^2)\Big|_{x\sim 0{,}5} \approx 1 + \frac{4\tilde{U}_A}{\Delta E} \sim 0.75 - 0.85; \tag{2.16}$$

here we have used $\overline{U}_A = 40$ MeV and[13] $\Delta E = 0.5$–0.8 GeV.

This estimate agrees reasonably well with the observed value of R_A. A typical property of (2.7) and (2.16) is the factorization of the dependence of $R_A - 1$ on x and on A at $x < 0.7$, where the contribution of nucleon correlations amounts to a correction. This prediction[20] agrees reasonably well with experiment.[14] By way of contrast, we note that R_A itself depends in a complicated way on x and A (Fig. 4). Furthermore, expression (2.16), which uses measurements of $\overline{U}_A$ at various nuclei,[104] and a calculation[105] of $\overline{U}_{^4\mathrm{He}}$ lead to an A dependence of $R_A - 1$ which agrees well with recent measurements at the Stanford Linear Accelerator Center[14] (Fig. 5).

To incorporate the deformation of nucleons and the Fermi motion simultaneously for $F_{2A}(x,Q^2)$ we should replace $F_{2N}(x,Q^2)$ by $F_{2N}^{\mathrm{eff}}(x,Q^2)$—the structure function of the deformed nucleon—in the impulse approximation for F_{2A}. As a result we find

$$F_{2A}(x,\ Q)^2 = \int F_{2N}^{\mathrm{eff}}\left(\frac{x}{\alpha},\ Q^2\right) \rho_A^N(\alpha,\ p_t)\ \frac{d\alpha}{\alpha}\ d^2p_t; \tag{2.17}$$

here $x = -Aq^2/2q_0m_A$ and ρ_A^N is the density matrix of the nucleons in the nucleus.[94] To estimate $\eta(k,x) \equiv F_{2N}^{\mathrm{eff}}(k,x,Q^2)/F_{2N}(x,Q^2)$ we can use the arguments above, according to which we would have $\eta(k,x) \approx \delta(k)(\approx 1)$ at $x \geqslant 0.6$ $(x \leqslant 0.3)$. Equation (2.17) with $\delta(k)$ from (2.15), $\Delta E \sim (0.5\text{–}0.8)$ GeV, and ρ_A^N from Ref. 38 gives a satisfactory description of the data of Ref. 14 (Fig. 6) at $x > 0.35$.

The suppression of compressed configurations should evidently be manifested in scattering by the deuteron also. Recent data[106] on the reaction $e + D \to e + p + n$ seem to indicate such an effect (Subsection 4.1).

These effects have little influence on the average characteristics of the nucleons in a nucleus, since the poarizability of a nucleon basically reduces to a change in the probability for compressed configurations from $P \sim 0.1\text{–}0.2$ to $P\langle\delta\rangle$. As a result, the radius of a bound nucleon in a nucleus increases 1–3% because of an increase in the probability for intermediate configurations from $1 - P$ to $1 - P + (1 - \langle\delta\rangle)P$ (Ref. 20). A similar estimate results from correcting the bag model for a decrease in the interaction with decreasing radius of the nucleon bag[20]:

$$\frac{r_{\text{bound}}}{r_{\text{free}}} \approx 1 - \frac{2}{3}\ \frac{\overline{U}}{M_N} \approx 1.03. \tag{2.18}$$

This estimate of $r_{\text{bound}}/r_{\text{free}}$ is based on a specific assumption regarding the dynamics of the interaction of nucleons in an average quark-gluon configuration, and it is model-dependent. For example, in the model of chiral bags the incorporation of an increase in the pion field in the nucleus leads to a compression of the bag in comparison with the case of a free nucleon[69,70] It is nevertheless interesting to note that the success of the equations of the Glauber approximation in describing the total cross section for the scattering of hadrons by nuclei (see the review by Alkhazov *et al.*[81]) does not contradict a significant level of compressed configurations in either the incident hadron[100] or the nucleon bound in the nucleus[20] [cf. (2.18)].

The admixture of nonnucleon components, $P_{\Delta,N^*,N\pi}$, due to the supression of compressed configurations in a bound nucleon also turns out to be small: $4U_A P_{CK} \Delta E \sim 2$–4%. The suppression of compressed configurations in a bound nucleon also gives rise to a small correction to the binding energy of a nucleon in a nucleus,[20] $\Delta\varepsilon \sim 1$ MeV. The deformation of the average configurations in a nucleon may cause a comparable increase in the binding energy, but the corresponding calculations have so far run into difficulties.[14)]

Although the expansion of nucleons leads to only small corrections in the phenomena which are usually discussed in nuclear physics,[15)] the attraction between nucleons intensifies with increasing density of the nuclear matter. As a result, there are increases in the radius of a bound nucleon, (2.18), and thus the relative weight of the excited states in the

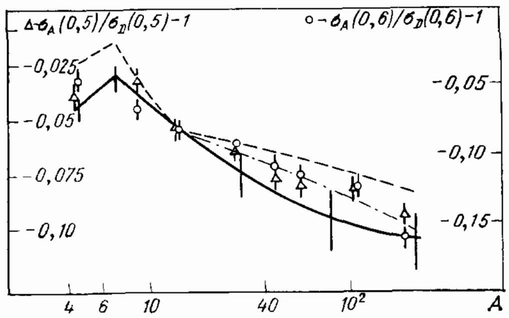

FIG. 5. Comparison of the data of Ref. 14 for $(2/A)\sigma_A(x)/\sigma_D(x) - 1$ with predictions of the mechanism involving a suppression of pointlike configurations (solid line; the error bars on this curve reflect the uncertainties in the value of $\overline{U}_A$) and mechanisms involving the deformation of two-nucleon and three-nucleon correlations (dashed and dot-dashed lines).

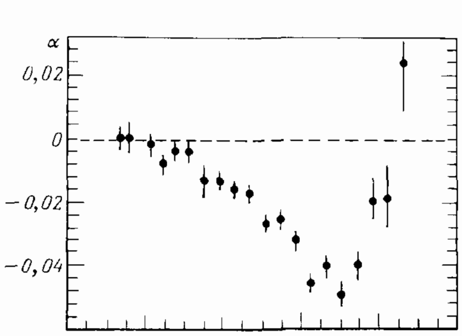

FIG. 4. Dependence of $(2/A)\sigma_A(x)/\sigma_D(x)$ on A parametrized in the form $c(x)A^{\alpha(x)}$ (Ref. 14).

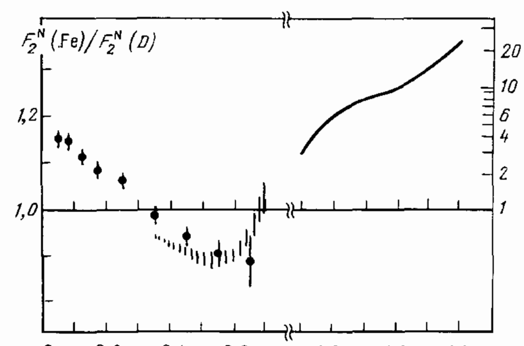

FIG. 6. Comparison of the data of Ref. 14 with calculations from Eq. (2.17).

nucleon. In other words, the local structure of the nuclear matter changes. We see an analogy between the properties of superdense nuclear matter and those of a nonideal exciton gas, where the internal structure of the excitons leads to similar changes in the properties of the excitons.[21] This analogy suggests that the EMC effect might be thought of as indicating a possible phase transition in superdense nuclear matter, analogous to the phase transition from an exciton gas to an electron-hole liquid.[21,22] However, the excitation energy of the nucleon, $m_{N^*} - m_N \gtrsim 0.3$ GeV, is large, so that in a relativistic situation this effect amounts to only a correction. With increasing density of the nuclear matter, the probability for the exchange of quarks and gluons between nucleons increases. As a result, the properties of the nuclear matter come to be determined by the nature of the excitations in the two nucleon system (dibaryon resonances?). The quantitative analysis of these questions is still in a primitive stage.

It was mentioned in Ref. 29 that the EMC effect can be described in a semiquantitative way, at all values of x, as a result of a difference between the radiative effects in QCD at a nucleus and at a nucleon if it is assumed that the color screening radius in a nucleon is 10–20% greater than that at a nucleon. Jaffe *et al.*[92,109] have suggested that a mixing of quarks between nucleons at small relative separations[34] might be responsible for an increase in the color screening radius. They concluded that $\sim 80\%$ of the nucleons in a heavy nucleus belong to a nonnucleon component. In models like that proposed by Low,[95] a 10–20% increase in the color screening radius should result in a comparable increase in the cross section for the scattering of hadrons by bound nucleons of a nucleus, e.g., in p^4He scattering. The Glauber approximation, however, agrees within $\sim 2\%$ with experimental data.[81]

To explain the EMC effect at $x > 0.3$ e must first refine the existing limitations on the admixture of an exotic component in the nuclear wave function (Subsection 3c). It is also necessary to study directly the compressed component in the wave function of the nucleon. Apparently the clearest approach here would be to study the multiplicity of hadrons associated with the production of a Drell-Yan lepton pair in hA scattering[19] (see Refs. 38, 101, and 102 for some other suggestions).

c) The EMC effect. The region $x < 0.2$. The EMC group[12,15] found that at $Q^2 \geqslant 9$ GeV^2 and $x < 0.2$ the ratio $R(x,Q^2)$ is significantly greater than 1 and depends weakly on Q^2 (Fig. 1). The experimental data from SLAC[13,14,110] are contradictory: Recent measurements[14] failed to reveal a Q^2 dependence of $\tilde{R}_A = \sigma_{eA}/\sigma_{eD}$ at $x \sim 0.1$ in the range $2 \leqslant Q^2 \leqslant 5$ GeV^2, but, in contrast with Ref. 12, they revealed $\tilde{R}_A \approx 1$ at $x \leqslant 0.3$. In contrast, some earlier data[13b,110] at $Q^2 \sim 2$ GeV^2 and $x \sim 0.1$ are compatible with Ref. 12, and at the same time they indicate that R_A is lower than in Ref. 12 at $x < 0.1$. As has been mentioned in several papers,[111] the discrepancy between the results of Ref. 14, on the one hand, and Refs. 12 and 110, on the other, might be due to a significant value $\sigma_L/\sigma_T|_{Fe} - \sigma_L/\sigma_T|_D \approx 0.15$ at $x \sim 0.1$ and $Q^2 \sim 2$–5 GeV^2. Using the standard formulas from a QCD perturbation theory (see Ref. 11, for example), we can show that a significant difference of this sort could apparently arise only from effects which fall off in a power-law fashion with increasing Q^2 (higher twists). This difference would lead to a qualitative difference between F_{1A}/F_D and F_{2A}/F_{2D} in the SLAC kinematics. For example, we would have $F^{(N)}_{1Fe}/F^{(N)}_{1D} < 1$ and $F^{(N)}_{2Fe}/F^{(N)}_{2D} > 1$ at $x \sim 0.1$ and $Q^2 \sim 2$ GeV^2. Since these ratios should be essentially identical in the scaling limit, a strong dependence of σ_L/σ_T on A (if confirmed by further experiments) would seriously complicate an interpretation of the SLAC data in terms of the distribution of quarks in a nucleus.

We begin our theoretical discussion of the origin of the EMC effect at $x < 0.2$ with the exact sum rules which follow from QCD. For this analysis it is convenient to write F_{2A} in the form $F_{2A}(x,Q^2) = V_A(x,O^2) + O_A(x,Q^2)$, where V_A—the distribution of valence quarks in the nucleus—carries information on the quantum numbers of the nucleus, and O_A is the contribution to the cross section from the scattering of γ^* by a quark-antiquark "sea" in the nucleus. A study of ν, $\bar{\nu}$, and A scattering will make it possible to distinguish experimentally V_A and O_A, since the neutrinos interact in different ways with quarks and antiquarks.

Using the exact sum rule for the baryon charge,[38,23]

$$\int_0^A \left[\frac{1}{A} V_A(x,\ Q^2) - \frac{1}{2} V_D(x,\ Q^2)\right] \frac{dx}{x} = 0, \qquad (2.19)$$

we can easily show that if we assume $(2/A)V_A(x,Q^2)/V_D(x,Q^2) \geqslant 1$ at $x < 0.05$ the observed increase in $F_{2A}(x,Q^2)$ at small x must be due to an increase in the sea in Fe by a factor of 1.4–1.6 (Ref. 23). On the other hand, to reconcile sum rules (2.19) with the observed EMC effect it is sufficient to have $(2/A)V_A(x,Q^2)/V_D(x,Q^2) \approx 2/3$ at $x \lesssim 10^{-2}$ as a consequence of Glauber screening. This value is close to the value observed for $(1/A)F_{2A}/F_{2N}$ in Ref. 112 at $x \sim 10^{-2}$ and $Q^2 \approx 2$ GeV^2. We wish to emphasize that we do not have an adequate theory for the screening in $\sigma_{\gamma^* A}$. All that is known is that screening is present in QCD in the leading term in Q^2 (Refs. 19 and 101). We might also note that, as is shown by an analysis of the sum rules for the total momentum of the nucleus, the assumption that screening is important only for $O_A(x,Q^2)$ at $x \sim 10^{-2}$ (Ref. 30) does not explain the existing EMC data at $x \sim 0.1$.

As x decreases, the longitudinal distances which are important in the interaction of γ^* with the nucleons of the nucleus increase: $z \simeq 1/m_N x$. At $x \sim 0.1$, they are comparable to the average internucleon distances in the nucleus. As a result, the impulse approximation in terms of quark distribution is not equivalent to the impulse approximation in terms of the nucleon (cluster) degrees of freedom in the nuclear wave function. Examples are diagrams of the type in Fig. 7, which shows an exchange of partons between different nucleons (clusters) of a nucleus.[113] At $x > 0.1$, where the momenta of the clusters in the nucleus can be ignored in comparison with the momenta of the partons in a cluster, the change in the parton distributions is of the form $\delta p_A(x) \sim \lambda p^2_{cl}(x)/(R_A m_N x)(R^2_A k^2_t)$, where R_A is the radius of the nucleus, k^2_t is the average square transverse momentum of a

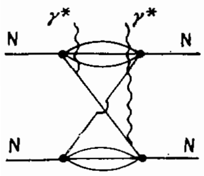

FIG. 7.

parton, $p_{\rm cl}(x)$ is the distribution of partons in a nucleon (cluster), and λ is a numerical factor ~ 1. The basic lessons to be drawn from this calculation are that at small values of x diagrams of the type in Fig. 7—fluctuations in the density of partons, pions, etc.—are important, but they decrease rapidly with increasing x, becoming an insignificant correction at $x > 0.2$.

On the basis of general considerations we can calculate only the functional dependence of $R_A(0.1)$ on the atomic number, since this dependence is governed by the circumstance that γ^* interacts with both nucleons of the nucleus:

$$R_A\,(x \sim 0.1) \sim \int \rho_{\rm A}(r_1)\,\rho_{\rm A}(r_2)\,{\rm d}^3r_1\,{\rm d}^3r_2\,\theta\left(\frac{1}{m_N x} - |z_1 - z_2|\right) \times \theta\,(1\,{\rm fm} - |r_{1t} - r_{2t}|); \qquad (2.20)$$

here $\rho_A(r)$ is the density of nucleons in the nucleus. The radius $R_A(0.1)$ is expected to be a slightly stronger function of A than $R_A(0.5)$ (Ref. 20).

Another mechanism corresponds to the coalescence of partons of different nucleons of the nucleus.[32] In QCD perturbation theory we would naturally expect[20] that diagrams of the type in Fig. 8, which correspond to the coalescence of quarks and gluons, would play a leading role at $x \sim 0.1$, since at $x > 0.05$ we have $V_{\rm N}(x) < O_{\rm N}(x)$. It is difficult to deal with this mechanism in a quantitative way (it is difficult even to determine the sign of the effect at $x \sim 0.1$), since it is necessary to solve balance equations in a region in which the virtualities at the parton interaction vertices are small.

The increase in R_A at $x \leqslant 0.2$ has been interpreted in several papers[23,31–33] as an inherently nuclear effect—an intensification of the nonstatic pion field in a nuclear medium. Calculations carried out using the model of a Fermi liquid show that a precondensation of the pion field in a nucleus leads to an increase in $O_A(x,Q^2)$ which is compatible with the observed value $\Delta = (1/A)F_{2\rm A}(x,Q^2) - (1/2)F_{2\rm D}(x,Q)^2$ at $x < 0.2$. The equations of Ref. 32, however, violate the exact (in the limit $Q^2 \to \infty$) sum rule which follows from QCD for the nuclear momentum carried off by partons. (In the formalism used in Ref. 32, the fraction of the nuclear momentum which is carried off by a pion, $\alpha_\pi \approx k_3 m_{\rm A}$, can be either positive or negative!) Incorporating the sum rules through the formalism of the nuclear wave function on the light cone leads to a much smaller (by a factor $\sim 1/3$) increase in $R_A - 1$ (Ref. 33). The calculations of Refs. 32 and 33, however, ignored the scattering of ν and $\bar\nu$ by mesons which are exchanged by nucleons (diagrams of the type in Fig. 7, where the exchange of quarks is replaced by an exchange of mesons). On the whole, the theoretical analysis of a possible role of a meson field in a nucleus is incomplete.

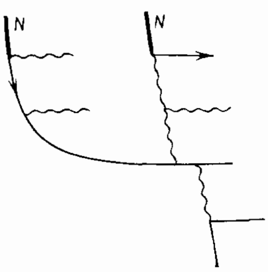

FIG. 8.

We wish to emphasize that the experimental data on ν, $\bar\nu$ scattering by nuclei[15] so far do not reveal any significant difference between $O_{\rm A}(x,Q^2)/O_{\rm N}(x,Q^2)$ and 1. In particular, preliminary data of the CDHS collaboration on the ratio of the antiquark distributions in the proton and in iron at $\langle Q^2\rangle = 7\ {\rm GeV}^2$ can be used to estimate the maximum possible contribution of the pion field of $\bar{\rm u}$,$\bar{\rm d}$ quarks to $\Delta(x,Q^2)$ It turns out to be much smaller than the value observed for Δ in Ref. 12, and this result is obviously inconsistent with the expectations of Refs. 23, 25, 27, 31, and 32. As was mentioned in Ref. 111d, the introduction of a longitudinal cross section $\sigma_{\rm L}$ which increases with A does not change this conclusion.

For completeness we note that a small increase in $F_{2\rm A}(x,Q^2)$ at small values of x, by a factor $(1 - \langle\delta\rangle P)/(1 - P) \sim 1.03$, results from the increase (discussed above) in the probability of the intermediate configurations in a bound nucleon. It is difficult to evaluate a possible contribution from the deformation of intermediate configurations.

It can be seen from an analysis of these mechanisms that distinguishing between the contributions of $V_{\rm A}$ and $O_{\rm A}$ is crucial for reaching an understanding of the nature of the EMC effect at $x < 0.2$. Such experiments are presently being carried out, on the basis of ν,$\bar\nu$, A scattering and also on the basis of the reaction $\mu + {\rm A} \to \mu + {\rm h} + {\rm X}$, with the hadron h being detected in the γ^* fragmentation region (the EMC group). It is also necessary to determine, theoretically and experimentally, the Q^2 dependence of the screening $\sigma_{\gamma^*{\rm A}}$ at $x \lesssim 10^{-2}$.

3. THE DEVELOPMENT OF IDEAS REGARDING NUCLEAR STRUCTURE

In this section we identify and discuss the theoretical problems that arise in an effort to derive a microscopic theory of the nucleus on the basis of quantum chromodynamics.

a. The meson hypothesis of the nature of nuclear forces. Study of the scattering of nucleons with energies below the threshold for meson production has revealed that over a broad range of the internucleon distance r the nuclear forces are similar to van der Waals forces (Fig. 9) but differ in that they fall off exponentially at large r. A qualitative understanding of the exponential decay of nuclear forces with increasing r. A qualitative understanding of the exponential decay of nuclear forces with increasing r is based on Yukawa's idea that the nuclear potential reduces to the

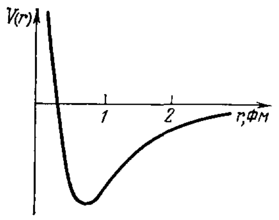

FIG. 9.

Fourier transform of diagrams with an exchange of particles:

$$V(r) \sim \int e^{iqr} \frac{d^3q}{M^2+q^2} \sim \frac{e^{-Mr}}{r}, \tag{3.1}$$

where M is the mass of the exchanged particle. The asymptotic form of the potential at large r is thus described by an exchange of a particle with a minimum mass: a π meson. Experimental confirmation of this idea comes from the agreement with experiment of calculations of the peripheral phase shifts for NN scattering on the basis of π exchange (Pomeranchuk). At present, all the data in nuclear physics indicate that the approximation of one-pion exchange gives a quantitatively correct description of the long-range part of the nuclear potential at $r \gtrsim 1/m_\pi$.

Attempts were later undertaken to calculate the potential $V(r)$ at short and intermediate range on the basis of ideas regarding the exchange of heavier mesons ($\rho, \omega, 2\pi, \ldots$) the OBEP model; see the review by Erkelenz[114]). A quantitative description of the NN-scattering phase shifts has emerged from work in this direction. However, some of the parameters of the meson-nucleon interaction which have been used are seriously at odds with the expectations based on the experimental physics of elementary particles. For example, the description of the nuclear core would require a value for the constant $g_{\omega NN}$ much larger than follows from the vector dominance model. There are many phnomena which can be interpreted as the interactions of the incident particle with the meson potential: the so-called meson currents (reviewed by Brown and Jackson[78]). In those cases in which the exchange of heavy mesons is important, however, the theoretical interpretation is ambiguous, since the contribution of the meson exchange currents depends strongly on the essentially unknown form factors for vertices of the $\rho\pi\gamma$ type.

The meson models of quantum field theory which have been offered to date are qualitatively at odds with QCD, since these models predict an increase in the invariant charge with decreasing distance at short range instead of asymptotic freedom (in QCD). On the other hand, the field-theory interaction of vector mesons with nucleons violates unitarity even in the "tree" (Born) approximation.[16)]

Because of the composite nature of nucleons and mesons, it is not obvious that the idea of meson exchanges is a suitable approximation at internucleon distances

$$r \leqslant 2r_N + 2r_M \sim 2 \text{ fm}, \tag{3.2}$$

at which a meson "cannot fit" between nucleons [r_N (r_M) is the radius of the nucleon (meson)]. From the standpoint of calculating the ladder quark-gluon diagrams corresponding to the exchange of a meson, the exchange of quarks and gluons between nucleons would appear to be a more natural source of the short-range nuclear forces. Furthermore, it is not clear whether the OBEP approximation is compatible with the experimental absence of a significant intensification of the nonstatic meson field in a nucleus (Subsection 2c).

The OBEP approximation has been applied not only to the NN system but also to the $N\bar{N}$ system. As was noted by Sakurai as early as 1960, the negative charge parity of vector mesons means that the exchange of ρ and ω in an $N\bar{N}$ system will give rise to large short-range attractive forces.[115] At the same time, in the diagrams of QCD perturbation theory we do not see a simple relationship between the NN and $N\bar{N}$ interactions (see Refs. 38 and 116, for example).

b) QCD and nuclear forces. As often happens in science, QCD automatically erased the seemingly inescapable difficulties of the meson hypothesis of nuclear forces. Instead of the "charge zero" we have asymptotic freedom in QCD. Furthermore, in contrast with the meson field-theory models the simplest theoretical description turns out to be the description of the processes which occur over short time intervals. Here we can use perturbation theory, by virtue of the asymptotic freedom in QCD, so that the scattering can be described in terms of quarks and gluons. These are the phenomena which are the present focus of attempts to test the QCD predictions. We wish to stress, however, that a brute-force attempt to derive a quantitative theory of nuclear forces on the basis of QCD runs into obvious difficulties. The kinematic region in which perturbation theory can be used is quite far from the region characteristic of the basic nuclear phenomena. [In particular, we have an invariant charge $\alpha(k^2) \gtrsim 1$ at $k \lesssim 0.2$–0.3 GeV/c.] To illustrate this assertion, we calculate the asymptotic behavior of the electromagnetic form factor of the deuteron, $F_D(q^2)$, in QCD perturbation theory. At sufficiently large values of q^2, the Feynman diagrams of the so-called democratic-chain type[35] are dominant. A calculation of these diagrams leads to a behavior $F_D(q \sim 1/(q^2)^5$ for the form factor.[35,91] This approximation, however, is valid only if all the quark and gluon propagators for the wave function of the deuteron are large; otherwise, perturbation theory is qualitatively incorrect. This condition can be written in the graphic form[34]

$$\frac{q}{6} > \langle k_q \rangle;$$

where $\langle k_q \rangle$ is the average (typical) momentum of the quark in a nucleon, ~ 0.3 GeV/c. In other words, the behavior $(1/q^2)^5$ is not justified theoretically at $|q^2| \lesssim 4$ GeV/c^2. In this region, the momenta of the nucleons in the deuteron would be $\lesssim 1$ Gev/c!

Consequently, perturbation-theory methods are almost never applicable for the scale times and scale distances characteristic of nuclear phenomena. Nevertheless, there is a basis for a restrained optimism. Significant progress has been achieved in recent years in the circumvention of the confinement problem on the basis of dispersion QCD sum rules through calculations of the static characteristics of hadrons determined by effects ignored by perturbation theory.[5] A recent calculation[118] has been successful in finding a dynam-

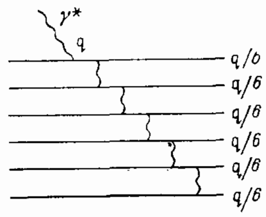

FIG. 10.

ic characteristic of a hadron—the pion form factor in the interval $0.5 < |q^2| < 4\ \mathrm{GeV}^2$—in agreement with experiment. It runs out that in this q^2 interval the form factor is determined by the interaction of quarks with vacuum fluctuations, and the one-gluon exchange diagram, which is dominant in the limit $|q^2| \to \infty$, is only a minor correction in this region.[17)]

Theoretical work has revealed several hadron properties which are specific to QCD. These properties should be considered in the construction of realistic models of the nucleus.

1. Since the masses of the u, d, and s quarks ($m_u \approx 4$ MeV, $m_d \approx 7$ MeV, and $m_s \approx 150$ MeV are small in comparison with the range of the strong interaction, $\sim m_\rho$, the QCD equations have an approximate chiral symmetry. ("Chiral symmetry" is the symmetry under independent transformations of the left-hand and right-hand components of the quarks u, d, and s.) In nature, the QCD equations have a solution in which the chiral symmetry is spontaneously broken. In the limit of vanishing current masses of the u and d quarks, the π mesons are massless Goldstone bosons. When the masses of the u and d quarks are taken into account, the pion acquires a small mass, much smaller than the masses of other hadrons. As a result, over space-time intervals $1/m_\pi \gg \sqrt{\Delta x^2} \gg 1 m_\rho$ the description of the hadron interactions in terms of a chiral Lagrangian containing only hadronic degrees of freedom works well. The chiral Lagrangian of QCD is successful in calculations not only on the low-energy interaction of π and K mesons but also the properties of the $\eta(560)$ and $\eta'(960)$ mesons (see the reviews in Refs. 5 and 117). Furthermore, we do not rule out the possibility that the nucleon can be described as a chiral soliton.[119] Accordingly, it does not year appear necessary to appeal to quark-gluon degrees of freedom in order to describe the phenomena that are determined by the average internucleon distances in a nucleus.

2. The interaction of slow pions with hadrons is proportional to the momentum of the pion. For vanishing current masses of the u and d quarks, this circumstance is a manifestation of chiral symmetry in QCD: the Goldstone nature of the π meson (the proof reduces to an examination of the Ward identities for the axial current). As a result, at a low pion energy E_π the $\mathrm{NN} \to \mathrm{NN}\pi$ reaction is dominated by the emission of pions from the external nucleon lines, while emission from the center is suppressed by a factor $\sim E_\pi / m_\rho$. A numerical analysis shows that the direct production of slow pions is slight. This theoretical result explains the experimental observation that the inelasticities in NN scattering are caused primarily by the production of baryon resonances, primarily Δ isobars. The scale value determined by the admixture of nonnucleon components is thus $m_\Delta - m_\mathrm{N} \sim 0.3$ GeV, while in the deuteron we would have 2 $(m_\Delta - m_\mathrm{N})$, because of the isoscalar nature of the deuteron. It is this QCD property which explains the dominance of the baryon degrees of freedom in the nuclear wave function and the existence of relativistic nuclear physics.

3. The presence of a condensate of quark-antiquark pairs and of gluons has the consequence that the energy of the nonperturbative vacuum is substantially lower than that of a perturbative vacuum.

4. The color forces between quarks and gluons are screened at large distances. Today, this is a plausible assumption: Theoretically, confinement has been seen only in two-dimensional models (the Schwinger model, etc.). From the discussion in Point 1 above we have $r_\mathrm{scr} < 1/m_\pi$.

Phenomenological models of the nucleus motivated by QCD have recently attracted considerable interest in nuclear physics. A fundamentally new element introduced in these models is the question of the role played by color degrees of freedom in the nuclear wave functions. The contribution of these degrees of freedom is determined by the energetics of the color excitations and the color screening radius. At present, the theory tells us little about such problems (see Points 1 and 2). It is in the treatment of these questions that the models discussed below differ from both the nonrelativistic nuclear theory and from each other.

A first group consists of the "bag" models, in which the confinement of the quarks and gluons is dealt with by introducing an external boundary which is impenetrable to the quarks and gluons, and asymptotic freedom by the fact that inside a bag the quarks and gluons are free particles (see the review by Thomas[71]). In the MIT bag model[120] and in the chiral bag model in which the pion field is present both inside and outside the bag[71] the radius of the bag is[71] 0.8–1 fm. Calculations based on these models have not yet led to a description of the interaction in a region in which bags overlap. For both models, there are comparatively light multiquark states, e.g., $M_{6q} - 2M_\mathrm{N} \sim 0.3$ GeV. (In the chiral models[119a] motivated by the analysis[119b] of the low-energy QCD Lagrangian, this difference is ~ 1 GeV and has been interpreted by the authors as a nuclear-core mechanism.) Since a multiquark system may experience a lowering of energy as a result of nonperturbative effects, forming white subsystems, the bag states apparently are not manifested as resonances.[77] On the other hand, they lead to poles in the P matrix for NN scattering and may determine the short-range dynamics of the NN interaction.[72,77]

These qualitative considerations have been incorporated in a model[72,73] of composite quark bags. It is assumed in this model that at $r \lesssim b \sim 1.1$–1.5 fm the nucleons make an abrupt transition into a quark bag, while at $r > b$ they interact through a phenomenological NN potential, which is basically a one-pion exchange. The introduction of a composite-quark-bag potential corresponds to an NN-interaction which depends on the energy; it simultaneously describes a short-range repulsion and an attraction at intermediate internucleon distances. This model has been successful in calculating the lowest NN-scattering phase shifts at $T_N < 1$

GeV in terms of the bag radius b, the residue of the NN→(composite quark bags) transitions, and the composite-quark-bag energy E_ν. Consequently, in contrast with the earlier studies (see Refs. 34–36 and 55–59, for example), it becomes possible to incorporate 6q bags in the wave functions of the deuteron and nuclei in a manner consistent with the NN-scattering phase shifts. Incorporating the strong coupling of composite quark bags and NN channels has the consequence that the admixture of nonnucleon degrees of freedom is small even though it determines the entire dynamics of the NN interaction at $r < b$ (Ref. 72):

$$P_{ex}\ (\mathrm{D}) \sim 2\ \%, \quad P_{ex}\ (A \gg 2) \sim 5\ \%. \tag{3.4}$$

While the wave function of a six-quark composite quark bag is described by the MIT model, its exotic component is due primarily not to an admixture of baryon resonances but to states constructed from two color three-quark clusters.[36] This conclusion may turn out to be qualitatively wrong since it ignores the differences between the characteristic excitation energies of white and color degrees of freedom.

An alternative to the bag model is the model of constituent quarks, in which it is assumed that the properties of the hadrons are determined by the interaction between quarks, while the interaction with the vacuum condensates is dealt with by introducing an effective quark mass $m_u \approx m_d \approx 300$ MeV. The constituent quark model is successful in describing the basic characteristics of hadrons and even high-energy collisions of hadrons (see the reviews in Refs. 121 and 122). The achievements of the constituent quark model in describing experimental data are stimulating the development of potential models of hadrons, in which it is assumed that the constituent quarks interact with each other through a binary potential which increases in a power-law fashion with distance (see the review and citations of earlier papers in Refs. 74–76). This description is incorrect at large distances between the nucleons, since it ignores the production of $q\bar{q}$ pairs, which converts the power-law decay of the NN potential into an exponential decay. Under the assumption that the potential is determined at $r > 2$ fm by one-pion exchange, Maltam and Isgur[76] showed that at $r < 2$ fm the NN potential calculated by a variational method from the model of Ref. 76 is similar to the phenomenological soft-core potentials. The attraction at $r \sim 1.4$ fm arises primarily from the presence of spatially separated color three-quark clusters in this model. In the core region, according to Ref. 76, white three-quark clusters are dominant, while Refs. 74 and 75 have a six-quark $^4S^2P$ configuration dominant here. These models are difficult to reconcile with the hypothesis of quark and gluon confinement since they contain van der Waals color forces, which provide the basic attraction between nucleons in Ref. 76. Furthermore, the interaction of color entities at distances > 1 fm is difficult to reconcile with QCD lattice calculations, which show that color-screening effects (the production of pairs from vacuum) become important at these distances.[18)]

On the whole, although the use of QCD to derive a microscopic theory of the nucleus has so far been at a semi-quantitative level, the existence of a significant high-momentum component in the nuclear wave function and the existence of a large number of phenomena (Section 4) in which this component can be studied quantitatively raise the hope that relativistic nuclear physics will eventually become a method for studying in detail nonperturbative effects in QCD.

c) Possible search for exotic degrees of freedom in the wave function of a nucleus. From the standpoint of the physics of strong interactions, a description of a nucleus in terms of nucleon degrees of freedom alone can be only approximate. Fundamental questions here are how important these exotic (i.e., nonnucleon) components are and what their nature is.

In the prequark era (before the November Revolution of 1974: the discovery of the ψ meson), interest focused on an admixture of baryon resonances in the nuclear wave function, particularly an admixture of a ΔΔ component in the deuteron.[19)] The probability for this configuration, $I_{\Delta\Delta}$, can be estimated on the basis of dimensional considerations, through a comparison with the probability that the nucleons in the deuteron are in a configuration with an energy non-conservation of the order of the difference between the masses of the ΔΔ system and the deuteron[115]:

$$k_{\mathrm{N}}^{\mathrm{vi}} \geqslant k_0 = \sqrt{m_\Delta^2 - m_{\mathrm{N}}^2} \approx 0.8\,\mathrm{GeV}/c, \tag{3.5}$$

i.e.,

$$I_{\Delta\Delta} = \int \psi_{\mathrm{D}}^2 (k)\, d^3k\, \theta\,(k - k_0) \approx 5 \cdot 10^{-3}. \tag{3.6}$$

Here realistic wave functions of the deuteron, with a core, have been used for ψ_{D}. The incorporation of overlap integrals in the model of one-gluon exchange leads to[115] an additional suppression of $I_{\Delta\Delta} : I_{\Delta\Delta} \sim 10^{-3}$. A similar value is expected in the model of composite quark bags.[73] This estimate is consistent with the results of an experimental search for isobars in the deuteron[123]:

$$I_{\Delta\Delta} = (1 \pm 2) \cdot 10^{-3}.$$

The experimental search for an admixture of isobars in the deuteron thus appears to be a very difficult task. Special conditions must be arranged; e.g., a search must be made for cumulative isobars in deep inelastic scattering of leptons in order to avoid background processes. For heavier nuclei, the theoretical restrictions are far less severe, since the NN→NΔ transitions are suppressed to a lesser extent by both energy denominators and the combinatorices than in the case of the NN→NΔ transitions. It would thus be interesting to measure the ratio of inclusive cross sections for reactions at high energies[20)]:

$$\frac{\alpha\, d\sigma^{h+A\to\Delta(N^*,\, N\pi\ldots)+X}}{d^2k_t}(\alpha,\ k_t)\left(\frac{\alpha\, d\sigma^{h+A\to N+X}}{d\alpha\, d^2k_t}\right)^{-1}$$

$$\text{for}\quad \alpha_\Delta > 1.$$

In order to suppress the possible production of N* and Δ in secondary processes, it would be necessary to use the lightest nuclei, with $A \leqslant 6$. To the best of our knowledge, absolutely no such data are presently available. Limitations at a level of even 1/3 would be informative.

Quark physics has raised the question of six-quark configurations in the deuteron in which the quarks belonging to the proton and the neutron are mixed.[34–36] As Feynman[124] has pointed out, however, any "white" state can be described in terms of both hadronic and quark-gluon degrees of freedom. Just which description is simpler is a dynamic problem. If the color clusters are not spatially separated, then a description in terms of white hadronic degrees of freedom would be just as convenient as the other description. If, on the other hand, there are spatially separated color groups of three quarks in the deuteron, then such a state could be described more conveniently in terms of quark degrees of freedom. There are clear advantages in using the two descriptions simultaneously. For example, the absence of significant inelasticities in the NN interaction with a zero isospin suggests that it might be possible to describe the deuteron in terms of nucleon degrees of freedom over a rather broad range [see (3.5)]. In terms of quarks, in contrast, it is a simpler matter to calculate the momentum dependence of the so-called hard nuclear processes.[125] We recall that it was this approach which proved successful in describing the properties of hadrons on the basis of dispersion sum rules.[5]

Here is a list of the existing limitations on an exotic baryon component in the wave function of a nucleus (cf. the discussion in Section 2).

a) The nonrelativistic theory of the deuteron at nucleon momenta $k \leqslant 0.2$–0.3 Gev/c agrees within $\sim 10\%$ with experiment, while at $k \geqslant 0.3$ Gev/c ("hard" nuclear process; Section 4) the two-nucleon approximation describes cross sections with $\sim 30\%$. An exotic component should therefore give us simply corrections to the contribution of the nucleon component. If we assume that a transition to an exotic component is possible only at distances $\leqslant 1/m_\pi$ or $k \geqslant 200$ MeV/c, then we have the following limitation on $P_{ex}(\mathrm{D})$:

$$P_{ex}\ (D) < 3\ \%. \tag{3.7}$$

A definite limitation is also imposed by the agreement of the nonrelativistic theory with the static characteristic of the deuteron.

b) For ^{3}He, an analysis of p^{3}He scattering[126] yields a limitation on the admixture of light nonnucleon components, primarily the Δ isobar (per nucleon):

$$P_{ex}\ (^3\mathrm{He}) < 1.7\ \%. \tag{3.8}$$

c) In the case of carbon, an analysis of (e,e′) reactions near the quasi-elastic peak[82] leads to an integral estimate: The probability for finding a nucleon in a nucleus with a momentum less than the Fermi momentum is[82] $\sim 75\%$. On the other hand, the probability for finding a nucleon with a momentum > 300 MeV/c in a nucleus is $\sim 30\%$ (Section 4). Taking into account the uncertainties in these numbers, we find the following conservative upper estimate of the magnitude of the exotic component:

$$P_{ex}\ (^{12}\mathrm{C}) < 10\ \%. \tag{3.9}$$

As in the case of the deuteron, the exotic component is only a correction to the high-momentum nucleon component over a broad range of nucleon momenta. An estimate close to (3.9) can also be derived for Fe, but it is less reliable in that case because of the lack of data on lepton reactions at Fe at[21] $x > 1$.

Correlation experiments hold promise for a search for exotic configurations. Of primary interest here are lepton reactions at $Q^2 \geqslant 1$ GeV^2 with several particles detected in the final state, for the lightest nuclei ($A \leqslant 6$): $l + A \to l' + \mathrm{N}(\Delta, N\pi) + \mathrm{X}$ with $\alpha_\mathrm{N}(\alpha_\Delta, \alpha_{\mathrm{N}\pi}) > 1$. In such reactions, however, we can search for only products of the decay of multiquark (primarily 6q) configurations. In order to detect the ejection of a six-quark as a whole we would naturally turn to processes involving a study of three-particle correlations, corresponding to three nearly coincident nucleons. If, as two nucleons close on each other, a six-quark is formed, then as three nucleons close on each other the momentum of one may be balanced by a six-quark. Accordingly, when this nucleon is ejected a six-quark may simultaneously be liberated. A necessary condition here is to select configurations in which scattering by triple correlations is predominant. The simplest example is the reaction $\mathrm{e} + {}^3\mathrm{He} \to \mathrm{e} + \mathrm{p} + \mathrm{X}$ with $\alpha_\mathrm{p} \geqslant 1.5$ at $E_\mathrm{e} \sim 5$ GeV. A selection of events with $(p_t)_\mathrm{p} \approx -(p_t)_\mathrm{e} \geqslant 0.5$ Gev/c sharply increases the contribution from the elastic scattering of the electron by fast nucleon of the nucleus. To search for manifestations of 6q configurations we should study the composition of the forward-emitted system (2N,2Nπ,...) and seek peaks in the mass spectrum of this system. Another possibility is to study the corresponding reaction with a cumulative pion, to study deep inelastic scattering, and to replace the lepton by a hadron (see also the discussion in Ref. 38 and in Subsection 4a).

Another important question in the effort to learn about the dynamics of the NN interaction is whether the dibaryon resonances which have recently been discussed extensively actually exist (see the review by Makarov[127]). At this point, the theoretical and experimental situation here is unsatisfactory, since the positions of the assumed resonances are close to the thresholds for the production of isobars. Calculations show (Ref. 128, for example) that if a reaction involving isobars came into play it would cause irregularities similar to those observed experimentally. Since the dibaryons have primarily inelastic decay modes, according to the existing phase-shift analyses, it becomes necessary to study accurately inelastic reactions in which the dibaryon contribution should be enhanced: $\mathrm{NN} \to \mathrm{N}\Delta$, $\mathrm{N} + \mathrm{N} \to \mathrm{N} + \Delta$, $\pi\mathrm{D} \to \mathrm{N}\Delta$, ($\pi$NN), $\mathrm{N} + \mathrm{d} \to \mathrm{NNN}$, $\mathrm{NN} \to \pi\mathrm{D}$. In such reactions, it is also a comparatively simple matter to measure the polarization of Δ, and there are some rather simple ways to measure the polarization of the fast deuteron.[129]

4. STUDY OF SHORT-RANGE CORRELATIONS IN NUCLEI IN THE SCATTERING OF HIGH-ENERGY PARTICLES

To the best of our knowledge, the methods of high-energy physics were first used in a purposeful way to measure the high-momentum component of a nuclear wave function by Lederman *et al.*,[130] who studied the production of antiprotons below the threshold. In recent years, as extensive data on hadron reactions[43–50] and then lepton reactions[37,39–42] with a large energy transfer to the nucleons of a nucleus have become available, it has become possible to take up the study of not only the magnitude but also the structure of short-range nucleon correlations in a nuclear wave function.

a) Near-threshold (e,e') *reactions at large* Q^2. The inelastic scattering of electrons by the nuclei D (Refs. 3 and 4), He, and Al has been studied at SLAC at[39,40] $Q^2 \geqslant 1$ GeV^2:

$$e + A \to e' + X. \tag{4.1}$$

This study has been carried out near the kinematic boundary, i.e., under conditions such that the mass of the product system X is small, $W - M_A \lesssim 0.3$ GeV, with $x = Q^2/2q_0 m_N > 1$. In this kinematics the production of pions is strongly suppressed, so that X is a system of nucleons and nuclear fragments. A natural mechanism for reaction (4.1) (see Ref. 131, for example, and the bibliography there) is the elastic scattering of the e by a nucleon with a large momentum **k**, *directed opposite* **q** *(in the rest frame of the nucleus A)*, since only in this case can $W - M_A$ be small. In other words, the contribution from the high-momentum component of the nuclear wave function is kinematically enhanced in such processes.

$A = 2$. In the region $W - M_D \geqslant 50$ MeV, where the final state interaction is slight, the process is described by the sum of diagrams in Fig. 11, which correspond to the impulse approximation and to an interference, respectively. The cross section is expressed in terms of the deuteron wave function on the light cone, which is essentially equal in the two-nucleon approximation to the nonrelativistic wave function of the deuteron[22)] (Ref. 52). The results calculated with the help of the standard wave functions of the deuteron with a core[125] agree reasonably well with the SLAC data[23)] (Ref. 39). An interesting way to check the dominance of the diagram in Fig. 11a was proposed by West[131] in a nonrelativistic approach, where at large values of Q^2 the ratio

$$\frac{\partial q_0}{\partial y}\,\frac{d\sigma^{e+D\to e'+x}}{d\Omega_{e'}\,dE_{e'}}\,(\sigma_{\mathrm{Mott}}\,|f^2(Q^2)|)^{-1} = F(y) \tag{4.2}$$

should depend only on the variable y determined by energy conservation:

$$q_0 + m_D = \sqrt{m^2 + (q_3 - y)^2} + \sqrt{m^2 + y^2} \tag{4.3}$$

This is so-called y scaling; here $|f(Q^2)|^2$ is the square of the corresponding nucleon form factor; and σ_{Mott} is the cross section for scattering by a point nucleon. The analysis by Bosted *et al.*[136] has shown that y scaling agrees with the data of Refs. 38 and 106. (A relativistic motion of the nucleons[25] does not significantly disrupt y scaling.[137]) This is an independent test of the suggestion that the wave function of the deuteron can be measured in reaction (4.1). On the whole,

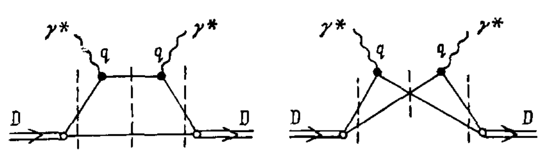

FIG. 11.

analysis of (e,e') reactions at the deuteron has shown that the two-nucleon approximation works reasonably well at $Q^2 \leqslant 2$ GeV^2 up to $k \leqslant 0.6$ GeV/c and that we have[24)] $\int \psi_D^2(k) d^3k \cdot \theta(k - \mathrm{GeV}/c) \sim 4$–$5\%$. This integral is dominated by the D wave, in agreement with the observed behavior of the elastic cross section for hD scattering in the momentum-transfer region 0.3–0.4 GeV^2 (Ref. 138, for example) and with data on elastic pD scattering[25)] (Ref. 139).

It is reasonable to expect that at $Q^2 > 2$ GeV^2 the form factor of the nucleon [like $F_{2N}(x,Q^2)$ at $x \geqslant 0.6$] will be dominated by scattering by compressed configurations in the wave function of the nucleon. Since the probability for compressed configurations is suppressed in a bound nucleon [cf. (2.15')], the two-nucleon approximation should overestimate the cross section for the reaction eD→eNN at large Q^2 and at $k \gtrsim 0.5$ Gev/c by a factor of[20] $\delta_D(k)$. Furthermore, at *$k \gtrsim 0.5$ Gev/c we would expect an additional decrease in the* cross section (beyond that stemming from the EMC effect), since a deformed nucleon cannot undergo fragmentation because of the limited phase volume to the states allowed by the quantum numbers. Refined data furnished us by Arnold *et al.*[106] are consistent with the existence of such an effect (Fig. 12). We wish to emphasize that since the standard deuteron wave functions (the Reid and soft-core wave functions) successfully describe the cross section at $Q^2 \leqslant 2$ GeV^2 (Fig. 12), it is a deviation from the two-nucleon approximation which we observe experimentally. We need detailed measurements of the cross section for the reaction eD→e + p + n, since this reaction at $Q^2 \geqslant 4$ GeV^2 can furnish direct information on the six-quark component in the deuteron (cf. Subsection 3a).

$A \geqslant 3$. At $A > 2$, the large momentum (k) of the nucleon by which the electron is scattered may be balanced by not only a single nucleon with a momentum of $-k$, as in the deuteron (a binary correlation), but by two nucleons, with $k_1 \simeq k_2 \sim -k/2$ (a ternary correlation), etc. It is thus convenient to write the cross section for reaction (4.1) as the sum of the cross sections for scattering by two-nucleon, three-nucleon, etc., correlations[38]:

$$\sigma_A(x, Q^2) = A \sum_{j=2}^{A} a_j(A)\,\sigma_j(x, Q^2). \tag{4.4}$$

We have used the circumstance that at large values of x the wave functions of a j-nucleon correlation and thus the cross section for scattering by it are determined exclusively by the local properties of the nuclear matter; the only property that depends on the nucleus is the relative probability for a j-nucleon correlation, $a_j(A)$. It is thus obvious that (4.4) must hold in cases in which electrons are scattered by 6q,9q,..., clusters[26)] (Ref. 140). From the kinematics we have

$$\sigma_j(x, Q^2) = 0 \quad \text{at} \quad x > j. \tag{4.5}$$

It is convenient to normalize σ_2 in such a way that at $x > 1$ we have $\sigma_2(x,Q^2) = \sigma_D(x,Q^2)$. It follows from condition (4.5) that at $x \geqslant 2$ we have $\sigma_3(x,Q^2) \sim \sigma_{^3He}(x,Q^2)$. Since the nucleus is a rather low-density system, the contribution of the j-nucleon correlation should dominate at $j > x > j - 1$. The ratio

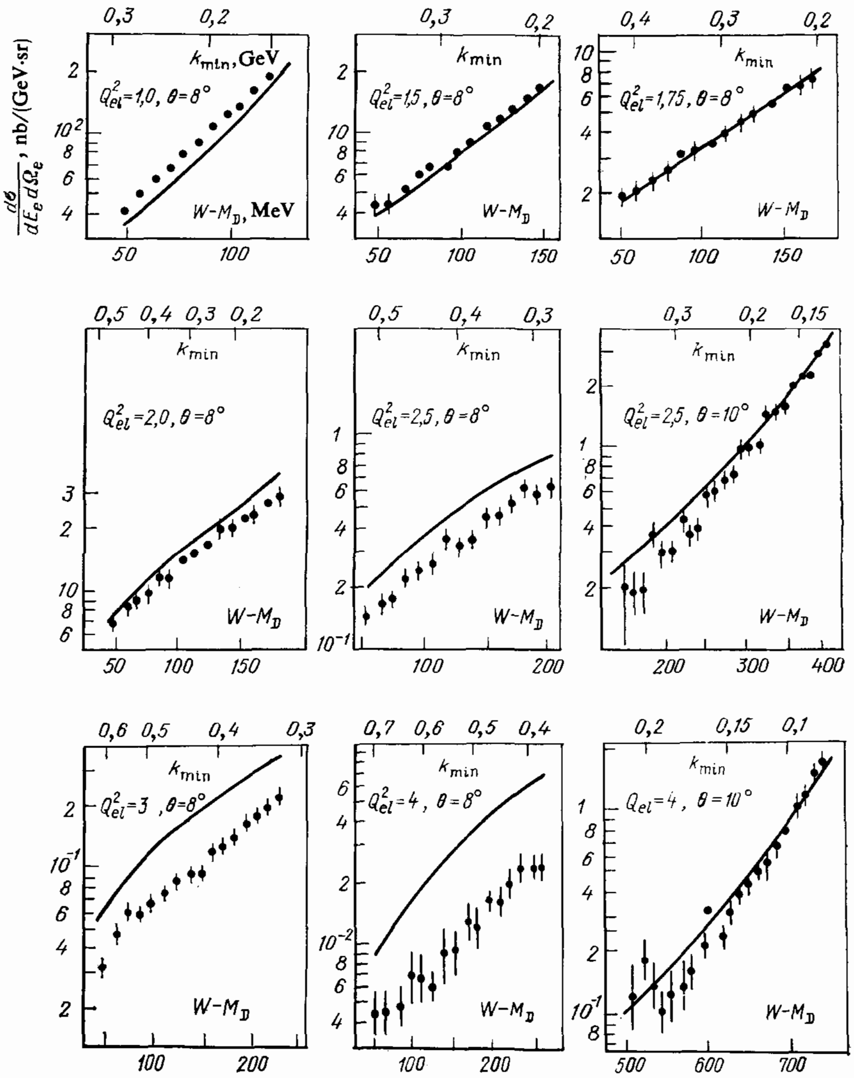

FIG. 12. Cross section for the reaction[106] $e + D \to e + X$ as a function of the invariant mass of the product system, W, and of the minimum momentum of the nucleon in the deuteron wave function k_{min}, allowed for the given values of Q^2 and W. Here Q^2_{el} is Q^2 for the elastic scattering of the electron by the deuteron (or nucleon) at $\theta = 8°$ (10°). The curves show calculations in the two-nucleon approximation with the deuteron wave functions (Reid, soft core). As in Ref. 136, a limitation $W - M_D \geqslant 50$ MeV must be imposed in order to suppress the final state interaction.

$$\frac{\sigma_A(x,\ Q)^2}{\sigma_{A=j}(x,\ Q^2)}\bigg|_{j-1<x<j} \approx \frac{A}{j}\,a_j \tag{4.6}$$

should therefore depend only weakly on[27] x and Q^2 (Ref. 38). Relation (4.6) explains the regularities observed in the experimental data on near-threshold (e,e′) reactions at ^{3,4}He and Al:

a) $\sigma_{^3He}(x,Q^2)\sigma_D(x,Q^2)$ is essentially independent of x for $1 < x < 2$ (Fig. 13). The magnitude of this ratio agrees with the expectations of Ref. 38, based on an analysis of the yields of cumulative p and π from ^{4}He (more on this below).

b) $R(x,Q^2) = \sigma_{^4He}(x,Q^2)/\sigma_{^3He}(x,Q^2)|_{2<x<3}$ is also slightly dependent on Q^2 (Fig. 14), although $\sigma_{^3He}$ and $\sigma_{^4He}$ fall off by a factor of 100 over this Q^2 interval. The observed value of $R \sim 3$ corresponds to a doubling of the probability for ppn and nnp trios in ^{4}He in comparison with the correlation probability in ^{3}He. This result is in agreement with a combinatorial estimate[38] based on the near equality of the radii of ^{3}He and ^{4}He.

c) Comparison of the spectrum from an empty aluminum target, reported in Ref. 40b, with data on ^{4}He indicates $(1/27)\sigma_{Al}/(1/4)\sigma_{^4He} \sim 2$ and that this cross section depends only weakly on x in the interval[28] $2 \leqslant x \leqslant 2.7$.

Day *et al.* and Rock *et al.*[40] report that the data on $\sigma_{^3He}(x,Q^2)$ exceed by a factor ~ 5 the theoretical results based on a solution of the Schrödinger equation for ^{3}He with a binary NN interaction. To a significant degree, the reason for this discrepancy is that the calculations used in Ref. 40 discarded the high-momentum component of the spectral function with momenta $k \geqslant 0.5$ Gev/c. Analysis of recent calculations[141] of the ^{3}He wave functions seems to indicate that the discrepancy with experiment is not large in the binary-correlation region. Specifically, the calculations lead to $a_2(^3\text{He}) \simeq 1.3$, while the data of Refs. 40 and 106 imply $a_2(^3\text{He}) \simeq 1.7$. The difference between these numbers is of the same order of magnitude as the decrease in the average density of nucleons in ^{3}He in the calculations of Ref. 141.

On the whole, the data on (e,e′) reactions indicate a significant absolute value of the high-momentum component of the nuclear wave function. The theoretical interpretation of the data, on the other hand, is ambiguous: The results are consistent both with an expansion of the cross section in contributions from binary, ternary, etc., nucleon correlations [see (4.4)] and with models which assume that the mo-

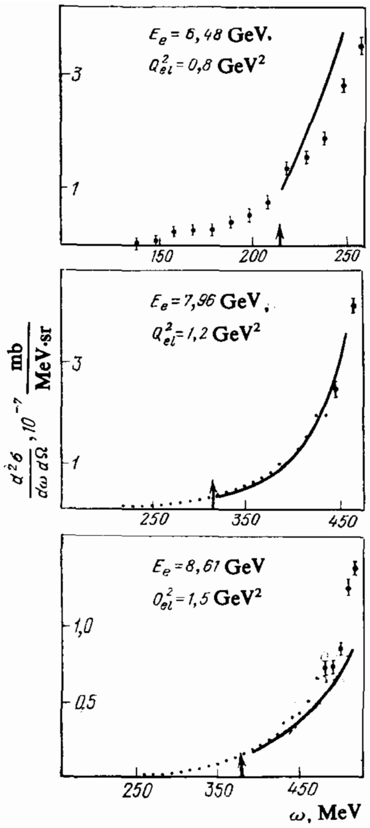

FIG. 13. The differential cross section for the (e,e') reaction on ^{3}He as a function of the energy transfer $\omega = E_i - E_f$ (Ref. 40a). Solid lines—Calculated in the binary-correlation approximation; arrows—boundary of the region kinematically allowed for scattering by the deuteron.

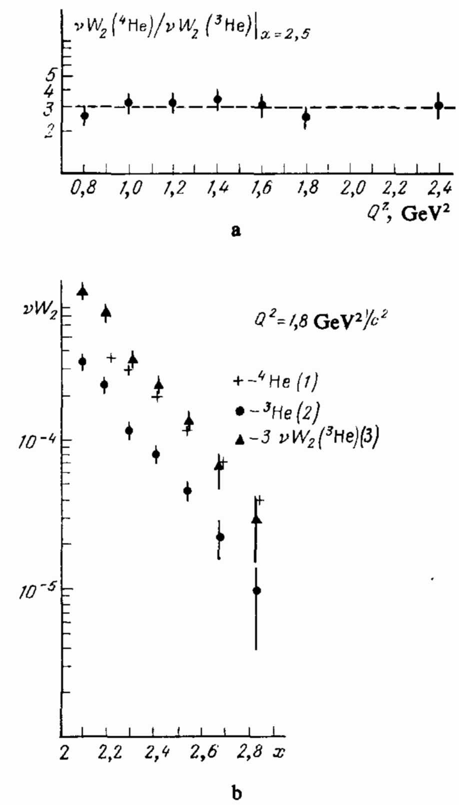

FIG. 14. a: Ratio of the structure functions for scattering by ^{4}He and ^{3}He at the fixed value $x = 2.5$ as a function of Q^2 (data from Ref. 40b). b: Comparison of the x dependence of the structure functions of ^{4}He and ^{3}He at $Q^2 = 1.8\,\text{GeV}^2$. 1—$F_{2\,^4\text{He}}(x,Q^2)$; 2—$F_{2\,^3\text{He}}(x,Q^2)$; 3—$3F_{2\,^3\text{He}}(x,Q^2)$ (data from Ref. 40b).

mentum of the fast nucleon is balanced by a system of nucleons with a mass[131] $\sim M_{A-1}$ (the mean field approximation).

b) Deep inelastic lA scattering at $x \gg 1$. The behavior of the nuclear structure functions $-F_{2A}(x,Q^2)$ at $x>1$ and their significant value were predicted in Refs. 132 and 142 on the basis of Eq. (2.17). Baldin[143] worked from the assumption that the x_F dependence for the spectra of cumulative particles agrees with the x distribution of quarks to predict $F_{2A}(x) \sim \exp(-ax)$ with $a \approx 7$. In the model of few-nucleon correlations, the similarity of the distributions of cumulative pions and quarks over a broad range $x>1$ arises at $A>2$ as a consequence of the similarities in the functional dependence on x at $x<0.8$ for the spectra of pions in hN scattering and the quark distribution in the nucleon. On the other hand, in the case of the deuteron the spectrum of cumulative pions is expected[125] to decrease as $x \to 2$ much more slowly than $F_{2D}(x,Q^2)$ [see (4.12)]. Preliminary results of the measurements of $F_{2C}(x,Q^2)$ at $x \gg 1$ in μ^{12}C scattering at $Q^2 \sim 100$ GeV² have recently appeared (Fig. 15).[41] The experimental slope of the spectrum, $a \approx 8$–10, agrees reasonably well with the expectations of Refs. 132, 142, and 143. On the other hand, as can be seen from Fig. 15, nuclear models which do not contain a significant high-momentum component, the Fermi-gas model, and the oscillator model are in qualitative contradiction of experiment. The calculations of Refs. 132 and 142 in the binary-correlation approximation, with $\rho^p_A(\alpha,k_t) = \lambda Z \rho^p_D(\alpha,k_t)$ at $k \gg 0.3$ GeV/c and $\lambda = 6$–8 (Ref. 53) and with the same ρ^p_D as in (e,e') reactions, agree reasonably well with experiment at $x \sim 1$. At large x, the data are consistent[144] with the calculations of Refs. 125 and 145, which use a density matrix for nucleons in a nucleus on the light cone, $-\rho^N_A(\alpha,k_t)$, extracted from a theoretical analysis of processes involving the production of cumulative particles. The need to use the wave functions on the light cone is discussed in the Appendix. The relation between ρ^N_A and the nonrelativistic nuclear wave function (in the region in which there is such a relationship) was studied in Ref. 38. The contribution of ternary and quaternary correlations causes an increase in the ratio of structure functions $F_{2A}(x,Q^2)/F_{2D}(x,Q^2)$ at $x>1$ (Fig. 6). The function $F_{2D}(x,Q^2)$ calculated in this manner agrees[52] with the scaling limit of $F_{2D}(x,Q^2)$ measured in Ref. 39, where the mass of the product hadron state was small. The estimate of the coefficient λ in Ref. 53 agrees reasonably well[68,38] with the existing analyses of the photo-disintegration of nuclei, where the number of triplet pn pairs in the ^{12}C nucleus with momenta $k \sim 0.3$–0.4 Gev/c has been measured. This value of λ corresponds to the following value of the high-momentum component:

$$\int n(k)\,\theta\,(k - 0{,}4\ \text{GeV}/c)\,d^3k \sim 0.15. \tag{4.7}$$

We wish to emphasize that in this definition λ incorporates

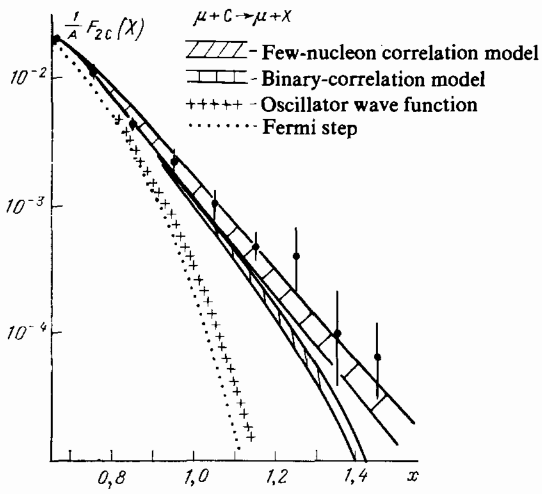

FIG. 15. Comparison of the data of Ref. 41 with calculations of $F_{2C}(x,Q^2)$ using several nuclear wave functions.

the contribution of higher-order correlations. For example, it can be expected on the basis of an analysis of the cross section for the production of cumulative nucleons[38] that the coefficient for specifically binary correlations would be $\lambda \sim 4$–6.

The data of Ref. 41 have also been analyzed in the fluctuon model[29)] (Refs. 27 and 146). In contrast with the few-nucleon correlation model, the fluctuon model predicts that the ratio $(1/6)F_{2C}(x,Q^2)/F_{2D}(x,Q^2)$ will depend only weakly on x at $1<x<1.5$ and will have a value ~ 2. This number is significantly different from the experimentally observed ratio of the inclusive cross sections for the production of cumulative pions from ^{12}C and D (per nucleon); this ratio is ~ 5–6 at $\alpha_\pi > 1$ and increases with increasing α_π (Ref. 87).

c) Production of cumulative nucleons in the scattering of (anti)neutrinos. In principle, the most effective way to study the structure of short-range correlations is to study deep inelastic processes and to detect the additional particles. By virtue of the kinematics, a nucleon emitted into the rear hemisphere with a large momentum $-p$ (a cumulative nucleon) may be produced as a result of the disruption of the correlation only if the ν is scattered by a nucleon which is correlated with an ejected nucleon which is moving forward with a momentum p' (Fig. 16); this is the so-called spectator mechanism.[30)] In this case we should have a sort of Doppler effect: the mean value of x for events in which a cumulative nucleon with a momentum p is detected should be smaller than in the average event.[53] In the case of $\bar{\nu}$Ne, where the effects of secondary interactions are apparently slight, a decrease in $\langle v\rangle_p = \langle xy\rangle_p$ compatible with the predictions of Refs. 38 and 53 has been observed.[42] Those predictions were based on a dominance of binary and ternary nucleon correlations in the high-momentum component of the nuclear wave function (Fig. 17a); here $y = E_\mu/E_{\bar{\nu}}$. The coefficients a_2 and a_3 have been determined independently, from an analysis of hadronic processes involving the production of cumulative protons[38] (see Ref. 38 for a discussion of the data of Ref. 148 for heavier nuclei with $A \sim 10^2$, where secondary-interaction effects are intensified). The absolute value of the cross sec-

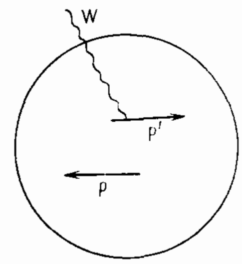

FIG. 16.

tion also agrees (within $\sim$20–30%) with the value extracted for the high-momentum component of the nuclear wave function from the value of $F_{2^{12}C}(x,Q^2)$ at $x \geqslant 1$ (Ref. 143). On the other hand, $\langle y\rangle$ did not increase (Fig. 17b). This is not a trivial point, since in the average event at a nucleon a decrease in x is accompanied by an increase in y (the contribution from the scattering of $\bar{\nu}$ by sea quarks increases).

d) Production of cumulative particles in high-energy hadronic processes. Extensive information has now been accumulated on the production of cumulative particles in high-energy processes[43–50,87,149]

$$a + A \to b + X, \text{ where } a = (\gamma, \nu, \bar{\nu}, \pi, p, \text{light nuclei}, \ldots), \; b = N, \pi, \tag{4.8}$$

i.e., in the limiting fragmentation region of the nucleus, which is kinematically forbidden for scattering hydrogen:

$$\alpha = (E_b - p_{b3})\, m_N^{-1} > 1. \tag{4.9}$$

The direction of the momentum of the incident particle, a, is selected as axis 3.

The study of processes (4.8) was stimulated by Baldin's suggestion[51] that the production of cumulative pions might be a limiting fragmentation process governed by the interaction of fast particles with fluctuations of the nuclear density and also by the formulation by Bayukov *et al.*[43] of the basic behavior in the production of cumulative protons in hadronic processes. The reader is referred to Refs. 38, 45–49, 58, 59b, and 68 for a detailed account of the extensive experimental results and a bibliography.

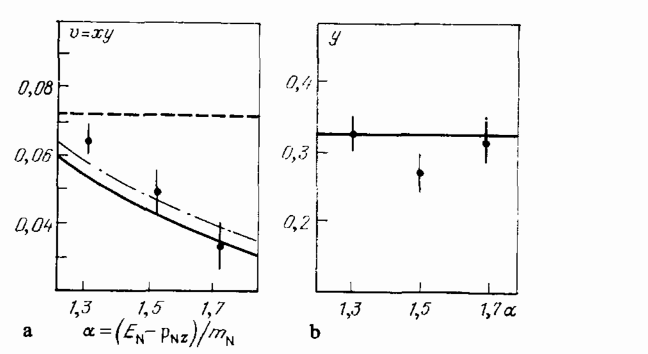

FIG. 17. The behavior of $v(\alpha)$ and $y(\alpha)$ for events with cumulative nucleons in the reaction $\bar{\nu} + \text{Ne} \to \mu^+ + p + X$. The data are from Ref. 42. Solid line—Prediction of the few-nucleon correlation model; dot-dashed line—effect of a 20% contribution of secondary interactions; dashed line—mean value of v for the sample.

In typical inelastic high-energy hadronic processes ($E_a \geqslant 10$ GeV) the average momentum transfer to the ejected nucleon of the nucleus is not small,[10] ~ 1 GeV/c so that again in the case of hadronic processes it is possible in principle to study short-range nucleon correlations.

$A = 2$. The basic contribution of the high-momentum component of the deuteron wave function to the cross section for the production of cumulative nucleons is the so-called spectator mechanism[52,53] The amplitude for this process corresponds to the scattering by a forward-moving nucleon (Fig. 18a) and to Glauber screening (Fig. 18b). The contribution of the direct mechanism—the scattering by a nucleon moving opposite to h (Fig. 18c)—is a rather small correction. At low values of p, Glauber screening cancels out for this mechanism.[140] The final result for the inclusive cross section

$$G_h^{D/N}(\alpha,\ p_t) \equiv \frac{\alpha\, d^3\sigma^{h+D\to N+x}}{d\alpha\, d^2 p_t},$$

i.e., for the cross section summed over all processes and over all parameters of the product particles except those under study, is[144]

$$G_h^{D/p}(\alpha,\ p_t) = \varkappa_h \sigma_{in}^{hN} \rho_D^N(\alpha,\ p_t), \tag{4.10}$$

where the factor $\varkappa_h = 0.7$–0.9 incorporates the Glauber screening and the contribution of the direct mechanism (see Refs. 129 and 150 for the expression for $\varkappa_h$).[31)] The cross section for the production of cumulative π mesons, etc., is described exclusively by the direct mechanism and is given by[52]

$$G_h^{D/\pi}(\alpha,\ p_t) \doteq \sum_{N=p,\,n} \int \rho_D^N(\beta,\ k_t)\, G_h^{N/\pi}\left(\frac{\alpha}{\beta},\ p_t + \frac{\alpha}{\beta} k_t\right) \frac{d\beta}{\beta} d^2 k_t. \tag{4.11}$$

The possibility that the production of cumulative pions from a deuteron is an effect of small internucleon distances in the deuteron was first raised by Baldin,[51] on the basis of the idea that two nucleons interact as a single hadron in this case, i.e., that the ratio $G_h^{D/\pi}(\alpha)/G_h^{N/\pi}(\alpha/2)$ is independent of α.

Calculations from Eq. (4.10) with realistic deuteron wave functions agree well with the data of Refs. 46 and 49 at $\alpha \lesssim 1.6$–1.7 ($k \leqslant 0.8$ Gev/c). At larger values of α the contributions of the more complex quark-gluon configurations in the deuteron[125] and the three-reggeon contribution,[38,132] corresponding to the regime $G_h^{D/p}(\alpha, p_t) \sim (2-\alpha)^2$, are parametrically enhanced. The possibility of a substantial production of nucleons through a final state interaction by means of color forces has also been discussed.[67] A distinctive feature of the latter mechanism is the small associative multiplicity; another is the suppression of the production of cumulative nucleons in deep inelastic scattering of leptons by the deuteron.

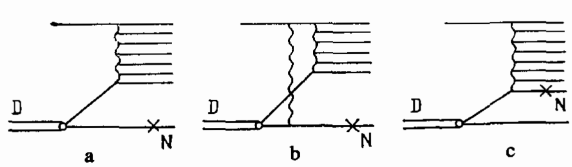

FIG. 18.

Equation (4.11) describes the cross section for the production of cumulative pions[46,50,87] up to $\alpha \sim 1.3$. At larger values of α (where the contribution from momenta $k \gtrsim 1$ GeV/c is dominant in the deuteron wave function in the impulse approximation), the contributions of more-complex quark-gluon configurations in the deuteron and three-reggeon contributions may prove important, as in the case of the production of cumulative nucleons. In the limit $\alpha \to 2$, these contributions are parametrically enhanced and lead to[125]

$$G_h^{D/\pi}(\alpha)|_{\alpha\to 2} \sim (2-\alpha)^5. \tag{4.12}$$

This result should be compared with $F_{2D}(x) \sim (2-x)^{10}$, expected in QCD perturbation theory in the limit $x \to 2$ (Ref. 90). [The behavior of $F_{2D}(x)$ is consistent with the scaling limit of the data of Ref. 39, where, however, the mass of the hadron state which is produced is small.] At present we have data on $G_p^{D/\pi}$ only at[46,50,87] $\alpha \leqslant 1.35$; at $\alpha > 1$, $G_p^{D/\pi}$ falls off approximately in proportion to $(2-\alpha)^9$. We wish to emphasize, however, that the use of the completeness approximation in the kinematics of the experiments of Refs. 46 and 87, especially at $\alpha \geqslant 1.3$, is a dubious procedure because of the limited phase volume (the invariant recoil mass is small).

Comparison of the data on hadron reactions (4.8) at the deuteron with lepton processes shows that reaction (4.8) incorporates direct information on the high-momentum component of the deuteron wave functions over a broad kinematic region.

$A > 2$. We begin by listing the basic features of reactions (4.8) which indicate that we are dealing with a nontrivial phenomenon. The data are reviewed in more detail in Ref. 38.

a) The α dependence for the spectrum of cumulative p's and π's from ^{4}He and from heavier nuclei is essentially the same for $1.3 \leqslant \alpha \leqslant 2$; at $1.3 \leqslant \alpha \leqslant 1.6$, these spectra are similar to the spectra of p's and π's from the deuteron.[46,47]

b) The α dependence for the spectrum of cumulative pions is similar in shape to the x dependence for the structure function of the ^{12}C nucleus.

c) The normalized cross section $G_a^{A/p}(\alpha, p_t)\sigma_{in}^{aN}$ has an identical A dependence for any incident particle a, and within ~ 20–30% it is independent of the species of particles a(a $= \gamma, \pi$,p,ν).

d) The quantity $(1/A)G_a^{A/p,\pi}(\alpha, p_t)$ increases with increasing A, and up to $A \lesssim 12$–20 this increase is identical for p's and π's. At larger values of A, the quantity $(1/A)G_a^{A/p}$ increases significantly more rapidly.[46,47,159]

e) The cross sections for the production of two, three, etc., cumulative nucleons are larger; for the most part, two nucleons are emitted from remote points of the nucleus.[47,154]

f) The relation $G_a^{A/p}(\alpha, p_t) \geqslant 10^2 G_a^{A/\pi}(\alpha, p_t)$ holds, but the α dependences of $G_a^{A/p}$ and $G_a^{A/\pi}$ are similar.[46,47,149]

The spectator mechanism for the production of cumulative nucleons[53] and the theoretical conclusion[10,144] that the Glauber screenings cancel out in the inclusive cross section are crucial to an understanding of the basic behavior here, as we will demonstrate in the particular case of the violation by an initial hadron h of only the binary correlations in a nucleus. The incident hadron h successively disrupts all the nucleon correlations at the impact parameter (Fig. 19). In the average event, the hadron h interacts inelastically with $\nu = A\sigma_{hN}^{inel}/\sigma_{hA}^{inel}$ nucleons. Consequently, the inclusive cross section, i.e., the cross section summed over the emission of spectator nucleons from arbitrary points of the nucleus, is[32)]

$$E_N \frac{d^3\sigma^{h+A\to N+X}}{d^3p_N} \approx \nu\sigma_{hA}^{inel}\varphi_A^2(p_N) = \sigma_{hN}^{inel} A\varphi_A^2(p_N), \qquad (4.13)$$

where φ_A^2 is the probability for a binary nucleon correlation. The reggeon-diagram technique can be used to generalize this result to arbitrary correlations and to incorporate the Glauber screenings caused by scattering by correlations.[38,144] Qualitative differences from the "actual" momentum approximation are seen in the significant emission from heavy and intermediate nuclei of several cumulative nucleons and also in an increase with A ($\sim A^{1/3}$) of the longitudinal dimensions of the nucleon emission region. This prediction[10] is in agreement with the data of Refs. 47 and 154. The picture of a sequential ejection of nucleons also explains the angular dependence of the sticking coefficient for the production of fast deuterons from nuclei.[47]

As a result, the cross section for reactions (4.8) can be written in a form analogous to that for (e,e′) reactions:

$$G_a^{A/h}(\alpha, p_t) = A\sum_{j=2}^{A} a_j(A)\,\sigma_a^{j/h}(\alpha, p_t), \qquad (4.14)$$

with the same coefficients (with an accuracy to within isotopic effects) as in (4.4). It obviously follows from (4.14)—by analogy with the case of reaction (4.1)—that the spectra are similar (Point a in the list above) and that the A dependences of $G_a^{A/p,\pi}$ are identical for light nuclei (Point c). The A dependences of $G_a^{A/p,\pi}(\alpha, p_t)$ and $\sigma^{(e,e')}(x,Q^2)$ are directly related at $\alpha\sim\alpha>1$ [Eqs. (4.13) and (4.8)]. For example, the value $a_2(^4He)\approx 4$ extracted from region (4.1), agrees with the value $a_2^N(^4He)/a_2^p(^3He)\approx 9/5$ estimated in the Wigner model in Ref. 38. We wish to emphasize that (4.14) holds in all models in which reaction (4.8) results from a scattering by short-range nucleon correlations [the models of few-nucleon correlations, 6q(9q,...) bags, and fluctuons]. It does not hold, however, in the mean field models, where it is assumed that the momentum of the fast nucleon is balanced by the momentum of the entire residual nucleus, or in models in which reaction (4.8) is caused by secondary interactions.

For more-detailed calculations we can use the approximation of a few-nucleon correlations for the high-momentum component of the nuclear wave functions. In this approach, the cross sections $G_a^{A/p,\pi}$ have the same form as for the deuteron [expression (4.11)] after the substitution $\rho_D^N(\alpha, p_t)\to\rho_A^N(\alpha, p_t)$. Generalizing (4.13) to incorporate the ejection of nucleons from arbitrary correlations (not simply binary correlations!), we conclude that the inclusive spectrum of cumulative nucleons, $G_a^{A/N}/\sigma_{in}^{aN}$, must have an identical dependence on the atomic number of the nucleus, A, for any particles $a = (\gamma,\nu,\bar{\nu},e,p,\pi)$ (Ref. 10) and even for nucleus-nucleus scattering.[156] In contrast, the empirical hypothesis of nuclear scaling[47,48] predicts a universal A dependence of the multiplicity of cumulative nucleons: $G_a^{A/N}/\sigma_{in}^{aN}$. A universal A dependence of the inclusive spectrum has been found experimentally not only for p, π (Refs. 47 and 48) and γ (Ref. 49) scattering by nuclei but also for nucleus-nucleus scattering.[157] Because of the Glauber screening in the ejection of a cumulative nucleons from a few-nucleon correlation we would expect a completely definite dependence of $G_a^{A/N}/\sigma_{in}^{aN}$ on the species of the particle a (Refs. 145 and 156).

The absolute value of $G_p^{A/\pi}$ and its dependence on α and p_t can be calculated from(4.10), (4.11), and the data on $G_p^{A/p}$; the results agree reasonably well with experiment.[38]

To calculate the dependence of $G_a^{A/N}$ on α and p_t it is convenient to expand ρ_A^N in a series in the contributions of two-nucleon and three-nucleon correlations,[38,94]

$$\rho_A^N(\alpha, p_t) = A\sum_{j=2}^{A} a_j(A)\,\rho_j(\alpha, p_t), \qquad (4.15)$$

and to approximate ρ_j as the convolution of $j-1$ binary correlations. For $\rho_2(\alpha, p_t=0)\sim(2-\alpha)^m$ we then find[38,94] (experimentally and in QCD perturbation theory,[125] we have $m\simeq 3$ over a broad range of α)

$$\rho_j(\alpha, 0) \sim \left(1-\frac{\alpha-1}{j-1}\right)^{m(j-1)+j-2},$$
$$\rho_j(\alpha, p_t) \approx \frac{1}{j-\alpha} F_j\left(\frac{m_N^2(1-\alpha)^2+p_t^2}{\alpha(j-\alpha)}\right). \qquad (4.16)$$

As a result, we calculate an absolute value for $G_h^{A/p}$ which is in reasonable agreement with experiment. We might note that since the nucleus is a rather dense system, the contribution of the j-nucleon correlation to ρ_A^N in the region $\alpha\to 1$ (i.e., at those values of α for which the description of this correlation in terms of nucleon degrees of freedom appears to be qualitatively wrong) is masked by a contribution from a higher order correlation. As a result, internucleon distances $\gtrsim$ 1–1.2 fm are important in most of the phenomena discussed in this section of the paper.[38]

The value $a_2(^{12}C)\sim$ 6–8 found from this analysis agrees with the value found for $a_2(^{12}C)$ from the deep inelastic μA scattering and from photonuclear reactions (see the discussion in Subsection 4b). Analysis of the experimental data on the production of cumulative particles in hadronic processes, above-threshold (e,e′) reactions with large Q^2, etc., shows that the probability (P_j) for a j-nucleon correlation falls off rapidly with j, $P_j\sim(P_2)^{j-1}\sim(1/5)^{j-1}$, reflecting the small value of the probability for an additional nucleon to enter the small volume.

Several alternative models have been proposed for the

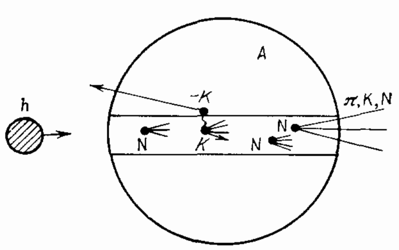

FIG. 19.

production of cumulative particles. These models have been compared in detail with experiment elsewhere.[38,68] Here we will briefly discuss the basic physical ideas. The first group of models are models which assume that the cumulative particles arise from a violation of short-range nucleon correlations in a hard hadronic process; i.e., the impulse approximation is assumed to be valid. These models ignore the spectator mechanism for the production of cumulative nucleons. Examples are the fluctuon model[33)] and the models of 6q, 9q, and 12q bags.[34,45,46,51,55,56,59,158] These models fail to explain (a) the large absolute value of the yield of one, two, etc., cumulative nucleons; (b) the result $p/\pi \gtrsim 10^2$ in the case of cumulative particles; (c) the large value observed experimentally for R_A (the radius of the region from which the two cumulative protons are emitted)[154]; and (d) the large absolute value of the yield of cumulative pions under the assumption that the process is hard. This conclusion is indicated by a comparison of $G_p^{A/\pi}$ at $\alpha_\pi > 1$ with data on $F_{2A}(x,Q^2)$ and on the production of leading pions in the scattering of leptons by hadrons.[58] If we take the Glauber rescatterings and the spectator mechanism into account—so far, this has been done only in the model of few-nucleon correlations—we find that the specific difference between these models reduces to the magnitude of the admixture of nonnucleon degrees of freedom in the nuclear wave functions (cf. Subsections 2a and 3c). The large absolute value of the cross section for the production of cumulative protons seriously limits the possibility of introducing an admixture of nonnucleon components in nuclear wave functions (Subsection 3c). It follows from the set of data discussed above that in the case of light nuclei the nucleon degrees of freedom are responsible for at least half of the high-momentum component of the wave function. In order to resolve this situation; it is necessary to search for the production of cumulative baryon resonances.[161]

It has been assumed in several papers (Refs. 65 and 162, for example) that the production of cumulative nucleons is the result of the absorption of slow secondary pions (with momenta ~ 100–300 MeV/c) by correlated pairs of nucleons in the nucleus. The role played by this mechanism was overestimated in a qualitative way in Refs. 162 and 65, since it is heavy resonances, for the most part, rather than pions which propagate in a nucleus. Many experiments have shown[163] that pions with a Feynman variable $x < 0.2$ are primarily products of the decay of Δ, ω, and ρ. These resonances essentially do not have time to decay inside the nucleus, since the distance traversed by a resonance without decay, the mean free path of the resonance, and the radius of the region in which the resonance interacts with the nucleons of the nucleus (the resonance "formation" region in the elementary reaction) are all ~ 1–1.5 fm. That this mechanism may play a role in the case of nuclei with $A \geqslant 12$ may be indicated by the observation in the reaction $p + A \to \pi + X$ of a dip in the spectrum of pions emitted into the rear hemisphere with momenta ~ 200–300 MeV/c (with $x > 0.4$)[87] (see Ref. 86 for another explanation of this dip, in the spirit of the EMC effect). The cascade model, based on classical mechanics, overestimates the role of cascade processes in a qualitative way. In addition to the important role of heavy resonances, discussed above, that model ignores (a) the formation time of fast hadrons and (b) the fact that quantum mechanics leads to an interaction of slow particles which is qualitatively different from that in classical mechanics (Ref. 164, for example). In particular, for a slow particle (with a momentum $k < k_F$), the probability for an interaction with a nucleon of the nucleus, i.e., the probability to enter a volume of radius $\sim 1/k_F$, is suppressed (in comparison with a calculation based on classical mechanics) because of the wave nature of the particle. As a result, the cross section for the absorption of slow pions by correlated pairs of nucleons in a nucleus is suppressed by a factor k_π^2/k_F^2 in comparison with the cross section found from cascade calculations.[165]

In general, there has not been an adequate theoretical study of a possible role of secondary interactions with the formation of cumulative particles or of the problem of the passage of cumulative particles through nuclear matter. If, for example, cumulative pions and nucleons are produced in a pointlike configuration (cf. Section 2), their secondary interaction with the nucleus will be strongly suppressed. Consequently, the only processes that can be used for a reliable study of nuclear structure at this point are the processes in which cumulative particles are produced in the lightest nuclei, with $A \leqslant 6$–9. To determine the boundary of the A region in which it is possible to study nuclear structure, rather than the spatial picture of the strong interaction, it will be necessary to carry out correlation experiments (see Refs. 10, 38, and 68, for example).

In summary, the methods of high-energy physics have proven their effectiveness in research on short-range correlations in nuclear wave functions. In a next step in the research, these methods may make it possible to study in detail the structure of these correlations. For theoretical nuclear physics, the presence of even crude experimental information on the magnitude and properties of the short-range correlations and on the properties of nuclear forces should have important practical consequences.

The problem of calculating the effective nucleon-nucleon interaction potential in terms of the interaction of nucleons in vacuum has been simplified slightly. Most of the progress in the quantitative description of the consequences of the small internucleon distances in a nucleus has resulted from the use of an expansion of the high-momentum component of the nuclear wave function in the sum of the contributions of short-range correlations, which are of a universal form, independent of the particular nucleus. It is this property which has made it possible to conceal the short-range nucleon correlations in the parameters of a universal effective potential describing the interaction of quasiparticles. In other words, the information which poses the most serious difficulties to calculations by the methods presently available can be extracted from experimental data in a universal form.

In high-energy inelastic processes, one studies the short-range correlations, primarily near the center of the nucleus.[34)] It is thus possible to make a slightly more direct comparison with the theory of infinite nuclear matter here than in low-energy processes.

5. CONCLUSION

1. The present QCD theory of strong interactions indicates a need for a qualitatively new interpretation of the short-range part of nuclear forces on the basis of the quark-gluon structure of hadrons.

2. The derivation of a quantitative theory of nuclear forces on the basis of QCD is a difficult problem, since we lack suitable calculation methods for dealing with the interaction of particles at distances $\geqslant 1$ fm, which are usually important for most nuclear phenomena (the methods of QCD perturbation theory cannot be used). The long-range part of the nuclear forces, for $r \gtrsim 1/m_\pi$, can be described by the exchange of a π meson and, perhaps, the exchange of two π mesons. The existence of a nucleus as a relativistic system of nucleons is a consequence of the presence of a spontaneously broken chiral symmetry in QCD.

3. Recent studies of high-energy processes with a large momentum transfer to the nucleons of a nucleus have proved to be an effective tool for studying short-range nucleon correlations in nuclei. As new numerical methods for solving the many-particle Schrödinger equation are developed, and as progress is made in the analysis of hard nuclear reactions, a detailed test of the basic assumptions on which nonrelativistic nuclear theory is based will become a practical possibility.

4. The methods of conventional nuclear physics, which make use of the concept of quasiparticles, are applicable to phenomena in which the contribution of a high-momentum component is not especially enhanced. In this approach, the effects of small internucleon distances are "concealed" in the parameters of the effective potential. In this sense, QCD does not discredit the existing methods in nuclear theory.

5. The next steps in the research will be (a) to study the quark-gluon wave functions of the nucleus over a broad range of x in deep inelastic scattering of leptons by nuclei, (b) to measure the A dependence for the probabilities of short-range correlations in measurements in (e,e′) reactions at large Q^2, (c) to study the structure of correlations and to seek exotic degrees of freedom in nuclei with the help of (e,e′N(Δ,N*)) reactions and hadronic processes, and (d) to study in detail two nucleon correlations in experiments with polarized deuterons.

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APPENDIX

Geometry of high-energy processes and short-range nucleon correlations

To a large extent, the study of short-range nucleon correlations in nuclei by the methods of high-energy physics is based on the use of a scattering geometry familiar from the Glauber approach. However, the multiparticle nature of high-energy processes and the resulting significant momentum transfer to the nucleons of the nucleus in a typical inelastic process make the external manifestations of short-range nucleon correlations qualitatively different from the typical physics at energies of the order of hundreds of MeV. Our purpose in this appendix is to explain the typical ideas and discussions in the analysis of high-energy processes. A more detailed discussion can be found in the review of Ref. 38.

The simplest example, and the one most important for relativistic nuclear physics, is the deep inelastic scattering of leptons by nuclei:

$$l + \mathrm{A} \to l' + \mathrm{X}. \tag{A.1}$$

All the notation corresponds to Fig. 2. To describe process (A.1) at large Q^2, and at values of x which are not too small, we can use the impulse approximation in a model in which the nucleus is described as a system consisting exclusively of nucleons. The form of the equations, however, is slightly different than in nonrelativistic nuclear physics because of the need to consider recoil. In our case there is no nonrelativistic Schrödinger equation to describe the scattering process, so that equations must be extracted directly from the diagrams of the time-dependent noncovariant perturbation theory. An attempt to calculate $F_{2\mathrm{A}}$ by the method customary in the nonrelativistic theory, in terms of the Schrödinger wave functions of the nucleus, leads to a paradox. To illustrate the situation, we consider scattering by a deuteron. Formally, the impulse approximation corresponds to the diagram in Fig. 20. The amplitude of the elementary process enters off the mass shell in this case. It is useful to calculate the nonconservation of the invariant energy in the upper block between the final state and the intermediate state, Δ:

$$\Delta = [s_f = (q + p_{\mathbf{D}} - p_{\mathrm{sp}})^2] - [s_i = (p + p_{\mathrm{int}})^2]; \tag{A.2}$$

here p_{sp} is the 4-momentum of the spectator nucleon, given by $p_{\mathrm{sp}} = (\sqrt{m_{\mathrm{N}}^2 + k^2}, -k)$; $p_{\mathrm{int}} = (\sqrt{m_{\mathrm{N}}^2 + k^2}\, k)$ is the momentum of the nucleon with which the interaction occurs; s_{f} is the invariant energy calculated from the final state; and s_{i} is that calculated from the intermediate state. Simple calculations yield

$$\Delta = 2q_0\left(M_{\mathbf{D}} - 2\sqrt{m_{\mathrm{N}}^2 + k^2}\right) + M_{\mathrm{D}}^2 - 2M_{\mathrm{D}}\sqrt{m_{\mathrm{N}}^2 + k^2} \to \infty. \tag{A.3}$$

Since the amplitude for the elementary process—the wave function—falls off with increasing Δ, we conclude that this contribution is zero (even in the limit $k \to 0$). Despite the paradoxical nature of this result, it is correct, and it reflects a specific feature of relativistic effects and multiple production. A detailed discussion of these problems goes beyond the scope of the present paper; the reader is referred to Ref. 38. The formal reason for the paradox is that we have not

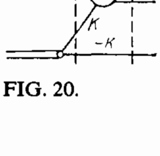

FIG. 20.

considered all the diagrams but only those which are customary in the nonrelativistic theory, where energy is not conserved in the intermediate state. On the other hand, to bring out the important diagrams it is useful to study the equivalent system of diagrams in which Σp_i^- but not $\Sigma p_i^+ (p_i^\pm = \sqrt{m_N^2 + p_t^2} \pm p_{iz})$ is conserved in the intermediate state. We direct the z axis along the momentum of the photon. It is convenient to introduce the light-cone variables $\alpha_i = Ap_i^-/p_A^-$ and k_{it}. The condition for the conservation of $\Sigma_i p_i^-$ is then equivalent to $\Sigma\alpha_i = A$. Here the sum is over all the particles in the intermediate state into which the nucleus A goes. An analogous calculation of the difference $s_f - s_i = \Delta$ yields the following result in this case in the limit $q_0 \to \infty$:

$$\Delta \to M_D^2 \left(1 - \frac{\alpha}{2}\right) - (m_N^2 + k_t^2) \left(\frac{\alpha}{2}\right)^{-1}. \qquad \text{(A.4)}$$

We see that $\Delta_{q0 \to \infty}$ is finite. Furthermore, in the energy denominators which characterize the time evolution of the system we can ignore Δ in comparison with W^2 (Bjorken and Feynman scaling).[35)] We thus find, from simply an analysis of the simplest physical phenomenon, that the wave functions of the nucleus on the light cone must be used in order to describe high-energy scattering. Here we will not discuss the more general arguments for this assertion, which are based on the Gribov-Feynman space-time picture of the strong interaction.[38] Our reward for the correct description of the scattering is a simple relation between the hadronic and quark-gluon degrees of freedom in the nuclear wave function.[36)] This circumstance is familiar in the theoretical work, from the saturation of the sum rules in current algebra and from the success of the dispersion sum rules. The use of nuclear wave functions on the light cone makes it possible to derive directly from Feynman diagrams an impulse-approximation expression for the nuclear structure function F_{2A} in terms of the nucleon structure function $F_{2N}(x,Q^2)$ and $\rho_A^N(\alpha,k_t)$, the density matrix of the nucleons in the nucleus.[132] It is simple to show that at small values of $|1-\alpha|$ and k_t the diagrams which are of importance for $\rho_A^N(\alpha,k_t)$ are the same as the diagrams of nonrelativistic nuclear theory.[38] In this kinematic region, $\rho_A^N(\alpha,k_t)$ can thus be calculated by the standard nonrelativistic nuclear theory.

The correspondence between the dynamics on the light cone and the nonrelativistic nuclear theory can be explained in the example of the interaction of two spinless nucleons. The restiction to nucleon degrees of freedom leads to a Weinberg equation for the amplitude for $N\bar{N}$ scattering off the mass shell:

$$\begin{aligned} T(\alpha_1\ k_{t1},\ \alpha_3,\ k_{t3}) &= V(\alpha_1,\ k_{t1},\ \alpha_3,\ k_{t3}) \\ &+ \int V(\alpha_1,\ k_{t1},\ \alpha_5,\ k_{t5})\,(M_{5,6}^2 - M_{3,4}^2)^{-1} \\ &\times \frac{d\alpha_5}{\alpha_5(2-\alpha_5)}\ \frac{d^2k_{t5}}{(2\pi)^3}\ T(\alpha_5,\ k_{t5},\ \alpha_3,\ k_{t3}); \end{aligned}$$

here $[d\alpha/\alpha(2-\alpha)]d^2k_t/(2\pi)^3$ is the phase volume for the two-nucleon system, α_i and k_t are the light momenta of the nucleons in the initial and final states, $1/(M_{5,6}^2 - M_{3,4}^2)$ is the energy denominator for the two-nucleons state, and

$$\frac{1}{4}M_{5,6}^2 = \frac{m^2 + k_{t5}^2}{\alpha_5(2-\alpha_5)},\quad \frac{1}{4}M_{3,4}^2 = \frac{m^2 + k_{t3}^2}{\alpha_3(2-\alpha_3)}$$

are the invariant masses of the systems of nucleons 5,6 and 3,4. The potential V is the sum of the diagrams which do not have a two-nucleon intermediate state.

The Lorentz invariance of T and V on the mass shell dictates the form of V off the mass shell: $V = V(k^2, k''^2, (k-k'')^2)$. Here k (k'') is the momentum of the nucleon at the center of mass of the initial (intermediate) state, and

$$\alpha_5 = 1 + \frac{k_z''}{\sqrt{m^2 + k''^2}},\quad k_{t5} = k_t''.$$

(To find this restriction it is convenient to transform to the center-of-mass frame of the colliding nucleons, in which T and V are rotationally invariant on the mass shell. Consequently, the terms of the perturbation-theory series for the amplitude for NN scattering in the potential V must be rotationally invariant. From this condition it is a simple matter to derive the general form of the potential V which is given above.) As a result of incorporating the restrictions on the shape of the potential V, we find that the Weinberg equation takes the form similar to that of the nonrelativistic Schrödinger equation:

$$T(k,\ k') = V(k,\ k') + \int \frac{V(k,\ k'')}{(k''^2 - k'^2)}\ \frac{d^3k''}{\sqrt{m^2 + k''^2}}\ \frac{T(k'',\ k')}{(2\pi)^3}. \qquad \text{(A.5)}$$

This equation was discussed in Refs. 38 and 166 in connection with the light dynamics of the deuteron. The restriction derived above on the shape of the potential is a nontrivial (for quantum mechanics on the light cone) manifestation of the rotational invariance of the theory. The simple form of the angular-momentum operator arises as a consequence of the restriction to nucleon degrees of freedom. (The same form was proposed for the angular momentum by Terent'ev[167] in the context of quark models.) In quantum field theory in perturbation theory, because of the nonconservation of the number of particles in the intermediate state, the angular-momentum operator is much more complicated. The reader is directed to Ref. 38 for a discussion of the various questions associated with the many-particle Weinberg equation and its relationship with nonrelativistic nuclear theory.

[1)]All the notation corresponds to Fig. 2; $x = -q^2/2m_N q_0$ is the standard Bjorken variable, q^2 is the square of the 4-momentum, and q_0 is the energy transferred to hadrons.

[2)]A similar change in the characteristics of excitons is expected in a nonideal exciton gas of sufficiently high density.[21] The analogy between a superdense nuclear medium and the physics of the phase transition of a gas of excitons into electron-hole droplets[21,22] may thus prove valuable.

[3)]This value of the momentum component of the wave function of the nucleus is consistent with the calculations which have been carried out in the Brueckner model,[8] where its value is 75–85%.

[4)]As Baldin has pointed out,[86] a qualitatively similar behavior is observed in high-energy hadronic processes for the cross section for fragmentation of pions from nuclei. The dependence on the atomic number for the EMC effect, however, is considerably weaker than that in the case of fragmentation of pions from nuclei, where the cross section at $0.6 < x < 1$ can be parametrized as $A^{2/3 + x/3}$ (Ref. 86; cf. Fig. 4 of the present paper). We do not rule out the possibility that this difference stems from the absorption of pions in the nucleus (Ref. 87, for example). Furthermore, the fragmentation $^4He \to \pi$ has so far revealed an increase in the yield of pions, rather than a decrease, as in Ref. 14.

[5)]Calculations using realistic nuclear wave functions show that (2.7) is a reasonable approximation of (2.5) at $x < 0.6$. We emphasize that the uncertainties in the magnitude of the correction for Fermi motion at $x < 0.6$ are small, since $\langle k_A^2 \rangle$ is known from experiments on the scattering of electrons by nuclei in the quasielastic region.

[6)]Figure 1 incorporates a small correction for the nonidentical p and n numbers in Fe.

[7)]Analysis of experimental data on hydrogen indicates[88a] that at $x > 0.65$ a high-twist contribution will be important in the Q^2 region studied in Ref. 14. At $x > 0.65$, the measurement of $F_{2A}(x,Q^2)$ in Ref. 14 is not expressed directly in terms of the quark distribution in the nucleus, and the Fermi-motion effects according to the estimate of Ref. 88b are substantially weaker than at $Q^2 \sim 100$ GeV^2.

[8)]An admixture of Δ isobars in the nuclear function, at the level ordinarily considered in nuclear physics, does not explain the effect: When charge independence is taken into account, we find that we need a $\sim 30\%$ admixture of Δ isobars.[89]

[9)]This limitation can be relaxed by introducing not only a multiquark component but also a pion component in the nuclear wave function (Refs. 25 and 27, for example).

[10)]Other experimental indications of compressed configurations in hadrons have been discussed in the literature. It has also been suggested that experiments be carried out to search directly for compressed configurations.[98–102,19,20,38]

[11)]This result agrees with the hypothesis of vector dominance, which is ordinarily applied to the γNN vertex, in the sense that as a nucleon is compressed the relative weight of compressed configurations in the photon will increase. We recall that compressed configurations have been observed in the photon in the process $\gamma\gamma \to$hadrons (Ref. 103, for example).

[12)]As J. de Kam and H. J. Pirner have pointed out (1982), effects of this type are important for stabilizing a bag with a pion field.

[13)]This is the range of ΔE values which is required for explaining diffraction processes at nucleons and nuclei.[99,100]

[14)]The swelling of quasiparticle nucleons in nuclei has been under discussion for a long time in nuclear physics; see, for example, the papers by Noble[107] and Kopeliovich and Niedermayer (1981). The characteristics r^* and m^* of this entity[107] are qualitatively different ($r^*/r_{\rm free} \approx m_N/m^* \sim 1.3$, $m_N \approx 0.25$ GeV, from those for a bound nucleon in a nucleus (2.18). We wish to stress that the cross section for deep inelastic scattering of leptons by nuclei cannot be expressed in any simple way in terms of the cross section for scattering by quasiparticles. If we ignore the qualitative difference between a quasiparticle and a bound nucleon in a nucleus, if we use Noble's assumption $r^* \sim 1/m^*$, and if we use (2.18) without the factor of 2/3 and (2.8), then we can fit the EMC data.[108] The cost of ignoring the dependence of the interaction on the radius of the quark-gluon configuration is the value found in Ref. 108 for the nuclear binding energy (per nucleon): $\varepsilon = m_N - m^* - \langle k_A{}^2 \rangle/2m_N = -U_A - 3k_F^2/10m_N \approx 30$ meV.

[15)]The correction to the magnetic moments of the nucleons bound in a nucleus for swelling of a nucleon leads, in the MIT bag model for the nucleon, to an increase in $\mu_{\rm theo}$ by a factor of $r_{\rm bound}/r_{\rm free}$ (Ref. 20) and significantly improves the agreement between the theoretical calculations of μ_d and $\mu_{^3He} + \mu_{^3H}$ with experiment.

[16)]As readers from the field of particle physics are well aware, the renormalizability of the interaction of massive vector mesons requires trilinear and quarternary interactions of vector mesons, the introduction of a scalar particle, and the introduction of a spontaneous-breaking mechanism, which gives vector mesons a mass.

[17)]Although calculations of the absolute values of the hadron form factors in QCD perturbation theory have been unsuccessful, dimensionality estimates of the $Q^2(x)$ dependence of the elastic (inelastic) form factors of hadrons agree fairly well with experiment.

[18)]We wish to thank Yu. M. Makeenko for a discussion of this question.

[19)]Because of the zero isospin of D, it has no ΔN component.

[20)]Here α/A is the fraction of the momentum of the nucleus which is carried off by the Δ isobar in the frame of reference in which the nucleus is fast [see Eq. (4.9)].

[21)]The successful description of the many characteristics of heavy nuclei by the conventional nuclear approaches is difficult to use to find limitations on P_{ex}, since this description is based on an effective NN interaction in whose parameters a large nonnucleon component might in principle be "hidden."

[22)]This result follows from the requirement that the two-nucleon approximation for the wave function of the NN system on the light cone hold for an arbitrary direction of the quantization axis. If the potential in the Weinberg equation violates this condition, then it is not possible to achieve Lorentz invariance for the amplitude for NN$\to$NN scattering on the mass shell.

[23)]The calculations of Ref. 132 with the same wave functions are also successful in describing the elastic form factors of the deuteron[37] up to $Q^2 \sim 4$ GeV^2. In this case, however, the cross section depends strongly on the electric form factor of the neutron, about which little is known experimentally.[38,132–135]

[24)]Bosted *et al.*[136] report that a fit of experimental data on the (e,e') reaction at the deuteron requires a high-momentum component of the deuteron wave function larger by a factor of two or three than in the standard deuteron wave functions. Giofi degli Atti has shown that this conclusion is basically a result of some incorrect approximations used in the data analysis in Ref. 136.

[25)]Since W^2 is bounded in this reaction, a final state interaction can be significant. The impulse approximation using plane waves to describe the final state, as a consequence of the attraction between nucleons, underestimates the cross section. This question requires further study.

[26)]It was assumed in Ref. 140 that the process results from the scattering by 6q and 9q configurations. Since the kinematics of the process is near the threshold, the use of the scaling formulas in Ref. 140 and thus the approximation of completeness are unjustified. It is necessary to take into account the suppression of the cross section which results from the small value of the overlap integral with the two-nucleon, three-nucleon, etc., systems.

[27)]Corrections for the finite separation energy of the nucleons show that relation (4.6) should begin to hold at large values of ν for heavy nuclei.

[28)]R. G. Arnold has informed us that data on electron scattering by aluminum at $x > 1$ are presently being subjected to a special analysis.

[29)]The behavior of $F_{2A}(x,Q^2)$ at $x \gg 1$ had been analyzed previously[59b] in the fluctuon model. The expression derived for $F_{2A}(x,Q^2)$ at $x \sim 1$ in Ref. 59b leads to $a \approx 0.4$. The cumulative number[79] used in the analysis[27,146] of the production of cumulative π mesons depends on $E_{\rm inc}$ and thus leads to a violation of limiting fragmentation.

[30)]By virtue of the kinematics of $\bar{\nu}$N scattering, the relative contribution to the scattering by the nucleon moving opposite to ν, accompanied by the production of a cumulative nucleon, is less than 1% (Ref. 147). This result poses a serious difficulty for mean field models,[60–63] where only this mechanism for the production of cumulative nucleons is considered. The production of cumulative nucleons from correlations not violated by $\bar{\nu}$ is sharply suppressed by the final state interaction (see, for example, Ref. 10).

[31)]With decreasing initial energy, the final state interaction becomes increasingly important, in particular because of two-step processes involving the production of a Δ isobar in an intermediate state.[150,151] This contribution to the cross section is also determined by the probability for binary nucleon correlations, since it is proportional to $|\psi_D(k \sim 0.3$ Gev$/c|^2$. A possible role of this process is indicated by data on pD scattering at $E_p = 1$ GeV (Refs. 152 and 153), where the contribution of the spectator mechanism is a small fraction of the cross section (see Ref. 54, for example). At a higher energy it is possible to avoid the corrections due to these processes by selecting events with the average multiplicity.

[32)]In the case of heavy nuclei, a substantial contribution to the production of cumulative nucleons (but not of pions!) comes from the disruption of short-range correlations by secondary particles.[38] This effect gives rise to an additional factor $\gamma(A)$ on the right side of (4.13). Although the absolute yield of cumulative p and π from heavy nuclei can be described when these phenomena are taken into account, the extent to which the analysis is ambiguous is not clear. With decreasing initial energy, expression (4.13) becomes qualitatively incorrect, since in this case the hadron can produce cumulative nucleons only from the region near the surface of the nucleus. See Refs. 38, 94, and 155 for attempts to estimate the contribution of few-nucleon correlations to the production of cumulative nucleons in pA scattering at $E_p \leqslant 1$ GeV.

[33)]On occasion it is assumed in the fluctuon model[45,51,57,58] and also in the coherent-tube model[159] that partons of the fast nucleus which belong to nucleons at the same impact parameter are coherent. So far, no dynamic mechanism has been proposed for the establishment of such a coherence. In QCD perturbation theory, there is no such effect (see the discussion in Ref. 69), since the Lorentz compression of the fast nucleus in the fragmentation region of the nucleus is cancelled exactly by Lorentz time dilation. We recall that the coherent-tube model runs into qualitative difficulties in a description of hadron-nucleus scattering (Ref. 160, for example).

[34)]In low-energy processes, the distribution of the nucleons near the Fermi surface is studied, while at high energies in elastic hA scattering the distribution of nucleons near the surface of the nucleus is studied.

[35)]The logarithmic violation of scaling which occurs in QCD is inconsequential here.

[36)]The recipe for incorporating $N\bar{N}$ pairs in nuclear wave functions by means of Bethe-Salpeter wave functions—a polular recipe in theoretical nuclear physics—violates the exact sum rules which follow from QCD and makes the Glauber approximation unsuitable for describing high-energy processes.[38]

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