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EVOLUTION OF THE NUCLEON STRUCTURE IN LIGHT NUCLEI

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The evolution of the EMC effect as a function of atomic mass A is considered for the first time for the lightest nuclei, D, ^{3}He and, ^{4}He, with an approach based on the Bethe–Salpeter formalism. We show that the pattern of the oscillation of the ratio $r^A(x) = F_2^A / F_2^{N(D)}$ with respect to the line $r^A(x) = 1$ varies with A, unlike the pattern for nuclei with masses $A > 4$, where only the amplitude of the oscillation changes. It is found that the shape of the structure function distortions, which is typical for metals, is being reached in ^{3}He.

The investigation has been performed at the Bogoliubov Laboratory of Theoretical Physics, JINR.

Эволюция структуры нуклона в легких ядрах

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Эволюция эффекта ЕМС как функция атомной массы A впервые рассмотрена для легчайших ядер D, ^{3}He и ^{4}He в рамках единого подхода, в основе которого — формализм Бете–Солпитера. Показано, что форма осцилляции $r^A(x) = F_2^A / F_2^{N(D)}$ относительно линии $r^A(x) = 1$ преобразуется с ростом A, отличаясь тем самым от случая с ядрами $A > 4$, где изменяется только амплитуда осцилляции. Получено, что типичная для металлов форма искажений структурной функции достигается в ядре ^{3}He.

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Strictly speaking, the EMC effect, first observed as a distortion of the free nucleon structure function in an iron nucleus [1], is rather the phenomenon of differences between the deuteron $F_2^D(x)$ and helium $F_2^{^4\mathrm{He}}(x)$ structure functions. Indeed, as demonstrated in [2,3], well established by the experiments of SLAC [4] and NMC [5] the oscillation pattern of the ratio $r(x) = F_2^{^4\mathrm{He}} / F_2^D$ in the range $x < 0.9$ is retained in heavier nuclei. The pattern of the oscillations for nuclei with masses $A > 4$ is fixed by the positions of the three cross-over points $x_i, i = 1–3$, in which $r^A(x) = 1$ independently of A [3]. The evolution of the EMC effect in the range of masses from $A = 4$ to $A \sim 200$ manifests itself in the increase of the oscillation amplitude, $a_{\mathrm{EMC}} = 1 - r^A_{\mathrm{min}}$, by a factor of $\sim$ 3. It is well understood as a nuclear density effect if the surface nucleons are excluded from consideration. There does not exist any data on the EMC effect in the range of $A < 4$. Most challenging therefore is to find how the effect evolves in the range of the lightest nuclear masses, $r^D(x) \to r^{A=3}(x) \to r^{A=4}(x)$.

Below we perform derivations of the relative changes in the nuclear structure function $F_2^A(x)$ with respect to the isoscalar nucleon one, $F_2^N(x) = \frac{1}{2}(F_2^p(x) + F_2^n(x))$, where p and n denote the free proton and the free neutron respectively. On the other hand, the comparison with experimental data can be done only in terms of $r^A(x)$, obtained with the deuteron structure function $F_2^D(x)$ as a reference. Therefore our final results will be presented for both cases. In the considered range of x $(0.3 < x < 0.9)$ the experiments (see Ref. 4) are consistent with no Q^2 dependence of $r^A(x)$. We perform numerical calculations for the fixed $Q^2 = 10$ GeV^2.

In the present paper, we consider a model independent calculation of $F_2^A(x)$. Our approach originates from the Bethe–Salpeter formalism [6], which allows one to treat nuclear binding effects by using general properties of nucleon Green functions. A basic assumption which dictates the entire approximation is that the nuclear fragments have a small relative energy. It allows one, as has been shown previously [7], to derive the deuteron structure function in the form:

$$F_2^D(x_D) = \int \frac{d^3k}{(2\pi)^3} \left(F_2^N(x_N) \left(1 - \frac{k_3}{E} \right) - \frac{M_D - 2E}{E} x_N \frac{dF_2^N(x_N)}{dx_N} \right) \Psi^2(\mathbf{k}), \qquad (1)$$

where E is the on-mass-shell nucleon energy $E^2 = \mathbf{k}^2 + m^2$, and k is the relative four momentum of the bound nucleons, M_D is the mass of a deuteron. The nucleon Bjorken variable is defined as $x_N = x_D m/(E - k_3)$. The function $\Psi^2(\mathbf{k})$ is an analog of the three dimensional momentum distribution and is directly related to the Bethe–Salpeter vertex function for the deuteron $\Gamma^D(P, k)$:

$$\Psi^2(\mathbf{k}) = \frac{m^2}{4E^2 M_D (M_D - 2E)^2} \left\{ \Phi^2(M_D, k) \right\}_{k_0 = E - \frac{M_D}{2}},$$

$$\Phi^2(M_D, k) = \bar{\Gamma}^D(M_D, k) \left(\sum_s u^s(\mathbf{k}) \bar{u}^s(\mathbf{k}) \right) \otimes \left(\sum_s u^s(-\mathbf{k}) \bar{u}^s(-\mathbf{k}) \right) \Gamma^D(M_D, k). \qquad (2)$$

Expression (1) shows that the relative time dependence of the amplitude of lepton deep inelastic scattering (DIS) off a bound nucleon results in a depletion of the ratio F_2^D/F_2^N for $0.2 < x < 0.7$. In this paper we briefly describe the extension of this approach for light nuclei, $A = 3, 4$.

Using the unitarity condition one can relate the hadronic part of the DIS amplitude (hadronic tensor) with the forward Compton scattering amplitude,

$$W^A_{\mu\nu}(P, q) = \frac{1}{2\pi} \, \mathrm{Im} \, T^A_{\mu\nu}(P, q), \tag{3}$$

which is defined as a product of electromagnetic currents averaged over nuclear states,

$$T^A_{\mu\nu}(P, q) = i \int d^4x e^{iqx} \langle A \,|\, T(J_\mu(x) \, J_\nu(0)) \,|\, A \rangle. \tag{4}$$

Starting from the field theory framework one can define the matrix element in terms of solutions of the n-nucleon Bethe–Salpeter equation and n-nucleon Green functions with insertion of the T-product of electromagnetic currents:

$$T^A_{\mu\nu}(P, q) = \int \frac{d^4k_1}{(2\pi)^4} \cdots \frac{d^4k_{n-1}}{(2\pi)^4} \frac{d^4k'_1}{(2\pi)^4} \cdots \frac{d^4k'_{n-1}}{(2\pi)^4} \, \overline{\Gamma}(P, k_1, \ldots, k_{n-1}) \, S_{2n}(P, k_1, \ldots, k_{n-1}) \times$$

$$\times \, \overline{G}_{2(n+1)_{\mu\nu}}(q; P, k_1, \ldots, k_{n-1}, k'_1, \ldots, k'_{n-1}) \, S_{2n}\,(P, k'_1, \ldots, k_{n-1}) \, \Gamma(P, k'_1, \ldots, k'_{n-1}). \tag{5}$$

The Bethe–Salpeter vertex function $\Gamma(P, k_1, \ldots, k_{n-1})$ satisfies the homogeneous Bethe–Salpeter equation

$$\Gamma(P, k_1, \ldots, k_{n-1}) = - \int \frac{d^4k_1}{(2\pi)^4} \cdots \frac{d^4k_{n-1}}{(2\pi)^4} \times$$

$$\times \, \overline{G}_{2n}(P, k_1, \ldots, k_{n-1}; k'_1, \ldots, k'_{n-1}) \, S_{2n}(P, k'_1, \ldots, k'_{n-1}) \, \Gamma(P, k'_1, \ldots, k'_{n-1}). \tag{6}$$

Here k_i are Jacoby relative momenta of the nucleons inside a nucleus, and P is the total momentum of the nucleus. The $S_{2n}(P, k'_1, \ldots, k'_{n-1})$ is a direct product of n nucleons propagators. The $\overline{G}_{2n}(P, k_1, \ldots, k_{n-1}, k'_1, \ldots, k'_{n-1})$ term denotes the irreducible truncated n-nucleon Green function.

All irreducible interaction corrections to the imaginary part of the Compton amplitude are suppressed by powers of $1/(-q^2)$ [7]. This justifies consideration of the zeroth order term of $\overline{G}_{2(n+1)_{\mu\nu}}$:

$$T^A_{\mu\nu}(P, q) = \int \frac{d^4k_1}{(2\pi)^4} \cdots \frac{d^4k_{n-1}}{(2\pi)^4} \, \overline{\Gamma}(P, k_1, \ldots, k_{n-1}) \, S_{2n}(P, k_1, \ldots, k_{n-1}) \times$$

$$\times \, \overline{G^{(0)}_{2(n+1)_{\mu\nu}}}(q; P, k_1, \ldots, k_{n-1}) \, S_{2n}(P, k_1, \ldots, k_{n-1}) \, \Gamma(P, k_1, \ldots, k_{n-1}), \tag{7}$$

where the Green function $\overline{G^{(0)}}_{2(n+1)_{\mu\nu}}$ is defined via the truncated nucleon Compton amplitude $G_{4\mu\nu}(q;\, P,\, k_i)$:

$$\overline{G^{(0)}_{2(n+1)_{\mu\nu}}}(q;\, P,\, k_1,\dots,\, k_{n-1}) =$$

$$= \sum_i G_{4\mu\nu}(q;\, P,\, k_i) \otimes S^{-1}_{2n-1}(k_1,\dots,\, k_{i-1},\, k_{i+1},\dots,\, k_{n-1}). \tag{8}$$

Substituting this expression into Eq.(7) one could get the Compton amplitude of a nucleus in terms of the nucleon one:

$$T^A_{\mu\nu}(P,\, q) = \int \frac{d^4k_1}{(2\pi)^4} \cdots \frac{d^4k_{n-1}}{(2\pi)^4} \bar{\Gamma}(P,\, k_1,\dots,\, k_{n-1}) \sum_i (S(P,\, k_i)\, G_{4\mu\nu}(q;\, P,\, k_i) S(P,\, k_i)) \otimes$$

$$\otimes\, S_{2n-1}(P,\, k_1,\dots,\, k_{i-1},\, k_{i+1},\dots,\, k_{n-1})\, \Gamma(P,\, k_1,\dots,\, k_{n-1}). \tag{9}$$

However, unlike the deuteron case, where singularities in the Bethe–Salpeter vertex function can be neglected, in this case there are singularities connected with nucleon-nucleon bound states, which lie in the range of low relative momenta. Having obtained exact solutions of the Bethe–Salpeter equation one could obtain the Compton amplitude in terms of the nucleon one. Presently there are no solutions for Eq.(6) with $n \geq 3$, so we investigate a possibility of expressing the nuclear amplitude in terms of amplitudes of the physical nuclear fragments and three dimensional momentum distributions. To do this we have to remove the singularities explicitly.

One can remove the singularities by introducing the «bare» BS vertex function γ, which is regular with respect to the relative nucleon momenta:

$$\Gamma(P,\, k_1,\dots,\, k_{n-1}) = -\int \frac{d^4k_1'}{(2\pi)^4} \cdots \frac{d^4k_{n-1}'}{(2\pi)^4} \times$$

$$\times\, G_{2n}(P,\, k_1,\dots,\, k_{n-1};\, k_1',\dots,\, k_{n-1}')\, S_{2n}(P,\, k_1',\dots,\, k_{n-1}')\, \gamma(P,\, k_1',\dots,\, k_{n-1}'). \tag{10}$$

In case of ^{3}He we have a pole in G_4 connected with a bound deuteron and nucleon-nucleon continuous spectrum g_4:

$$G_4\left(\frac{2P}{3} + k,\, k_1,\, k_1'\right) = \frac{\Gamma^D\left(\frac{2P}{3} + k,\, k_1\right)\bar{\Gamma}^D\left(\frac{2P}{3} + k,\, k_1'\right)}{\left(\frac{2P}{3} + k\right)^2 - M_D^2} + g_4\left(\frac{2P}{3} + k,\, k_1,\, k_1'\right). \tag{11}$$

For ^{4}He one has additionally the ^{3}He and ^{3}H poles. For example, for the neutron-proton-proton Green function one has the ^{3}He pole and three-nucleon continuous spectrum g_6:

$$G_6\left(\frac{3P}{4}+k,\, k_1,\, k_1',\, k_2,\, k_2'\right)=$$

$$=\frac{\Gamma^{{}^3\mathrm{He}}\left(\frac{3P}{4}+k,\, k_1,\, k_2\right)\bar{\Gamma}^{{}^3\mathrm{He}}\left(\frac{3P}{4}+k,\, k_1',\, k_2'\right)}{\left(\frac{3P}{4}+k\right)^2-M^2_{{}^3\mathrm{He}}}+g_6\left(\frac{3P}{4}+k,\, k_1,\, k_1',\, k_2,\, k_2'\right). \quad (12)$$

Substitution of these expressions into Eq.(7) gives $T^A_{\mu\nu}$ for ^{3}He, and by applying the unitarity condition we find corresponding expression for the hadronic tensor:

$$W^{{}^3\mathrm{He}}_{\mu\nu}(P,\, q)=\int\frac{d^4k}{(2\pi)^4}\frac{d^4K}{(2\pi)^4}\frac{d^4k'}{(2\pi)^4}\frac{d^4k_1}{(2\pi)^4}\times$$

$$\times\left[W^D_{\mu\nu}\left(\frac{2P}{3}+K,\, q\right)\bar{\gamma}(P,\, K,\, k)\frac{\Gamma^D\left(\frac{2P}{3}+K,\, k\right)\bar{\Gamma}^D\left(\frac{2P}{3}+K,\, k'\right)}{\left(\left(\frac{2P}{3}+K\right)^2-M_D^2\right)^2}\otimes S\left(\frac{P}{3}-K\right)\gamma(P,\, K,\, k')+\right.$$

$$+\frac{W^N_{\mu\nu}\left(\frac{P}{3}-K\right)}{\left(\frac{P}{3}-K\right)^2-m^2}\bar{\gamma}(P,\, K,\, k)\left[G_4\left(\frac{2P}{3}-K,\, k,\, k_1\right)S\left(\frac{P}{3}+\frac{K}{2}+k_1\right)\otimes S\left(\frac{P}{3}+\frac{K}{2}-k_1\right)\times\right.$$

$$\left.\left.\times G_4\left(\frac{2P}{3}+K,\, k_1,\, k'\right)\right]\otimes S\left(\frac{P}{3}-K\right)\gamma(P,\, K,\, k')\right].$$

Assuming small relative energy of nuclear fragments one can integrate this expression over the zeroth component of different fragments relative momenta and obtain the ^{3}He, ^{3}H, and ^{4}He hadronic tensors respectively in terms of physical amplitudes of the fragments. Using the projection operator $g_{\mu\nu}$ one gets

$$\lim_{Q^2\to\infty} g^{\mu\nu}W^{N(A)}_{\mu\nu}(P,\, q)=-\frac{1}{x}F_2^{N(A)}(x).$$

Introducing now Bjorken variables $x_A=\frac{Q^2}{2P_A\cdot q}$ and $x_N=\frac{Q^2}{2P_N\cdot q}$, we find F_2^A for ^{3}He and ^{3}H in the form:

$$F_2^{{}^3\mathrm{He}}(x_{{}^3\mathrm{He}})=\int\frac{d^3k}{(2\pi)^3}\left[\frac{E_p-k_3}{E_p}F_2^p(x_p)+\frac{E_D-k_3}{E_D}F_2^D(x_D)+\right.$$

$$\left.+\frac{\Delta_p^{{}^3\mathrm{He}}}{E_p}x_p\frac{dF_2^p(x_p)}{dx_p}+\frac{\Delta_p^{{}^3\mathrm{He}}}{E_D}x_D\frac{dF_2^D(x_D)}{dx_D}\right]\Phi^2_{{}^3\mathrm{He}}(\mathbf{k}), \quad (13)$$

$$F_2^{^3\mathrm{H}}(x_{^3\mathrm{H}}) = \int \frac{d^3k}{(2\pi)^3}\left[\frac{E_n - k_3}{E_n} F_2^n(x_n) + \frac{E_D - k_3}{E_D} F_2^D(x_D) + \right.$$

$$\left. + \frac{\Delta_n^{^3\mathrm{H}}}{E_n} x_n \frac{dF_2^n(x_n)}{dx_n} + \frac{\Delta_n^{^3\mathrm{H}}}{E_D} x_D \frac{dF_2^D(x_D)}{dx_D}\right] \Phi^2_{^3\mathrm{H}}(\mathbf{k}), \tag{14}$$

and for ^{4}He in the form:

$$F_2^{^4\mathrm{He}}(x_{^4\mathrm{He}}) = \int \frac{d^3k}{(2\pi)^3}\left[\frac{E_p - k_3}{E_p} F_2^p(x_p) + \frac{E_{^3\mathrm{H}} - k_3}{E_{^3\mathrm{H}}} F_2^{^3\mathrm{H}}(x_{^3\mathrm{H}}) + \right.$$

$$+ \frac{\Delta_p^{^4\mathrm{He}}}{E_p} x_p \frac{dF_2^p(x_p)}{dx_p} + \frac{\Delta_p^{^4\mathrm{He}}}{E_{^3\mathrm{H}}} x_{^3\mathrm{H}} \frac{dF_2^{^3\mathrm{H}}(x_{^3\mathrm{H}})}{dx_{^3\mathrm{H}}} + \frac{E_n - k_3}{E_n} F_2^n(x_n) + \frac{E_{^3\mathrm{He}} - k_3}{E_{^3\mathrm{He}}} F_2^{^3\mathrm{He}}(x_{^3\mathrm{He}}) +$$

$$\left. + \frac{\Delta_n^{^4\mathrm{He}}}{E_n} x_n \frac{dF_2^n(x_n)}{dx_n} + \frac{\Delta_n^{^4\mathrm{He}}}{E_{^3\mathrm{He}}} x_{^3\mathrm{He}} \frac{dF_2^{^3\mathrm{He}}(x_{^3\mathrm{He}})}{dx_{^3\mathrm{He}}}\right] \Phi^2_{^4\mathrm{He}}(\mathbf{k}), \tag{15}$$

where $\Delta_N^A = -M_A + E_N + E_{A-1}$ is the binding energy of the corresponding nuclear fragment. The three-dimensional momentum distributions $\Phi_A^2(\mathbf{k})$ are defined via the «bare» Bethe–Salpeter vertex functions. For example for ^{3}He one has:

$$\Phi^2_{^3\mathrm{He}}(\mathbf{k}) = \frac{mM_D}{4E_pE_DM_{^3\mathrm{He}}(M_D - E_p - E_D)^2}\left\{\int \frac{d^4k_1}{(2\pi)^4}\frac{d^4k_1'}{(2\pi)^4}\bar{\gamma}^{^3\mathrm{He}}(P, k, k_1)\, S_2\left(\frac{2P}{3} + k, k_1\right)\times\right.$$

$$\times \Gamma^D\left(\frac{2P}{3} + k, k_1\right)\bar{\Gamma}^D\left(\frac{2P}{3} + k, k_1'\right)S_2\left(\frac{2P}{3} + k, k_1'\right)\otimes$$

$$\left.\otimes \left(\sum_s u_\alpha^s(\mathbf{k})\, \bar{u}_\delta^s(\mathbf{k})\right) \gamma^{^3\mathrm{He}}(P, k, k_1')\right\}_{k_0 = k_{0_p}}, \tag{16}$$

where $k_{0_p} = \dfrac{M_{^3\mathrm{H}}}{3} - E_p$. Since presently there are no realistic solutions of the Bethe–Salpeter equation for a bound system of three or more nucleons, one has to use phenomenological momentum distributions. It is reasonable to assume that the momentum distributions in Eqs.(13), (14), and (15) can be related with the distributions extracted from experimental data. In the numerical calculations we make use of the distributions available from [8] and [9].

The obtained result reduces to the one obtained within the x-rescaling model [10] and for $A = 3$ becomes:

$$F_2^{^3\mathrm{He}}(x_{^3\mathrm{He}}) =$$

$$= \int dy\, d\varepsilon \left\{ F_2^p\left(\frac{x_{^3\mathrm{He}}}{y - \frac{\varepsilon}{M_{^3\mathrm{He}}}}\right) f^{p/^3\mathrm{He}}(y, \varepsilon) + F_2^D\left(\frac{x_{^3\mathrm{He}}}{y - \frac{\varepsilon}{M_{^3\mathrm{He}}}}\right) f^{D/^3\mathrm{He}}(y, \varepsilon) \right\}, \qquad (17)$$

where $\varepsilon = \Delta_p^{^3\mathrm{He}}$ has the meaning of a nucleon (deuteron) separation energy and $f^{p(D)/^3\mathrm{He}}(y, \varepsilon)$ are the ^{3}He spectral functions for a bound proton (deuteron):

$$f^{p(D)/^3\mathrm{He}}(y, \varepsilon) = \int \frac{d^3k}{(2\pi)^3}\, \Phi^2_{^3\mathrm{He}}(\mathbf{k})\, \frac{m}{E_{p(D)}}\, y\delta\left(y - \frac{E_{p(D)} - k_3}{m} \right) \delta(\varepsilon - (E_p + E_D - M_{^3\mathrm{He}})).$$

Input structure functions $F_2^{p(n)}(x)$ are introduced via parameterizations based on the measurements of the proton and the deuteron structure functions by DIS experiments. We used the most recent parameterization of $F_2^p(x, Q^2)$ found in [11] and fixed the value of Q^2 to 10 GeV^2. The structure function $F_2^n(x)$ is evaluated from $F_2^p(x)$ and from the ratio $F_2^n(x)/F_2^p(x)$ determined in [12]. We have verified that the uncertainties in $F_2^{p(n)}(x)$ are suppressed in the obtained ratio $r^A(x)$ and thus can be neglected in the considered kinematic range.

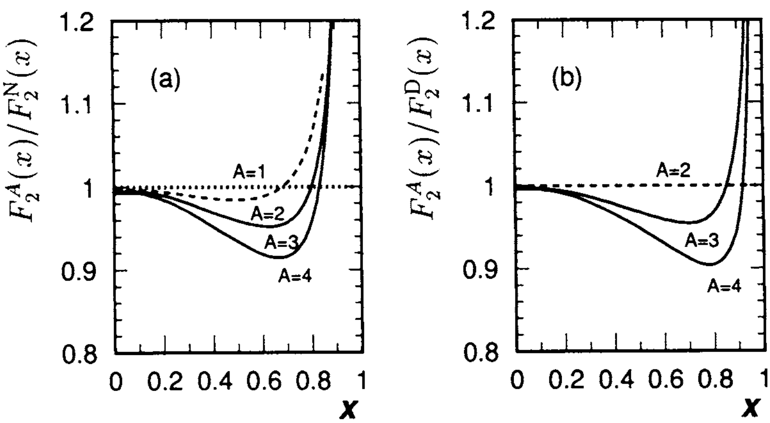

Fig.1. (a) The ratio of the nuclear structure function $F_2^A(x)$ to the isoscalar nucleon one $F_2^N(x)$. (b) The ratio of the nuclear structure function $F_2^A(x)$ to the deuteron structure function $F_2^D(x)$ ($A = 4$) and to the combination of structure functions $(2F_2^D(x) + F_2^p(x))/3$ ($A = 3$). The dashed curve in Fig.(a) shows the result of calculations, described in the text, for $A = 2$. The results for $A = 3$ and $A = 4$ are shown with the solid curves

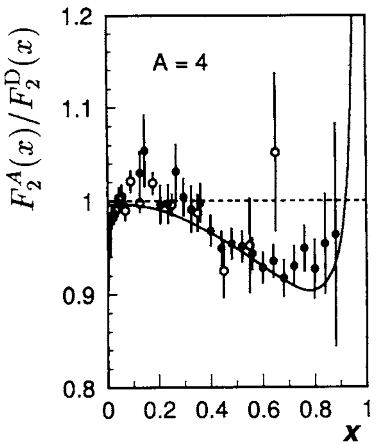

Fig.2. The ratio of the ^{4}He structure function $F_2^{\mathrm{He}}(x)$ to the deuteron structure function $F_2^D(x)$. Results of calculations are shown with the solid curve. The data are from Ref. 4 (filled circles) and Ref. 5 (empty circles)

The results of the numerical calculations, which show how the free nucleon structure function $F_2^N(x)$ $(A = 1)$ evolves to the deuteron $(A = 2)$ and helium ($A = 3$ and 4) structure functions, are presented in Fig.1(a). The evolution, which starts from $F_2^D(x)$, is shown in Fig.1(b). Contrary to what is observed for nuclei with masses $A > 4$, the pattern of the oscillation of $r^A(x)$ changes its shape in the range of $A \leq 4$. The rate at which the changes occur is consistent with the fast buildup of the short range binding forces.

We compare our results for the ratio $F_2^{^4\mathrm{He}}(x)/F_2^D(x)$ with the available data from [4,5] in Fig.2. The position of the cross-over point, obtained from our calculations as $x_3 = 0.919$, is in reasonable agreement with the extrapolated data. On the other hand, the corresponding point for $A = 3$, $x_3 = 0.855$, falls within the interval 0.84—0.86, which is where the ratios $r^A(x)$ $(A > 4)$ cross the line $r^A = 1$ (cf. Ref. 4). This means that the pattern of the EMC effect observed in such dense nuclei as metals is being reached at $A = 3$. The effect of saturation of binding forces in light nuclei is shown illustratively in Fig.3 by comparing our results for ^{3}He with the experimental data obtained on iron. The larger value of x_3 at $A = 4$ must be related with the anomalous binding energy of ^{4}He.

It is important to note the following. The expression (1) shows that the integral

$$I_D = \int_0^1 \frac{dx}{x}\,(2F_2^p(x) - 2F_2^D(x)),$$

which is generally used for the experimental tests of the Gottfried sum rule (cf. Ref. 13), is equal to the Gottfried sum integral to within a correction proportional to $F_2^N(x = 0)$:

$$I_D = \int_0^1 \frac{dx}{x}\,(F_2^p(x) - F_2^n(x)) - 2\frac{\langle M_D - 2F_D\rangle_D}{m}\,F_2^N(x = 0).$$

Apparently, such tests cannot be performed rigorously because the value of $F_2^N(x)$ at $x = 0$ is unknown. On the other hand, if one uses Eqs.(13) and (14) to calculate the integral

$$I = \int_0^1 \frac{dx}{x}\,(F_2^{^3\mathrm{He}}(x) - F_2^{^3\mathrm{H}}(x)),$$

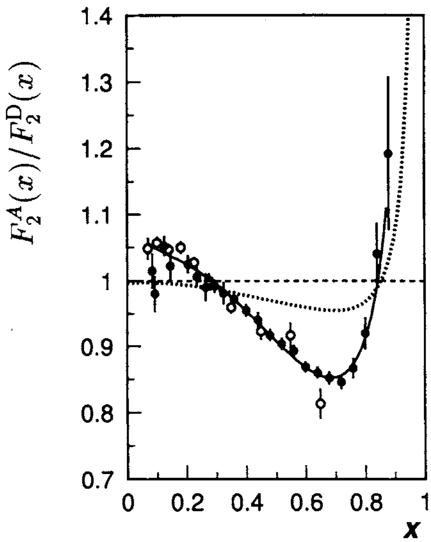

Fig.3. The ratio of the nuclear structure function $F_2^A(x)$ to the deuteron structure function $F_2^D(x)$. Results of calculations for ^{3}He, described in the text, are shown with the dotted curve. The data obtained on iron are from Ref. 4 (filled circles) and Ref. 14 (empty circles). The data points are approximated with the solid curve

one sees that the binding corrections cancel and the integral I is equal to the Gottfried sum. An experiment, which used ^{3}He and ^{3}H targets would therefore be able to verify the Gottfried sum rule independently of the model uncertainties in the binding corrections.

In conclusion, the method for the model-free calculations of the evolution of the structure functions in the lightest nuclei has been developed as the extension of an approach based on the Bethe–Salpeter formalism. The method allows one to express $F_2^A(x)$ in terms of structure functions of nuclear fragments and three-dimensional momentum distributions. As a result, $F_2^A(x)$ have been evaluated numerically without finding solutions of Eqs. (6) and (10).

The obtained pattern of distortions of the nucleon structure function proves that the EMC effect in the lightest nuclei, D, ^{3}He, and ^{4}He, is basically the manifestation of the short range binding forces in the nucleon parton distributions. The quantitative predictions for ^{3}He and ^{4}He nuclei, which have to be verified in future experiments at HERA or CEBAF, indicate that the EMC effect in heavy nuclei can be naturally understood as distortions of the nucleon parton distributions in ^{3}He, which are modified by the nuclear density effects.

Finally, the obtained results prove that in the EMC effect range ($0.3 < x < 0.9$) the two-nucleon interactions can be considered as the dominant mechanism for the description of the nuclear binding forces.

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