European Muon Collaboration effect in deuteron and in three-body nuclei

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We investigate the European Muon Collaboration (EMC) effect in deuteron and three-body nuclei. The binding and the Fermi motion effects of nucleons are studied in detail. They reduce the structure function ratios from unity by less than 5%. Roles of excess pion are also discussed in connection with the momentum sum rule, which is violated by the binding effect.

I. INTRODUCTION

The European Muon Collaboration (EMC) effect[1] is that the nuclear structure functions per nucleon are significantly different from those of deuteron. It has attracted much attention because it might imply that the quark distribution in a nucleon is considerably distorted inside nuclei.[2] Recently, Akulinichev and Birbrair *et al.*[3] have presented a plausible explanation: The binding and the Fermi motion of nucleons in nuclei give most of the EMC effect. They failed, however, to reproduce its dependence on the nuclear mass number A observed in the SLAC experiment.[4] This is due to the saturation property of middle and heavy nuclei that the averaged value of separation energy is almost independent of A. They also could not explain the behavior of the structure function ratio at $x \lesssim 0.3$.[2] These are the main problems in dispute on the EMC effect. A naive inclusion of the binding effect of nucleons violates the energy-momentum sum rule for the quark distributions. How to restore the sum rule is another problem to be discussed. The most likely candidate to do so is excess pion in nuclei. The pion may contribute to the enhancement of the structure function ratio at small x.

In order to check the conventional nuclear effect quantitatively and to explore a possibility of the other mechanisms, we had best start with the light nuclei since their structure is well understood. In this paper, we study the EMC effect in deuteron and three-body nuclei, because excellent wave functions for the nuclei are available. We take the following conventional picture. Nuclei can be well understood in terms of nucleons and, if needed, meson degrees of freedom. Their properties do not change inside nuclei. Hence, we start from a study on the contributions of single nucleon: The binding and the Fermi motion effects are considered on the same kinematical basis. Next we include pion to restore the momentum sum rule. If they cannot give a satisfactory explanation of the EMC effect, then we will proceed to the other nuclear effects such as final-state interactions and/or additional mechanisms such as the Q^2 rescaling[5] and the explicit appearance of quark degrees of freedom.[6]

In the next section, the binding and the Fermi motion effects are calculated by using realistic nuclear wave functions. In Sec. III we discuss the role of excess pion as one of the binding quanta. By requiring it to recover momentum sum rule, which is violated by the binding effect, we estimate pionic contributions. Section IV is devoted to the summary and remarks. Unfortunately, there is no conclusive experimental data on the EMC effect in deuteron and three-body nuclei. We hope that such experiments are performed soon, especially for ^{3}He.

II. THE BINDING AND THE FERMI MOTION EFFECTS

At first, we will study the binding and the Fermi motion effects of nucleons in nuclei. We adopt the incoherent impulse approximation:[3,7] One nucleon with a momentum $\mathbf{p}$ and separation energy ϵ_f is knocked out off the nucleus A, and the spectator remains in the corresponding excited state $(A-1)^f$ (see Fig. 1). Then, the nuclear electromagnetic tensor is given by the convolution of the spectral function $S(p)$ and the nucleon tensor:

$$W^A_{\alpha\beta} = \int d^4p\, S(p) \frac{\nu}{\nu'} W^N_{\alpha\beta} \quad \text{with } m\nu' = pq \;, \qquad (1)$$

where m is the nucleon mass, and $q=(\nu,\mathbf{q})$ is the momentum transfer to the nucleus in the laboratory frame. The

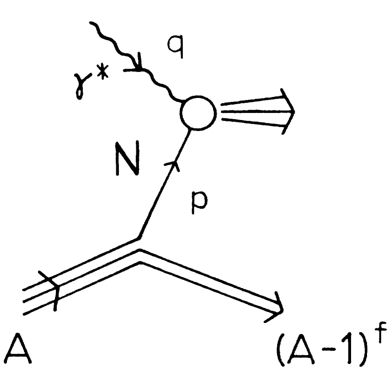

FIG. 1. Single nucleon contribution to the nuclear deep-inelastic scattering.

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nucleon tensor has the gauge invariant form:[7]

$$W^N_{\alpha\beta}=W^N_1(-g_{\alpha\beta}+q_\alpha q_\beta/q^2)+W^N_2/m^2[(p_\alpha-pqq_\alpha/q^2)(p_\beta-pqq_\beta/q^2)]$$

with two independent Lorentz-invariant functions W^N_1 and W^N_2. The nuclear tensor has a similar expression. In Eq. (1) we have introduced the factor ν/ν' to insure the parton-number sum rule in the Bjorken limit[8] (see Sec. III). The spectral function gives the four-momentum distribution of a nucleon in the nucleus:[9]

$$S(p)=\sum_f |\phi_f(\mathbf{p})|^2\delta[p^0-E_f(\mathbf{p})] , \tag{2}$$

where

$$E_f(\mathbf{p})=M_A-M^f_{A-1}-\frac{\mathbf{p}^2}{2M^f_{A-1}}=m-\epsilon_f-T_R ,$$

and $|\phi_f(\mathbf{p})|^2$ is the probability of finding a nucleon with given momentum $\mathbf{p}$ and the spectator residue in the state f. Furthermore, M_A and M^f_{A-1} are masses of the target and spectator nucleus, respectively. The separation energy is defined as $\epsilon_f=M^f_{A-1}+m-M_A$, and T_R is the recoil energy. Here we adopt the nonrelativistic kinematics. The spectral function is normalized to the mass number A with respect to the four-momentum p.[10]

By comparing specific components on both sides of Eq. (1) with the expressions of $W^{N,A}_{\alpha\beta}$, and by defining the nuclear structure function per nucleon as $F^N_2(A)\equiv\nu W^A_2/A$, we obtain

$$F^N_2(x,Q^2;A)=\frac{1}{A}\int d^4p\, S(p)\frac{1}{\omega'\tilde{\omega}}\left[\left[\omega'+\frac{2p_3}{q_3}\right]^2+\frac{2(\mathbf{p}^2-p_3^2)}{Q^2}\frac{\nu^2}{q_3^2}\right]\times F^N_2(\tilde{x},Q^2) , \tag{3}$$

where

$$\omega'=\frac{2m\nu'}{Q^2},\quad \tilde{\omega}=\frac{2m\tilde{\nu}}{Q^2}=\frac{1}{\tilde{x}} ,$$

$$\tilde{\nu}=\nu'-\frac{m^2-p^2}{2m},\quad Q^2\equiv-q^2 ,$$

and x is the Bjorken variable. We have chosen the third axis as parallel to $\mathbf{q}$. The off-shell structure function of interacting nucleon is assumed to be the same as the on-shell one with same Q^2 and final-state mass.[11]

For three-body nuclei, we use the spectral function given by the Faddeev wave function calculated by the Sendai group[12] with the Reid soft-core (RSC) potential. Then, $\phi_f(\mathbf{p})$ in Eq. (2) is given as

$$\phi_f(\mathbf{p})=\sqrt{3}\langle\psi^f_2,\mathbf{p}|\Psi_3\rangle ,$$

where $|\Psi_3\rangle$ is the ground-state wave function, and $|\mathbf{p},\psi^f_2\rangle$ is the state where the interacting nucleon has a relative momentum $\mathbf{p}$ to the correlated two-nucleon subsystem. The suffix f specifies that the spectator is deuteron or in some scattering state. For $\phi_f(\mathbf{p})$, the charge symmetry between ^{3}He and ^{3}H is assumed. We fix the binding energies of the nuclei to the experimental values: $E_3=-7.718$ MeV for ^{3}He and -8.482 MeV for ^{3}H, because the separation energies play a crucial role to determine the structure function ratio.[3] All results are calculated in the finite domain of the spectral function where $\epsilon_f+E_3\le 200$ MeV and $|\mathbf{p}|\le 6.5\ \text{fm}^{-1}$. In this domain, it is renormalized to $A=3$. In the case of deuteron, the spectator is a single nucleon. Hence, the spectral function consists of only one term. For it, we use the momentum distribution given by the same two-body potential as the above. As the nucleon structure function, the analytic function of Glück *et al.*[13] is employed in the following results.

In Fig. 2 we show the results for deuteron at $Q^2=5$, 10, and 100 GeV^2. We have defined the EMC effect of deuteron by the structure function ratio to the free nucleon (averaged over proton and neutron). The data points, in the region of $Q^2\lesssim 20$ GeV^2, are taken from the compilation by Bodek and Simon.[14] Our results are compatible with the experiments. The reduction of the structure function is almost due to the recoil energy term in Eq. (2) (see Table I).

Next, the results for ^{3}He are shown in Fig. 3. The structure function ratios are corrected for the isoscalarity:

$$\frac{F^N_2(^3\text{He})}{F^N_2(d)}\equiv\frac{3F^N_2(^3\text{He})}{2F^p_2+F^n_2}\frac{F^p_2+F^n_2}{2F^N_2(d)} . \tag{4}$$

For reference, the SLAC data[4] for ^{4}He (averaged over the Q^2 region of $2\le Q^2\le 15$ GeV^2) are shown in the figure. The deviation of ratios from unity is at most 3% at $x\lesssim 0.8$. This is about one-third of the effect observed in ^{4}He. In the figure, we also give the result for ^{3}H at $Q^2=10$ GeV^2. The difference between ^{3}He and ^{3}H at the same Q^2 comes from the difference between proton and neutron structure functions, which remains even after the isoscalarity correction.

It should be noted that all results in the above are rather insensitive to the nucleon structure function and the

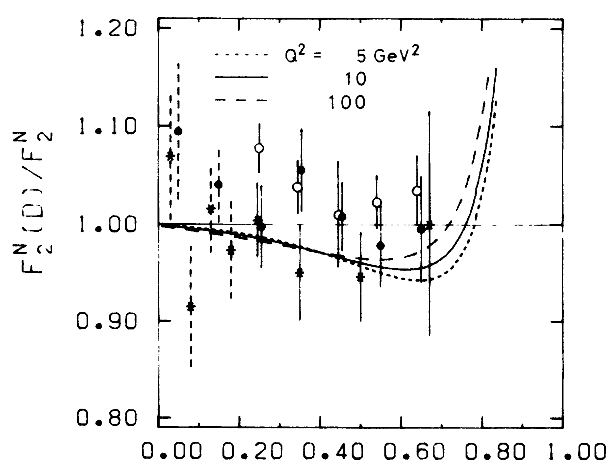

FIG. 2. Binding and Fermi motion effects in deuteron. The experimental data are taken from the compilation of Ref. 14.

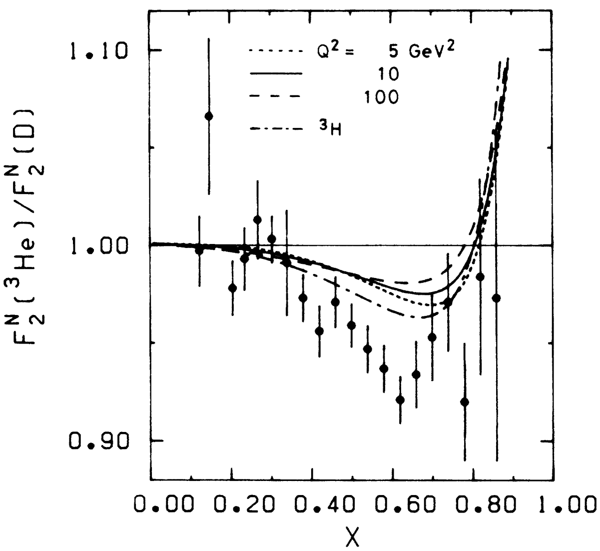

FIG. 3. The same as Fig. 2, but for ^{3}He. The result for ^{3}H is given at $Q^2=10$ GeV^2. For reference, the SLAC data (Ref. 4) for ^{4}He are shown.

two-body potential to calculate the wave functions.

Before closing this section, we will compare the above results with an estimation of the shell model. It is instructive to make clear what we have done with the realistic wave functions, though the single-particle model has little meaning for deuteron and ^{3}He. In the shell model, final states of the spectator nucleus are limited to one-hole states. Hence, we calculate the spectral function of ^{3}He in the following way. For proton knock-out process, the spectator is deuteron: The spectator energy is $\epsilon_f=E_3-E_2$, where E_2 is the binding energy of deuteron. For neutron knock-out process, we take $\epsilon_f=E_3$ because the spectator of two protons has no bound state. Moreover, the $\phi_f(\mathbf{p})$ is taken to be a harmonic oscillator wave function with oscillator length $b=1.97$ fm, which reproduce the experimental value of charge RMS radius.[15] For deuteron, the separation energy is $\epsilon_f=E_2$ and we take $b=2.116$ fm.[16]

The structure function ratio of ^{3}He to deuteron at $Q^2=10$ GeV^2 is shown in Fig. 4. The result is quite similar to the one given by the realistic wave functions. In particular, they are almost identical for $x \lesssim 0.6$. The reason why the difference is so small is as follows. If we take the large-Q^2 limit [see Eq. (5)] and neglect the tiny effect of the Fermi motion for small x, the nuclear structure function Eq. (3) can be well approximated as

$$F_2^N(x,A)=F_2^N(x)-\frac{\langle\epsilon\rangle+\langle T_R\rangle}{m}\frac{\partial}{\partial x}F_2^N(x) ,$$

where $\langle\epsilon\rangle$ and $\langle T_R\rangle$ are the expectation values of the separation energy and recoil energy per nucleon, respectively, such as

$$\begin{bmatrix}\langle\epsilon\rangle\\ \langle T_R\rangle\end{bmatrix}=\int d^4p\, S(p)\begin{bmatrix}\epsilon\\ T_R\end{bmatrix} .$$

Therefore, the deviation of the ratio from unity is proportional to the difference of $\langle\epsilon\rangle+\langle T_R\rangle$ between ^{3}He and deuteron. In Table I we give the values of $\langle\epsilon\rangle$, $\langle T_R\rangle$ and their sums for the nuclei with the realistic and harmonic oscillator wave functions. We can see that the values of $\langle\epsilon\rangle+\langle T_R\rangle$ play a decisive role in the binding effect on the EMC effect. Our realistic wave function of ^{3}He certainly includes some amplitudes where the spectator is in the higher-excited states. Accordingly, the value of $\langle\epsilon\rangle$ is larger than that of the shell model. But, finally, the difference of $\langle\epsilon\rangle+\langle T_R\rangle$ between ^{3}He and deuteron by the realistic wave functions is almost equal to that by the harmonic oscillator ones. This is due to the effect of recoil terms.

At large x, the Fermi motion effect becomes also important. High-momentum components of wave functions contribute to the structure function in a complicated way: The binding effect reduces it through the recoil term, while the Fermi motion effect increases it by smearing over a wide range of x. A little faster rising of the ratio in Fig. 4 is a result of this competition.

III. RESTORATION OF THE MOMENTUM SUM RULE AND PIONIC CONTRIBUTION

Next, we will study roles of pion in connection with the momentum sum rule. In the Bjorken limit, the nuclear structure function given by Eq. (3) is reduced to the convolution form:

$$F_2^N(x;A)=\int_0^{M_A/m} dz\, f_N(z)F_2^N(x/z) , \tag{5}$$

where z is the light-cone momentum fraction of the nucleus carried by a nucleon, and its distribution $f_N(z)$ is given as

$$f_N(z)=\frac{1}{A}\int d^4p\, S(p)\delta\left[z-\frac{p^0-p^3}{m}\right] . \tag{6}$$

TABLE I. Values of $\langle\epsilon\rangle$ and $\langle T_R\rangle$ of deuteron and ^{3}He with the realistic (RSC) and harmonic oscillator (HO) wave functions (in units of MeV).

	$\langle\epsilon\rangle$	$\langle T_R\rangle$	$\langle\epsilon\rangle+\langle T_R\rangle$
D (RSC)	2.225	11.05	13.28
^{3}He (Faddeev + RSC)	11.35	6.324	17.67
		difference	4.394
D (HO)	2.225	2.605	4.830
^{3}He (HO)	6.235	2.674	8.909
		difference	4.079

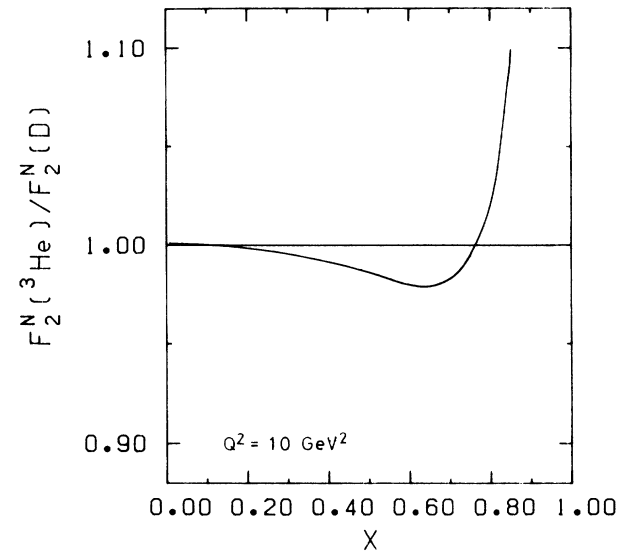

FIG. 4. Shell-model estimation of the EMC effect for ^{3}He. $Q^2 = 10$ GeV^2.

From this equation, we obtain

$$\int_0^{M_A/m} dz\, f_N(z) = 1 \tag{7a}$$

and

$$\int_0^{M_A/m} dz\, z f_N(z) \equiv \langle z \rangle_N = 1 - \frac{\langle \epsilon \rangle + \langle T_R \rangle}{m} . \tag{7b}$$

Equation (7a) insures the parton-number sum rule. However, Eq. (7b) means that the momentum sum rule is violated:

$$\langle z \rangle_N \neq \frac{M_A}{Am} = 1 + \frac{E_A}{Am} , \tag{8}$$

where E_A is the binding energy of the nucleus. This remark leads us to introduce quanta responsible for the nuclear binding. They are expected to carry a momentum fraction η:

$$\eta = -\frac{\langle V \rangle}{m} , \tag{9}$$

where $\langle V \rangle$ is the potential energy per nucleon. Including this fraction, we can recover the momentum sum rule:

$$\langle z \rangle_N + \eta = 1 + \frac{E_A}{Am} . \tag{10}$$

Among these quanta, we should first study roles of pion field in nuclei. In the remains of this section, we simply estimate its contributions to the structure functions of deuteron and ^{3}He.

Following Eq. (5), the pionic contribution is assumed to be given by

$$\delta F_2^N(x;A) = \int_0^{M_A/m} dz\, f_\pi(z) F_2^\pi(x/z) . \tag{11}$$

The pion structure function $F_2^\pi(x)$ has been observed in the experiment of J/ψ and dimuon productions. We use the analytic function given by Owens[17] (his second set). The momentum distribution $f_\pi(z)$ is constructed as follows. In a free-nucleon case, the contribution of pion cloud is given by Eq. (11) and[18,19]

$$f_\pi^N(z) = \frac{3g^2}{16\pi^2} z \int_{t_0}^{\infty} dt \frac{t\, |F(t)|^2}{(t+m_\pi^2)^2} \quad \text{with } t_0 = m^2 \frac{z^2}{1-z} , \tag{12}$$

where $-t$ is the four-momentum squared of pion, and g ($=13.5$) is the coupling constant. The form factor $F(t)$ at the πNN vertex is taken as

$$F(t) = \exp\left[-\lambda \frac{t+m_\pi^2}{m_\pi^2}\right] .$$

The cutoff parameter λ plays the most substantial role. When the nucleon is embedded in a nucleus, several modifications such as the polarization of nuclear medium occur. They may be expressed by an effective change of λ. In deuteron and three-body nuclei, this change is expected to be small. Then, $f_\pi^A(z)$ is given as

$$\begin{aligned} f_\pi^A(z) &= f_\pi^N(z;\lambda-\Delta\lambda) - f_\pi^N(z;\lambda) \\ &\simeq -\frac{\partial f_\pi^N(z)}{\partial\lambda}\Delta\lambda \\ &= \frac{3g^2}{16\pi^2}\Delta\lambda z \left[\frac{1}{\lambda}\exp\left[-2\lambda\frac{t_0+m_\pi^2}{m_\pi^2}\right] \right. \\ &\qquad \left. + \tfrac{1}{2} Ei\left[-2\lambda \frac{t_0+m_\pi^2}{m_\pi^2}\right]\right] , \end{aligned} \tag{13}$$

where

$$Ei(-z) = -\int_z^\infty dt \frac{e^{-t}}{t} .$$

We take $\Delta\lambda$ (>0) so that

$$\int_0^{M_A/m} dz\, z f_\pi^A(z) = \eta_\pi , \tag{14}$$

where the momentum fraction of pion η_π is determined by the momentum balance Eq. (10): Our wave functions give $\eta_\pi = 0.0130$ for deuteron and 0.0155 for ^{3}He.

The results with and without the pionic contributions at $Q^2 = 10$ GeV^2 are shown in Figs. 5 and 6. We have taken $\lambda = 0.026$,[19] but our results are insensitive to the value of λ. We can see that the pionic correction is at most a few percent effective and that it is limited to the region of $x \lesssim 0.4$. In particular, it is very small in ^{3}He, because the pionic contributions to the structure functions of deuteron and ^{3}He cancel each other out in the ratio.

In our estimation, excess pion in the nuclei gives a slight enhancement of the structure function ratios at small x. However, recent experimental data[20] show a complicated structure in the region of $x \lesssim 0.3$, namely, a small enhancement around $x = 0.15$ and a reduction near $x = 0$. This behavior might imply that the nuclear shadowing[21] works at small x. Hence, detailed information on the nuclei studied here, especially at small x, becomes more and more important to investigate the mechanism.

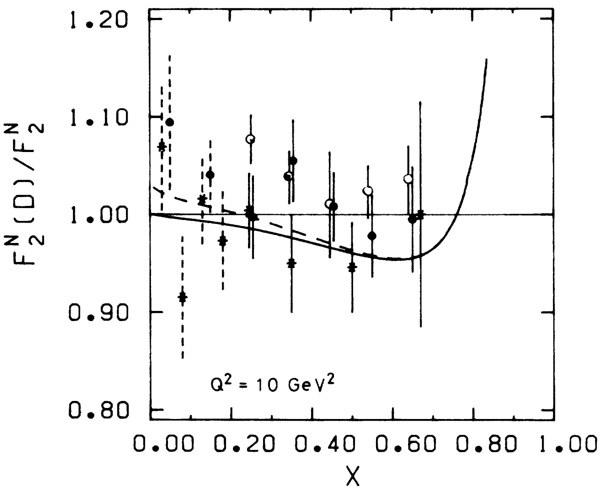

FIG. 5. Pionic effect in deuteron. The solid and dashed lines show the results without and with the pionic contribution, respectively.

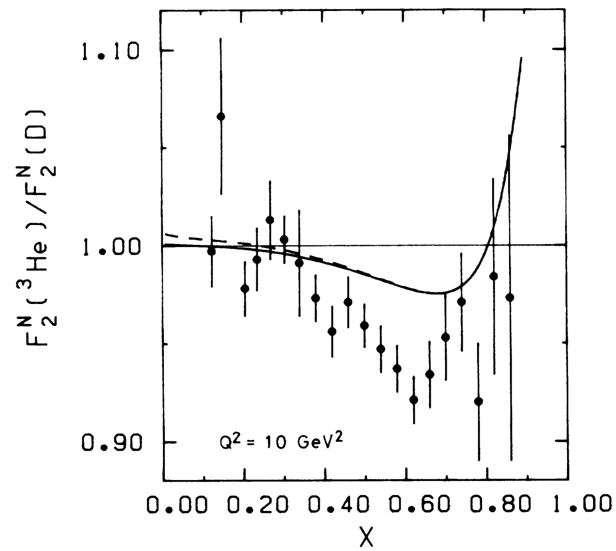

FIG. 6. The same as Fig. 5, but for ^{3}He. The SLAC data (Ref. 4) for ^{4}He are shown for comparison.

IV. SUMMARY AND REMARKS

We have carefully evaluated the binding and the Fermi motion effects of nucleons on the EMC effect in deuteron and three-body nuclei by using their realistic wave functions. The result for deuteron is compatible with the present data. The results for ^{3}He and ^{3}H show about half of the effect observed in ^{4}He. We also have studied the roles of pion. Although it is important to restore the momentum sum rule, its effect is fairly small. Until now, there is no experimental data to be compared with our results for ^{3}He. We hope that such experiment is performed soon.

In this paper, we have corrected the structure function ratios for the isoscalarity such as in Eq. (4). The results are almost independent of the used nucleon structure function. However, experiments give uncorrected ratios as $F_2^N(d)/F_2^p$ and $F_2^N(^3\mathrm{He})/F_2^N(d)$ because we cannot directly measure the neutron structure function. Therefore, we need to reduce the ambiguity due to the isoscalarity correction, that is, the uncertainty in the ratio of neutron structure function to proton one. The large errors of deuteron data in Fig. 2 just come out from it. For this reduction, we invoke neutrino and antineutrino scattering experiments. Moreover, we had better measure both the ratios of ^{3}He and ^{3}H to deuteron.

Finally, we would like to stress that few-body nuclei are very suited to investigate the origin of the EMC effect. We believe that our results are useful for this study. Of course, further investigations remain to be done. For example, the other nuclear effects such as final-state interactions, off shellness of nucleon structure function, and relativistic effect will contribute to the nuclear structure function. And, the Drell-Yan process[22] and nuclear structure functions beyond $x=1$ (Ref. 23) should be also explored. It seems premature to discuss the explicit quark degrees of freedom in the EMC effect.

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$$\int d^4p\, \alpha S_{\rm rel}(p)\theta(\alpha)\theta(A-\alpha)=A$$

$$\text{with } \alpha=(p^0-p^3)/M_A/A\ .$$

The step functions can be replaced by 1 to good approximation. The nonrelativistic version of the sum rule is given as

$$\int d^4p\, S(p)=A\ .$$

It seems, therefore, natural to relate the nonrelativistic spectral function as $S(p)=\alpha S_{\rm rel}(p)$. On the contrary, Frankfurt and Strikman have identified it as $S(p)=S_{\rm rel}(p)$ and argued that the contribution of the nuclear binding to the EMC effect is largely reduced. We think that their reduction is spurious.

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