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THE NUCLEON MASS AND THE EOS OF NUCLEAR MATTER*

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We show in the simple model independent of photon momentum transfer the density evolution for the parton distribution in nuclear matter. The correction to the Fermi energy from term proportional to the pressure are very important.

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The nuclear EMC effect, quite strong as witnessed by Fig. 1 (where it is shown for mass number $A = 56$), is reflection of the influence of nuclear field on the partonic structure of nucleons. We investigated this effect using extended Relativistic Mean Field model (RMF) [1] together with Monte Carlo method and calculated parton distributions in nuclei [2] for Björken variable $0 < x < A$. In this region the change of the nuclear virtual pion cloud (connected with the existence of exchanged mesons originated from the nuclear forces) turns out to be crucial. In order to reproduce observed behavior of experimental data in that region we adjusted in our model [3] the value of the parameter which determines the relative number of intermediate pions (mediating the nucleon–nucleon interaction [4]) in the function of Björken x. We argue therefore that results on deep inelastic e–A scatterings show partial deconfinement [4] of nucleons inside the nuclear matter (NM) enhancing thus the role played by the partonic degrees of freedom. The magnitude of the nuclear Fermi motion is sensitive to the residual interactions between partons, influencing both the nucleon Structure Function (SF) and the value of the nucleon mass M in NM [5]. The sea parton distributions are described by allowing for additional virtual pions in hadron in the quantity which reproduces both the nuclear lepton pair production data and saturates the energy-momentum sum rule. The extension of this model [6] to the dense medium with the nonzero pressure will be discussed.

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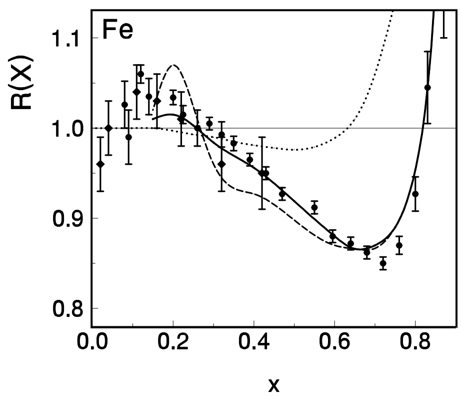

Fig. 1. Ratio of the nuclear (iron) SF $F_2(x)$ to deuterium in the deep inelastic electron–nucleus scattering.

1. The nuclear deeply inelastic limit

The nuclear quark SF F_2^A in a nucleus with the mass number A is constructed in the convolution model (CM) from the free nucleon SF $F_2^N(x)$ in the nucleon and the nucleon distribution function $\rho^A(y_A)$ in the nucleus. The Björken variable x in this model is equal to the fraction of longitudinal momentum carried by quark in the nucleon and is equal to $x = (k_0 + k_3)/M$ in the nucleon rest frame. It was shown that in fact the nuclear SF in CM depends on the contributions from both scalar $U_{\rm S}$ and vector $U_{\rm V}$ potentials. The final result in the relativistic Fermi gas model [7] is:

$$\rho^A(y_A) = \frac{4}{\rho} \int \frac{d^4p}{(2\pi)^4} S_N(p_0, \boldsymbol{p})(1 + p_3/E^*(p))\delta(y - (p_0 + p_3)/\mu)\,, \qquad (1)$$

where $\overline{p} = (p_0, \boldsymbol{p})$ is the nucleon four momentum, $E^* = \sqrt{\boldsymbol{p}^2 + (m - U_{\rm S})^2}$ and the factor $(1 + p_3/E^*)$ corrects [8] the nonrelativistic expression. The nucleon energy is equal to the chemical potential $p_0 = \mu$ at the Fermi surface. Here the nucleon spectral function was taken in the RMF approach: $S_N = n(p)\delta(p_0 - (E^*))$.

In the nuclear medium, characterized by ϵ and Fermi energy $E_{\rm F}$, the rest energy of the nucleon $M_{\rm B} = \sum_i k_{Ni}^+$, in the Björken on shell limit, takes effectively different value [6] than its free nucleon mass M. It can be thought as the sum of the corresponding partonic energies k_{Ni}^0 expressed in the rest frame of nucleon (notice that they differ from k_{Ai}^0). Such $M_{\rm B}$, accounts therefore effectively for the Fermi motion of nucleons inside the nucleus. This is, in addition to the standard Fermi smearing on a nuclear level, the influence of the Fermi motion emerging from a nucleonic (x) level.

2. RMF models, EOS and parton SF

The Equation of State (EOS) for nuclear matter has to match the saturation point for $T = 0$ but then the behavior for higher densities is different for different RMF models — see Fig. 2 taken from [9]. We have here stiff

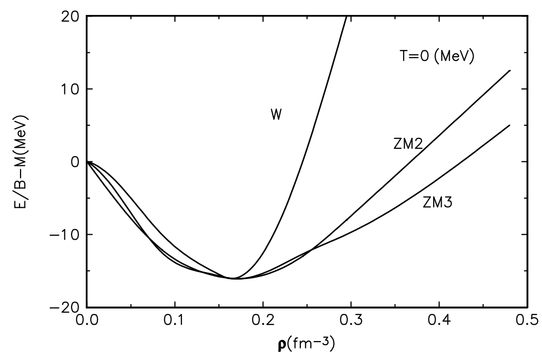

Fig. 2. The nucleon energy in function of nuclear matter density in the RMF approach for Walecka and ZM models [9].

model of Walecka [10] and two Zimanyi–Moszkowski (ZM) [11] models ZM2, ZM3. There are however non-equilibrium correction to nuclear SF in the RMF. For non zero pressure in nuclear matter the Fermi energy is not equal to the average binding energy but there are the well known correction (for example) [12]:

$$E_{\mathrm{F}} \;=\; \frac{d}{d/\varrho}\left(\frac{E}{\Omega}\right), \tag{2}$$

$$E_{\mathrm{F}} \;=\; \left(\frac{E}{B}\right) + \varrho\frac{d}{d/\varrho}\left(\frac{E}{B}\right), \tag{3}$$

where $B/\varrho = \Omega$ give volume. Equivalently we can use the pressure $p = -\left(\frac{\partial E}{\partial \Omega}\right)_B = \varrho^2 \frac{d}{d/\varrho}\left(\frac{E}{B}\right)$ to obtain:

$$E_{\mathrm{F}} = E/B + p/\varrho\,. \tag{4}$$

In formula (1) the nucleon distribution is simplified in the RMF to the form:

$$\rho^A(y_A) = \frac{3}{4v_A^3}(v_A^2 - (y_A - 1)^2)\,, \tag{5}$$

where $v_A = (p_{\mathrm{F}}/\mu)$, and y takes the values given by inequality $0 < (1 - p_{\mathrm{F}}/\mu) < y < (1 + p_{\mathrm{F}}/\mu)$. Thus all the nuclear dependence is hidden in nucleon chemical potential $\mu = E_{\mathrm{F}}$ (Hugenholtz–van Hove theorem) when the pressure is absent in the saturation point. Taking only first term in (3) we present comparison between nuclear SF calculated for density $3\rho_0$ with Walecka model [10] and softer (ZM2) model [11] with compressibility slightly below 200 MeV in Fig. 3. In Walecka model, which was chosen rather for

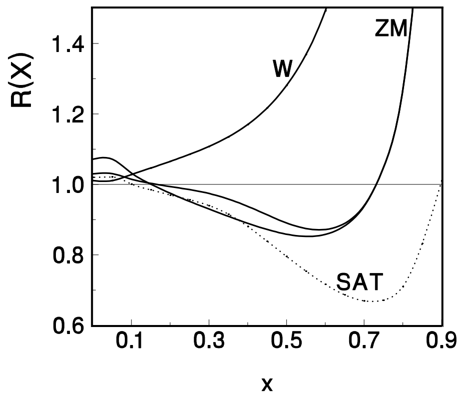

Fig. 3. Results for the ratio $R(x) = F_2^{\rm NM}(x)/F_2^N(x)$ which shows the evolution of the nucleon SF for Walecka (W) and ZM models of RMF for density $\rho = 0.51$ fm^{-3}. Results for equilibrium density (SAT) $\rho_0 = 0.17$ fm^{-3} the same for both models are shown for reference.

reference, we have also consistency between non zero value of quark condensates and unaffected nucleon SF for high density. The main reason is the dominance of the strong repulsive vector field in this region. In contrary the nonlinear coupling of the meson field in the RMF model looks much more realistic with respect to possible chiral phase transition for (3,4) ρ_0 connected with the change of the nucleon SF inside dense medium. The final calculation have to include the changes in energy (presented already in Fig. 3) and finite pressure. The influence of second term in Eq. (3) (or Eq. (4)) to the Fermi energy is indeed very strong, it increases $E_{\rm F}$ by 8% for density $3\rho_0$. The nucleon structure function is changed and in our model this corresponds to the smaller (by 80 MeV) nucleon mass in the deep inelastic limit. In this case we have to shift Björken x in order to satisfy the nuclear momentum sum rule. Fig. 4 shows the ratio of the SF calculated for ZM2 model for different densities where the pressure p is negative for $\rho < \rho_0$ or positive otherwise. Generally, the positive pressure increases Fermi energy, therefore we have to balance it by the nucleon momentum scaling the Björken x. This corresponds to decrease of nucleon mass in the deep inelastic process. In that way we obtain the deep-inelastic prescription how to decrease the nucleon mass along with the increase of pressure in the NM. For negative pressure these corrections are small and are connected with the scattering on exchanged mesons.

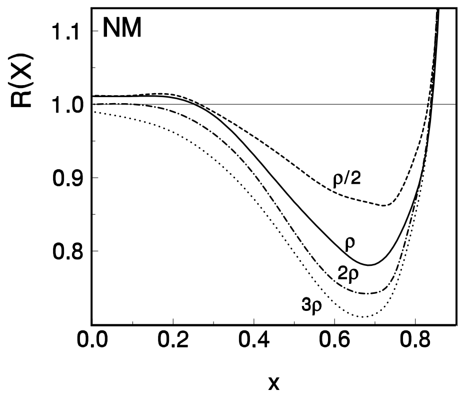

Fig. 4. Results for the ratio $R(x) = F_2^{\rm NM}(x)/F_2^N(x)$ which shows the evolution, without nuclear sea, of the nucleon SF for ZM2 model of RMF for densities $\rho = \rho_0/2, \rho_0, 2\rho_0, 3\rho_0$. $\rho_0 = 0.16\ fm^{-3}$ is the equilibrium density for NM.

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