Nuclear-structure effects and non-nucleonic degrees of freedom in deep inelastic lepton–nucleus scattering

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A systematic analysis is given of the contributions to deep inelastic lepton–nucleus scattering made by meson exchange currents, high-momentum components, nucleon off-shell effects, and multiquark nuclear components. The problem of separating the nuclear structure factors in deep inelastic scattering by the lightest nuclei when these nuclei are polarized is investigated.

INTRODUCTION

It is now clear that the notion of nuclei as systems consisting solely of nucleons has restricted validity. Nucleons are composite objects, consisting, according to modern ideas, of quarks that interact through non-Abelian gauge fields: gluons. The idea of spontaneous breaking of chiral symmetry at the quark level has as a necessary consequence the existence of meson fields. The mesons most probably generate the forces that bind the nucleons in the nuclei. The mesonic degrees of freedom must be manifested in nuclei and in nuclear processes at relatively large internucleon separations.

At relatively small nucleon–nucleon separations, the nucleons lose their individuality on account of the identity of the quarks, and we are confronted with virtual multiquark systems. If the states of such systems are represented as superpositions of colorless three-quark clusters (nucleons, Δ isobars, N resonances, etc.), then one may speak of the existence in nuclei of non-nucleonic baryonic components. Thus, we see that the picture of the nucleus is rather complicated and different from the traditional proton–neutron model.

Although the nucleonic component makes the main contribution to the total normalization of the nuclear wave function, the non-nucleonic components may in a number of cases make a contribution (and even the dominant one) to observable processes. In fact, if we wish to obtain reliable information about nuclear "exotics," we must take into account correctly the nucleonic components themselves, that is, we must take into account the part played by the Fermi motion of the nucleons, the interaction of nucleons in the final state, etc. However, nuclear reactions at high energies are often investigated with an unjustifiably simplified treatment of the nucleonic part of the problem, and this leads to incorrect conclusions being drawn about the other degrees of freedom.

In this review, we investigate manifestations of mesonic exchange currents, high-momentum nucleonic components, and nucleon off-shell effects in deep inelastic lepton–nucleus scattering; the part played by nuclear structure in the deep inelastic scattering of polarized leptons; and the manifestation of multiquark components in deep inelastic scattering by nuclei.

Section 1 gives a brief introduction to the problem of deep inelastic lepton–nucleus scattering. It is of the nature of a review and provides the basis for the following exposition.

A self-consistent method for taking into account mesonic exchange currents in deep inelastic lepton–nucleus scattering is considered in Sec. 2. Self-consistency is understood in the sense that the same mesons (with their constants that characterize the meson–nucleon interaction) determine both the Schrödinger equation with an OBE potential for the nuclear wave function and the basic properties of deep inelastic reactions, namely, the structure functions, the contribution of the mesonic exchange currents, binding effects, and the energy sum rule. The problem of taking into account high-momentum nucleonic components in deep inelastic scattering is also investigated.

In Sec. 3 we consider the nuclear-structure effects of the lightest nuclei in deep inelastic scattering of polarized leptons by polarized nuclei. In the final section, we investigate the manifestation of multiquark structure of nuclei in deep inelastic processes.

1. DEEP INELASTIC SCATTERING OF LEPTONS BY NUCLEI

Phenomenology of the $l+T\to l'+X$ reaction

We consider a process $l+T\to l'+X$, where T is a hadronic target [nucleus A, nucleon (proton, neutron) N (p, n), etc.], and X is an unobserved hadronic final state (Fig. 1). Such a process may be due to either an electromagnetic or a weak interaction. To be definite, we shall consider deep inelastic scattering of leptons in an electromagnetic interaction; the generalization to the case of electroweak processes does not present fundamental difficulties. In the lowest order in the electromagnetic coupling constant, the cross section of such a process is determined by an invariant amplitude M:[1–3]

$$|M|^2=L_{\mu\nu}W^{\mu\nu}, \tag{1}$$

where the leptonic tensor $L_{\mu\nu}$ is calculated exactly, while the hadronic tensor $W^{\mu\nu}$, given by the expression

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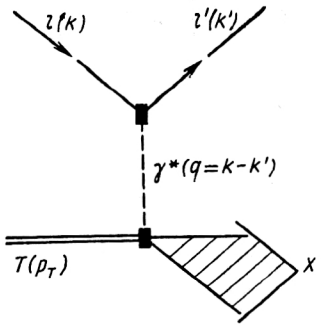


FIG. 1. Deep inelastic scattering of lepton l by a hadronic target T.

$$W^{\mu\nu}(p_T,q)=\sum_X (2\pi)^4\delta^4(p_T+q-p_X)\langle p_T|J^\mu(0)|X\rangle \times\langle X|J^\nu(0)|p_T\rangle, \quad (2)$$

can be parametrized in terms of Lorentz-invariant structure functions $F_{1,2}$ and g_1 (here, we write out only those structure functions that will be discussed in what follows):

$$W^{\mu\nu}(p_T,q,s)=\left[-g_{\mu\nu}+\frac{q^\mu q^\nu}{q^2}\right]\frac{F_1^T(x_T,Q^2)}{M_T} +\left[p^\mu-q^\mu\frac{p_Tq}{q^2}\right]\left[p^\nu-q^\nu\frac{p_Tq}{q^2}\right]\frac{F_2}{\nu M_T^2} +i\varepsilon^{\mu\nu\alpha\beta}q^\alpha s^\beta\frac{g_1^T(p_T,q)}{p_Tq}, \quad (3)$$

where M_T is the mass of the target, p_T and s are its momentum and polarization, $\nu = q_0$ in the target rest frame (laboratory system), and the variable x is determined by

$$x_T\equiv\frac{-q^2}{2p_Tq}=\left(\frac{-q^2}{2M_T\nu}\right)_{\text{lab}} Q^2=-q^2\geqslant 0. \quad (4)$$

It follows from kinematic restrictions that $0\leqslant x\leqslant 1$ for any target, i.e., $F_{1,2}^T(x) = 0$ for $x\leqslant 0$ and $x\geqslant 1$. The two structure functions $F_{1,2}^T$ in (3) are related to the total cross sections for absorption of transverse (helicity ± 1) and longitudinal (helicity 0) photons:

$$\sigma_{\pm 1,0}\approx\varepsilon^{*\mu}_{\pm 1,0}W^{\mu\nu}\varepsilon^\nu_{\pm 1,0}, \quad (5)$$

where $\varepsilon^\nu_{\pm 1,0}$, the polarization vectors of virtual intermediate photons, satisfy the condition $(q\cdot\varepsilon)=0$ and $q_\mu W^{\mu\nu}(p,q) = 0$. Already the early experiments on deep inelastic electron scattering[4–6] showed that Bjorken scaling[7] holds, i.e.,

$$F_{1,2}^T(x_T,Q^2)\to F_{1,2}^T(x_T)\ (x_T,\ \text{fixed and}\ \nu,\ Q^2\to\infty), \quad (6)$$

and that in this limit the well-known Callan–Gross relation is fulfilled:

$$F_2^T(x_T)=2x_TF_1^T(x_T). \quad (7)$$

The main subject of our investigation is the nuclear structure functions, expressed in terms of the structure functions of the constituents (nucleons, mesons,...); the latter are assumed to be given, and in actual calculations are taken either from experiment or from known parametrizations of them.

It follows from the naive picture of the nucleus as a collection of almost free nucleons that

$$F_{1,2}^A(x_A)=\frac{1}{A}\left[ZF_{1,2}^p(x_N)+(A-Z)F_{1,2}^n(x_N)\right], \quad (8)$$

where $x_A = -q^2/2M_A\nu$, $x_N = -q^2/2m\nu \approx x_A\cdot A$, A is the atomic mass of the nucleus, Z is its charge, and m is the nucleon mass. The factor $1/A$ is introduced for convenience ("normalization to one nucleon"); in what follows, we shall always use such a normalization unless stated otherwise. Until 1982, it was assumed that deviations from the formula (8) can be expected at the edges of the kinematically allowed region with respect to x_N, i.e., at very small ($x_N\to 0$), large ($x_N\to 1$), and very large ($x_N\geqslant 1$) values of the scaled variable x.

At very small values, $x\leqslant 0.05$, nuclear screening was predicted:[8] the nucleon structure function measured in the nucleus must be smaller than the structure function measured on an isolated nucleon: $(F_2^{N/A}/F_2^N)_{x\leqslant 0.05}\leqslant 1$.

At large x, when $0.8\leqslant x$, it was assumed[9,10] that there would be a manifestation of effects associated with motion of the nucleons in the nucleus, and also cumulative effects, so that values $x\geqslant 1$ would become possible, and the structure function would not tend to zero as $x\to 1$, and therefore

$$\left(\frac{F_2^{N/A}}{F_2^N}\right)_{x\geqslant 0.8}\geqslant 1. \quad (9)$$

The theoretical predictions in the region of very small x_N recently received experimental confirmation.[11] In the regions neighboring the boundary of the single-nucleon kinematics ($x_N\sim 1$), firm experimental data are not yet available.

The experimental data in the intermediate region $0.1\leqslant x\leqslant 0.8$, obtained by the European Muon Collaboration (EMC),[12] were unexpected. The first qualitative comparison of the structure function for a heavy nucleus (iron) and an "isoscalar" nucleon (in reality, the deuteron) demonstrated an approximately 15% difference in the region of intermediate x_N, where the Fermi-motion effects could have been expected to be unimportant. The "EMC effect" at $0.3\leqslant x_N\leqslant 0.8$ was confirmed by SLAC data[13] and later in other experiments,[11,14–17] including ones made with other nuclei.

In the intervening period, quite extensive material has been accumulated on the x, Q^2, and A dependences of deep inelastic lepton–nucleus scattering (see Refs. 15, 18, and 19). During the last five years, these experimental results have stimulated active theoretical study of non-nucleonic degrees of freedom in nuclei.

Parton model: application to deep inelastic scattering by nuclei

The parton model[7,20] gives a very simple and transparent picture of deep inelastic scattering. In this model, the partons possess the same quantum numbers as the quarks

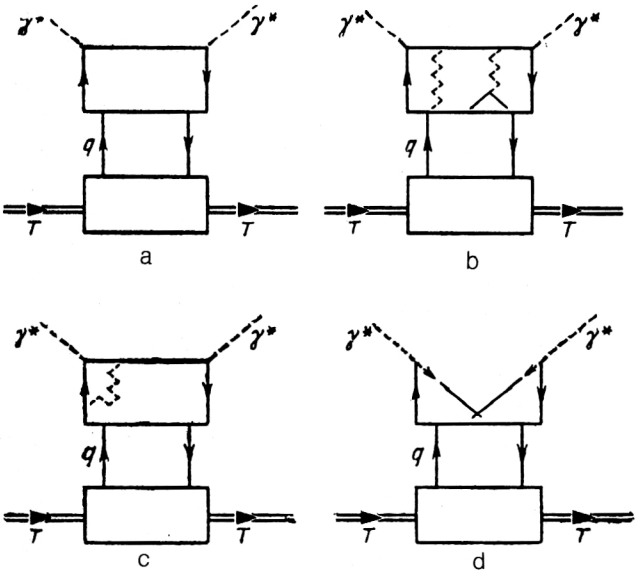


FIG. 2. Some contributions to the amplitude of forward Compton scattering in QCD: a) parton-model diagram; b) final-state interaction; c) vertex corrections; d) interference.

and are usually identified with them; the model is therefore often also called the quark–parton model. The model has been justified in QCD in the framework of the operator expansion on the light cone[1,2,21] (in the limit of free quark fields).

In the quark–parton model, it is convenient to investigate in place of the tensor $W_{\mu\nu}$ the amplitude of virtual Compton scattering $T_{\mu\nu}$, which is related to $W_{\mu\nu}$ by the optical theorem[19,20,22] [see also (5)]. Characteristic diagrams that contribute to $T_{\mu\nu}$ are shown in Fig. 2. They can be classified in accordance with the emission and interaction of quark lines (Fig. 2). In the parton model, it is assumed that at large Q^2 the currents (but not the states!) can be treated in the free theory. Therefore, the final-state interaction (Fig. 2b) and the vertex corrections (Fig. 2c) can be ignored. The contributions of the type shown in Fig. 2b are usually called the corrections of the "higher twists" (in the light-cone operator expansion, they appear as the result of allowance for operators of higher twists[2,19]). The contribution from the diagram in Fig. 2d decreases rapidly, as a power in Q^2, and is therefore ignored. There remains only the contribution of the diagram in Fig. 2a, which represents the incoherent elastic scattering by each quark of a target T. It is important to remember that the parton model is approximate, and it follows from QCD that quantities that depend only on the scaled variable x may also depend logarithmically on Q^2 if the quark–gluon interaction is taken into account by perturbation theory.

In the parton model, the hadronic tensor $W_{\mu\nu}$ (2) can be found from the Compton amplitude (Fig. 2a) with on-shell particles in the intermediate state. If the mass of the partons and their transverse momenta are ignored, then for the structure functions $F_{1,2}^T(x_T)$ we obtain the expressions

$$F_1^T(x_T)=\tfrac{1}{2}\sum_q e_q^2 q^T(x_T);\quad F_2^T(x_T)=\sum_q e_q^2 x_T q^T(x_T), \tag{10}$$

where in the parton model x is the fraction of the longitudinal momentum of the target T carried by the parton. The function $q_i^T(x_T)$ is the density of the probability that the parton momentum is x_T. The index of summation ranges over all parton species ($q=u,\ \bar{u},\ d,\ \bar{d},\ s,\ \bar{s},\ c,\ \bar{c},\ldots$), and e^2 is the square of the parton charge in units of the electron charge.

We note that in the parton model the Callan–Gross relation (7) holds exactly.

At the present time, the quark distributions $q(x)$, the existence of which is postulated in the parton model, cannot be exactly calculated in QCD. There exist only model approaches, of which the best known are the so-called statistical models such as those of Refs. 23–25. Information about $q(x)$ can also be obtained from sum rules of various types, asymptotic dynamical estimates, analysis of the symmetries of the interaction, and, of course, experiments.

The quark distribution functions q can be separated into the contribution q_v of the valence quarks and the contribution q_s of the continuum (sea) of quark–antiquark pairs, i.e., $q(x)=q_v(x)+q_s(x)$. Regge analysis and quark counting rules[2,3,26] permit estimates of the behavior of q_v and q_s at the boundaries of the kinematic region of x ($x\to 0$, $x\to 1$). For example, for the nucleon (N) and meson (M), this analysis gives

$$\left.\begin{aligned} &q_v^N(x)\sim x^{-1/2}(1-x)^3,\ q_s^N\sim x^{-1}(1-x)^7,\\ &q_v^M(x)\sim x^{-1/2}(1-x)^1,\ q_s^M\sim x^{-1}(1-x)^4. \end{aligned}\right\} \tag{11}$$

The asymptotic estimates (11) are approximate and are to be regarded as a first (qualitative) approximation. It follows from (11), for example, that in the region of intermediate and large x the valence quarks are dominant, but in the region of small x the "sea" quarks dominate. Group-theoretical analysis of the quantum numbers of the hadrons and quarks makes it possible to establish the quark composition of the hadrons and, therefore, the quark structure of the parton distributions. The quark distribution functions must satisfy certain exact relations that follow from current algebra.[3,27] For example, for valence quarks, protons, and neutrons (p, n):

$$u(x)\equiv u^p(x)=d^n(x),\quad d(x)\equiv d^p(x)=u^n(x), \tag{12}$$

$$\int_0^1 dx u_v(x)=2,\quad \int_0^1 dx d_v(x)=1, \tag{13}$$

where the relations (12) are a direct consequence of isotopic invariance, and (13) expresses the law of charge conservation. Integral relations of the type (13) are called sum rules for the quark distributions.

From the physical meaning of the parton model, it is possible to write down an exact sum rule for the total momentum of the hadron T:

$$\langle x\rangle_T\equiv\sum_q\langle x_q\rangle_T+\langle x_g\rangle_T, \tag{14}$$

$$\langle x_q\rangle_T \equiv \int_0^1 dx x q^T(x), \quad \langle x_g\rangle \equiv \int_0^1 dx x g^T(x), \quad (15)$$

where $g^T(x)$ are the distribution functions of the gluons in the target T (not to be confused with the spin structure functions); the integrals (15) have the meaning of the total fraction of the target momentum carried by the partons.

To what extent does everything said above about the parton model apply to deep inelastic scattering by nuclei? So far, we have not yet in any way particularized the type of hadronic target, and therefore all the above results are also fully valid for nuclei. In particular, the structure functions of nuclei in deep inelastic lepton scattering are, in accordance with the basic expressions (10) of the parton model, determined by the parton distribution functions in the nucleus, and these have properties of the type (12)–(15). Operating with the "naive proton–neutron model" [see (8)], one could suppose that the quark distributions in the nucleus are obtained by simple combination of the distributions in the free nucleons. The validity of such an assumption was put into doubt by the discovery of the EMC effect, which in the language of the parton model means that, compared with free nucleons, the quarks in the nucleons of nuclei possess large values of x (the region of valence quarks) with a reduced probability; conversely, there is a somewhat enhanced probability for them to have a small value of x (the region of sea quarks). "Transfer" of some of the momentum of the nuclear nucleons also occurs in the region of large x (due to nonrelativistic Fermi motion of the nucleons), but this effect has little influence in the region of intermediate x because of the smallness of the absolute value of the structure functions at large x.

The currently existing theoretical approaches to the description of the nuclear structure functions F_2^A and, hence, of the parton distributions, with the explanation of their A dependence, can be nominally divided into two directions.[19] The first assumes a change of the quark distributions of nuclear nucleons due to the possible change in the nucleus of the conditions of Q^2 evolution of the nucleonic structure functions (Q^2 rescaling[28–30]). At the same time, the intranuclear motion of the nucleons is effectively ignored. At the present time, the rescaling parameter can be neither calculated theoretically nor determined from independent experiments. It is therefore chosen by the condition of best fitting of experiments (EMC, BCDMS, SLAC,...). In contrast, the second direction explains the difference of F_2^A for different A by correct allowance for the internuclear motion of the nucleons and the fact that they are off shell, leaving the quark distributions of the nucleons themselves in the nucleus unchanged. At the same time, allowance is made for the fact that the properties of the nuclear quasiparticles, the nucleons, differ from those of free nucleons. In particular, bound nucleons possess an effective mass $m^* \neq m$, which depends on the energy of the shell. This leads to a renormalization of the scaled variable $x_N \to m^*/m x_N$ (x rescaling[31–33]). In principle, there are no free parameters in this approach—all the parameters are fixed by nonrelativistic nuclear theory.

It appears very probable that the two approaches to the description of deep inelastic scattering by nuclei are "genetically" related, i.e., describe the same phenomenon—the change in the properties of a nucleon in a nuclear medium—but in "different languages." One thing is not in doubt—the nuclear environment significantly influences the properties of nucleons. The advance toward understanding of the nature of these effects is from two sides: from the side of nuclear-physics ideas and from the side of field-theoretical ideas about the interaction of quarks and gluons in hadrons and nuclei.

In this review, deep inelastic scattering of leptons by nuclei is considered from the position of nuclear theory, and therefore we discuss here in more detail the main ideas of approaches of this type. Assuming that we are dealing with an isoscalar nucleus with charge $Z = A/2$ (it is indeed such nuclei with which we shall mainly deal in what follows: ^{2}D, ^{12}C, ^{56}Fe), and proceeding from the probability interpretation of the parton model, we can write down the following relations for the distribution of the quarks in the nucleus with respect to the fraction of the longitudinal momentum:

$$\frac{1}{A} q^A(x_N) = \frac{M_A}{m} \int_0^1 dz \int_0^{y_{\max}} dy q^N(z) f^{N/A}(y) \delta(x_N - zy)$$

$$= \frac{M_A}{m} \int_{x_N}^{y_{\max}} \frac{dy}{y} q^N\left(\frac{x_N}{y}\right) f^{N/A}(y), \quad (16)$$

where we have introduced the notation

$$\frac{1}{A} q^A(x_A) = \frac{1}{A} q^A\left(\frac{m}{M_A} x_N\right) \Rightarrow \frac{1}{A} q^A(x_N). \quad (17)$$

The factor M_A/m in (16) arises because of the difference between the normalizations of the fluxes in the cross sections of deep inelastic scattering by a nucleus and by the nucleon. In (16), the function $f^{N/A}(y)$ is the probability of finding in a nucleus A a nucleon N with fraction y of the longitudinal momentum of the nucleus; the δ function ensures momentum conservation. Under the redefinition (17), the kinematically allowed domain of definition of the functions $q^A(x_A)$, which is $0 \leqslant x \leqslant 1$, goes over to $0 \leqslant x \leqslant y_{\max} = M_A/m$. In what follows, for brevity, we shall denote the variable x_N simply by x. Formally, sum rules of the type (13)–(15) can be rewritten in the form (the notation is obvious)

$$\frac{1}{A} \langle q^A\rangle = \langle q^N\rangle \langle f^{N/A}\rangle, \quad \frac{1}{A} \langle x^A\rangle = \frac{m}{M_A} \langle x\rangle_N \langle y\rangle_{N/A}. \quad (18)$$

From this there follow normalization conditions for the momentum distribution of the nucleons:

$$\langle f^{N/A}\rangle = 1, \quad \frac{Am}{M_A} \langle y\rangle_{N/A} = 1. \quad (19)$$

In studies of Vagradov's group,[31,32] it was shown that to calculate the nuclear structure functions it is necessary to take into account the fact that the nucleons are bound (off-shell effects), especially the contribution of deep quasi-

one-particle levels of collective nature. In accordance with (16), the structure function $F(x)$ is determined by the expressions

$$F_2^A(x)=\int_x^{y_{\max}} dyF_2^N\left(\frac{x}{y}\right)f^{N/A}(y), \tag{20}$$

$$f^{N/A}(y)=\int d^4pS_A(p_0,p)\left(1+\frac{p_3}{m}\right)\delta\left(y-\frac{pq}{m\nu}\right), \tag{21}$$

where $S_A(p_0,p)$ is a "spectral function," which depends on the momentum p and the nucleon energy p_0. It is obvious that the distribution function $f^{N/A}$ satisfies the condition of normalization (19) to the total charge. The momentum conservation law for the nucleus has the form

$$\int_0^{y_{\max}} dyyf^{N/A}(y)=\int d^4pS_A(p_0,p)\left(1+\frac{p_3}{m}\right)\frac{p_0+p_3}{m}$$

$$=1+\frac{\langle\varepsilon\rangle}{m}+\frac{\langle p_3^2\rangle}{m^2}, \tag{22}$$

where $p_0 = m + \varepsilon$ is the energy of the nucleon on which the deep inelastic scattering process occurs; ε is the nucleon separation energy (here, we have ignored the kinetic energy of a heavy residual nucleus; in the case of the lightest nuclei, such as the deuteron, this contribution must be taken into account); and $\langle ...\rangle$ denotes averaging with the spectral function $S_A(p_0,p)$. The averaged quantities on the right-hand side of (22) depend on the choice of the spectral function. For example, calculation of $\langle\varepsilon\rangle$ with allowance for deep quasi-one-particle levels gives[31,33,34]

$$\langle\varepsilon\rangle\approx -40 \text{ MeV}. \tag{23}$$

For $\langle p_3^2\rangle/m$ we have the estimate $\langle p_3^2\rangle/m = \langle p^2\rangle/3m = (2/3)T_k$, where T_k is the kinetic energy of the nucleons ($T_k \approx 18$ MeV). Substituting all quantities in the expressions (19) and (22), we find that $f^{N/A}$ as defined by (21) with allowance for the off-shell effects does not satisfy the normalization condition (19), the scale of the discrepancy being 2–3%.

This violation of the energy sum rules is merely apparent. What we must have here is only an incomplete description of the nucleus. Indeed, in this description the important thing is the strong binding of the nucleons in the nucleus ($\langle\varepsilon\rangle_s \approx -40$ MeV), but as yet we have said nothing at all about the nature of the binding itself. At the present time, it is widely accepted that the NN potential is due to the exchange of mesons.[35,36] Thus, the nucleus is a system of nucleons and other virtual constituents of the nucleus (mesons), which must be taken into account in the nuclear structure functions.

In the parton model, the contribution of the mesonic degrees of freedom to the nuclear structure functions can be expressed in the form

$$\delta F_2^M(x)=\sum_M\int_0^{y_{\max}} dyF_2^M(x/y)f^{M/A}(y), \tag{24}$$

where F_2^M is the structure function of the meson M; $f^{M/A}$ is the distribution of the mesons M in the nucleus A with respect to the fraction of the longitudinal momentum; and the summation in (24) is over all meson species. Since the mesons do not carry baryonic charge, the normalization (18) of the nuclear structure function to the total charge is unchanged when the mesons are taken into account. In the energy sum rule there appears a meson correction (the momentum of the nucleus is carried by all the constituents), and (19) becomes

$$\frac{Am}{M_A}[\langle y\rangle_{N/A}+\langle y\rangle_{M/A}]=1, \tag{25}$$

where $\langle y\rangle_{M/A}=\Sigma\int_0^{y_{\max}}dyyf^{M/A}(y)$.

In the early studies,[37–42] the contribution of the mesonic degrees of freedom were essentially taken into account phenomenologically in the description of the experiments on deep inelastic scattering by nuclei, since the parameters of the models were determined from the same experiments. Therefore, although the introduction of mesonic degrees of freedom did make it possible to establish (in principle) an energy sum rule, the question of the absolute magnitude of the mesonic corrections to the cross section for deep inelastic scattering and of their ability to recover the energy sum rule remained open in such an approach.

A qualitative jump in the understanding of the part played by nuclear structure in deep inelastic scattering was made in Refs. 43 and 44, which established a connection between the nucleon off-shell effects and the meson corrections, the contribution of the mesonic currents to the nuclear structure function was directly calculated, and the sum rules (18) and (25) were analyzed. This approach will be considered in Sec. 2.

To complete this section, we discuss nuclear sum rules at large values of x. The nucleon–meson picture of the nucleus considered above determines the behavior of the nuclear structure functions in a restricted range of x—from zero to unity and, perhaps, slightly above unity because of the Fermi motion of the nucleons within the nucleus. However, it follows from (4) that, in principle, the nuclear structure functions are defined in the kinematic range $0\leqslant x\leqslant A$. Does the Fermi motion of the nucleons exhaust the contributions to the structure functions at $x\geqslant 1$? Do there exist other sources of a high-momentum component? Investigation of cumulative hadron–nucleus processes[45,46] gives an indication of the existence of a non-nucleonic high-momentum component.

If it is assumed that not only nucleons and mesons can exist in a nucleus but also clusters with a large number of quarks,[47–51] then the nuclear structure function can be represented in the form[48,52]

$$F_2^A(x)=\left[1-\sum_k P_k\right]\bar{F}_2^A(x)+\sum_k P_kF_2^{(3kq)}(x/k), \tag{26}$$

where $\bar{F}_2^A(x)$ is the nuclear structure function with allowance for all degrees of freedom (nucleonic, mesonic,...) except the multiquark degrees; P_k is the probability of find-

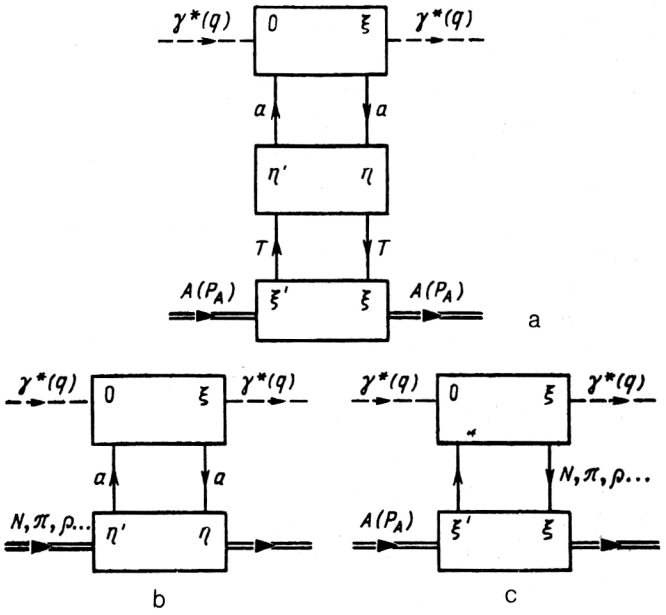


FIG. 3. Convolution model of deep inelastic scattering by nuclei: a) two stages of the convolution model (the nucleus emits constituent T, which then emits a parton a); b) structure function of target in the parton model; c) deep inelastic scattering by a nucleus in the convolution model (the upper block corresponds to deep inelastic scattering by a nuclear constituent).

ing a configuration of $3k$ quarks in the nucleus; and $F_2^{(3kq)}$ is the structure function of the $3k$-quark configuration.

Since the mechanism of electromagnetic scattering of leptons is well known, experimental studies of deep inelastic lepton–nucleus scattering can assist us to a much more accurate knowledge about the nature of the high-momentum component of the quark–parton structure function of nuclei. Some possible effects in different deep inelastic scattering processes to which the existence of multiquark configurations in nuclei can lead are considered in Sec. 4.

Convolution model

There is an intuitively transparent (and easy to use) approach according to which deep inelastic scattering by a nucleus is treated as incoherent scattering by nuclear constituents T (nucleons, pions, Δ isobars, etc.). It is known in the literature as the convolution model[19,22,53] (Fig. 3):

$$F_2^A(x)=\sum_T \int_0^{y_{\max}} dy F_2^T(x/y) f^{T/A}(y). \tag{27}$$

Generally speaking, the convolution approximation is not obvious. *A priori*, there are no reasons for neglect of the nonconvolution contributions (see, for example, Fig. 4), namely, a large q^2 "flows" only through the upper vertex, and therefore only ξ^2, and not $(\eta-\eta')^2$ and $(\xi-\xi')^2$, tends to zero. However, it was shown in Refs. 19 and 53 that the convolution picture does have a dynamical justification, and that deep inelastic scattering of leptons by nuclei can be regarded as a sum of incoherent scatterings on the constituents of the nuclear target and that nonconvolution corrections of the higher twists (Fig. 4) can be ignored. In any case, we shall here adhere to just such a point of view, though other approaches do exist.[22,53]

To calculate the graphs in Fig. 4a, we need explicit expressions for the amplitudes for scattering of virtual photons by the nuclear constituents. They can be obtained by using the philosophy of the light-cone operator expansion.[54]

The hadronic tensor $W_{\mu\nu}^T(p_T,q)$ is related to the amplitude of virtual Compton scattering $T_{\mu\nu}^T(p_T,q)$, which is the matrix element of a certain operator $\hat{t}_{\mu\nu}^T(q)$ with respect to the target state $|p_T\rangle$. To investigate the properties of this operator $\hat{t}_{\mu\nu}^T(q)$, we write the expression for it in terms of the commutator of two hadronic currents,[19,22] obtaining this expression from (2) by means of the completeness condition and the conservation laws in the intermediate states $|X\rangle$:

$$\hat{t}_{\mu\nu}^T(q)=\mathrm{i}\int d^4x e^{\mathrm{i}qx}[J_\mu^T(x),J_\nu^T(0)]. \tag{28}$$

The nucleonic operator $\hat{t}_{\mu\nu}^T(q)$ is averaged over the spin with the wave function of the nucleus and therefore de-

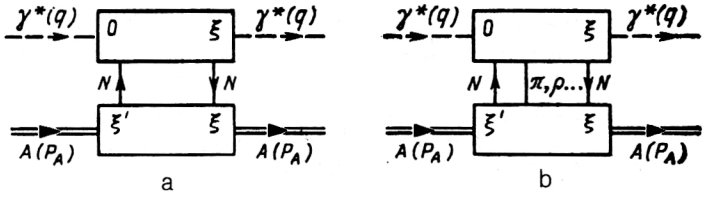


FIG. 4. Convolution (a) and nonconvolution (b) contributions to deep inelastic scattering by a nucleus.

pends on spinor indices.[43,55,56] The operator expansion makes it possible to represent $\hat{t}^T_{\mu\nu}(q)$ as a quadratic form in the nucleonic fields:

$$\hat{t}^N_{\mu\nu}(q)=\sum_n C_n(q^2)\frac{2^n q_{\mu 1}\dots q_{\mu n}}{(-q^2)^n}\{\bar{N}(0)\gamma^{\mu 1}\overleftrightarrow{\partial}^{\mu 2}\dots\overleftrightarrow{\partial}^{\mu n}N(0)\}, \tag{29}$$

where $\{\dots\}$ denotes symmetrization and subtraction of all traces over $\mu_1\dots\mu_n$ in the corresponding operators. The coefficients of the expansion (29) are related to the unpolarized amplitude $T^N_{\mu\nu}(p,q)$:

$$\sum_n C_n(q^2)\frac{\omega^n}{m}=T^N_{\mu\nu}(p,q),\quad \omega=\frac{2pq}{-q^2}. \tag{30}$$

The expression (29) is obtained in the approximation of the leading twist, and its spinor structure ($\sim q_\alpha\gamma^\alpha$) is identical to the result of Ref. 43 for a polarized nucleon:

$$\langle N|\hat{t}^N_{\mu\nu}(q,s)|N\rangle=\frac{m}{E}T^N_{\mu\nu}(p,q)\frac{\bar{u}q_\alpha\gamma^\alpha u}{pq}m,\quad \bar{u}u=1. \tag{31}$$

Going over in (29) to the momentum representation, and substituting (30), we obtain

$$\begin{aligned}\hat{t}^N_{\mu\nu}(q)=\sum_s\int\frac{dp_1dp_2}{(2\pi)^6}\sqrt{\frac{m^2}{E_1E_2}}T^N_{\mu\nu}(p_2,q)\\ \times\frac{m}{p_2q}[a^+(p_1,s_1)\bar{u}(p_1,s_1)\\ +d(p_1,s_1)\bar{v}(p_1,s_1)]\hat{q}[a(p_2,s_2)u(p_2,s_2)\\ +d^+(p_2,s_2)v(p_2,s_2)],\end{aligned} \tag{32}$$

where a and a^+ are the nucleon annihilation and creation operators. The expression for the pions has the form

$$\begin{aligned}\hat{t}^\pi_{\mu\nu}=\int\frac{dk_1dk_2}{(2\pi)^6}\sqrt{\frac{m^2}{\omega_1\omega_2}}[b^+(k_1)b(k_2)\\ +b^+(k_2)b(k_1)]T^\pi_{\mu\nu}(k_2,q).\end{aligned} \tag{33}$$

Similar expressions can be obtained for the other mesons.

2. MESONIC DEGREES OF FREEDOM AND FERMI MOTION OF NUCLEONS IN DEEP INELASTIC LEPTON–NUCLEUS SCATTERING

We consider a self-consistent method of allowing for mesonic exchange currents in deep inelastic scattering of leptons by nuclei. Self-consistency is understood in the sense that the same mesons (with their constants that characterize the meson–nucleon interaction) determine both the Schrödinger equation with an OBE potential (for example, in the potential of the Bonn group,[35] these are the π, ω, σ,... mesons) for the wave function of the nucleus and the basic characteristics of deep inelastic scattering: the structure functions, the contribution of the mesonic exchange currents, binding effects, and the energy sum rule.

Nonrelativistic reduction of the equations of motion of the meson–nucleon system

A procedure of nonrelativistic reduction of the system of equations of interacting fields and separation of the contribution of mesonic exchange currents in various elastic and quasielastic processes has been developed by many authors.[57,58] Despite the differences between the approaches and employed approximations, the final results are similar. The general prescription for calculating the matrix elements of any operator $\hat{O}$ between the states of the nucleus is as follows: a) passage to the nonrelativistic limit in the equations of motion of the interacting mesonic and nucleonic fields and the derivation of an equation for the wave function of the nucleus; b) representation of the covariant relativistic expression for the operator $\hat{O}$ in terms of the nucleonic and mesonic fields; c) nonrelativistic reduction of the operator $\hat{O}$ and calculation of the corresponding matrix element.

As an example, we consider the method developed in Ref. 59 to describe elastic and quasielastic electron–deuteron scattering. For simplicity, we restrict ourselves to inclusion of the contribution of pseudoscalar isovector mesons (pions), which is fundamentally important for deep inelastic scattering by the deuteron. The calculations for the other mesons (vector and scalar) can be made similarly. The classical equations of motion of the nucleons interacting with the pions have the form

$$(i\hat{\partial}-m)N(x)=ig\gamma_5\Phi(x)N(x); \tag{34}$$

$$(\Box+\mu_\pi^2)\Phi_i(x)=-ig\bar{N}(x)\gamma_5\tau_iN(x), \tag{35}$$

where $N(x)=\binom{f(x)}{\phi(x)}$ and $\Phi(x)$ are the nucleon and pion fields, respectively [$\Phi(x)\equiv\tau_i\Phi^i(x)$, $i=1,2,3$]; m and μ are the corresponding masses. The matrix equation (34) makes it possible to express the small component $\Phi(x)$ of the Dirac bispinor $N(x)$ in terms of the large component $f(x)$ in the leading order of the expansion in powers of $1/m$:

$$\Phi(x)=-\frac{i}{2m}(\vec{\sigma}\overleftarrow{\partial}+g\Phi(x))f(x). \tag{36}$$

Using the nonrelativistic relation (36) to eliminate the small components $\Phi(x)$ in the nucleon spinors $N(x)$, we arrive at an equation of motion for the large components $f(x)$:

$$\begin{aligned}if(x)=-\frac{\Delta}{2m}f(x)+ig(\vec{\sigma}\cdot\vec{\partial}\Phi(x))f(x)\\ +\frac{g^2}{4m^2}\Phi^2(x)f(x),\end{aligned} \tag{37}$$

this equation determining, in principle, the nonrelativistic Hamiltonian of the system. The functions $f(x)$ do not have a probability interpretation[60] because of the violation of the normalization conditions and the sum rules for the baryon charge, which must hold if we wish to interpret $f(x)$ as a canonical field. Indeed, the covariant expression for the conserved baryon number and conserved isospin of the meson–nucleon system has the form

$$Q^\alpha = Q_N^\alpha + Q_\pi^\alpha = \int d^3x N^+(x)\tau^\alpha N(x) + \varepsilon_{\alpha\beta\gamma}\cdot\int d^3x \Phi^\beta(x)\Phi^\gamma(x), \tag{38}$$

where $\alpha = 0, 1, 2, 3$, $\tau^\alpha = \hat{I}$, and $\varepsilon_{0ik} = 0$. In accordance with Eq. (36), the nucleon contribution Q_N^0 must be rewritten in the form

$$\begin{aligned} Q_N^\alpha &= \int d^3x[f^+(x)\tau^\alpha f(x) + \phi(x)\tau^\alpha\phi(x)] \\ &= \int d^3x\Big\{f^+(x)\tau^\alpha f(x) + \frac{1}{4m^2}[f^+(x) \\ &\times(\overleftarrow{\sigma\partial})\tau^\alpha(\overrightarrow{\sigma\partial})f(x) + g\Phi(x)f^+(x)((\overleftarrow{\sigma\partial})\tau^\alpha \\ &+ \tau^\alpha(\overrightarrow{\sigma\partial})) + g^2 f^+(x)f(x)\Phi(x)\Phi(x)]\Big\}, \end{aligned} \tag{39}$$

whereas in the nonrelativistic case Q_N^α must be expressed solely by the first term in (39). Therefore, we redefine $f(x)$ in terms of a new field

$$\psi(x) = (\hat{I} + \hat{F})f(x), \tag{40}$$

which must satisfy the usual normalization condition for canonical second-quantized fields:

$$Q_N^\alpha = \int d^3x\psi^+(x)\tau^\alpha\psi(x). \tag{41}$$

The condition (41) determines the operator $\hat{F}$, and by direct calculation we find

$$\hat{F} = -\frac{\Delta}{8m^2} - \frac{g}{4m^2}\vec{\sigma}\cdot\vec{\partial}\Phi(x) + \frac{g^2}{8m^2}\Phi^2(x). \tag{42}$$

The equation of motion for $f(x)$ (37) and the definitions of $\psi(x)$ in accordance with (40) and (42) lead to an equation of motion for the field $\psi(x)$, which in the leading order of the expansion in powers of the operator $\hat{F}$ is identical to the Schrödinger equation with a Yukawa potential,

$$i\dot{\psi}(x) = -\frac{\Delta}{2m}\psi(x) - \frac{g^2}{4m^2}\int d^3y\partial_x Y \times(x-y)\partial_y(\psi^+(y)\tau_i\sigma\psi(y))\tau^i\sigma\psi(x), \tag{43}$$

where the single-time Green's function $Y(x-y)$ is determined from (35) with the condition of absence of pion fields in the initial and final states:

$$Y(x) = \frac{\exp(-\mu_\pi x)}{4\pi x}. \tag{44}$$

In the second-quantization representation, $\psi^+(x)$ and $\psi(x)$ in (43) are nucleon creation and annihilation operators [in contrast to the exact relativistic equations (34) and (35), antinucleons are here absent], and any many-nucleon state can be obtained by means of them. Analogous (relativistic) creation and annihilation operators also exist for the meson fields. Hitherto, we have nowhere particularized the nucleus that is considered. Thus, everything said above applies equally to all nuclei. The generalization to the case of isoscalar, scalar, and vector mesons is analogous. To determine the ground state $|A\rangle$ of the nucleus, we need, in the first place, the nonrelativistic Hamiltonian of the system: $H = H_0 + H_{\text{int}}$. Using the well-known covariant expressions and the relations (36), (40), and (42), we obtain the required Hamiltonian:

$$\begin{aligned} H_0^\pi &= \frac{1}{2}\int d^3x\{\partial\Phi^i(x)\partial\Phi^i(x) + \dot{\Phi}^i(x)\dot{\Phi}^i(x) \\ &+ \mu_\pi^2\Phi^i(x)\Phi^i(x)\}; \\ H_0^N + H_{\text{int}} &= \int d^3x\left\{\frac{1}{2m}\partial\psi^+(x)\partial\psi(x) + m\psi^+(x)\psi(x)\right\} \\ &- \int d^3x\left\{\frac{g}{2m}\psi^+(x)(\sigma\partial\Phi(x))\psi(x) \right. \\ &\left. + \frac{g^2}{2m}\psi^+(x)\psi(x)\Phi^2(x)\right\}. \end{aligned} \tag{45}$$

We note that for this Hamiltonian the Fock vacuum is the true vacuum and that this Hamiltonian commutes with the pion number operator. Because of this, the physical mesons are (in our model) identical to the free mesons, and we can expand $|A\rangle$ with respect to eigenvectors of the pion number operator (Tamm–Dancoff method). The corresponding expansion coefficients are wave functions that determine the probability amplitude for finding in the nucleus A "bare" nucleons (without a pion cloud) and a certain number of pions:[61]

$$|A\rangle = \sqrt{1-Z}\alpha_0^A|N_1N_2...N_A\rangle + \alpha_1^A|N_1N_2...N_AM\rangle + \alpha_2^A|N_1N_2...N_AMM\rangle + \dots, \tag{46}$$

where the renormalization constant Z is determined from the condition $\langle A|A\rangle = 1$, and the expansion coefficients α_i^A in (46) are determined from the equation

$$(H_0 + H_{\text{int}})|A\rangle = M_A|A\rangle, \tag{47}$$

where H_{int} is the usual nonrelativistic Hamiltonian of the meson–nucleon interaction, and M_A is the mass of the nucleus. Projecting (47) onto the various Fock states (46), we can, in principle, find the explicit form of the wave functions α_i^A. In the case $A = 2$ (deuteron), no additional assumptions are needed in order to find α_i^A. At the same time, α_0^D satisfies a Schrödinger equation with a one-boson exchange potential, and $\alpha_{1,2}^D$ are functions of α_0^D and the matrix elements of H_{int}:

$$(2E-M_D)\alpha_0+|\langle NN\pi|H_{\rm int}|NN\rangle|^2\alpha_0/\Delta_1=0;$$

$$\alpha_1=\langle NN\pi|H_{\rm int}|NN\rangle\alpha_0/\Delta_1+O(g^3);$$

$$\alpha_2=\langle NN\pi\pi|H_{\rm int}|NN\pi\rangle\times\langle NN\pi|H_{\rm int}|NN\rangle\alpha_0/\Delta_2+O(g^3), \quad (48)$$

where

$$\Delta 1=\mathrm{MD}-\mathrm{E1}-\mathrm{E2}-\omega,$$

$$\Delta_2=M_D-E_1-E_2-\omega_1-\omega_2;$$

$$E=m+p^2/2m; \quad \omega=\sqrt{\mu_\pi^2+k^2}. \quad (49)$$

In what follows, we shall relate the function α_0 to the ordinary deuteron wave function $\Psi_D(p)$, which is the solution of a Schrödinger equation with a realistic meson-exchange NN potential (for example, the OBE potential of the Bonn group). Then the matrix elements of the operators relating to deep inelastic scattering reactions will be completely determined by the parameters of the corresponding mesons (coupling constants, masses, vertex form factors). For heavier nuclei, it is not possible to avoid the use of approximate methods to construct the many-particle wave functions α_i^A. Indeed, the nucleon wave function α_0^A satisfies a many-particle Schrödinger equation with a nonrelativistic two-particle meson-exchange potential:

$$\left(Am-M_A+\sum_{i=1}^{A}T_i\right)\alpha_0^A+\frac{1}{2}\sum_M\sum_{i,j}\int\frac{d^3k}{(2\pi)^3}\times V_M(k)\alpha_0^A(...p_i+k,...p_j-k,...)=0, \quad (50)$$

where T_i is the single-particle kinetic energy; here, $V_M(k)$ is the meson-exchange NN potential. To solve Eq. (50), we can use the Hartree approximation. Then, as in the deuteron case, the functions $\alpha_{1,2}^{AH}$ can be expressed in terms of α_0^{AH} and the matrix elements of $H_{\rm int}$, and α_0^{AH} itself is a product of single-particle functions, the equations for which are obtained from (50) by iteration. Thus, it is possible to obtain both the structure of a nucleus and the deep inelastic scattering by this nucleus in the framework of the same assumptions.

Amplitude for deep inelastic scattering by a nucleus with allowance for mesonic exchange currents

We recall that the basic quantity measured in deep inelastic scattering is the nuclear tensor $W_{\mu\nu}^A$, which is related by the optical theorem to the imaginary part of the amplitude $T_{\mu\nu}^A$ for forward elastic scattering of a virtual photon by the target nucleus [see (5)], which, in its turn, is the matrix element of a certain operator $t_{\mu\nu}^A$ with respect to the ground state of the nucleus. Strictly speaking, an explicit expression of the total operator $\hat{t}_{\mu\nu}^A$ for the system of interacting fields is unknown. It can be obtained by using the connection between $W_{\mu\nu}^A$ and the energy–momentum tensor:

$$\int F_2(x)dx=\frac{1}{m}\langle A|\Theta_{\lambda\lambda}|A\rangle, \quad (51)$$

where $F_2(x)=-xg^{\mu\nu}W_{\mu\nu}m$, and $\Theta_{\lambda\lambda}$ is the trace of the energy–momentum tensor, which for a system of interacting spinor scalar and vector fields (respectively, ψ, ϕ, and V^α) has the form[62]

$$\Theta_{\lambda\lambda}=m\bar{\psi}\psi+\mu^2\varphi\varphi+M^2V_\alpha V^\alpha, \quad (52)$$

i.e., is a sum, diagonal with respect to all fields, in which each term is the $\Theta_{\lambda\lambda}$ of the corresponding field. It is obvious that the relation (51) is also valid for the free fields. Thus, in our case the operator $\hat{t}_{\mu\nu}^A$ can be written as a sum of operators, $\hat{t}_{\mu\nu}^A=\hat{t}_{\mu\nu}^N+\Sigma_M\hat{t}_{\mu\nu}^M$, which describe Compton scattering by a nucleon and mesons, respectively:

$$\left.\begin{aligned}\langle A|\hat{t}_{\mu\nu}^A|A\rangle&=\frac{m}{M_A}T_{\mu\nu}^A(p_A,q);\\ \langle N|\hat{t}_{\mu\nu}^N|N\rangle&=\frac{m}{E}T_{\mu\nu}^N(p,q);\\ \langle m|\hat{t}_{\mu\nu}^M|M\rangle&=\frac{m}{\omega}T_{\mu\nu}^M(k,q).\end{aligned}\right\} \quad (53)$$

The normalization in (53) is chosen in such a way that the hadronic tensors have the same dimension. The relativistic expressions for the operators $\hat{t}_{\mu\nu}^{N,M}$ are determined by (32) and (33). The nonrelativistic reduction of the operator $\hat{t}_{\mu\nu}^N$ in the form (29) gives in accordance with the above procedure

$$\hat{t}_{\mu\nu}^N(q)=\sum_n C_n(q^2)\frac{2^n q_0 q_{\mu_2}...q_{\mu_n}}{(-q^2)^n}\left\{\psi^+(0)\left[\overleftrightarrow{\partial}^{\mu_2}...\overleftrightarrow{\partial}^{\mu_n}-\frac{i\overleftrightarrow{\partial}_z}{2m}\overleftrightarrow{\partial}^{\mu_2}...\overleftrightarrow{\partial}^{\mu_n}+\frac{g^2}{8m^2}(2\Phi(0)\overleftrightarrow{\partial}^{\mu_2}...\overleftrightarrow{\partial}^{\mu_n}\Phi(0)-\Phi^2(0)\overleftrightarrow{\partial}^{\mu_2}...\overleftrightarrow{\partial}^{\mu_n}-\overleftrightarrow{\partial}^{\mu_2}...\overleftrightarrow{\partial}^{\mu_n}\Phi^2(0))+\frac{ig\sigma_z}{2m}(\Phi(0)\overleftrightarrow{\partial}^{\mu_2}...\overleftrightarrow{\partial}^{\mu_n}-\overleftrightarrow{\partial}^{\mu_2}...\overleftrightarrow{\partial}^{\mu_n}\Phi(0))\right]\psi(0)\right\}, \quad (54)$$

where we have used the approximation usually employed in deep inelastic scattering, $q_\alpha\gamma^\alpha\approx q^0(\gamma_0+\gamma_z)$, and the Z axis has been taken in the direction opposite to that of the vector q. In our case, the last two terms do not contribute to the matrix element (53) and will therefore be omitted in what follows. Substituting (30) in (54), we obtain an explicit expression for the nonrelativistic operator $\hat{t}_{\mu\nu}^N(q)$:

$$\hat{t}_{\mu\nu}^N(q)=\int\frac{d^3p_1d^3p_2}{(2\pi)^6}\frac{m}{p_{20}+p_{2z}}(1+p_{2z}/m)T_{\mu\nu}^N(p_2,q)\times[a_\sigma^+(p_1)\chi_\sigma^+a_\rho(p_2)\chi_\rho], \quad (55)$$

where $a_{\sigma,\rho}^+$ and $a_{\sigma,\rho}$ are the nucleon creation and annihilation operators, respectively, σ and ρ are the spin–isospin indices, and $\chi_{\sigma,\rho}$ are the spin–isospin functions. For the pion operator $\hat{t}_{\mu\nu}^\pi(q)$, the relativistic form (33) is pre-

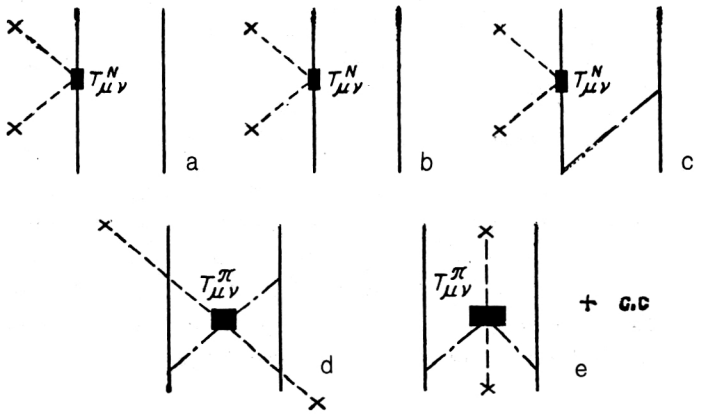


FIG. 5. Nonrelativistic diagrams of Compton scattering of a virtual photon by a nucleus.

served. From Eqs. (33), (46), (53), and (55), we can obtain the amplitude $T^A_{\mu\nu}$ in the second order in the constant g:

$$T^A_{\mu\nu}=\frac{M_A}{m}\,(\langle AN|\alpha_0^+\hat{t}^N_{\mu\nu}(q)\alpha_0|AN\rangle - Z\langle AN|\alpha_0^+\hat{t}^N_{\mu\nu}(q)\alpha_0|AN\rangle + \langle AN|\alpha_1^+\hat{t}^N_{\mu\nu}(q)\alpha_1|AN\rangle$$
$$+\langle ANM|\alpha_1^+\hat{t}^M_{\mu\nu}(q)\alpha_1|ANM\rangle + \langle ANMM|\alpha_2^+\hat{t}^M_{\mu\nu}(q)\alpha_0|AN\rangle + \langle AN|\alpha_0^+\hat{t}^M_{\mu\nu}(q)\alpha_2|ANMM\rangle);$$
$$(|AN...\rangle\equiv|N_1N_2...N_A...\rangle), \tag{56}$$

where the various terms correspond to the nonrelativistic diagrams in Fig. 5: impulse approximation (a), renormalization of the wave function (b), recoil (c), and exchange diagrams (d and e). We can now calculate the contribution of the various diagrams to the nuclear tensor $W_{\mu\nu}$. In (33) and (56), we have not yet determined the "elementary" amplitudes $T^{N(M)}_{\mu\nu}$ of Compton scattering of a photon by bound nucleons and virtual mesons. In what follows, we shall assume that the imaginary parts of these amplitudes, which are the corresponding elementary hadronic tensors $W_{\mu\nu}$, differ insignificantly from the free quantities, and we shall take them from experiment. In a self-consistent description, the matrix elements in (56) are determined by the parameters of the mesons (coupling constants, masses, vertex form factors) that generate the NN potential in Eqs. (49) and (50). Moreover, as is usual in the nonrelativistic approach, we ignore terms of self-energy type,[63,64] and in the expressions (49), (50), and (56) we take into account only the exchange contributions.

Impulse approximation in deep inelastic scattering of leptons by nuclei

Instead of the tensors $W_{\mu\nu}$, it is more convenient to use the structure functions $F^A_{1,2}(x)$, which are directly measured in experiments. The main contribution to (56) is made by the impulse-approximation diagram in Fig. 5a. We recall that the standard calculation of the nuclear structure function in the impulse approximation gives an incorrect behavior in the complete range of variation of the kinematic variable x.[9,10,65] The updating of the impulse approximation consists of correctly taking into account the fact that the nucleons are bound. Then, assuming the existence in nuclei of deep quasi-single-particle levels, we can satisfactorily describe the nuclear structure functions in the region of intermediate x ($x\geqslant 0.3$).[31,33] The contribution of the impulse approximation can be most accurately calculated in the case of the simplest nucleus, the deuteron:

$$F_2^{N/D}(x_D)=2\int\frac{d_p^3}{(2\pi)^3}\frac{1}{3}\sum_M|\Psi_D^M(p)|^2$$
$$\times\left(1+\frac{p_z}{m}\right)F_2\left(\frac{mx}{p_0+p_z}\right), \tag{57}$$

where, for convenience, we have replaced the deuteron wave function α_0 by the new, more customary notation Ψ_D^M (M is the projection of the total angular momentum of the deuteron); the nucleon off-shell effects are taken into account in $p_0=M_D-E$ and pp. To construct many-particle wave functions in the case of heavier nuclei, we use the Hartree approximation, in which the explicit expression for the structure function $F_2^{N/A}$ has the form

$$F_2^{N/A}(x_A)=\sum_\lambda \int \frac{d_p^3}{(2\pi)^3}\frac{1}{3}|\Psi_\lambda^A(p)|^2 \times\left(1+\frac{p_z}{m}\right)F_2\left(\frac{mx}{p_{\lambda 0}+p_z}\right), \tag{58}$$

where Ψ_λ^A are the single-particle wave functions of the nucleus (λ labels the states). In accordance with the definition of the model, $p_{\lambda 0} = m + \varepsilon_\lambda$, where ε_λ is the energy of the single-particle state. We note that our expressions (57) and (58) differ from the corresponding expression of Refs. 31 and 33 by a factor $(1 + p_z/m)$, which in the considered approach appears automatically because of the nonrelativistic reduction of the operator $\hat{t}_{\mu\nu}^N$ [see (54)]. The need for inclusion of this factor in calculations of $F_2^A(x)$ was discussed in Ref. 66, and also in Refs. 43, 54–56, and 67. This factor does not change the normalization of the function Ψ^A and, therefore, does not violate the conservation law for the baryon charge; however, its inclusion somewhat reduces the depth of the "pocket" at intermediate x in the ratio $R^{A/D}$ of the heavy-nucleus and deuteron structure functions.[34,67,68] The expressions (57) and (58) can be rewritten in the convolution form

$$\frac{1}{A}F_2^{N/A}(x)=\int_z^{M_A/m} d\xi F_2^N\left(\frac{x}{\xi}\right)f^{N/A}(\xi), \tag{59}$$

where the function $f^{N/A}(\xi)$ is the probability of the distribution of nucleons in the nucleus with momentum fraction x:

$$\begin{aligned} f^{N/D}(\xi)&=\int \frac{d^3p}{(2\pi)^3}\frac{1}{3}\sum_M |\Psi_D^M(p)|^2\left(1+\frac{p_z}{m}\right) \times\delta\left(\xi-\left(\frac{p_0+p_z}{m}\right)\right), \\ f^{N/A}(\xi)&=\frac{1}{A}\sum_\lambda \frac{d^3p}{(2\pi)^3}|\Psi_\lambda^A(p)|^2\left(1+\frac{p_z}{m}\right) \times\delta\left(\xi-\frac{p_{\lambda 0}+p_z}{m}\right). \end{aligned} \tag{60}$$

We can now determine the mean momentum carried by the quarks in the nucleus and compare it with the corresponding momentum of deep inelastic scattering by a nucleon. From (60), we obtain

$$\left.\begin{aligned} &\frac{2m\int_0^1 dxF_2^{N/D}(x)}{M_D\int_0^1 dxF_{2N}(x)}=\frac{2m\langle\xi\rangle}{M_D}; \\ &\langle\xi\rangle\equiv\int f_{N/D}(\xi)\xi d\xi\approx 1+\frac{\varepsilon_D}{m}-\frac{\langle p^2\rangle}{6m^2}\leqslant 1; \\ &\frac{Am\int_0^1 dxF_2^{N/A}(x)}{M_A\int_0^1 dxF_{2N}(x)}=\frac{Am\langle\xi\rangle}{M_A}; \\ &\xi\equiv\int f_{N/A}(\xi)\xi d\xi\approx 1+\frac{\langle\varepsilon\rangle}{m}-\frac{\langle 2T\rangle}{3m}\leqslant 1, \end{aligned}\right\} \tag{61}$$

where the averaging is with respect to the ground state of the target nucleus; ε_D is the deuteron binding energy; $\langle\varepsilon\rangle = (1/A)\Sigma\langle\varepsilon_\lambda\rangle = \langle V\rangle + \langle T\rangle$ (V and T are the potential and kinetic energies of the nucleon in the nucleus). It can be seen from the expressions in (61) that allowance for the fact that we have off-shell nucleons in the nucleus has the consequence that the mean momentum carried by quarks in a nuclear nucleon is less than in a free one. In the "nucleon–meson" model of the nucleus, the mesons responsible for the binding of the nucleus also carry a definite fraction of the momentum. To determine this fraction of the momentum and compare it with the missing momentum in (61), it is necessary to calculate the mesonic corrections to the impulse approximation (the diagrams of Figs. 5b–5e).

Mesonic currents in deep inelastic scattering of leptons by nuclei

The diagrams in Figs. 5b and 5c are usually called the renormalization and recoil diagrams in the literature. Their explicit form is given in Ref. 54, in which the total contribution of these graphs to the deuteron sum rules is also estimated. In what follows, we shall ignore the corrections from these diagrams, since, on the one hand, their total contribution to the sum rule (61) is proportional to g^4, and our treatment is restricted to the corrections up to g^2. On the other hand, an analysis of their contribution to elastic and quasielastic processes made by various authors[59,69,70] has shown that it is either exactly zero (the diagrams "cancel") or vanishing in the limit of small momentum transfers. The remaining diagrams (Figs. 5d and 5e) are due to mesonic currents and have an exchange nature:

$$\begin{aligned} \frac{1}{2}F_2^{M/D}&=\int_x^{M_{D/m}} d\eta f_\mu(\eta)F_2^\mu(x/\eta); \\ f_\pi(\eta)&=\frac{g^2}{2m^2}\int\frac{d^3p_1d^3p_2}{(2\pi)^6}\frac{1}{3\Omega^4}\sum_M \Psi_D^{M+}(p_1)(\boldsymbol{\sigma}_1\cdot\mathbf{k}) \times(\boldsymbol{\sigma}_2\cdot\mathbf{k})(\boldsymbol{\tau}_1\cdot\boldsymbol{\tau}_2)\Psi_D^M(p_2)(k_0+k_z) \times\delta(\eta-(k_0+k_z)/m)\Theta(k_0+k_z) \\ &\equiv\int d^3k n_\pi(k)(k_0+k_z)/\Omega \times\delta(\eta-(k_0+k_z)/m); \end{aligned} \tag{62}$$

$$\begin{aligned} f_\sigma(\eta)&=2g_\sigma^2\int\frac{d^3p_1d^3p_2}{(2\pi)^6}\frac{1}{3\Omega_\sigma^4}\sum_M \Psi_D^{M+}(p_1)\Psi_D^M(p_2) \times(k_0+k_z)\delta(\eta-(k_0+k_z)/m)\Theta(k_0+k_z); \\ f_\omega(\eta)&=-2g_\omega^2\int\frac{d^3p_1d^3p_2}{(2\pi)^6}\frac{1}{3\Omega_\omega^4}\sum_M \Psi_D^{M+}(p_1)\Psi_D^M(p_2) \times(k_0+k_z)\delta(\eta-(k_0+k_z)/m)\Theta(k_0+k_z), \end{aligned} \tag{63}$$

where the variables of integration p_1 and p_2 are the momentum of the nucleon before and after the interaction;

$k = p_1 + p_2$ is the meson momentum; $\Omega^2 = (k_0^2 - \omega_\mu^2)$; $\mu = \pi, \omega, \sigma$. In calculations of the mesonic corrections to deep inelastic scattering by a heavier nucleus, it is necessary to choose a definite meson-exchange NN potential, i.e., to choose the set of mesons that are subsequently to be taken into account. In this, the nuclear case differs fundamentally from the deuteron case, for which the mesons must determine the ordinary one-boson NN potential. For our purpose, the most suitable potential is the generalized Miller–Green OBE potential,[71] which will be used in what follows. In its form, this is an ordinary one-boson exchange potential and includes the scalar (σ and ε), vector isoscalar (ω), and vector (ρ) mesons. The parameters of the ε and ω mesons are such that their contributions to the static potential cancel, while the contribution of the ρ meson can be ignored for the description of spherical nuclei. We note that the parameters of the Miller–Green potential are chosen in such a way as to describe the basic properties of spherical nuclei in the Hartree approximation, as a consequence of which (and also because of the absence of the ρ meson, which in this approximation does not contribute to the single-particle potential for spherical nuclei) this potential is not a realistic NN potential, i.e., it does not describe NN scattering and the properties of the deuteron. An argument in support of such a modification is provided by the observation that the NN interaction in vacuum and in a medium (nucleus) are different; in addition, such a parametrization may take into account effectively nucleonic correlations, which are explicitly excluded in the Hartree approximation. For comparison, we also give the results of calculations with the realistic Bonn potential (σ and ω mesons). The contribution of the scalar mesons to the nuclear structure function in the Hartree approximation has the form (see also Refs. 72 and 73)

$$F_2^s(x) = 2g_s^2 \int \frac{d^3k}{(2\pi)^3} F_2^M(mx/k_3) k_3/\omega_s^4 \rho(\mathbf{r}_1)\rho(\mathbf{r}_2) \times e^{i\mathbf{k}(\mathbf{r}_1 - \mathbf{r}_2)} d\mathbf{r}_1 d\mathbf{r}_2, \tag{64}$$

where $\omega = k^2 + \mu_s^2$; $\rho(r)$ is the ordinary scalar nuclear density. To the accuracy adopted in our calculations ($\sim g^2$), the contribution of the vector mesons is determined by a formula analogous to (64). The conservation law for the total momentum carried by the quarks and gluons in the nucleus (for brevity, we again only give the expressions for the deuteron) now takes the form

$$\langle x_q \rangle_D = \frac{2m}{M_D} \{\langle x_q \rangle_N \langle \xi \rangle + \langle x_q \rangle_\pi \langle \eta_\pi \rangle + \langle x_q \rangle_\sigma \langle \eta_\sigma \rangle + \langle x_q \rangle_\omega \langle \eta_\omega \rangle + \ldots\};$$

$$\langle x_g \rangle_D = \frac{2m}{M_D} \{\langle x_g \rangle_N \langle \xi \rangle + \langle x_g \rangle_\pi \langle \eta_\pi \rangle + \langle x_g \rangle_\sigma \langle \eta_\sigma \rangle + \langle x_g \rangle_\omega \langle \eta_\omega \rangle + \ldots\}, \tag{65}$$

where the meaning of the notation that we have introduced is obvious. Adding the right- and left-hand sides of (65), and using the equation $2m/M_D \approx 1 - \langle \varepsilon \rangle / 2m$, we obtain a sum rule for the mean values of the momenta of the deuteron constituents:

$$1 = 1 + \delta_N + \langle \eta_\pi \rangle + \langle \eta_\sigma \rangle + \langle \eta_\omega \rangle + \ldots;$$

$$\delta_N = \frac{\langle \varepsilon \rangle}{2m} - \frac{\langle p^2 \rangle}{6m^2}. \tag{66}$$

Let us consider how (66) is satisfied in our case. For example, let us suppose that there are only pions in the deuteron. Then instead of (66) we must obtain the equation $\langle \delta_N \rangle = -\langle \eta_\pi \rangle$, where in accordance with (63)

$$\langle \eta_\pi \rangle = \frac{g^2}{2m^2} \int \frac{d^3p_1 d^3p_2}{(2\pi)^6} \frac{1}{3\Omega^4} \Sigma_M \Psi_D^{M+}(p_1)(\boldsymbol{\sigma}_1 \cdot \mathbf{k})(\boldsymbol{\sigma}_2 \cdot \mathbf{k}) \times (\tau_1 \cdot \tau_2) \Psi_D^M(p_2)(k_0 + k_z)^2 / m\Theta(k_0 + k_3). \tag{67}$$

Indeed, bearing in mind that $-g^2(\sigma_1 k)(\sigma_2 k)(\tau_1 \tau_2)/4m^2\omega^2$ is precisely the pion–nucleon interaction potential $V(k)$,[74] and using the relations

$$V(k)k^2/\omega^2 = V(k)\left(1 - \frac{\mu_\pi^2}{\omega^2}\right) = \left(1 - \frac{1}{2} k \frac{d}{dk}\right) V(k) \tag{68}$$

and $k_2^2 \sim k^2/3$ in the static limit ($k_0 \approx 0$, $\Omega^2 \cong \omega^2$), we can rewrite (67) in the coordinate representation, obtaining

$$\langle \eta_\pi \rangle = -\frac{1}{3m} \int d^3r \frac{1}{3} \Psi_D^{M+}(r) \times \{\tfrac{3}{2} V(r) + \tfrac{1}{2} r dV(r)/dr\} \Psi_D^M(r). \tag{69}$$

Then in accordance with the virial theorem, $\langle r dV(r)/dr \rangle = \langle 2T \rangle$, we obtain

$$\langle \eta_\pi \rangle = -\frac{1}{3m} \left(\frac{3}{2} \langle V \rangle + \langle T \rangle\right), \tag{70}$$

where T is the kinetic energy of a nucleon in the deuteron. Substituting $\langle V \rangle = \langle \varepsilon \rangle - \langle T \rangle$ in (70), we see that

$$\delta_N = -\langle \eta_\pi \rangle, \tag{71}$$

i.e., the sum rule holds exactly. One can show that (71) is also satisfied for the remaining mesons [for the scalar mesons, the relation (71) is obtained in Ref. 43], the only difference being that for each of them the corresponding part of the potential will occur in the virial theorem, and in the sum we again obtain (74). For a nucleus with $A \geqslant 2$, the treatment is again as in Refs. 43 and 75, and Eq. (71) also holds. We merely note here that the determination of the mesonic contributions to the nuclear structure function on the basis of (51), (52), and (64) contains a certain arbitrariness, since the relation (51) is an integral relation. Thus, it was shown in a study of deep inelastic scattering by a nucleus in the framework of Walecka's model[72] that the contribution of the vector fields may contain other corrections besides (64). However, in the analysis of the sum rule (71) these corrections exactly cancel, and, in integral form, (64) is identical to the result of Ref. 72 for the vector mesons. Thus, in the static limit allowance for the meson-exchange corrections to the structure functions in deep inelastic scattering completely restores the energy sum

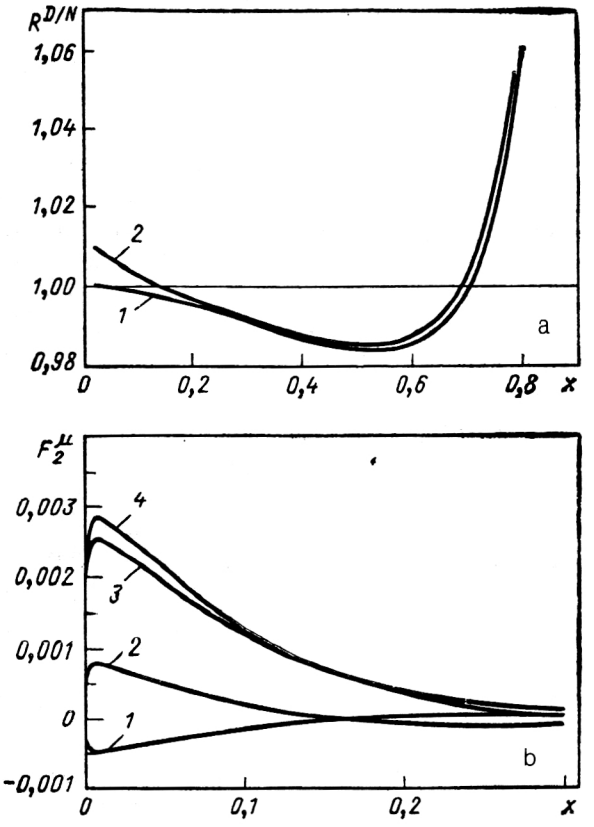


FIG. 6. Contribution of mesonic exchange currents to deuteron structure function: a) ratio of structure functions of the deuteron and the free isoscalar nucleon: 1) impulse approximation; 2) impulse approximation with allowance for mesonic exchange currents; b) absolute contributions of mesonic currents to deuteron structure function: 1)–3) contributions of ω, σ, and π mesons, respectively; 4) total contribution of all mesons.

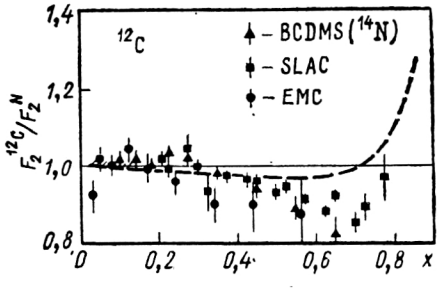


FIG. 7. Ratio of structure functions of the ^{12}C nucleus and the deuteron. The curve is calculated in the Hartree approximation with Skyrme I forces; the narrow hatched region [*sic*] demonstrates the contributions of the mesonic corrections calculated under different assumptions (see the text).

rule, which is violated in the impulse approximation because of the binding of the nucleons. In a real case ($k_0 \neq 0$), this need not be so. Note also that (71) has been obtained under the assumption that g is a constant.

In Fig. 6, we give the results of a calculation of the ratio $R^{D/N}(x)$ of the deuteron structure function and the structure function of an isoscalar nucleon. Curve 1 corresponds to the impulse approximation, and curve 2 to calculation with the deuteron structure function in accordance with the expressions (57) and (62), (63). The deuteron wave function was determined by solution of the Schrödinger equation with the Bonn potential (Ref. 35); the parameters that determine the contribution of the ρ, ω, and σ mesons in the meson-exchange diagrams are taken to be the same as in the determination of the NN potential. Thus, the problem is self-consistent overall. A small failure of self-consistency has a general nature and results from the fact that the main diagrams of the deep inelastic scattering, and also the Schrödinger equation for the function Ψ_D are obtained in the approximation up to g^2, whereas the wave function Ψ_D in (57) and (63) actually contains all powers of g. In addition, in the realistic NN potentials phenomenological meson–nucleon vertex form factors are used, $g^2 \Rightarrow g^2\Lambda^2/(\Lambda^2 + k^2)$. Thus, the contribution of the impulse-approximation diagram already depends on Λ. Therefore, it should also be taken into account in the exchange diagrams. These circumstances may lead to violation of the relation (71).

It can be seen from Fig. 6a that the qualitative behavior of the ratio $R^{D/N}$ is the same as for heavy nuclei, but the depth of the minimum at $x \sim 0.6$ is here appreciably less because of the relative weakness of the nuclear effects in the deuteron. The main contribution to the structure functions is made by the pion-exchange diagrams (see Fig. 6b). The contribution of the heavier mesons is suppressed by their greater masses, which occur in the corresponding propagators, and by the effects of the scalar (σ meson) and vector (ω meson) coupling. A numerical calculation of the sum rule (66) for the deuteron gave the result $\delta_N = -0.0036$, and as the total contribution of all mesons $\Sigma\langle\eta_\nu\rangle \approx 0.0022$. Thus, it can be seen that numerically the mesonic corrections only partly ($\sim 60\%$) restore the energy sum rule violated in the impulse approximation.

We now consider a heavier nucleus: ^{12}C. In our case (only in the Hartree approximation!) the contributions of the impulse approximation, (58), and of the mesonic currents, (64), are expressed in terms of the ordinary nuclear densities, which, in their turn, can be taken from independent sources—from experiment or some realistic model. Figure 7 gives the result of calculation of the ratio $R^{A/D}(x)$ of the structure function of the ^{12}C nucleus in accordance with the expressions (58) and (64) and the deuteron in accordance with the expressions (57), (62), and (63). Note that the contribution of the mesonic currents to $R^{A/D}(x)$ is small—the curves calculated with different mesonic contributions (Fig. 8) differ in the narrow hatched region at $x \leqslant 0.2$. It can be seen from Fig. 7 that, in contrast to Refs. 31, 33, and 34, $R^{A/D}(x)$ does not describe the data in the region $x \sim 0.5$–0.7. As was noted above, this is due to the fact that in the approach developed here we do not take into account collective effects, which lead to the formation of deep quasi-single-particle levels. In this respect, our results agree with those of Ref. 67.

In Fig. 8, we give the results of calculations of the mesonic contributions to the structure function. Curves 1 and 2 (continuous and broken) correspond to two different sets of mesons (Miller–Green and Bonn potentials) and

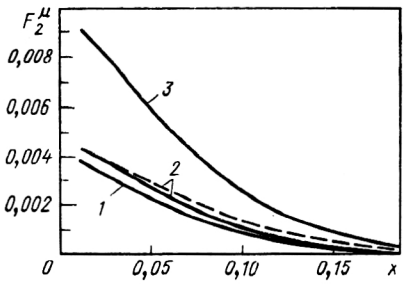


FIG. 8. Absolute contribution of mesonic currents to the carbon structure function: 1) calculation with Woods–Saxon nuclear density and Miller–Green mesons; 2) with Skyrme I density and two sets of mesons (continuous curve for Miller–Green mesons, broken curve for σ and ω mesons of the Bonn potential). Curves 1 and 2 are calculated with cutoff parameters Λ from the corresponding potentials; 3) calculation with Skyrme I density and Miller–Green mesons in the limit $\Lambda \to \infty$.

different nuclear densities (Skyrme I and Woods–Saxon). It can be seen that the results of these calculations are practically identical, indicating that the assumptions made in choosing the meson potentials and nuclear densities were correct. Numerical calculation of the sum rule (66) (Skyrme I) gave the result $\delta_N = (Am/M_A) \times (\xi - 1) = -0.0046$ (see for comparison Ref. 67), and the total contribution of all the mesons (Miller–Green) is $\Sigma\langle\eta_\mu\rangle \approx 0.0023$. Thus, it can be seen that, as in the deuteron, the mesonic corrections only partly ($\sim 50\%$) restore the energy sum rule. Elimination of the form factors from the exchange diagrams ($\Lambda \to \infty$) completely restores the sum rules [$(\Sigma\langle\eta_\mu\rangle) \approx 0.005$; see also Fig. 8], but the self-consistency of the method is then evidently lost.

For complete restoration of the sum rule, it will be necessary in what follows to investigate: a) The contributions of the following terms of the expansion in g^2, since such contributions are contained in the deuteron wave function. In particular, we must consider in more detail the systematic construction, with controlled accuracy in g^2, of the "elementary" operators $\hat{t}^M_{\mu\nu}$ and $\hat{t}^N_{\mu\nu}$. b) The contribution of other non-nucleonic (N^*, Δ, multiquark) degrees of freedom in the deuteron. Here, it must be borne in mind that this contribution is partly taken into account phenomenologically in the realistic deuteron wave function. c) The contribution from heavier mesons (by which the k^2 dependence of the vertex form factors can be imitated). In particular, we could here have mesonlike objects neutral with respect to the electromagnetic interaction (like glueballs[76]), which could contribute to the potential but not be "detected" in deep inelastic scattering.

Realistic nucleon momentum distributions and different schemes for taking into account Fermi motion

It was noted above that for satisfactory description of the structure functions of intermediate and heavy nuclei at $x \leqslant 0.2$–0.3 it is necessary to take into account scattering by deep levels having a more complicated collective nature. In the impulse approximation, all information about nuclear structure is contained in $S^A(\varepsilon,k)$ [see (20) and (21)] and, accordingly, in $f^{N/A}$. Exact relativistic expressions for $S^A(\varepsilon,k)$ and $f^{N/A}$ are unknown. However, $S^A(\varepsilon,k)$ must satisfy the normalization conditions (19)–(22) and must describe well the basic properties of nuclei and collective effects such as the energy widths and centroids, the cross sections for the scattering of particles and ions by the considered nucleus in various reactions, the momentum distributions, etc. There are several approaches that permit the construction of such a spectral function.[31,33,68] One of them is the coherent density-fluctuation model (CDFM).[77–79] In this method, the spectral function $S^A(\varepsilon,k)$ has the comparatively simple form

$$S^A(\varepsilon,k) = \frac{16\pi r_0^3}{3}\frac{\alpha}{2|k|}\frac{|f(r_0)|^2}{\sqrt{\mu(\varepsilon - E_f)}}, \tag{72}$$

where $r_0 = \alpha/(|k|\sqrt{(\varepsilon - E_f)/\mu})$, $\alpha = (9\pi A/8)^{1/3}$. The values of the parameters E_f and μ are taken from Refs. 77 and 78: $\mu = -50$ MeV and $E_f = 8$ MeV. The function $|f(r_0)|^2$ is related to the nuclear density by the equation

$$|f(r_0)|^2 = -\frac{4\pi r_0^3}{3A}\left(\frac{d\rho(r)}{dr}\right)_{r=r_0}. \tag{73}$$

The nuclear density $\rho(r)$ in (73) can be extracted from electron–nucleus scattering experiments; for example,[80]

$$|f(x)|^2 = \frac{4\pi x^3}{3Ab}\rho_0\left\{\frac{e^{-(x+R)/b}}{(1+e^{-(x+R)b})} - \frac{e^{(x-R)/b}}{(1+e^{(x-R)/b})^2}\right\}, \tag{74}$$

where $\rho_0 = 3A/\{4\pi R^3[1 + (\pi b/R)^2]\}$. The spectral function constructed in this manner describes electron $(e, e'p)$ reactions well, including the data on deep hole states. In accordance with (21), the distribution function $f^{A/N}(y)$ of the nuclear nucleons with respect to the longitudinal momentum is determined by

$$f^{A/N}(y) = \int \frac{d^3k}{(2\pi)^3}\, d\varepsilon S(\varepsilon,k)(1 + k_z/m) \times \delta(y - k_+/m),$$
$$k_+ = m + \varepsilon + k_z. \tag{75}$$

Accordingly, the mean momentum $\langle y\rangle$ has the form

$$\langle y\rangle \cong 1 + \frac{\langle\varepsilon\rangle}{m}, \quad \langle\varepsilon\rangle = \int \frac{d^3k}{(2\pi)^3}\, d\varepsilon S(\varepsilon,k)\varepsilon. \tag{76}$$

Figures 9 and 10 show the results of calculations of F_2^A for the nuclei ^{12}C and ^{56}Fe in accordance with the expressions (20), (21), and (72)–(75) using the following density parameters (74):

$$\left.\begin{aligned} R = 2.214 \text{ fm}, \ b = 0.488 \text{ fm} \quad &\text{for } ^{12}\text{C}; \\ R = 4.054 \text{ fm}, \ b = 0.600 \text{ fm} \quad &\text{for } ^{56}\text{Fe}. \end{aligned}\right\} \tag{77}$$

It can be seen that the theoretical results of the CDFM (continuous curves) agree well with the experimental data in the region $0.3 \leqslant x \leqslant 0.7$. For methodological purposes, we also give the results of calculations without the flux factor $(1 + k_z/m)$ (broken curve). Figure 9 also shows the

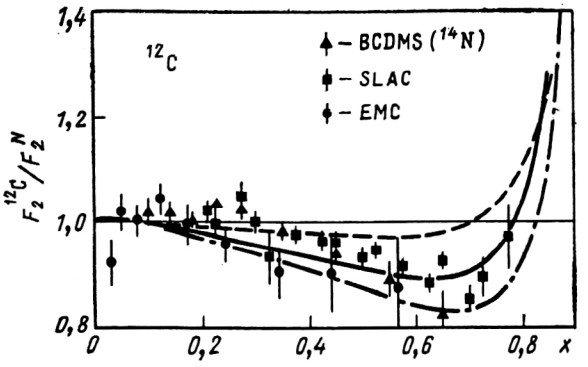


FIG. 9. Ratio $F_2^{^{12}C}/F_2^N(x)$ calculated in the coherent density-fluctuation model (with and without flux factor: continuous and chain curves, respectively) and in the Hartree approximation with Skyrme forces (broken curve).

results for the nucleus ^{12}C obtained using the single-particle spectral function in the Hartree approximation with Skyrme forces. The behaviors of the ratios in the two cases (CDFM and Hartree approximation) are qualitatively similar. However, the depths of the minimum are different, since the behavior of the structure function at intermediate x depends strongly on $\langle\varepsilon\rangle$ (76). Indeed, expanding $F_2^N(x/y)$ in (20) in the neighborhood of $\langle y\rangle$ [i.e., at the point of the sharp peak of $f^{N/A}(y)$],

$$\frac{1}{A}F_2^A(x)\cong F_2^N(x/\langle y\rangle)+\frac{1}{2}(\langle y^2\rangle-\langle y\rangle^2)\frac{\partial^2}{\partial\langle y\rangle}\times F_2^N(x/\langle y\rangle)+\ldots \tag{78}$$

and substituting $\langle y\rangle$ (76), we obtain an estimate for $F_2^A(x)$:

$$F_2^A(x)\approx F_2^N(x/(1+\langle\varepsilon\rangle/m). \tag{79}$$

It is clear from this that the disagreement in the results of the CDFM and of the single-particle Hartree approximation is a consequence of the different values of $\langle\varepsilon\rangle$ obtained in these approaches—the CDFM gives $\langle\varepsilon\rangle\approx -38$ MeV, in agreement with the results of Refs. 31 and 33, whereas the Hartree approximation gives $\langle\varepsilon\rangle\approx -20$ MeV (see also Ref. 67).

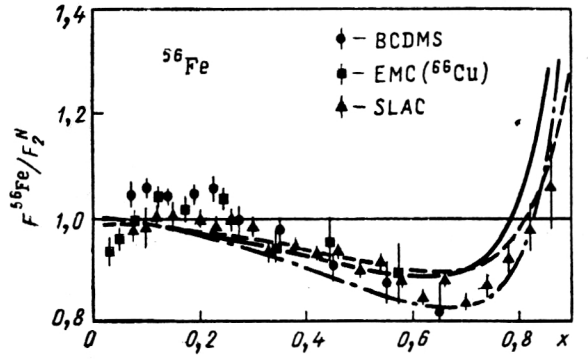


FIG. 10. Ratio $F_2^{^{56}Fe}/F_2^N(x)$ calculated in the coherent density-fluctuation model (with and without flux factor: continuous curve and broken curve, respectively). The ratio $F_2^{^{56}Fe}/F_2^D(x)$ (chain curve) is shown for comparison.

3. NUCLEAR-STRUCTURE EFFECTS FOR THE LIGHTEST NUCLEI IN DEEP INELASTIC SCATTERING OF POLARIZED LEPTONS BY POLARIZED NUCLEI

There is great interest currently in investigations of spin structure in deep inelastic scattering. This is largely due to the results of measurement of the proton spin structure function. It has been found that the fraction of the nucleon spin carried by quarks is, to within the experimental errors, zero, in sharp contradiction with the picture provided by the parton model. This work has been based on models of the relative distribution of the proton and neutron. Since direct information about the neutron structure function can be obtained from nuclear data, it is necessary to investigate the part played by nuclear structure in deep inelastic scattering in polarization experiments. These questions are discussed below.

Effects of Fermi motion in deep inelastic scattering by polarized deuterons

The cross sections σ^+ and σ^- (with mutually parallel and antiparallel orientation of the beam and target spins along the beam direction) in the deep inelastic scattering of polarized leptons by a polarized hadronic target T are determined by the structure functions $F_{1,2}^T(x)$ (3) and the so-called spin functions $g_{1,2}^T(x)$.[3,81] They depend on the target polarization and characterize the spin distribution of the nucleon with respect to the quarks of the different species. In the parton model, $g_{1,2}^T(x)$ are given by

$$g_1^T(x)=\sum_\alpha e_\alpha^2(q_{\alpha/T\uparrow}^T(x_T)-q_{\alpha/T\downarrow}^T(x_T)),\quad g_2=0, \tag{80}$$

where the index α ranges over all species of charged partons; $q_{\alpha\uparrow/T\uparrow}^T(x_T)$ and $q_{\alpha\uparrow/T\downarrow}^T(x_T)$ are the densities of the probability of finding in the target T parton α with fraction x_T of the longitudinal momentum and spin parallel or antiparallel to the target spin.

The spin structure functions $g_{1,2}$ have been the subject of much less experimental and theoretical study than the functions $F_{1,2}$. In early studies[82,83] having methodological aims, the following assumptions were used to determine the spin structure functions: 1) a contribution to g_1^N is made by only the valence quarks (the "sea" is unpolarized), the x distribution for which was taken to be the same as in the unpolarized case; 2) the distribution of the valence quarks with respect to the spin orientations is determined by the nucleon wave function in the representation [56, $L_z=0$] of the $SU(6)$ classification:[3]

$$|p\uparrow\rangle=\sqrt{\tfrac{2}{3}}\,|u\uparrow u\uparrow d\downarrow\rangle+\tfrac{1}{3}\,|u\uparrow d\uparrow u\downarrow\rangle. \tag{81}$$

These assumptions lead to the following structure functions for the proton (p) and isoscalar nucleon (N):

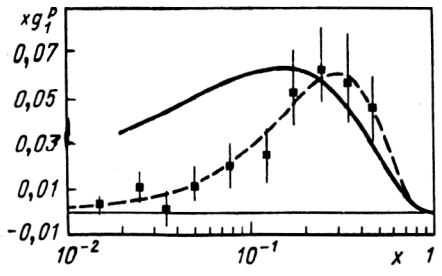


FIG. 11. Proton spin function. The experimental points and the broken curve are from Ref. 86; the continuous curve is $xg_1^p(x)[(SU(6)]$ (see the text).

$$g_1^p(x)=\tfrac{1}{27}\{8u_v(x)-d_v(x)\}; \tag{82}$$

$$g_1^N(x)=\tfrac{5}{54}\{2u_v(x)-d_v(x)\}, \tag{83}$$

where $u_v(x)$ and $d_v(x)$ are the distributions of the valence u and d quarks in the proton (here, these distributions are normalized to 1). In addition, we ignore electromagnetic corrections of higher orders in α and the corrections from neutral currents (including the possible interference of the electromagnetic and neutral currents), the contribution of which at the planned initial lepton energies and momentum transfers is small and does not exceed a few percent.[85] It can be seen from Fig. 11 that the functions (81) do not describe the experimental data at $x\leqslant 0.2$ and agree reasonably with them at large x. Application of the convolution model to the deep inelastic scattering of polarized leptons by polarized nuclei gives for the structure function $g_1^A(x)$ the expression[22]

$$g_1^A(x)=\sum_t \int_x^{y_{\max}} dy/y\, g_1^t(x/y)[\,f_{t\uparrow/A\uparrow}(y)-f_{t\downarrow/A\uparrow}(y)\,], \tag{84}$$

where t is a virtual hadronic constituent of nucleus A, $f_{t\uparrow/A\uparrow}(y)$ and $f_{t\downarrow/A\uparrow}$ are the distribution functions of constituents t with different spin orientations (spin 1/2) relative to the spin of nucleus A. Here, we consider only nucleonic degrees of freedom in the nucleus, and therefore $t=p,\ n$. Integration of (84) makes it possible to obtain integral sum rules for the nuclear structure functions g_1. An example of such a sum rule will be considered below. To reduce to a minimum the uncertainties associated with the determination of the elementary structure functions g_1^N, we shall consider relative quantities of the type $R_g^{D/N}(x)=g_1^D(x)/g_1^N(x)$, where g_1^D is determined by (84). Then to a good accuracy the uncertainties in the structure functions g_1^p and g_1^n (see Fig. 11) cancel (in any case when $x\geqslant 0.2$). Therefore, we shall regard the "elementary" g_1^p and g_1^n as given [for definiteness, we use the functions (82) and (83)] and concentrate our attention on the truly nuclear part of the convolution formula (84): $f_{t\uparrow/A\uparrow}(y)$ and $f_{t\downarrow/A\uparrow}$. It is clear that, in principle, the distribution functions $f_{t\uparrow/A\uparrow}(y)$ and $f_{t\downarrow/A\uparrow}$ must reflect not only the spin effects themselves but also all the same nuclear effects as in the study of unpolarized processes (the off-shell behavior of the nucleons, corrections of the type $\langle p\rangle/m$, mesonic currents,...). However, in the case of the lightest nuclei ($A\sim 2$–3) these effects are small compared with those of the spin–orbit structure. Therefore, for definiteness we shall here, when calculating the function $R_g^{D/N}(x)$ for $x\leqslant 1$, use the theory of bound systems in the light-front kinematics.[87,88] Taking into account the spin–orbit structure of the deuteron wave function in this formalism, we obtain

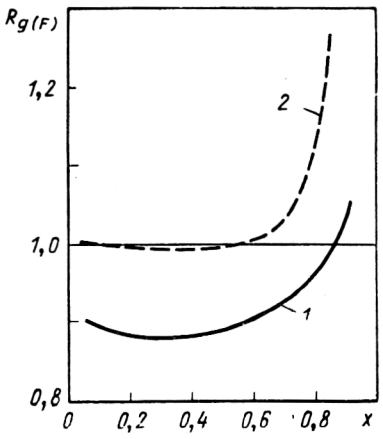


FIG. 12. Ratio of structure functions of the deuteron and isoscalar nucleon: 1) ratio of the spin structure functions $g_1(x)$; 2) ratio of the structure functions F_2. The calculations were made in light-front kinematics with the deuteron wave function in the Paris potential.

$$f_{t\uparrow A\uparrow}(y)=\int dk^2 \frac{f_{\uparrow/t\uparrow A\uparrow}(k)}{2y(1-y)},$$

$$f_{t\downarrow/A\uparrow}(y)=\int dk^2 \frac{f_{t\downarrow/A\uparrow}(k)}{2y(1-y)}; \tag{85}$$

$$f_{t\uparrow/A\uparrow}(k)=[4\pi(u(k)Y_{00}(\hat{k})-w(\hat{k}))Y_{20}(k)/\sqrt{10}];$$

$$f_{t\downarrow/A\uparrow}(k)=\sqrt{\tfrac{3}{5}}\,4\pi w(k)Y_{22}(\hat{k}), \tag{86}$$

where u, w, S, D are the components of the deuteron wave function. For actual calculations, we used the realistic deuteron wave function in the "Paris" potential[89] and the structure function (83) with distributions $u_v(x)$ and $d_v(x)$ from Ref. 39. Figure 12 gives the ratio $R_g^{D/N}$ (curve 1) and, for comparison, $R_F^{D/N}=F_2^D(x)/F_2^N(x)$, calculated in the same light-front kinematics (curve 2). It can be seen that the behaviors of $R_g^{D/N}$ and $R_F^{D/N}$ differ quantitatively. The destructive contribution of the D wave in $g_1^D(x)$ has the consequence that $R_g^{D/N}$ is significantly less than unity in almost the complete range of x (and less than $R_F^{D/N}$ at all x).

Bjorken sum rule and deep inelastic scattering by polarized nuclei

In this subsection, we discuss (Ref. 84): a) the problem of separating the first moments of the nucleon structure functions $g_1^{p,n}$ (in the context of extraction of additional information for investigation of the "spin crisis" from the nuclear structure functions in polarized deep in-

elastic scattering); b) the connection between the nuclear structure functions and ordinary nuclear physics (weak decays, nuclear magnetic moments, etc.). For this, we shall consider the spin structure functions $g_1^T(x)$ of different targets (nucleons and nuclei). We shall also discuss the Bjorken sum rule,[90] which is one of the basic relations used to analyze deep inelastic scattering of polarized particles (in addition, this sum rule makes it possible to relate the structure function g_1 to quantities and matrix elements often encountered in nonrelativistic nuclear physics):

$$\int_0^1 dx\left(g_1^p(x,Q^2)-g_1^n(x,Q^2)\right)=\frac{1}{6}\frac{g_A}{g_V}\left(1-\frac{\alpha_s}{\pi}\right). \quad (87)$$

This relation is also true for deep inelastic scattering by nuclear targets; in this case, we shall replace g_A/g_V by $(g_A/g_V)_A$ on the right-hand side of Eq. (87). Direct experimental verification of fundamental relations like the Bjorken sum rule and the more accurate determination of the nucleon spin structure functions is also important because analysis of the experimental EMC data using information on the weak decays of baryons casts doubt on the validity of perturbative QCD (!?) if the Bjorken sum rule is correct.[91] Alternatively, if one believes QCD, then the Bjorken sum rule must be violated. Using the convolution model and ignoring the off-shell corrections (which are unimportant in the case of the lightest nuclei) to integrated quantities such as (87), we integrate (80) and obtain for the nuclear structure functions

$$\int_0^1 dx g_1^A(x)=\sum_{t=p,n}\int_0^1 dx g_1^t(x)\left[p_{t\uparrow/A\uparrow}-p_{t\downarrow/A\uparrow}\right], \quad (88)$$

where $p_{t\uparrow/A\uparrow}$ and $p_{t\downarrow/A\uparrow}$ are the probabilities of finding in the ground state of nucleus A the constituent t with spin parallel and antiparallel to the nuclear spin, respectively. In the formalism of ordinary nuclear physics, $[p_{t\uparrow/A\uparrow}-p_{t\downarrow/A\uparrow}]$ can be calculated as a matrix element of the operator $\Sigma_i\sigma_z^t(i)$ (sum over all nucleons of species t) with respect to the ground-state nuclear wave function $|A\rangle$:

$$\int_0^1 dx g_1^A(x)=\left[\int_0^1 dx g_1^p(x)\right]S^{pA}+\left[\int_0^1 dx g_1^n(x)\right]S^{nA}, \quad (89)$$

where $S^{pA(nA)}=\langle A|\Sigma_{p(n)}\sigma_v^{p(n)}|A\rangle$ is the spin structure factor of the nucleus. Equation (89) is a direct consequence of the convolution approach, which is usually used to investigate deep inelastic scattering by nuclei. However, the interpretation of this relation requires some care in view of (a) the possible change of the nucleon properties in the nuclear medium and because (b) the simplified treatment of the nuclear structure can lead to incorrect results.

Close *et al.*[22] used the Bjorken sum rule (87) and relations of the type (89) to determine g_A/g_V for bound nucleons, using the structure functions g_1^A and $q_1^{A^*}$ of mirror nuclei, i.e., nuclei A and A^* in which the numbers of protons and neutrons and interchanged, with spin $J=1/2$. For mirror nuclei, we have $S^{pA}=S^{nA^*}$, $S^{nA}=S^{pA^*}$, and therefore

$$\int_0^1 dx(g_1^A(x)-g_1^{A^*}(x))$$
$$=\frac{1}{6}\left(\frac{g_A}{g_V}\right)_{AA^*}\left(1-\frac{\alpha_s}{\pi}\right)$$
$$=\frac{1}{6}\left(\frac{g_A}{g_V}\right)_{\text{bd. nucl}}\left(1-\frac{\alpha_s}{\pi}\right)[S^{pA}-S^{nA}]. \quad (90)$$

Following the logic of the convolution model, we should set $(g_A/g_V)_{\text{bd. nucl}}g_A/g_V$. In the context of our treatment, this would mean equality of the first moments of the structure functions $g_1(x)$ of free and bound nucleons. However, as is well known, the ratio g_A/g_V is "renormalized" in a nucleus into an effective constant $(g_A/g_V)_{\text{bd. nucl}}$, which can be determined from β decay and Gamow–Teller transitions of nuclei.[92] This discrepancy arises from the limitations of the convolution approach, which operates solely with nucleonic degrees of freedom, whereas mesonic and other non-nucleonic constituents (Δ isobars, multiquarks, etc.) can also contribute to the nuclear structure function. But if we consider only nucleons, then we must clearly use an effective constant, except, perhaps, for the case of the lightest nuclei, in which the non-nucleonic contributions can be expected to be slight. Further, $[S^{pA}-S^{nA}]$ is the matrix element of the operator σ_3 with respect to the wave function of nucleus A. It is precisely this matrix element that arises when one considers β-decay processes. Experiments on weak decays of mirror nuclei give the empirical quantity $(g_A/g_V)_{AA^*}$, which makes it possible to determine the renormalized ratio g_A/g_V of bound nucleons:

$$\left(\frac{g_A}{g_V}\right)_{AA^*}=\left(\frac{g_A}{g_V}\right)_{\text{bd. nucl}}\left\langle A\left|\sum_{i=1}^{A}\sigma_3(i)\tau_3(i)\right|A\right\rangle$$
$$\equiv\langle\sigma_3\tau_3\rangle\left(\frac{g_A}{g_V}\right)_{\text{bd. nucl}}. \quad (91)$$

Equations (90) and (91) illustrate the connection between deep inelastic scattering and weak processes. Using these relations and experimental information on β decay, and calculating the spin structure factors of the nuclei, we can extract information about the neutron structure function and about the ratio g_A/g_V from deep inelastic scattering by nuclei.

We now consider specific examples. We begin with the case of the simplest nucleus—the deuteron. It appears obvious that the properties of the nucleons in the deuteron differ least from the properties of free nucleons. Therefore, the deuteron is the most appropriate source of information about the neutron spin structure function, as in the unpolarized case.[92,94] The main problem in extracting the structure function of the "almost free" neutron from deuteron data is the determination and elimination of the nucleon Fermi motion (polarized or unpolarized). If we intend to work solely with the integrated quantities (87) and (89), the influence of the deuteron structure reduces to the factors $S^{pA,nA}$. (Polarization effects in deep inelastic scattering by deuterons were considered in Refs. 81 and 82. The spin momentum distributions of the nucleons in the deuteron

were obtained. However, the problem of extracting the neutron structure function and, therefore, the ratio g_A/g_V was not posed.) Direct calculations in (89) with a realistic deuteron wave function give

$$\int_0^1 dx g_1^D(x) = \int_0^1 dx[g_1^p(x) + g_1^n(x)]S^D, \tag{92}$$

$$S^{pD} = S^{nD} = S^D = 1 - \tfrac{3}{2} P_D, \tag{93}$$

where P_D is the weight of the D wave in the deuteron. Equations (92) and (93) provide a possibility for extracting the first moment of the structure function of a polarized neutron from the corresponding proton and deuteron structure functions. Ignoring the change in the properties of nucleons in the deuteron, we can readily find the required expression for $(g_A/g_V)_{pD} \cong g_A/g_V$:

$$\int dx(2g_1^p(x) - (S^D)^{-1} g_1^D(x)) = \frac{1}{6}\frac{g_A}{g_V}\left(1 - \frac{\alpha_s}{\pi}\right). \tag{94}$$

Equation (94) is a helpful relation for experimental verification of the Bjorken sum rule, or, if it is correct, for verification of QCD (for more details, see Ref. 91). Substituting the realistic value $P_D = 0.0425$ (Ref. 35) in (93), we find $S^D = 0.936$ a value that differs appreciably from unity because of the orbital motion of the nucleons in the deuteron (D wave). Allowance for this factor leads to corrections to the final result compared with the case when the deuteron is represented by a simple sum of a neutron and proton that are both at rest ($S^D = 1$). We now consider the example of mirror nuclei, for which we choose the simplest case—the pair ^{3}H–^{3}He. As was shown above, the difference between the structure functions of mirror nuclei is related to $(g_A/g_V)_{\text{bd. nucl}}$ by Eq. (90). The wave functions of the H and He nuclei are well known in few-body theory (see, for example, Refs. 95 and 96). The general form is

$$\Psi_A^{JM_J TM_T} = \sum_{L,m,\mu} R_{Lm}^{\nu}(r_1 r_2 r_3)(LMS\mu|JM_J)|S\mu, TM_T, \nu\rangle, \tag{95}$$

where J, T, M_J, M_T are the total angular momentum and total isospin of the nucleus and their projections corresponding to the considered nucleus; S and L are the spin and angular momentum; the index ν labels states with different symmetries; and $|S\mu,\ TM_T,\ \nu\rangle$ are spin–isospin states with symmetry ν. Ground states in the wave function (95) are the completely symmetric S state and the S' state of mixed symmetry with $L = 0$, and also the D state of mixed symmetry with $L = 2$. The wave function (95) permits explicit calculation of the spin factors:

$$\begin{aligned} S^{n^3\text{He}} &= S^{p^3\text{H}} = P_S + \frac{1}{3}P_{S'} - \frac{1}{3}P_D, \\ S^{p^3\text{He}} &= S^{n^3\text{H}} = \frac{2}{3}P_{S'} - \frac{2}{3}P_D, \\ S^{p^3\text{H}} - S^{n^3\text{H}} &= P_S - \frac{1}{3}P_{S'} + \frac{1}{3}P_D, \end{aligned} \tag{96}$$

where $P_{S'}$, P_S, P_D are the weights of the S, S', and D states in ^{3}He or ^{3}H, respectively. Using the typical values $P_S = 0.897$, $P_{S'} = 0.0017$, and $P_D = 0.086$, we find

$$S^{p^3\text{H}} - S^{n^3\text{H}} = 0.920, \tag{97}$$

and this value differs appreciably from the value of unity obtained in the standard model.[22] Further, the experimental value of (91) for $(g_A/g_V)_{A=3} \cong 1.2055$ (see Ref. 92) and the estimate (97) give

$$\left(\frac{g_A}{g_V}\right)_{\text{bd. nucl}} \approx 1.310, \tag{98}$$

in place of the 1.259 for free nucleons.[97] Because the ratio g_A/g_V is renormalized in a nucleus with $A = 3$ by about 4%, such nuclei are not very suitable sources of information about the neutron spin structure function. However, there is still a possibility to test the Bjorken sum rule. Substituting (97) and (98) in (90), we obtain

$$\begin{aligned} &\int_0^1 dx(g_1^{^3\text{H}}(x) - g_1^{^3\text{He}}(x)) \\ &= \left(1 - \frac{\alpha_s}{\pi}\right)\cdot(S^{p^3\text{H}} - S^{n^3\text{H}})\frac{1}{6}\left(\frac{g_A}{g_V}\right)_{\text{bd. nucl}} \approx 0.1836, \end{aligned} \tag{99}$$

which differs somewhat ($\sim$5%) from the result of the ordinary convolution model using the spin structure factor (97) (not unity!):

$$\begin{aligned} &\int dx(g_1^{^3\text{H}}(x) - g_1^{^3\text{He}}(x)) \\ &= \frac{1}{6}\frac{g_A}{g_V}\left(1 - \frac{\alpha_s}{\pi}\right)\cdot(S^{p^3\text{H}} - S^{n^3\text{H}}) \approx 0.1762. \end{aligned} \tag{100}$$

Allowance for these apparently slight ($\sim$5%) corrections is important because of the high sensitivity of the following analysis of the Bjorken sum rule and the spin content of the nucleons to small variations of the data used in this analysis. To end this section, we make some concluding remarks:

1. The analysis has shown that polarized deep inelastic scattering of leptons by nuclei encounters the same difficulties as the unpolarized case (EMC effect), namely, the usual convolution approach must be "corrected" to take into account the fact that (a) the properties of nucleons may be changed in the nuclear medium and (b) a simplified treatment of nuclear structure may lead to incorrect results. Above, we have taken into account the influence of the nuclear spin–orbit structure on the spin structure functions of the lightest nuclei. Further, using the "renormalized" constant g_A/g_V, we have also effectively taken into account other effects: the change of the nucleon properties in nuclei, non-nucleonic degrees of freedom (Δ isobars, multiquarks,...), etc.

2. To understand the possible role of mesonic exchange currents, we consider a "pedagogical" example. The operator averaged in (89) is very close to the usual operator of the nuclear magnetic moment:

$$\hat{\mu}_z=\mu_p \sum_p \sigma_z^p+\mu_n \sum_n \sigma_z^n+\sum_p l_z^p. \quad (101)$$

It is known that calculation of the magnetic moments of nuclei with allowance for only the nucleonic degrees of freedom leads to incorrect results. For example in the case of ^{3}He and ^{3}H the discrepancy with experiment is about 20%:

$$\mu(^3\mathrm{He})_{\mathrm{cal}}\cong 1.743; \quad \mu(^3\mathrm{He})_{\mathrm{exp}}\cong -2.128;$$

$$\mu(^3\mathrm{H})_{\mathrm{cal}}\cong 2.558; \quad \mu(^3\mathrm{H})_{\mathrm{exp}}\cong 2.978.$$

The deviation of μ_{cal} calculated using the P_S, $P_{S'}$, and P_D given above from the experimental value μ_{exp} can be naturally explained when allowance is made for mesonic exchange corrections.[98,99] We note that in the case of the neutron the deviation is only 0.2%, and it is obvious that here the mesonic corrections are small (therefore, renormalization effects can be ignored).

3. The relations established between the nuclear structure, the renormalized ratio g_A/g_V, and the spin structure functions of nuclei make it possible to predict the value of $(g_A/g_V)_A$ (87). In particular, according to the experimental values[92] of $(g_A/g_V)_A$, the Bjorken sum rule must have a nontrivial A dependence. On the other hand, it is possible to use (at least in principle) the expression (90) to extract the ratio (g_A/g_V) from deep inelastic scattering experiments.

4. The analysis indicates that extraction of the value of g_A/g_V from the combination of proton and deuteron data is just as good as extraction from data on deep inelastic scattering by mirror nuclei. In both cases the procedure is model-dependent, since it is impossible to avoid using model wave functions to calculate the nuclear factors $S^{pA,nA}$. In addition, the deuteron wave function has been better studied in nuclear physics than the other wave functions, and the effect of renormalization of (g_A/g_V) are minimal.

4. MANIFESTATION OF MULTIQUARK STRUCTURE OF NUCLEI IN DEEP INELASTIC PROCESSES

As a rule, the occurrence of a high-momentum component in the study of deep inelastic scattering is attributed to NN interaction at short distances leading to the formation of strongly correlated nucleon clusters that can participate as single objects in a reaction with short interaction time. In this section, we shall investigate the manifestation of multibaryon–multiquark correlations in nuclei in deep inelastic scattering. The problem of multiquark interactions was posed already in 1971.[100] Later, ideas were developed about the existence in nuclei of multiquark systems, a prototype of which is provided by Blokhintsev's nuclear density fluctuations.[101] Therefore, following the tradition of these studies,[21,49,102] we shall here often refer to multiquark systems as "fluctons." Fluctons are unstable objects that, in principle, can be detected by reactions with short interaction time and are manifested primarily in the high-momentum part of nuclear spectra.

Chromodynamic evolution of nuclear structure functions near the limit of single-nucleon kinematics

The role of multiquark states (high-momentum component) has been studied in detail in cumulative h reactions[48] at comparatively small momentum transfers ($Q^2 \leqslant 10$ GeV$^2/c^2$). In deep inelastic scattering of leptons, the momentum transfer is comparatively large ($Q^2 \sim$ 100–300 GeV$^2/c^2$), and the behavior of a multiquark distribution in the region $x\sim 1$ may differ appreciably from distributions for which the parameters are determined from hA reactions. A first obvious source of this difference is the logarithmic violation of scaling. A qualitative analysis of the Q^2 dependence of the nuclear structure functions for $x\geqslant 1$ without allowance for Fermi motion of the nucleons was given in Ref. 103. However, since in the boundary region for single-nucleon kinematics the contribution of the Fermi motion is important, to obtain quantitative estimates it is necessary to investigate the total contribution of the Fermi motion and the multiquark components to the structure function $F_2^A(x,Q^2)$. This problem is especially topical because the effect of scaling violation in the structure function of the free nucleon increases as $x\to 1$.

The structure function of a nucleus A with allowance for $6q$ admixtures at given Q^2 can be represented in the form

$$F_2^A(x)=(1-P_A)\int dy f^{N/A}(y)F_2^N(x/y) + P_A F_2^{6q}(x/2), \quad (102)$$

where F_2^{6q} is the flucton structure function (in the region $0.8\leqslant x\leqslant 1.0$ that we consider, the main contribution is made by $6q$ configurations; we ignore multiquarks of higher orders); P_A is the corresponding probability.

The distribution function $f^{N/A}(y)$, calculated in the usual shell model [see, for example, (58)], gives an unsatisfactory description of the data in the noncumulative region (Fig. 7), and it is therefore natural to use a function $f^{N/A}(y)$ that takes into account more complicated collective effects. For this, it is again convenient to use the relatively simple model of coherent fluctuations of the nuclear density [see Eqs. (72)–(77)].

In principle, the isoscalar nucleon structure function $F_2^N(x)$ in (102) must be determined from the condition of best description of the data on deep inelastic scattering of leptons by the proton and neutron (actually from a combination of the proton and deuteron data with allowance for the Fermi motion and the fact that the nucleons in the deuteron are bound) for given Q^2. Unfortunately, the experimental data are at present insufficient for reliable parametrization of $F_2^N(x,Q^2)$. Therefore, in Refs. 104 and 105 we used in the range $Q^2 \sim$ 15–25 GeV$^2/c^2$ the simple parametrization from Refs. 39 and 103, which gives a satisfactory description of the EMC data[93] and CDHS data.[17] The results of calculations with this function in the region $x\sim 0.8$–1.2 are practically identical to those using the function in Refs. 44, 54, and 75. The parameters $F_2^{6q}(x)$ were also determined for comparatively small Q^2 (Refs. 48 and 103):

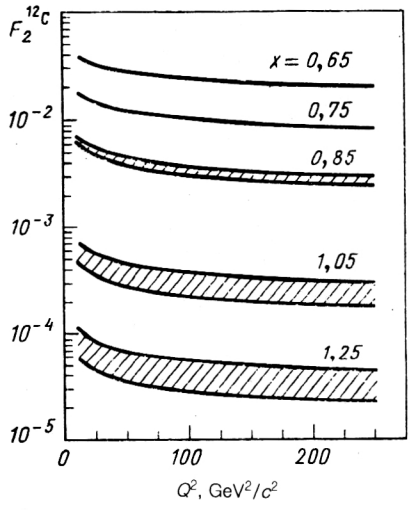


FIG. 13. The Q^2 evolution of the carbon structure function with allowance for the effects of the off-shell nucleons and a six-quark admixture for different x.

$$F_2^{6q}(x,Q_0^2)=\tfrac{5}{18}\,4.7x^{0.5}(1-x)^7+\tfrac{2}{3}\,0.23(1-x)^{11}. \qquad (103)$$

Thus, we can calculate the nuclear structure function $F_2^A(x,Q_0^2)$ with $Q_0^2 \sim 15$ GeV$^2/c^2$. To make a comparison with the BCDMS experimental data ($Q_0^2 \sim$ 100–300 GeV$^2/c^2$) we must, in the first place, take into account the usual logarithmic violation of scaling (Q^2 evolution). We solved the Lipatov–Altarelli–Parisi evolution equations[106,107] in the leading order with QCD dimensional parameter $\Lambda \cong 0.21$–0.23 GeV and initial conditions (104) at $Q_0^2 \cong 15$ GeV$^2/c^2$. It is obvious that these initial conditions inescapably contain a certain arbitrariness, which is mainly due to the errors in the determination of the nucleon structure function from the incomplete experimental data on the proton and neutron structure functions. A second source of arbitrariness is associated with the uncertainties in the determination of the contribution of the multiquark configurations to both $F_2^{6q}(x)$ and P_A. These uncertainties arise from the impossibility of a sufficiently correct allowance for secondary nuclear effects such as hadron rescattering in the initial and final states, nuclear absorption, etc. Figure 13 gives the results of numerical calculations of the Q^2 evolution for the structure function of carbon at different values of x with allowance for both binding of the nucleons in the nucleus and the contribution of the flucton component. The hatched regions in Fig. 13 reflect the uncertainties in the initial conditions mentioned above; the probability P_A of a $6q$ admixture in the carbon nucleus was varied from 5 to 10% up to $Q^2 = 150$ GeV$^2/c^2$, depending on the value of x. Allowance for the Fermi motion of the bound nucleons satisfactorily describes the experimental data in the noncumulative region. The contribution of the multiquark component is comparable with the contribution of the Fermi motion near the limit of the single-nucleon kinematics and begins to be dominant with increasing values of x. Inclusion of the effects of the off-shell behavior of the nucleons and the Q^2 evolution leads to a significant softening of the high-momentum component in the nucleus. This result should be taken into account when a comparison is made of experimental data on nuclear structure functions obtained in different processes: lepton–nucleus (large Q^2) or hadron–nucleus (intermediate and small Q^2).

Multiquark configurations in deep inelastic scattering by polarized deuterons

In Sec. 3, we have already considered the effects of the nuclear structure of the lightest nuclei in the deep inelastic scattering of polarized leptons by polarized nuclei. In particular, we demonstrated the qualitatively different behavior of the ratios of the structure functions of the deuteron and the "isoscalar" nucleon in the cases of the functions $g_1(x)$ and $F_2(x)$. We now consider the influence on the behavior of the spin structure functions of a possible admixture of the $6q$ component in the deuteron wave function. If it is assumed that a six-quark flucton is present in the deuteron with probability P_D, then the expression (84) for $g_1^D(x)$ is modified to

$$g_1^D(x)=(1-P_D)\int dy/y\,p_{\rm pol}^{N/A}(y)g_1^N(x/y) + P_D g_4^{6q}(x/2), \qquad (104)$$

where to estimate the spin structure function of the six-quark flucton we can use, for example, the model of Ref. 47 for a $6q$ bag with quantum numbers of the deuteron and quarks in the S shell; the spin distributions of the quark flavors (u, d) with respect to x are taken to be the same as in the unpolarized case. It is clear that, as in the case of nucleons, such an estimate has a very approximate nature, but it can be used for a qualitative analysis of the role of the multiquark admixture in the spin structure function of the deuteron. The explicit expression for $q_1^{6q}(x)$ can be written in the form

$$g_1^{6q}(x)=\tfrac{1}{9}\,(2u_v^{6q}(x)+d_v^{6q}(x)), \qquad (105)$$

where $u_v^{6q}(x)$ and $d_v^{6q}(x)$ and the probability P_D are taken from Refs. 39 and 103.

Analysis of $g_1^D(x)$ (104) permits the following conclusions: a) allowance for the contribution of the $6q$ component qualitatively changes the behavior of the ratio $R_g(x) = g_1^D(x)/g_1^n(x)$ in the complete range of x values; b) the ratio R_g is, compared with $R_F(x) = F_2^D(x)/F_2^N(x)$, more sensitive to the presence of a $6q$ admixture in the deuteron. We note that allowance for more complicated quark configurations of a $6q$ bag in the deuteron (for example, $S+P$ wave[108]) leads to a certain decrease of R_g in the region of small and intermediate x.

CONCLUSIONS

Thus, in high-energy nuclear physics the investigation of the different forms of "exotics" is intimately related to calculations of the ordinary nucleon structure of nuclei, which, of course, have their own "high-energy" specifics.

Multisided investigations of non-nucleonic degrees of freedom in nuclei extend our knowledge about nuclear structure and, simultaneously, enable us to understand more deeply the nature of the strong interaction of elementary particles.

At the present time, theoretical investigations of the effects of nuclear structure and non-nucleonic degrees of freedom also have a special topical importance in view of the large number of major experimental programs, both operating and planned. In deep inelastic scattering of leptons, there are, in the first place, the experiments, including polarization experiments, of the EMC, BCDMS, and SLAC collaborations; in hadron–nucleus and nucleus–nucleus cumulative processes there are experiments of the High Energy Laboratory of the Joint Institute for Nuclear Research, the Institute of Theoretical and Experimental Physics, and others. Finally, the program of the CE-BAF accelerator in the United States includes an important part relating to nuclear "exotics."

[1] C. Itzykson and J.-B. Zuber, *Quantum Field Theory* (McGraw-Hill, New York, 1980) [Russ. transl., Mir, Moscow, 1984].
[2] F. J. Ynduráin, *Quantum Chromodynamics: An Introduction to the Theory of Quarks and Gluons* (Springer-Verlag, New York, 1983) [Russ. transl., Mir, Moscow, 1986].
[3] F. E. Close, *An Introduction to Quarks and Partons* (Academic Press, London, 1979) [Russ. transl., Mir, Moscow, 1982].
[4] E. D. Bloom and F. J. Gilman, Phys. Rev. Lett. **9**, 1140 (1970).
[5] E. D. Bloom, R. L. Cottrell, D. H. Coward *et al.*, *Proc. of the 15th Intern. Conf. on High Energy Physics*, Vol. 1 (Kiev, 1970), p. 324.
[6] D. Miller, E. D. Bloom, E. D. Cottrell *et al.*, *Proc. of the 15th Intern. Conf. on High Energy Physics*, Vol. 1 (Kiev, 1970), p. 323.
[7] J. D. Bjorken, Phys. Rev. **179**, 1547 (1969).
[8] N. N. Nikolaev, Fiz. Elem. Chastits At. Yadra **12**, 162 (1981) [Sov. J. Part. Nucl. **12**, 63 (1981)].
[9] G. Berland, A. Dar, and G. Eilam, Phys. Rev. D **22**, 1547 (1980).
[10] A. Bodek and J. L. Ritchie, Phys. Rev. D **23**, 1070 (1981).
[11] J. Ashman, B. Badelek, G. Baum *et al.*, Phys. Lett. **202B**, 603 (1988).
[12] J. J. Aubert, G. Bassompierre, K. H. Becks *et al.*, Phys. Lett. **123B**, 275 (1983).
[13] A. Bodek, N. Ciokaris, W. B. Atwood *et al.*, Phys. Rev. Lett. **50**, 1431 (1983).
[14] A. C. Benvenuti, D. Bollini, G. Bruni *et al.*, Phys. Lett. **189B**, 483 (1987).
[15] G. Bari, A. C. Benvenuti, D. Bollini *et al.*, Phys. Lett. **163B**, 282 (1985).
[16] R. G. Arnold, P. E. Bosted, C. C. Vhang *et al.*, Phys. Rev. Lett. **52**, 727 (1984).
[17] H. Abramowicz, G. Hansl-Kozanecki, J. May *et al.*, Z. Phys. C **25**, 29 (1984).
[18] I. A. Savin, Preprint R1-85-315 [in Russian], JINR, Dubna (1985).
[19] R. P. Jaffe, *Relativistic Dynamics and Quark Nuclear Physics* (Wiley, New York, 1986).
[20] R. P. Feynman, *Photon–Hadron Interactions* (Benjamin, Reading, Mass., 1972) [Russ. transl., Mir, Moscow, 1975].
[21] A. V. Efremov and A. V. Radyushkin, Riv. Nuovo Cimento **3**, 1 (1980).
[22] F. E. Close, R. G. Roberts, and G. G. Ross, Nucl. Phys. **B296**, 582 (1988).
[23] J. Kuti and V. Wisskopf, Phys. Rev. D **4**, 3418 (1971).
[24] P. S. Isaev and S. G. Kovalenko, Hadronic J. **3**, 919 (1980).
[25] S. Date, Prog. Theor. Phys. **70**, 1682 (1983).
[26] V. A. Matveev, R. M. Muradyan, and A. N. Tavkhelidze, Lett. Nuovo Cimento **7**, 719 (1973).
[27] V. de Alfaro *et al.*, *Currents in Hadron Physics* (North-Holland, Amsterdam, 1973) [Russ. transl., Mir, Moscow, 1976].
[28] R. L. Jaffe, F. E. Close, R. G. Roberts, and G. G. Ross, Phys. Rev. D **31**, 1004 (1985).
[29] F. E. Close, R. G. Roberts, and G. G. Gross, Phys. Lett. **129B**, 346 (1983).
[30] J. Cleymans and R. L. Thews, Phys. Rev. D **31**, 1014 (1985).
[31] S. V. Akulinichev, G. M. Vagradov, and S. A. Kulagin, Pis'ma Zh. Eksp. Teor. Fiz. **42**, 105 (1985) [JETP Lett. **42**, 127 (1985)].
[32] S. V. Akulinichev, S. A. Kulagin, and G. M. Vagradov, Phys. Lett. **158B**, 475 (1985); J. Phys. G **11**, L245 (1985).
[33] B. L. Birbrair, E. M. Levin, and A. G. Shuvaev, Phys. Lett. **166B**, 119 (1986).
[34] A. N. Antonov, L. P. Kaptari, V. A. Nikolaev, and A. Yu. Umnikov, Rapid Communications, Vol. 2 [41]-90, JINR, Dubna (1990), p. 14.
[35] R. Machleid, K. Holinde, and C. Elster, Phys. Rep. **149**, 1 (1987).
[36] R. Dymarz and F. C. Khanna, Nucl. Phys. **A507**, 53 (1990).
[37] A. Yu. Umnikov, *Structure Functions of Nuclei in Deep Inelastic Processes* [in Russian] (Far Eastern State University, Vladivostok, 1987).
[38] C. H. Llewellyn Smith, Phys. Lett. **128B**, 107 (1983).
[39] A. I. Titov, Yad. Fiz. **40**, 76 (1984) [Sov. J. Nucl. Phys. **40**, 50 (1984)].
[40] E. L. Berger, F. Coester, and R. B. Wiringa, Phys. Rev. D **29**, 398 (1984).
[41] É. E. Sapershteĭn and M. Zh. Shmatikov, Pis'ma Zh. Eksp. Teor. Fiz. **41**, 44 (1985) [JETP Lett. **41**, 53 (1985)].
[42] M. Ericson and A. W. Thomas, Phys. Lett. **128B**, 112 (1983).
[43] B. L. Birbrair, E. M. Levin, and A. G. Shuvaev, Nucl. Phys. **A491**, 618 (1989).
[44] L. P. Kaptar', B. L. Reznik, A. I. Titov, and A. Yu. Umnikov, Pis'ma Zh. Eksp. Teor. Fiz. **47**, 428 (1988) [JETP Lett. **47**, 508 (1988)].
[45] A. M. Baldin, V. K. Bondarev, N. Ghiordanescu *et al.*, Preprint E2-82-472, JINR, Dubna (1982).
[46] V. B. Gavrilov and G. A. Leksin, Preprint 124 [in Russian], Institute of Theoretical and Experimental Physics, Moscow (1983).
[47] V. A. Matveev and P. Sorba, Lett. Nuovo Cimento **20**, 435 (1977).
[48] V. V. Burov, V. K. Luk'yanov, and A. I. Titov, Fiz. Elem. Chastits At. Yadra **15**, 1249 (1984) [Sov. J. Part. Nucl. **15**, 558 (1984)].
[49] V. K. Luk'yanov and A. I. Titov, Fiz. Elem. Chastits At. Yadra **10**, 815 (1979) [Sov. J. Part. Nucl. **10**, 321 (1979)].
[50] A. V. Efremov, Fiz. Elem. Chastits At. Yadra **13**, 613 (1982) [Sov. J. Part. Nucl. **13**, 254 (1982)].
[51] L. A. Kondratyuk and M. Zh. Shmatikov, Yad. Fiz. **41**, 222 (1985) [Sov. J. Nucl. Phys. **41**, 141 (1985)].
[52] A. V. Efremov and E. A. Bondarchenko, Preprint E2-84-124, JINR, Dubna (1984).
[53] R. L. Jaffe, F. E. Close, R. G. Roberts, and G. G. Ross, Phys. Lett. **13B**, 449 (1984).
[54] E. L. Bratkovskaya, L. P. Kaptari, A. I. Titov, and A. Yu. Umnikov, Nucl. Phys. **A512**, 684 (1990).
[55] S. A. Kulagin, Nucl. Phys. **A500**, 653 (1989).
[56] H. Jung and G. A. Miller, Phys. Lett. **200B**, 351 (1988).
[57] M. Chemtob and M. Rho, Nucl. Phys. **A163**, 1 (1971).
[58] J. F. Mathiot, Phys. Lett. **115B**, 174 (1982).
[59] M. Rosa-Clot and M. Testa, Nuovo Cimento **A78**, 113 (1983).
[60] A. I. Akhiezer and V. B. Berestetskii, *Quantum Electrodynamics* (Wiley, New York, 1965).
[61] H. A. Bethe and F. de Hoffmann, *Mesons and Fields* (Wiley, New York, 1955).
[62] C. G. Callan, S. Coleman, and R. Jackiw, Ann. Phys. (N.Y.) **59**, 42 (1970).
[63] J. D. Sullivan, Phys. Rev. D **5**, 1732 (1972).
[64] G. G. Arakelyan and K. G. Boreskov, Yad. Fiz. **41**, 416 (1985) [Sov. J. Nucl. Phys. **41**, 267 (1985)].
[65] L. L. Frankfurt and M. I. Strikman, Phys. Rep. **76C**, 215 (1981).
[66] L. L. Frankfurt and M. I. Strikman, Phys. Lett. **183B**, 254 (1987).
[67] G. L. Li, K. F. Liu, and G. E. Brown, Phys. Lett. **213B**, 531 (1988).
[68] C. Ciofidegli Atti and S. Luiti, Phys. Lett. **225B**, 215 (1989).
[69] E. A. Ivanov and E. Truhlík, Fiz. Elem. Chastits At. Yadra **12**, 492 (1981) [Sov. J. Part. Nucl. **12**, 198 (1981)].
[70] M. Gari and H. Hyuga, Z. Phys. A **277**, 291 (1976); Nucl. Phys. **A264**, 409 (1976).
[71] L. D. Miller and A. E. Green, Phys. Rev. C **5**, 241 (1972).
[72] B. L. Birbrair, E. M. Levin, and A. G. Shuvaev, Preprint LNPI-1296, Leningrad (1987).
[73] B. L. Birbrair, E. M. Levin, and A. G. Shuvaev, Phys. Lett. **222B**, 281 (1989).
[74] G. E. Brown and A. D. Jackson, *The Nucleon–Nucleon Interaction* (New York, 1976).

[75] L. P. Kaptar', A. I. Titov, and A. Yu. Umnikov, Yad. Fiz. **51**, 864 (1990) [Sov. J. Nucl. Phys. **51**, 549 (1990)].
[76] W. Zhu and J.-G. Shen, Commun. Theor. Phys. **9**, 51 (1988).
[77] A. N. Antonov, P. E. Hodgson, and I. Zh. Petkov, *Nucleon Momentum and Density Distribution in Nuclei* (Clarendon Press, Oxford, 1988).
[78] A. N. Antonov and I. Zh. Petkov, Nuovo Cimento **A94**, 68 (1986).
[79] A. N. Antonov, V. A. Nikolaev, and I. Zh. Petkov, Z. Phys. A **297**, 257 (1980).
[80] V. K. Luk'yanov and Yu. S. Pol', Fiz. Elem. Chastits At. Yadra **5**, 955 (1974) [Sov. J. Part. Nucl. **5**, 385 (1975)].
[81] P. Hoodbhoy, R. L. Jaffe, and A. Manohar, Nucl. Phys. **B312**, 571 (1989).
[82] L. P. Kaptar', A. I. Titov, and A. Yu. Umnikov, Preprint R2-86-26 [in Russian], JINR, Dubna (1986).
[83] I. I. Bazhanskiĭ, L. P. Kaptar', B. L. Reznik *et al.*, in *Proc. of the Eighth International Seminar on Problems of High Energy Physics*, D1-2-86-668, Vol. 1 [in Russian] (JINR, Dubna, 1987), p. 318.
[84] L. P. Kaptari and A. Yu. Umnikov, Phys. Lett. **240B**, 203 (1990).
[85] D. Yu. Bardin, Yad. Fiz. **33**, 707 (1981) [Sov. J. Nucl. Phys. **33**, 368 (1981)].
[86] J. Ashman, B. Badelek, G. Baum *et al.*, Phys. Lett. **206B**, 364 (1988).
[87] V. A. Karmanov and I. S. Shapiro, Fiz. Elem. Chastits At. Yadra **9**, 327 (1978) [Sov. J. Part. Nucl. **9**, 134 (1978)].
[88] V. R. Garsevanishvili and V. A. Matveev, Teor. Mat. Fiz. **24**, 3 (1975).
[89] M. Lacombe, B. Loiseau, J. M. Richard, and R. Vinh Mau, Phys. Rev. C **21**, 861 (1980).
[90] J. D. Bjorken, Phys. Rev. **148**, 1467 (1966).
[91] G. Preparata and J. Soffer, Phys. Rev. Lett. **61**, 1167 (1988).
[92] S. Raman, C. A. Houser, T. A. Walkiéwicz, and I. S. Towner, At. Data Nucl. Data Tables **21**, 567 (1978).
[93] J. J. Aubert, G. Bassompierre, G. Becks *et al.*, Nucl. Phys. **B293**, 740 (1987).
[94] A. C. Benvenutti, D. Bollini, G. Bruni, Preprint EP/89-171, CERN (1989).
[95] A. C. Phillips, Rep. Prog. Phys. **40**, 905 (1977).
[96] J. L. Friar, B. F. Gibson, and G. L. Payne, Ann. Rev. Nucl. Sci. **34**, 403 (1984).
[97] Review of Particle Properties, Particle Data Group, Phys. Lett. **204B** (1988).
[98] E. P. Harper, Y. E. Kim, and A. Tubis, Phys. Lett. **40B**, 533 (1972).
[99] W. Jaus, Nucl. Phys. **A271**, 495 (1974).
[100] A. M. Baldin, Kratk. Soobshch. Fiz. **1**, 35 (1971).
[101] D. I. Blokhintsev, Zh. Eksp. Teor. Fiz. **33**, 1295 (1957) [Sov. Phys. JETP **6**, 995 (1958)].
[102] D. I. Blokhintsev, A. V. Efremov, V. K. Luk'yanov, and A. I. Titov, in *International Seminar on Problems of High Energy Physics*, 1, 2-12036 [in Russian] (JINR, Dubna, 1978).
[103] A. I. Titov, Yad. Fiz. **38**, 1582 (1983) [Sov. J. Nucl. Phys. **38**, 964 (1983)].
[104] L. P. Kaptar', B. L. Reznik, A. I. Titov, and A. Yu. Umnikov, in *Proc. of the Ninth International Seminar on Problems of High Energy Physics*, D12-88-652, Vol. 2 [in Russian] (JINR, Dubna, 1988).
[105] L. P. Kaptari and A. Yu. Umnikov, Rapid Communication No. 6 [32]-88, JINR, Dubna (1988), p. 17.
[106] G. Altarelli, Phys. Rep. **81**, 1 (1982).
[107] L. P. Lipatov, Yad. Fiz. **20**, 181 (1974) [Sov. J. Nucl. Phys. **20**, 94 (1975)].
[108] V. G. Neudatchin, I. T. Obukhovsky, Yu. F. Smirnov, and E. Y. Tkaya, Z. Phys. A **313**, 357 (1983).

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