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A.I.Titov

QUARK DISTRIBUTIONS OF HEAVY ATOMIC NUCLEI

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1. INTRODUCTION

Recently, experimental data of European Muon Collaboration (EMC)[1] completed by the data of SLAC[2] have become available on the ratio R of structure functions of deep inelastic scattering of muons (electrons) on the iron and deuterium:

$$R(x) \equiv \frac{2\sigma^A}{A\sigma^D} = \frac{F_2^A(x)}{F_2^D(x)}; \qquad A = {}^{56}Fe, \tag{1}$$

where x is the Bjorken scale variable ($x = Q^2/2M\nu$; $t = -Q^2$ is the four-momentum transfer, $\nu = E_\mu - E_{\mu'}$ is the energy transfer and M is the nucleon mass), $F_2^{A(D)}(x)$ is the structure function of the nucleus (deuterium). If the nucleus is assumed to consist of A nucleons only (where A is the atomic weight), the R should be equal to unity. But the experiment shows that the R has a bump at small values of x ($x \simeq 0.05$, $R-1 \simeq 0.17$); a minimum at $x \simeq 0.65$ ($R-1 \simeq -0.15$) and a sharp increase at $x \sim 1$ ($x \simeq 0.9$, $R-1 \simeq 0.2$). In other words, the experiment shows that the idea of the atomic nucleus being a system of A -nucleons is not complete, and it is necessary to take into account the other nonnucleon components which we call conditionally "exotics".

In principle, the nuclear exotics is taken into account phenomenologically by including it into the structure function of a neutron, which is determined as the difference between the nuclear and hydrogen data. So, the result of the mentioned experiments can be reformulated as follows: the exotics of the heavy nucleus differs from the exotics of the deuteron. Really, the difference between the structure functions of the proton $F_2^p(x)$ and neutron $F_2^n(x)$ extracted from the deuteron data[3] gives the factor: $R-1 = -\epsilon$, where $\epsilon = \frac{A-2Z}{2A}\,\frac{F_2^p(x) - F_2^n(x)}{0.5(F_2^p(x) + F_2^n(x))} \approx 1.3\cdot 10^{-3}$, which is by an order of magnitude less than the experimental one and can be omitted.

In our paper we present theoretical investigations of the ratio R. We shall show that the observed increase of R at small and at large x comes from different reasons, namely: the increase at small x is caused by the long-range meson fields the contribution of which into the heavy nuclei is somewhat larger compared to the deuteron[4]. The increase at large x is connected with the short-range nucleon-nucleon dynamics and after all with the admixture of the multiquark configurations in the nuclei. Recall that at present there are available some experimental manifestations of the multiquark states in nuclei and nuclear processes. First, these include the reactions of the "cumulative" particle production in the hadron-nucleus[5,6] and lepton-nucleus collisions at large momentum transfer. These reactions show that an incident particle interacts with a system the mass of which is several times (up to 4) as large as the nucleon mass. The radius of the system is about $r \sim 1$ fm, so the nucleons of the system lose their individuality, and it is natural to consider that this group of nucleons is a multiquark system. The second example is the data on the deuteron charge formfactor at large Q^2 ($Q^2 \sim 10 (\mathrm{GeV}/c)^2$) which cannot be explained without the six-quark states admixture in the deuteron[12]. And, finally, there are many results on the nucleon-nucleon interactions at intermediate energies where some peculiarities of the partial waves have been found[13]. These peculiarities are connected with the existence of the dibaryons, direct candidates for six quark systems[14,15].

We begin in §2 by analyzing the influence of the multiquark configurations in nuclei on the behaviour of the ratio R, and in §3 we investigate the contribution of the pion degrees of freedom.

2. MULTIQUARK STATES IN NUCLEI

Let us suppose here that the nucleons and 6q-states which are the simplest multiquark states give the main contribution to the nuclear structure function. The mass of the 6q-system is $M_f \simeq 2M$, so the structure function of the 6q-system depends on the scale variable $x^f = Q^2/2M\nu \simeq x/2$. Assuming the quark distributions for u- and d-quarks in the 6-quark system to be the same we can write the expression for the structure function of a nucleus:

$$F_2^A(x) = (1 - P^A)\, F_2^N(x) + P^A F_2^{q^6}(x), \tag{2}$$

where $P^{A(D)}$ is the probability of the 6-admixture in the nucleus (deuteron), $F_2^N(x)$ is the structure function of the nucleon:

$$F_2^N(x) = \frac{1}{2}(F_2^p(x) + F_2^n(x)) = \frac{5}{18}(x u_v(x) + x d_v(x)) + \frac{12}{9} x u_s(x), \tag{3}$$

$F_2^{q^6}(x)$ is the contribution of the 6q-structure function:

$$F_2^{q^6}(x) = \frac{5}{18}\left[\frac{x}{2} u_v^f\left(\frac{x}{2}\right)\right] + \frac{6}{9}\left[\frac{x}{2} u_s^f\left(\frac{x}{2}\right)\right], \tag{4}$$

where

$$u_s = \bar{u}_s = d_s = \bar{d}_s = s_s = \bar{s}_s; \quad u_s^f = \bar{u}_s^f = d_s^f = \bar{d}_s^f = s_s^f = \bar{s}_s^f.$$

A theoretical estimation of the probability of the 6q-admixture predicts the P^D value in an interval of $(2\text{-}7)\cdot 10^{-2}$ [10,15-17]. The same values come from the data of nuclear reactions at large momentum transfers[10,18]. The analysis of the cumulative particle production reactions[10,8,9] and theoretical calculations[19] show that the probability of the 6q-systems in a heavy nucleus is larger than the probability in a deuteron by a factor of two, i.e., $P^A \simeq 2P^D$. This result is clear because the ratio P^A/P^D is approximately equal to $[\langle r_{NN}\rangle_D / \langle r_{NN}\rangle_A]^3$, where $4/3\pi \langle r_{NN}\rangle^3$ is the volume of a nucleon in the nucleus (in the deuteron). As far as $\langle r_{NN}\rangle_D > \langle r_{NN}\rangle_A$, $P^A > P^D$. Below we take $P^D = 7.10^{-2}$. The parameters of $F_2^N(x)$ are chosen to describe the deep inelastic μD-scattering and are tabulated in the table. The parameters of $F_2^{q^6}(x)$ were fixed by the conditions: $\int u_v^f(x)\,dx = 3$, $\int x u_s^f\,dx = \int x u_s\,dx$. The quark distribution u_v^f decreases as $(1-x)^7$ when $x \to 1$, that means an incomplete "defreezing" of all colour degrees of freedom of the 6q-system[8-10]. So, if $P^A \neq P^D$, then $R = F^A/F^D \neq 1$, i.e.,

$$R(x) = 1 + (P^A - P^D)(F_2^{q^6}(x) - F_2^N(x))/F_2^D(x). \tag{5}$$

As $x \to 1$, $F^N \sim 0$, $F^D \sim P^D \cdot F^{q^6}$ and

$$R(x \sim 1) - 1 \simeq P^A/P^D - 1 \simeq 1. \tag{6}$$

That agrees with the experimental data qualitatively at $x \sim 1$. However, eq. (6) cannot explain the experimental data in the whole range of x. The corresponding calculations are shown in fig. 1 by the curve "0". For example, at $x = 0$ $R_{theor} - 1 = -7\cdot 10^{-2}$ $x0.27 = -1.9\cdot 10^{-2}$, as far as $R_{exp} - 1 = 1.8\cdot 10^{-1}$. So, we can conclude that the increase of R at small x is caused by other reasons. Note that the strong increase of R as $x \to 1$ was predicted also in other models which took into account the short-

Table

Parameters of the quark distribution in the deuteron

$x q_v(x) = c_v \sqrt{x}(1 - x)^{\gamma_v}$; $x q_s(x) = c_s (1 - x)^{\gamma_s}$

	u_v	d_v	u_s	u_v^f	u_s^f
γ	3	4	7	7	11
c	2.25	1.55	0.15	4.6	0.23

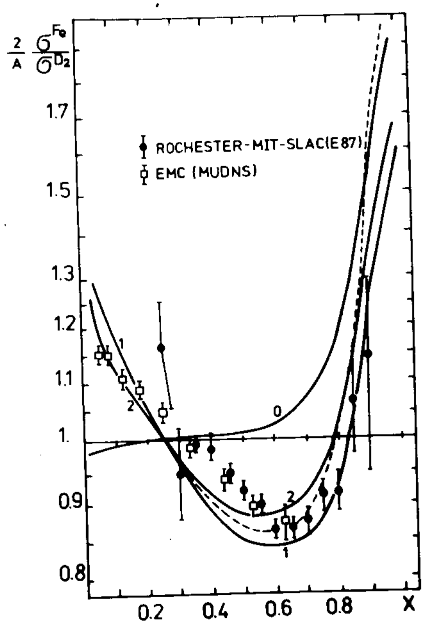


Fig.1. The calculation and comparison with experiment of the ratio R. Curves: "0" - only 6q - states; 1,2 - 6q -states and the meson fields. 1 - $\rho_c/\rho_0=3.9$, 2 - $\rho_c/\rho_0=4.2$, $P^A/P^D=2$. Dashed line - $P^A/P^D=2.3$ $\rho_c/\rho_0=3.9$. Experiment: ☐ - from /1/, ● - from /2/.

range nucleon-nucleon interactions, for example, "a few-nucleon correlation" model /20/. All of them cannot describe the increase of R at small x, moreover, they give a rather strong increase of R(x) at $x \sim 1$. The theoretical calculation /20/ predicts $R_{theor}(0.8) - 1 \simeq 1$, as far as $R_{exp}(0.8) - 1 \simeq 3.10^{-2}$. One can decrease the ratio R at $x \simeq 1$ by decreasing the contribution of the nucleon high-momentum components in a heavy nucleus. However, the decrease will give disagreement with experimental data on $F_2^A(x)$ at $x \geq 1$ /7/.

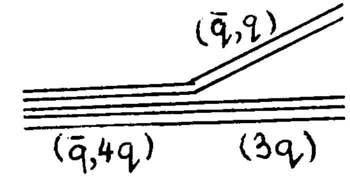
3. MESONIC DEGREES OF FREEDOM IN NUCLEI

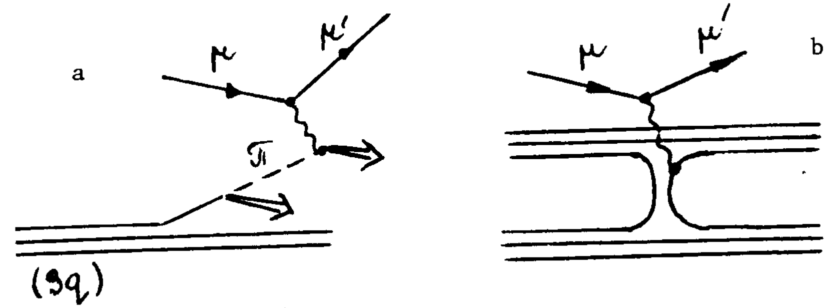
The calculation of the π-meson degrees of freedom leads to the increase of ratio R at small x. Qualitatively, this may be understood as follows. Suppose that there are n π-mesons in the nucleus. Then neglecting its internal motion in the nucleus one could expect that the respective contribution into the structure function will be proportional to $n \cdot F^\pi(x_\pi)$, where $F^\pi(x_\pi)$ is the pion structure function, x_π is the pion scaling variable: $x_\pi = Q^2/2m_\pi\nu = M/m_\pi x$. From the condition $0 \leq x_\pi \leq 1$ it follows that the range of the pion contribution into the structure function is restricted by small x: $0 \leq x \leq m_\pi/M \simeq 0.15$. The calculation of the pion internal motion will lead to a little extension of this region.

It is necessary to stress that here we condider soft pions. The soft pions are similar to fragmentational sea pions described by the diagram in fig.2. The quark counting rules predict the threshold behaviour of the pion distribution:

$$n(y) \sim (1-y)^5 \quad \text{at} \quad y \to 1. \tag{7}$$

Fig.2. Fragmentation of the soft pions.





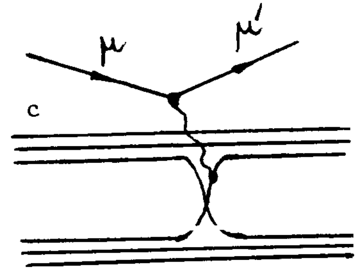


Fig.3. Contribution of the hard pions.

The soft pions must be different both from hard pions appearing in the fragmentation of fast quarks (Fig.3a) and from "exchange hard pions" described by exchange diagrams shown in Fig.3b,c. The first of them contribute into the structure function, this contribution decreases when $x \to 1$ as $(1-x)^{3\div 4}$, and is taken into account in the nucleon structure function. The second are taken into account in $F^{q^6}(x)$ because such diagrams contribute only at small nucleon-nucleon distances and imitate the short-range part of N-N forces /15,16/. The soft pions are responsible for the long-range part of N-N forces and they can not be described by exchange quark diagrams as in Fig.3b,c.

To estimate the effective number of pions in the nuclear matter, which (for simplicity) is supposed to be isotopically symmetric, we consider the amplitude of the process the diagram of which is shown in Fig.4:

$$S_{if} = \sqrt{\frac{M\mu}{E_N E_\mu E_{\mu'}}}\, \vec{\mathfrak{M}}^{\mu\pi}_{if} \sqrt{\frac{2m_\pi M}{E_{N'}}}\, D^\pi_A(k)\; g_0(k)\; \bar u(N')\,\gamma_5\,\vec\tau\, u(N), \tag{8}$$

where E_μ, $E_{\mu'}$ are the initial and final energies of the muon, μ is its mass, $\mathfrak{M}^{\mu\pi}_{if}$ is the amplitude of the deep-inelastic $\mu\pi$-scattering, $g_0(k)$ is the vertex function, $D^\pi_A(k)^{-1}$ is the pion propagator in the matter

$$D^\pi_A(k)^{-1} = D^\pi_0(k)^{-1} - \Pi^\pi_A(k), \tag{9}$$

$D^\pi_0(k)$ is the pion propagator in the vacuum $D^\pi_0(k)^{-1} = k^2 - m_\pi^2 + i\epsilon$, $\Pi^\pi_A(k)$ is the polarization operator. Averaging over initial and summing over final spin-isospin states we get the cross-section of this process

$$d\sigma^{\mu A}_\pi(x) = \int d\sigma^{\mu\pi}(x(k))\; n^\pi_A(k)\, d\vec{k}, \tag{10}$$

where $d\sigma^{\mu\pi}$ is the cross-section of the deep-inelastic $\mu\pi$-scattering, $n^\pi_A(k)$ is the effective number of pions in the nucleus

$$n^\pi_A(k) = \frac{3}{\pi^2 m_\pi} f^2_{NN\pi}(k)\; a^2(k) |D^\pi_A(k)|^2, \tag{11}$$

where

$$a^2(k) = (\vec{p}'/(E'_N + M) - \vec{p}/(E_N + M))^2 M(E_N + M)(E'_N + M)/E',$$

p, p' are momenta of nucleons N and N'; $E_N = \sqrt{p^2 + M^2}$ $E'_N = \sqrt{p'^2 + M^2}$ in the system of coordinates, where $\vec{p} = 0$ $a^2(k) = \vec{k}^2(2M^2/(E'(E'+M)))$. The vertex function $f_{NN\pi}$ is connected with g_0 by the relation

$$f_{NN\pi}(k) = f_{NN\pi}(0)\cdot v(k); \quad f^2_{NN\pi}(0) = \frac{g_0(0)^2}{4\pi}\left(\frac{m_\pi}{2M}\right)^2 \simeq 0.08, \tag{12}$$

where $v(k)$ is the form factor decreasing with growing k^2 which is introduced phenomenologically [4,22]

$$v(k) = \left(1 - \frac{k^2}{\Lambda^2}\right); \qquad \Lambda = 0.52\ \text{GeV/c}. \tag{13}$$

To calculate D^π_A, it is necessary to solve a coupled-equation system for vertex functions of the interaction of pions, nucleons, nucleon resonances, and Δ-isobars in nuclear matter. It is an independent problem and some special papers [4,22,23] are devoted to its solution. For our purpose it is convenient to use the analytic approximation $D^\pi_A(k)$ proposed in the paper [23]

$$D^\pi_A(k)^{-1} = m^2_\pi(1.6\bar{\omega}^2 - \omega_0^2 - 0.8(|\bar{k}| - k_0) + i\,0.5\bar{\omega}) \tag{14}$$

$$\bar{\omega} = \omega/m_\pi; \quad \bar{k} = \vec{k}/m_\pi; \quad \omega_0^2 = 0.5(\rho_c/\rho_0 - 1); \quad k_0^2 = 3 - 1.08\omega_0^2, \tag{15}$$

where $\rho = 0.18\ \text{Fm}^{-3}$ is the nornal nuclear debsity. The parameter ρ_c is the critical nuclear density, includes the whole information on nuclear matter and fundamental interactions of its constituents (pions, nucleons, Δ-isobars, etc.), and it is the only parameter in the equation (15).

The same parameter determines the effective number of pions in matter n^A_π:

$$n^A_\pi = \int n^A_\pi(\vec{k})\, d\vec{k}. \tag{16}$$

Note that an additional contribution into the structure function $F^A_2(x)$ comes from the difference $\delta n^A_\pi = n^\pi_A - n^\pi_0$, where n^π_0 corresponds to the contribution from the diagram of Fig.4 to the "vacuum", the "deuteron", which is taken into account in the nucleon structure function. As an illustration, the calculation of δn^A_π as a function of ρ_c/ρ_0 is shown in Fig.5. Note that theoretical estimates predict $\rho_c/\rho_0 \sim 3.2$ [22] that corresponds approximately to one pion/nucleon in the matter. Once the distribution $n(k)$ is known, one could determine the pion contribution into the nuclear structure function $\delta F^{A(\pi)}_2$ with the help of the "light cone" formalism [24] and the substitution:

$$\vec{k} \to \vec{k}(y); \quad dk_z = |dk_z/dy|\, dy; \quad d\vec{k}_\perp(y) = d\vec{k}_\perp,$$

where $k_z(y)$ is determined from the equation:

$$(\sqrt{\vec{k}^2 + M^2} + \sqrt{\vec{k}^2 + M^2})^2 = (m^2_\pi + k^2_\perp)/y + (M^2 + k^2_\perp)/(1-y)$$

$$k^2_z = \vec{k}^2 - k^2_\perp.$$

Then

$$\delta F^{A(\pi)}_2(x) = \int dy\; \delta n^\pi_A(y)\; F^\pi\!\left(\frac{x}{y}\right), \tag{17}$$

where

$$\delta n^\pi_A(y) = \int d\vec{k}_\perp\, |dk_z/dy|\; \delta n^\pi_A(\vec{k}), \tag{18}$$

$$\int dy\; \delta n^\pi_A(y) = \int d\vec{k}\; \delta n^\pi_A(\vec{k}) = \delta n^A_\pi. \tag{19}$$

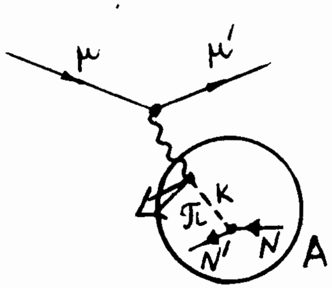


Fig.4. Deep-inelastic $\mu\pi$-scattering in the nucleus.

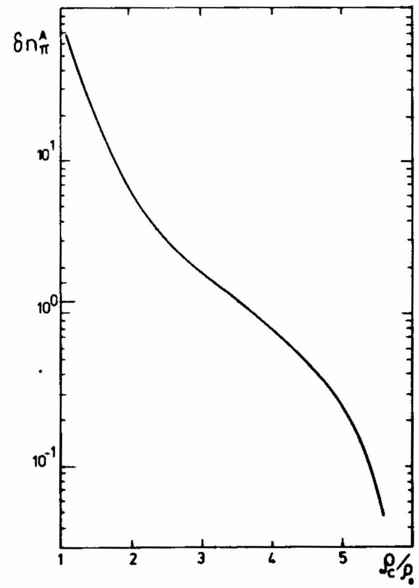


Fig.5. The effective number of pions in the nucleus.

Unfortunately, we cannot accomplish this scheme without introducing new parameters into the form factor $v(k^2)$ which would provide a rapid fall-off of $v(k^2)$ at large $|k^2|$ $(-k^2 > 1\,(\mathrm{GeV}/c)^2$. Note that the form factors we know from nuclear physics are not suitable for our purpose. They are determined at relatively small momenta $|k^2| \lesssim 0.5$ $(\mathrm{GeV}/c)^2$. And though this range of k^2 gives a main contribution to the total value of the pions δn_π^A the choice of $v(k^2)$ in (13) does not provide the "soft pion" condition (7). So, instead of introducing additional parameters in $v(k^2)$ we use considerations given at the beginning of this section and following (7) and (19) we take

$$\delta n_A^\pi(y) = 6\,\delta n_A^\pi (1-y)^5 . \tag{20}$$

Now before to calculate the nucleus structure function, we note one more fact: the pion contribution increases the momentum transferred by quarks and decreases the momentum transferred by gluons. However, we have no serious reasons to assume that such a redistribution really takes place, most probably it is natural to assume that the momentum transferred by gluons weakly depends on the hadron sort and practically does not depend on the atomic weight of the nucleus. Thus,

$$\langle xq^D(x)\rangle = \langle xq^A(x)\rangle = \langle x\tilde{q}^A(x)\rangle + \langle xq_A^\pi(x)\rangle , \tag{21}$$

where $\tilde{q}^A$ is the distribution of the nucleon and $6q$-quarks,

q_A^π is the distribution of pion quarks

$$q_A^\pi(x) = \int q^\pi\left(\frac{x}{y}\right) \delta n_A^\pi(y)\, dy , \tag{22}$$

q^π is the distribution of quarks in the pion. Besides, it is necessary to take into account the conservation of the baryon charge

$$\int \tilde{q}_v^D(x)\, dx = \int \tilde{q}_v^A(x)\, dx , \tag{23}$$

where q_v is the distribution of valence quarks. Conditions (21) and (23) lead to a small change of the quark distribution in a heavy nucleus, namely, valence quark distributions increase at small x ($x \leq 0.5$) and decrease at $x > 0.05$, sea-quark distributions decrease in the whole range of x:

$$q_{v_i}^A = c_{v_i}^A\, x^{\alpha_i} (1-x)^{\gamma_{v_i}} ; \quad xq_{s_i}^A = c_{s_i}^A (1-x)^{\gamma_{s_i}} , \tag{24}$$

where:

$$c_{v_i}^A = c_{v_i}^D \frac{\Gamma(0.5)\,\Gamma(\gamma_{v_i}+2-\alpha_i)}{\Gamma(1-\alpha_i)\,\Gamma(\gamma_{v_i}+1.5)} ; \quad \alpha_i = \frac{\gamma_{v_i}(2-\delta)-2\delta+3}{2\gamma_{v_i}-\delta+3} ; \quad c_{s_i}^A = \delta c_{s_i}^D \tag{25}$$

$$\delta = \frac{1-(1-a)P^D - \eta\,\delta n_A^\pi}{1-(1-a)P^A} ; \tag{26}$$

$$\eta = \frac{\int xq_A^\pi(x)\, dx}{\int xq^D(x)\, dx} \simeq 0.12 , \qquad a = \frac{\int xu^f(x)\, dx}{\int xq^D(x)\, dx} \sim 0.5 . \tag{27}$$

So the final expression for $F_2^A(x)$ is

$$F_2^A(x) = (1-P^A)\left[\frac{5}{18}(xu_v^A(x) + xd_v^A(x)) + \frac{12}{9} xu_s^A(x)\right] + \\ + P^A\left[\frac{5}{18}\left(\frac{x}{2} u_v^{A,f}\left(\frac{x}{2}\right) + \frac{6}{9}\left(\frac{x}{2} u_s^{A,f}\left(\frac{x}{2}\right)\right)\right] + \delta F_2^{A(\pi)}(x), \tag{28}$$

$$\delta F_2^{A(\pi)}(x) = \int dy\, F^\pi\left(\frac{x}{y}\right) \delta n_A^\pi(y) . \tag{29}$$

In concrete calculations for $F_2^\pi(x)$ we have used the parametrization from the paper [25]. The calculations show that $\delta F_2^{A(\pi)}(x)$ has the behaviour $(1-x)^{6\div 7}$ and imitates an increasing concen-

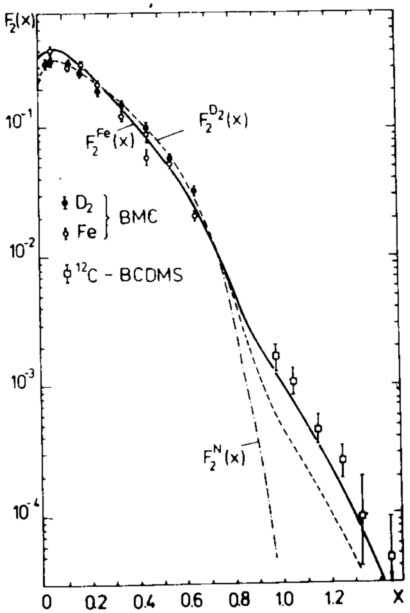


Fig.6. Calculation and comparison with experiment of the structure function of the heavy nucleus (solid line) and the deuteron (dashed line). Experiment: ● - from/3/, □ - from/7/, o - from/26/.

tration of sea-quarks in the nucleus. The contribution $\delta F_2^{A(\pi)}$ is proportional to δn_π^A and depends on the parameter ρ_c. The calculation of R is shown in Fig.1 for two values of ρ_c. Curve 1 is calculated for $\rho_c = 3.9\rho_0$ ($\delta n_A^\pi \simeq 0.83$); curve 2, for $\rho_c \simeq 4.2\rho_0$ ($\delta n_A^\pi \simeq 0.63$). The renormalization of quark distributions (25) leads to small decrease of R in the cumulative region. To keep this ratio equal to 2 at $x \sim 1$ /7,8/, it is necessary to somewhat increase P^A. The respective calculation is shown in Fig.1 (dashed line): $P^A = 2.3P^D$, $\rho_c = 3.9\rho_0$. It is seen that theoretical curves comply with experiment.

The comparison of structure functions $F_2^A(x)$ and $F_2^D(x)$ with experiment at $0 \lesssim x \lesssim 1.4$ is shown in Fig.6. It is seen that the calculated curves correctly describe the experiment. The calculation of F_2^D is predictable at $x \gtrsim 1$.

Note that some break of theoretical curves at $x \sim 1$ is due to the absolute ignorance of the Fermi motion of nucleons in the nucleus (deuteron), which contribute at $0.9 \lesssim x \lesssim 1.1$. Its calculation leads to smoothing both the curves in this region without changing all the results.

CONCLUSION

So, we have investigated possible reasons of the deviation observed from the unity of the ratio of structure functions of a heavy nucleus and the deuteron. The increase of R in the cumulative region is due to the dynamics of NN-interaction at small distances, namely, due to the admixture of the multiquark states in the nucleus. The increase of R in the range $0 \le x \le 0.2$ is due to admixture of the long-range pion fields which are enhanced in heavy nuclei. Thus, this increase is a pure nuclear effect*. The meson field rearrangement takes place in sufficiently heavy nuclei, and nuclei with $A \simeq 9$-16 are transitional in this sense.

Meson fields in the nucleus lead to the renormalization of quark distributions. The renormalization effect may be observed in the neutrino experiment in the measurement of the structure function $F_3(x)$. In particular, the approach developed here predicts the ratio of average values $\langle F_3 \rangle$ for a heavy nucleus and the deuteron at $0 \le x \le 1$

$$R_3 = \langle F_3^A \rangle / \langle F_3^D \rangle \simeq 0.90 \pm 0.05.$$

The pion quark distribution in the nucleus resembles the sea-quark distribution in the nucleon and six-quark systems. Therefore, this effect may be taken into account phenomenologically by a significant increase of the contribution of respective sea-quarks in the nuclear hadron. In this sense our approach substantiates this phenomenology and fills it by the physical content. In principle, the small increase of "the sea" in more complex six-quark systems may really occur. However, from our point of view this increase will be in nature of a correction to the basic effect.

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*As we have known the similar idea has been developed by M.Ericson and A.W.Thomas /27/.

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Титов А.И. E2-83-460

Кварковые распределения в тяжелых атомных ядрах

Анализируются недавние экспериментальные данные по отношению структурных функций глубоконеупругого рассеяния мюонов на железе и дейтерии. Показано, что наблюдаемое усиление структурных функций тяжелого ядра при значениях бьеркеновской масштабной переменной $0 \le x \le 0.25$ вызвано вкладом мезонных полей. Поведение структурных функций ядер в "кумулятивной" области $x \ge 1$ определяется примесью многокварковых состояний.

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Quark Distributions of Heavy Atomic Nuclei

Recent experimental data on the ratio of the structure functions of iron and deuterium are analysed. It is shown that an observed increase of the structure function of a heavy nucleus for the Bjorken-scaling variable x in the range $0 \le x \le 0.25$ comes from the contribution of the mesonic fields. The behaviour of the structure functions at large value of $x \ge 1$ is determined by the admixture of the multi-quark states in the nuclei.

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

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