Constituent-quark contribution to lepton-nucleus deep-inelastic scattering

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The distortion behavior of the nucleon structure function in the nuclear medium is extensively studied by the constituent-quark model. The nuclear medium makes the oscillator constants of the three constituent quarks in a bound nucleon weaken so as to soften the distribution of the constituent quarks inside a nucleon. In addition, this distribution is smeared by the Fermi motion among the bound nucleons. The sea quark distribution is hardened by the gluon recombination mechanism. The Mueller-Qiu mechanism is improved by taking the coexistence of nuclear shadowing and antishadowing into account and this formulation follows the momentum sum rule for the parton-quark system. A free parameter is chosen on the basis of the quantum chromodynamic vacuum model. The theoretical predications for a series of nuclei with $4 \leq A \leq 200$ within the kinetic region $0.001 \leq x \leq 0.9$ are in fair agreement with the existing lepton-nucleus data in the full x region. Especially, the model predicts that two crossover points of the cross-section ratios $R_\mu(x)$ with the horizontal axis $R_\mu(x)=1$ at low x move in the x direction as the mass number A increases.

I. INTRODUCTION

The European Muon Collaboration[1] (EMC) first found in the muon-nucleus deep-inelastic scattering (DIS) the difference of the structure function for a bound nucleon measured on a nuclear target as compared with that for a quasifree nucleon measured on a deuterium target. This phenomenon is called the EMC effect. It has motivated for years much experimental and theoretical interest in the features of the nucleons bound in a nucleus. Several models were proposed. Most of them can be divided into two major categories: one based on conventional nuclear physics[2-5] and another on the parton-quark model and perturbative quantum chromodynamics.[6-9]

Quite a few signs, however, show that there exist various kinds of origins for the EMC effect. We consider[10,11] that there is at least concern about two physical objects, the constituent quark and the current quark, on two different levels, respectively. A probe in the DIS process "sees" a current quark in a bound nucleon and so the nuclear effect reflects on the distortion of the parton distributions in a nucleon struck by the probe. On the other hand, since the nucleon, before struck by a probe, is bound in a nuclear target, it can be considered as a composite system consisting of three constituent quarks. It is generally agreed[12-17] that the constituent quark and the current quark have different physical contents and essences.

We need a model accounting for the structure of a free nucleon and will use it as a base for further studying its bound features. Such a model should first include the relationship between the constituent quark and the current quark. Second, it will be in a good agreement with the DIS data on a free nucleon. In what follows we employ the expression of the nucleon structure function for a free nucleon used by Hwa and Zahir;[12,13]

$$F_N(x,Q^2)=\sum_c \int_x^1 dy\, G_N^c(y)F_{c/N}^v(x/y,Q^2) + \sum_c \int_x^1 dy\, G_N^c(y)F_{c/N}^s(x/y,Q^2) \qquad (1.1)$$

in which the symbol v (s) refers to the valence (sea) quark. The distribution function $G_N^c(y)$ of the constituent quarks in a nucleon has nothing to do with a probe, but the distribution function $F_{c/N}^{v(s)}(z)$ of the partons in a constituent quark, which is also called the structure function of the constituent quarks, has much to do with a probe. The structure function of the constituent quarks is universal and its evolution is determined by the Altarelli-Parisi equations.[18] The formula (1.1) states the independence of the structure function $F_{c/N}^{v(s)}$, governed by perturbative QCD, with the distribution function G_N^c, governed by nonperturbative QCD, i.e., the implication of the factorization assumption.

Neglecting nuclear shadowing temporarily we generalized formula (1.1) in the case of a bound nucleon.[10,11] We find that the behavior of the ratio $R_\mu=F_2^A(x,Q^2)/F_2^N(x,Q^2)$ of the nucleon structure function for a nucleus to that for a free nucleon in the region of the Bjorken variable $x>0.3$ can be described by the constituent-quark model (CQM). The ratio R_μ in the region $x<0.3$ is found to have a change deviating somewhat from unity.[10,11]

The primary data measured by the EMC (Ref. 1) show a considerable enhancement in the ratio R_μ at low x and they are not consistent with the data measured by the SLAC (Refs. 19–21) and BCDMS (Ref. 22) collaborations. The systematic errors of the experiments were deduced by the European Muon Collaboration themselves recently and their new data[23,24] show that the ratio R_μ is less than unity at small x. The results remove the confusion on the data at low x and require theorists simultaneously to include a consistent explanation of nuclear shadowing in their own models.

In this paper we will investigate in detail the existing data about the EMC effect in terms of the constituent-quark model associated with the modified gluon recombination mechanism and with conventional nuclear physics. We find that the influence of the nuclear medium, which is concerned with nonperturbative QCD, can distort the distribution of constituent quarks in a bound nucleon. Most parts of the influences, e.g., the Fermi motion, the binding energy, and the nuclear surface effect, can be simply treated by conventional nuclear physics as we all have well known. The rest of the influence, which is involved in the modification of nonperturbative QCD, may be phenomenologically described by the weakening of the spring elastic coefficients among the constituent quarks. On the other hand, the QCD interaction of a struck parton with the partons of other bound nucleons makes the sea quark distribution and the gluon distribution change so as to produce nuclear shadowing and antishadowing and they can be described by the improved gluon recombination mechanism. It will be seen that our results are quantitatively consistent with the existing data about the EMC effect.

The organization of the paper is as follows. We address ourselves to the properties of the constituent quarks in a free nucleon in Sec. II and those of the constituent quarks in a bound nucleon in Sec. III. Nuclear shadowing and antishadowing are studied in Sec. IV. The lepton-induced DIS process in a nuclear target is examined in Sec. V. A conclusion is given in the last section.

II. PROPERTIES OF CONSTITUENT QUARKS IN A FREE NUCLEON

A free nucleon in the constituent-quark model[12,13] can be considered as a composite system consisting of three constituent quarks. Their distribution in a nucleon satisfies the normalized condition

$$\int_0^1 G_n^c(y)dy=1 \tag{2.1}$$

and the momentum sum rule

$$\int_0^1 G_n^c(y)y\,dy=\tfrac{1}{3}\ . \tag{2.2}$$

The sum rule means that each constituent quark carries, on an average, one-third of the longitudinal momentum of a nucleon. The form of the distribution G_N^c is determined by the interactions among the constituent quarks and is independent of a probe. As stated in Refs. 12 and 25 the experimental data of multiparticle production at low P_T or the nucleon electromagnetic form factor can be used to determine a simple phenomenological form for the distribution G_N^c:

$$\begin{aligned}G_N^c(y)&=B^{-1}(\xi_N/2+\tfrac{1}{2},\xi_N+1)y^{\xi_N/2-1/2}(1-y)^{\xi_N}\\&=\tfrac{105}{16}y^{1/2}(1-y)^2\quad\text{with } \xi_N=2\ ,\end{aligned} \tag{2.3}$$

where B is the beta function.

It is a good approach to derive the distribution G_N^c according to the QCD principles. However, this approach is concerned in complicated nonperturbative QCD and is beyond what we have understood as yet. But it does not prevent us from analyzing the form of the distribution (2.3) in terms of some phenomenological models. For example, it is interesting, as Ref. 26 stated, that the distribution of the constituent quarks in a free nucleon derived by the covariant harmonic oscillator model,[27,28] is substantially close to the expression (2.3) for $y>0.4$ determined by the experiments. The distribution acquired by the covariant harmonic oscillator model is

$$\begin{aligned}G_{\mathrm{osc},N}^c(y)=&(3m_N/\sqrt{2\pi\omega_N})\\&\times\exp[-m_N^2(1-3y)^2/2\omega_N]\end{aligned} \tag{2.4}$$

in which m_N is the nucleon mass and ω_N the oscillator constant $\omega_N=m_N^2/2$ for a free nucleon. The relation between ω_N and the spring elastic coefficient k_N is $\omega_N^2=mk_N$. The spring elastic coefficient k_N has much to do with interactions among the constituent quarks. An alternative statement for the structure function of a harmonic oscillator, equivalent to $G_{\mathrm{osc},N}^c(y)$, is

$$G_{\mathrm{osc},N}^{\prime c}(y)\sim\exp[-\tfrac{3}{2}m_N^2R_N^2(y-\tfrac{1}{3})^2]\ ,$$

where R_N is the scale of the nucleon.[29]

The inconsistency of $G_{\mathrm{osc},N}^c(y)$ with $G_N^c(y)$ for low y, as shown in the Appendix, originates from the fact that the distribution function for low y comes from the contribution of correlated constituent quarks at a large distance, where a hadron has so large a deformation, that the interactions among the constituent quarks do not mean the harmonic oscillator model is not appropriate to the full y region, but rather that the model is only adopted to describe phenomenologically the dynamics for the constituent quarks in the range $0.4\le y\le 0.9$.

The shortcoming of the distribution $G_{\mathrm{osc},N}^c(y)$ mentioned above, of course, does not prevent us from establishing the relation of two parameters ω and ξ for $0.45\le y\le 0.9$ according to the consistence between those two distributions in the high-y region. Using that relation one can understand how nuclear surroundings affect the parameter ξ. In addition, the distribution $G_N^c(y)$ restricted by the conditions (3.1) and (3.2), depends on the parameter ξ. The ξ change in nuclear surroundings also reflects how the distribution $G_N^c(y)$ in the full y region is distorted by nuclear medium.

The structure function $F_{c/N}^{v(s)}(z,Q^2)$ of the constituent quarks in formula (1.1) is dependent on a probe.[12] It is assumed in more detail that a probe "sees" a structureless constituent quark at a low value $Q^2=\mu^2$, that is, $F_{c/N}^{v}(z,\mu^2)=\delta(1-z)$ and $F_{c/N}^{s}(z,\mu^2)=0$. Starting from the initial conditions and making use of the QCD evolution in the leading logarithm approximation[18] (LLA), the structure function $F_{c/N}^{v(s)}(z,Q^2)$ of the constituent quarks at high value Q^2 can be derived.[12] The detailed calculations,[13,30] however, show that there are obvious deviations of the sea quark distribution and the gluon one from the data measured in recent years. Its cause should be ascribed to the fact that the leading logarithm approximation at low Q^2 is not a good approximation, in other words, the delta function in the LLA is not a good initial condition. In order to acquire a correct base for recent work we need the structure function of a free nucleon fitting the recent data. Let us substitute the structure function $F_N(z,Q^2)$ of a free nucleon at some Q^2 extracted from recent data[31] and the distribution G_N^c into both sides of expression (1.1). With the aid of their corresponding moment equations a simple parametrized structure function $F_{c/N}^{v(s)}(z,Q^2)$ can be established. It was shown in Ref. 11 that the dependence of the structure function $F_{c/N}^{v(s)}(z,Q^2)$ on Q^2 has a much weaker influence on the cross-section ratio $R_\mu(x)$. Hence we can take some value of Q^2 to examine the structure functions. For instance, our results are

$$F_{c/N}^{v}(z,Q^2)=0.509z^{0.5}(1-z)^{0.3} ,$$
$$F_{c/N}^{s}(z,Q^2)=0.08(1-z)^2 \quad \text{at } Q^2=10\ (\text{GeV}/c)^2 . \tag{2.5}$$

The comparison of our parametrized expressions with recent data[32] is shown in Fig. 1 in which the structure function of a free nucleon is determined by formulas (2.5), (2.3), and (1.1). We now acquire the model for the structure function of a free nucleon, which contains the constituent quark and the current quark on two different levels and is consistent with recent data about a free nucleon. On the basis of that, in what follows we can afford to address ourselves to discuss how the nuclear medium distorts the distribution of constituent quarks and how nuclear shadowing and antishadowing affects the distribution function of the current quarks under the DIS condition.

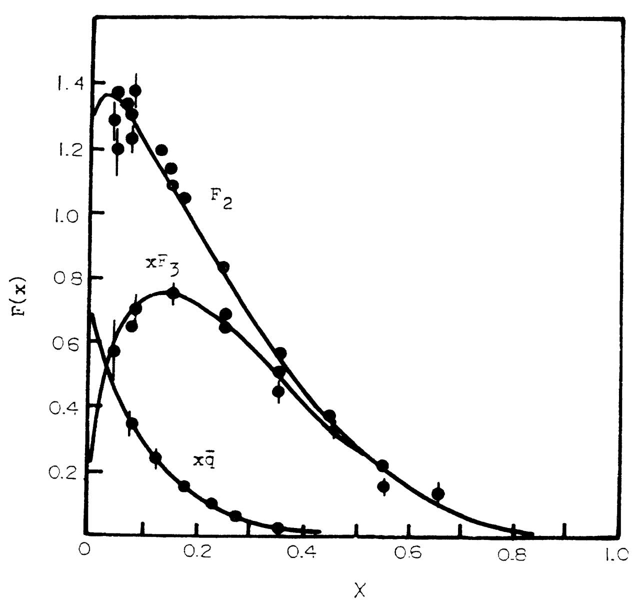

FIG. 1. The structure function $F_2(x,Q^2)$ for a free nucleon and its sea component $x\bar{q}(x,Q^2)$ and its valence component $xF_3(x,Q^2)$ and $Q^2=10$ (GeV/c)^2. They are determined by formulas (1.1), (2.3), and (2.5). The data are taken from Ref. 32.

III. FEATURES OF CONSTITUENT QUARKS IN A BOUND NUCLEON

It may be considered[11] that the interactions among bound nucleons merely change the distribution of the constituent quarks in a nucleon, i.e., $G_N^c \to G_A^c$. The average effective mass of the nucleons bound in a nucleus is $m^* = m_N - B_A$, where m_N is the nucleon mass. B_A here is the binding energy and is around 8 MeV. In comparison with the scale of about 1 GeV of the DIS process, the binding interactions among the nucleus in a nucleus can be neglected in the zeroth approximation and will be considered in the discussion of the following Fermi motion. Under the condition of this approximation, we think that a nuclear target in the DIS process can be considered as a fermion system consisting of quasi-independent A nucleons with the effective mass m^*. Under the condition of this approximation, we suppose that a bound nucleon still consists of three constituent quarks in the same manner as a free nucleon, i.e., the normalized condition is valid:

$$\int_0^1 G_A^c(y)dy = 1 . \tag{3.1}$$

In addition, not only do the binding interactions with a few MeV not change the longitudinal momentum carried by each bound nucleon in DIS, but also they do not change the average longitudinal momentum carried by each constituent quark. Therefore under the condition of the approximation, it seems to be reasonable to suppose that the momentum sum rule remains valid,

$$\int_0^1 G_A^c(y)y\, dy = \tfrac{1}{3} , \tag{3.2}$$

in the same manner as a free nucleon. The distribution of the constituent quarks in a bound nucleon, which satisfies the conditions (3.1) and (3.2) simultaneously, can in general be written as

$$G_A^c(y) = B^{-1}(\xi_A/2+\tfrac{1}{2}, x_A+1) y^{\xi_A/2-1/2}(1-y)^{\xi_A}$$
$$= B^{-1}(\delta_A+\tfrac{3}{2}, 2\delta_A+3) y^{\delta_A+1/2}(1-y)^{2(\delta_A+1)} , \tag{3.3}$$

where $\delta_A = \xi_A/2 - 1$ is called the distortion factor which reflects the influence of the nuclear medium on it.

The concrete influences of the nuclear medium on the nucleon structure function can be analyzed one by one as follows.

(1) We make use of the covariant harmonic oscillator model phenomenologically to describe[11] nonperturbative QCD which has not been understood as yet. The model decides the behavior of the constituent-quark distribution in the region $y > 0.4$. It is known from formula (2.4) that the change $k_n \Longrightarrow k_A$ (or $\omega_N \Longrightarrow \omega_A$) of the spring inelastic coefficient from the case of a free nucleon to that of a

bound nucleon, is a permissible approach for describing the nuclear influence on the distribution G_N^c. Denote the expression (2.4) distorted by the nuclear effects,

$$G_{\mathrm{osc},A}^c(y)=(3m_N/\sqrt{2\pi\omega_A})\times\exp[-m_N^2(1-3y)^2/2\omega_A] \;. \tag{3.4}$$

In the Appendix the QCD vacuum model is used to show that the nuclear medium makes the interactions among the constituent quarks inside a bound nucleon weaken so that $k_A<k_N$ or $\omega_A<\omega_N$. It means that the decrease of the distribution $G_{\mathrm{osc},A}^c(y)$ in a bound nucleon with y increasing is faster than that in a free nucleon, or equivalently the distortion factor δ_A in expression (3.3) is larger than zero.

The relation between the distortion factor δ_A and the oscillator constant ω (or its elastic coefficient k_A) can be established by comparing the distribution $G_{\mathrm{osc},A}^c(y)$ with the distribution $G_A^c(y)$ in the high y region since the distribution $G_{\mathrm{osc},N}^c$ at high y, derived by the covariant harmonic oscillator model, is close to the distribution G_N^c extracted from experimental data, as mentioned in the preceding section. Let us make a small change in ξ_N: $\xi_N\to\alpha\xi_N\equiv\xi_A$, where α is somewhat larger than 1, so that the distribution $G_N^c(y)$ for $y\geq 0.4$ also changes. Or alternatively one makes a small change in ω_N in the distribution G_{osc}^c: $\omega_N\to\omega_A=\beta\omega_N$ with β being somewhat smaller than 1, to gain the same object. Assume that the relation between α and β is

$$\alpha=c_1+c_2\beta+c_3\beta^2 \;. \tag{3.5}$$

The optimal fit for the parametrized set c_1, c_2, and c_3 is

$$\xi_A/\xi_N=5.72-8.12(\omega_A/\omega_N)+3.40(\omega_A/\omega_N)^2 \;. \tag{3.6}$$

It means the weakening of the oscillator constant ω results in the increase of the power for $(1-y)$, which decides the behavior of high y for the distribution $G_A^c(y)$. In addition, since there is only a parameter in the distribution (3.3) or (2.3), one can know how the change of high y affects that of low y for the distribution $G_A^c(y)$ although the dynamics mechanism for the low-y region is as yet not understood.

Next the dependence of ω_A (or k_A) on the mass number A is taken into account. The weakening of the spring elastic coefficient of a nucleon, distorted by the nuclear medium, should have much to do with the interaction between that nucleon and its neighboring nucleons. Hence the weakening dk_A of the elastic coefficient for a bound nucleon is enhanced by the increase in the number of its neighboring nucleons. The average number of the neighboring nucleons around a nucleon in the nuclear surface region is about half of that in the nuclear central region. Set $P_s(A)$ to be the probability of finding a nucleon inside the surface region with the mass number A and $k_A=k_N-dk_A$ to be the average elastic coefficient for a bound nucleon. Then it is reasonable to estimate

$$dk_A=P_s(A)dk_0/2+[1-P_s(A)]dk_0 \;, \tag{3.7}$$

in which k_0 is the relative reduced value of the spring elastic coefficient for a nucleon bound in the nuclear central region. $P_s(A)=1-V_c/V$ with V being the nuclear volume and $(V-V_c)$ being that of its surface region. The probability $P_s(A)$ can be obtained in terms of the nuclear density[33]

$$\rho_A(r)=\rho_0/\{1+\exp[(r-R)/b]\}$$

with $\rho_0=0.17$ nucleon/(fm)^3, $b=0.54$ fm, and $R=1.12A^{1/3}-0.86A^{-1/3}$. The surface thickness is given by $D=(4\ln 5)b$. As for light nuclei with the mass number $A\leq 8$, all the nucleons are inside the surface region, that is, $P_s(A)=1$. The value dK_A, however, should be continuously reduced with its neighboring nucleon number decreasing. In place of formula (3.6), a linear-approximation form is used such that

$$dk_A=(dk_0/2)(A-1)/7 \quad \text{for } A\leq 8 \;. \tag{3.8}$$

The relative reduced value dk_0 in formula (3.7) is considered as a parameter and its value can be estimated by the QCD vacuum model (see the Appendix).

(2) Another influence of the nuclear medium on bound nucleons is the Fermi motion of those nucleons in a nucleus. The Fermi motion is able to smear out the distribution of the constituent quarks in a bound nucleon[2]

$$G_A'^c(y)=\int_y^1 f_A^N(\xi)G_A^c(y/\xi)d\xi \;, \tag{3.9}$$

where

$$f_A^N(\xi)=\tfrac{3}{4}(m_N/k_F^A)^3[(k_F^A/m_N)^2-(\xi-\eta_A)^2] \tag{3.10}$$

for $\eta_A-k_F^A/m_N<\xi<\eta_A+k_F^A/N$; otherwise $f_A^N(\xi)=0$. Here $\eta_A=1-B_A/m_N$ is taken. The average effective mass of the bound nucleons is $m_N^*=m_N-B_A$. The Fermi momentum k_F^A and the binding energy B_A are taken as the values required by nuclear physics.

Equation (3.10) with $\eta_A=1-B_A/m_N$ means that the bound nucleon is nearly on shell and the influence of the nuclear binding effect on the bound nucleon may be negligible. The assumption about the nucleon on shell in the CQM does not imply that there are no interactions among the nucleons, but rather that the so-called binding effect does not conspicuously change the quark distributions. In order to account for this point, it is significant to compare the CQM with the nuclear binding model or the pionic enhancement model. We believe the following (i) Whether a bound nucleon is on-shell or off-shell, is model-dependent. For a bound nucleon, only a small amount of energy (approximately a few MeV) makes a contribution to the defect mass. Most of the interacting energy in the form of a mesonic cloud exists in a nucleus. A bound nucleon is off shell if the mesonic cloud is considered to be independent of it just as in the pion enhancement model.[34] In contrast with that, a bound nucleon is on shell if the mesonic cloud is considered as a part of a constituent quark as shown in Eq. (3.2). (ii) Whether a bound nucleon is on shell or off shell for a model is relevant to what extent the parton distribution in a nucleon is disturbed by the nuclear binding effect. (iii) A key point is that it is not understood yet from which partonic component of a nucleon the momentum and energy carried by its mesonic cloud come, primarily according to the information of conventional nuclear physics.

For the binding models[3-5] one takes $\eta'_A=(m_N-\langle \varepsilon \rangle)/m_N$ with $\langle \varepsilon \rangle$ the average separating energy. Replacing $(\xi-1)$ by $(\xi-\eta'_A)$ in Eq. (3.10), one gets the x rescaling for the bound nucleons: $x_A=x/\eta'_A$. As mentioned by Smith,[34] this procedure is only a guess. In fact, from the viewpoint of the parton picture, the above approach includes a hidden assumption: the x-rescaling coefficient η'_A is independent of the parton momentum and the parton flavor. This means that the valence quarks play an important role in the nuclear binding effect since the valence quarks carry larger momenta.

However, there is another possibility, we think–that the exchange of the wee partons[35] dominates the nuclear binding effect. In this case, the loss of the momentum fraction mainly arises in the sea quarks and the gluons. The change of momentum of the partons is

$$\langle \text{sea} \rangle_2 \rightarrow \langle \text{sea} \rangle_2 - \alpha(1-\eta_A) \equiv K_s \langle \text{sea} \rangle_2$$

and

$$\langle \text{gluon} \rangle_2 \rightarrow \langle \text{gluon} \rangle_2 - (1-\alpha)(1-\eta_A) \equiv K_g \langle \text{gluon} \rangle_2 ,$$

where $\langle\ \rangle_2$ denotes the average momentum fraction and $\alpha<1$. A straightforward way to change the momentum fraction carried by the sea quarks is to replace the primary distribution of sea quarks $F_s(x)$ by the distribution $K_sF_s(x)$ provided that the momentum loss of the sea quark is proportional to the quark density. Moreover, according to Eq. (3.2), in whatever form (pionic or other mesonic) the wee partons exchanged appear, the three constituent quarks are, on average, "dressed" in them. Since the momentum lost is transferred to the wee partons in the convolution form $\int f_\pi(h)F_\pi(x/y)dy$ with f_π the pion distribution and F_π the pion structure function.[34] The above convolution form with the average momentum fraction $(1-K_s)\langle \text{sea} \rangle_2$, is roughly similar to the distribution form $(1-K_s)F_s(x)$ provided that

$$F_\pi(0)\int f_\pi(y)dy \equiv (1-K_s)F_s(0) .$$

In this case, the probe cannot distinguish between the distribution of the original sea quarks and one of the compensative sea quarks. The interactions among bound nucleons might be considered to have nearly little interference on the nucleon structure function, i.e., the bound nucleon is on shell.

At the close of this section we want to emphasize that in the nuclear medium the normalized condition (3.1) and the momentum sum rule (3.2) are not trivial. The former, associated with the universality of the structure of the constituent quarks, plays an important role in the behavior of the cross-section ratio $R_\mu(x)$ which tends to unity as x approaches zero,[10,11] and the latter also plays a significant role in the momentum preservation of each species of the parton-quark system. We have postponed discussing them until the conclusion of the last section.

IV. NUCLEAR SHADOWING AND ANTISHADOWING

As stated in the preceding section, the influence of the nuclear medium on the nucleon structure function reflects on the distribution $G_A^c(y)$ of the constituent quarks and does not change the parton distribution in a constituent quark. Under the condition of the DIS process, however, because of the uncertainty relation, a parton with a low x, absorbing a virtual photon, is able to shadow some partons in other neighboring nucleons so that such a shadowing changes, in part, the parton distribution in a constituent quark. This is so-called nuclear shadowing.

Nikolaev and Zakharov[36] first qualitatively predicted the partonic shadowing effect mentioned above. Recently Mueller and Qiu[37] studied the interactions among shadowed partons and pointed out that since the number density of sea quark-antiquark pairs is much smaller than that of gluons, the contribution of the annihilation between sea quarks and antiquarks to the shadowing effect can be neglected. In this approximation the annihilation among the shadowed gluons from different nucleons not only degrades the gluon distribution function, but also indirectly degrades the sea quark distribution through the modified Altarelli-Parisi equations. Meanwhile most of the valence component is expected to remain the same. In Ref. 38 Qiu gave a simple parametrization of the shadowing factor $R_s(x,A)$ for the sea quark distribution for $0\le x\le x_n$. Here x_n is considered to be such a point at which the onset of shadowing would occur if the longitudinal size of overlapping gluons reached the longitudinal distance between two neighboring nucleons.[38]

It should be pointed out here that the gluon recombination only changes the distribution of the gluon number density, but does not change the average momentum of gluons. As in the case of gluons, the recombination process also affects the distribution of the number density of sea quarks, but does not affect their average momentum according to the modified Altarelli-Parisi equations. As a consequence the average momentum of each component (valence quarks, sea quarks, and gluons) in a bound nucleon can be preserved in the gluon recombination process. According to the momentum balance, the part of the gluon momentum lost in the nuclear shadowing process should be compensated in terms of new gluons with large x produced by the gluon annihilation. The process, required by the momentum balance, is also suitable for sea quarks. The enhancement of both the gluon distribution and the sea quark distribution is called nuclear antishadowing.

Nikolaev and Zakharov conjectured[36] that the shadowing and antishadowing belong to two different x regions. But now the case becomes more complicated: the fraction x carried by the enhanced gluons or the enhanced sea quarks (i.e., the contribution coming from the antishadowing just mentioned), can lie not only in the nonshadowing region $x>x_n$, but also in the shadowing region $0<x\le x_n$. For instance, the enhanced momentum fraction x carried by a newly produced gluon, which contributes to the antishadowing, could be smaller than x_n if two gluons with $x<x_n/2$ annihilated each other. On ac-

count of the coexistence (or admixture) of the shadowing and antishadowing effects, we cannot apply Qiu's formulation, which is merely suitable for the shadowing in the region $0<x\leq x_n$.

A convenient approach, without further considering the details about the coexistence or admixture of the shadowing and antishadowing, is to suppose that after considering the coexistence of both the effects at low x, the modified distribution of sea quarks has the following form extensively used in references: $F^s_{A,\mathrm{sa}}(x)=\alpha(1-x)^\beta$ for $0\leq x\leq x_n$. The subscript sa represents the inclusion of the shadowing and antishadowing so as to discriminate it from the sea quark distribution $F^s_A(x)$ in a bound nucleon, in which the two effects are not considered. The antishadowing can be omitted if x is close to zero. At that time the distribution of sea quarks degenerates into Qiu's shadowing formulation $F^s_{A,\mathrm{sa}}(0)=F^s_A(0)R_s(0,A)$, i.e., $\alpha=F^s_A(0)R_s(0,A)$, in which $R_s(0,A)=1-0.1(A^{1/3}-1)$ is a shadowing factor quoted by Qiu.[38] Here $(A^{1/3}-1)\equiv n_s$ denotes the effective nucleon number for which those nucleons are shadowed in a nucleus with the mass number A.

The antishadowing has remained until $x=2x_N$ and may disappear for $x\geq 2x_n$ since two gluons with $x=x_n$ can recombine into a new gluon with the enhanced momentum fraction $x=2x_n$, which makes a contribution to the antishadowing. The estimation, however, is approximate. In the infinite momentum frame the longitudinal size of the gluons with the momentum $k_z=xp$ is around $\Delta z\sim 1/xp$. The average distance between two neighboring nucleons along the longitudinal direction is $\Delta z_n\sim 2R_N m_N/p$. In fact, a nucleon has a definite size. Even though the momentum fraction x carried by gluons is larger than $x_n\equiv 1/2R_N m_N$, those gluons with the longitudinal size smaller than Δz_n, would have a small overlapping probability provided those gluons are close enough to each other. Therefore, strictly speaking, although the antishadowing rapidly decreases as x tends to $2x_n$, there is still a short antishadowing tail for $x>2x_n$.

Shadowing is responsible for the change of the sea quark distribution in the small-x range ($x<x_n$), whereas antishadowing is responsible for that of the distribution function $F^s_{A,\mathrm{sa}}$ in the larger-x range ($x>x_n$). Let us estimate the rate $f(x)$ for $F^s_{A,\mathrm{sa}}$ asymptotically tending to $F^s_A(x)$ as follows. The rate $f(x)$ was supposed to be the Gaussian form in a previous work.[39] Here we will estimate a plausible form for it. There are two factors involved in the rate. First, it is recognized in the gluon recombination mechanism that the gluon recombinations, on an average, proceed in two gluons with the same momentum fraction.[37] The magnitude of antishadowing at $x=x_0$ ($x_0>x_n$) is proportional to the square of the gluon density at $x=x_0/2$, the relative rate $f(x)\sim G^2(x_0/2)/G^2(x_n/2)$. Second, the two gluons with the momentum fraction $x_0/2$ fuse into a gluon with $x=x_0$. The recombination process for $x_0/2>x_n$ is solely restricted within the region of a spherical segment of cone with the height $h=1/m_N x_0$ between two neighboring nucleons with radius R. This restriction is

$$f(x)\sim 2\pi h^2(R-h/3)/(\tfrac{4}{3}\pi R^3)\ .$$

Based on the two factors mentioned previously one can write

$$F^x_{A,\mathrm{sa}}(x,A)=\begin{cases}\alpha(1-x)^\beta, & 0\leq x\leq x_n\ ,\\ [\alpha(1-x_n)^\beta-f^s_A(x_n)]f(x)+F^s_A(x), & x\geq x_n\end{cases} \tag{4.1}$$

where

$$f(x)=\begin{cases}[(1-x/2)/(1-x_n/2)]^{10}(x_n/x)^2, & x_n\leq x\leq 2x_n\\ [(1-x/2)/(1-x_n/2)]^{10}(x_n/x)^2[3/2(2x_n/x)^2(1-2x_n/(3x)], & x\geq 2x_n\end{cases}$$

where the gluon distribution is assumed to take the form suggested by the counting rule: $G(x)\sim x^{-1}(1-x)^5$. Here the modification of shadowing-antishadowing on the gluon distribution function $G(x)$, strictly speaking, should be taken into account. But most of the modification will be canceled by the ratio $G(x/2)/G(x_n/2)$.

The momentum balance for the parton-quark system demands that the distribution (4.1) has to obey the following condition:

$$\alpha/(\beta+1)-\alpha(1-x_n)^{\beta+1}/(\beta+1)+0.056[\alpha(1-x_n)^\beta-F^x_A(x_n)]-\int_0^{x_n}dx\,F^s_A(x)+\tfrac{5}{18}n_s\gamma=0\ , \tag{4.2}$$

where $n_s\gamma$ is the momentum fraction of sea quarks lost by the interactions of shadowed partons, predominantly due to the annihilations of sea quarks and antiquarks. The loss momentum is proportional to the effective nucleon number n_s. The value of β is able to be determined by the condition (4.2) if the value of γ is given.

In Mueller-Qiu's gluon recombination model[37,38] the shadowing among the sea quarks is neglected and shadowing is approximately considered to be entirely due to the recombinations among the gluons. The reason for this is that sea quark density is smaller than the gluon density in a nucleon. The parton momentum in that approach has no momentum transference between quarks and gluons; in other words, the gluon momentum and the quark momentum should be separately balanced. The following estimation, however, shows that it is more reasonable to include the annihilation of overlapping sea quarks in the parton-shadowing process. So a small part of the momentum due to the annihilation of overlapping sea quarks in the shadowing process is transferred into the gluons. It is sufficient for this modification to affect the ratio value R_μ within the range $0.1<x<0.2$, although it is not obvious for it to change the ratio R_μ of the structure functions except for $0.1<x<0.2$.

The value for the parameter γ in Eq. (4.2) can be estimated. Suppose that shadowing is primarily caused by the annihilations of the overlapping gluons and of the overlapping sea quarks. For that process the average momentum fraction $\langle \Delta x \rangle$ of the annihilation of overlapping partons can be estimated by Qiu's shadowing factor:[38]

$$\langle \Delta x \rangle = \int_0^{x_n} [xG(x)+\tfrac{18}{5}F_A^s(x)][1-R_s(x,A)]dx \ ,$$

where the factor R_g for the gluon shadowing is equal to that R_s for the sea quark shadowing.[38] In addition, since the shadowing probability is proportional to the square of the parton density, the event ratio

$$\sim \{G(x)/[\tfrac{18}{5}x^{-1}F_A^s(x)]\}^2$$

of the gluon recombination to the sea quark annihilation for $x<x_n$ is around 5:1–10:1. Here $G(x)$ is taken to be $2.8(1-x)^5/x$.[32] For iron it is really meant to be about 0.0034–0.005 of the momentum fraction loss for the sea quarks. The lost momentum is transferred to the gluons. It corresponds to $\gamma=0.0012$–0.0018 in Eq. (4.2). The comparison for $\gamma=0$, 0.001 27, and 0.0025 is given in the calculations.

The following edge effect has to be taken into account if calculating the effective number n_s for the shadowed nucleons in light nuclei. The geometric curvature of a nuclear surface becomes larger and larger with its mass number decreasing. At that time there is no shadowed nucleon in front of a struck nucleon lying in the convex region of the nuclear surface. Such an edge effect reduces the effective number of the shadowed nucleons in a light nucleus. Hence we cannot make use of $n_s=A^{1/3}-1$ for $A<8$ if calculating the effective number n_s of the shadowed nucleons in the shadowing factor $R_s(0,A)$. The values of that effective number for light nuclei may be estimated by a naive geometric consideration, as it stands, $n_s=0.3$ for helium and $n_s=0.5$ for lithium.

V. COMPARISON BETWEEN RESULTS AND EXPERIMENTS

The ratio of the structure function of a bound nucleon to that of a free nucleon in the muon-induced DIS is

$$R_\mu(x,Q^2)=[F_A^v(x,Q^2)+F_A^s(x,Q^2)]/[F_N^v(x,Q^2)+F_N^s(x,Q^2)] \ . \quad (5.1)$$

The original data measured by the European Muon Collaboration,[1] showed the following features: (1) a stronger enhancement ($\sim 15\%$) of the ratio $R_\mu(x)$ above unity for $x \lesssim 0.3$, (2) a decrease in the ratio $R_\mu(x)$ below

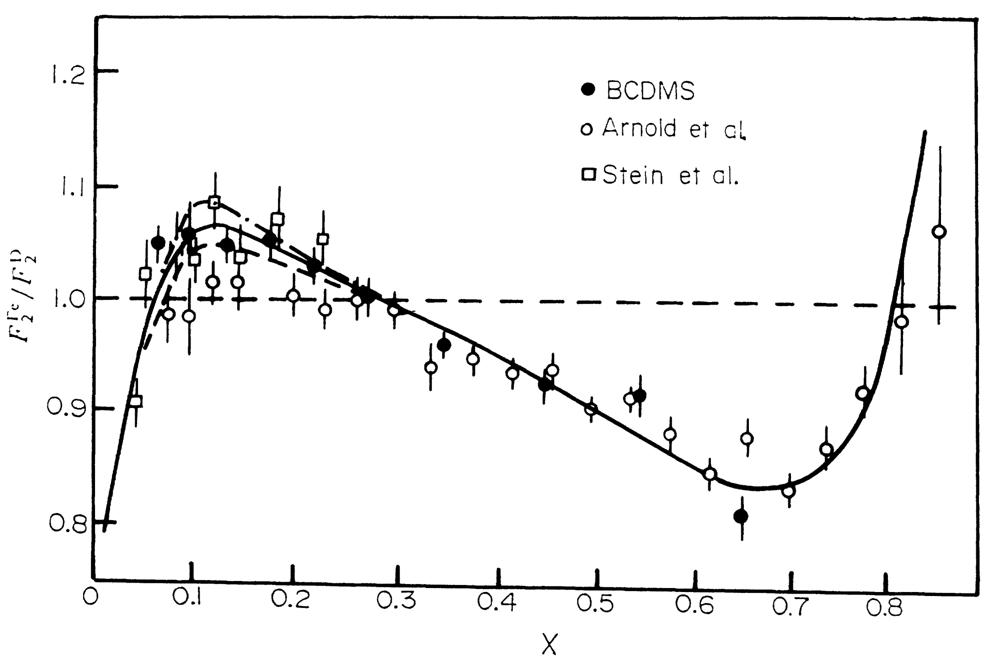

FIG. 2. The comparison of the BDCMS and SLAC data with the ratio of the structure function for iron to that for a free nucleon in the CQM. In Eq. (4.2), $\gamma=0$ (dashed-dotted line), $\gamma=0.001\,27$ (solid line), and $\gamma=0.0025$ (dashed line). The data are taken from Refs. 20, 21, and 22.

TABLE I. The values of the ratio ω_A/ω_N for two oscillator constants, the swelling coefficient R_A/R_N of nucleons, the distortion factor δ_A in the distribution (3.3), the power β in the distribution (4.1), and the first and the second crossover points x_{c1}, x_{c2} of the cross-section ratio $R_\mu(x)$ with the horizontal axis $R_\mu(x)=1$ for a series of nuclei with $4 \le A \le 200$. The input parameter is $dk_0=0.32k_N$.

A	4	6	12	20	27	40	56	100	120	200
ω_A/ω_N	0.965	0.939	0.902	0.897	0.893	0.888	0.883	0.876	0.873	0.868
R_A/R_N	1.018	1.032	1.053	1.056	1.058	1.061	1.064	1.068	1.070	1.073
δ_A	0.05	0.09	0.16	0.17	0.18	0.19	0.20	0.216	0.22	0.24
β	8.37	8.25	6.88	6.26	5.73	5.15	4.28	3.25	2.77	1.19
x_{c1}	0.018	0.030	0.051	0.056	0.058	0.063	0.065	0.072	0.075	0.079
x_{c2}	0.110	0.24	0.25	0.27	0.27	0.28	0.29	0.30	0.30	0.31

unity for $0.3 \lesssim x \lesssim 0.7$, and (3) a substantial independence of the ratio $R_\mu(x)$ on the four-dimensional transferred momentum square of a virtual photon.

The more recent data[23,24] measured by them, however, showed different features in the small x region. (1) The ratio in the region $0.1 \lesssim x \lesssim 0.25-0.3$ rises above unity and reaches a maximum enhancement of about 0.03–0.1 for those nuclei with $A=12-120$. (2) The ratio falls below unity at $x \lesssim 0.1$ for $Q^2 \gtrsim 1$ $(\mathrm{GeV}/c)^2$ and the first crossover point of the ratio $R_\mu(x)$ with the horizontal axis $R_\mu(x)=1$, evidently runs toward the x-increasing direction with the mass number A of nuclear targets rising. (3) The ratio has an A dependence in the small region, but no significant Q^2 dependence in that region. (4) There is no obvious change in the net momentum fraction carried by gluons and quarks within the experimental errors.

The relative reduced value of the spring elastic coefficient for a nucleon bound inside the nuclear central region is used, $dk_0=0.32k_N$, as input for the procedure and an explanation about the input is referred to that in the Appendix. The input means that $\omega_A/\omega_N=0.83$ or $R_A/R_N=1.10$ for a nucleon is bound inside the interior of a nucleus. For a series of nuclei, the ratios ω_A/ω_N of two spring elastic constants, calculated by the CQM, are listed in the first row of Table I. Meanwhile the nuclear distortion factor δ_A in the distribution (3.3) are also given in the third row of Table I.

The cross-section ratio for iron is presented in Fig. 2. The dashed-dotted line and the dashed line in Fig. 2 correspond, respectively, to the case with the parameter $\gamma=0$ and 0.0025. The parameter $\gamma=0.001\,27$, as shown by the solid line in Fig. 2, is chosen on the basis of the BCDMS data[22] and corresponds to that case in which the momentum fraction, lost by the sea quarks in iron owing to shadowing, is 0.0036. With the aid of Eq. (4.2), the power β in the distribution (4.1) can be determined and its values for various nuclei are shown in the fourth row of Table I. In addition the Fermi momentum k_F^A and the binding energy B_A concerned in the calculation are taken as the values required by conventional nuclear physics, and their values for a variety of nuclei are listed in Table II.

The ratios of the structure functions for some nuclei, carbon, calcium, and tin, to that for a free nucleon, are drawn in Fig. 3 in which the data are taken from Ref. 24. In a previous work[11] we neglected nuclear shadowing, and so the ratio for $x<0.3$ is nearly close to unity. It can be seen from Fig. 3 that it is necessary to take the

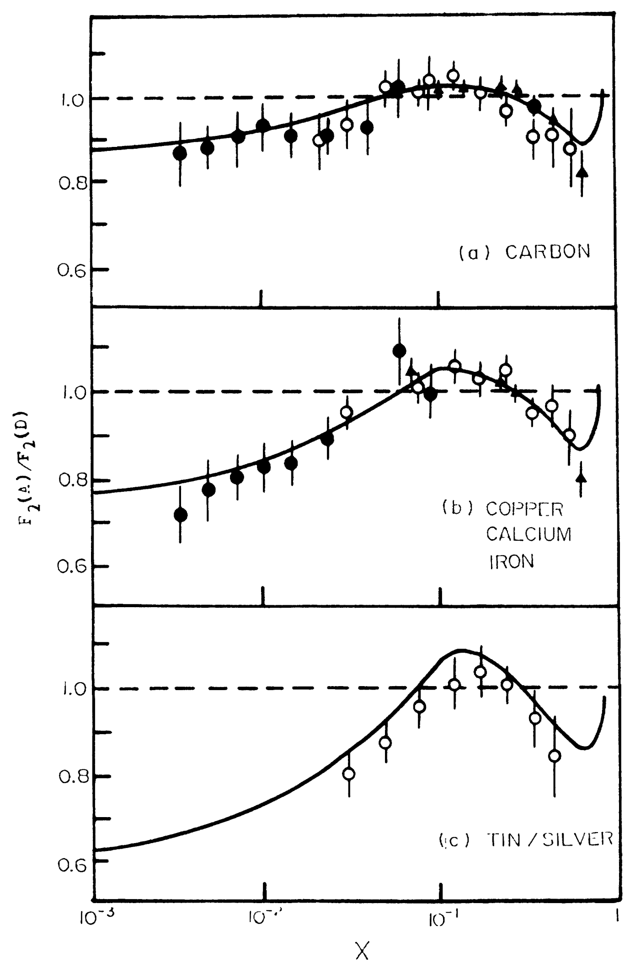

FIG. 3. The prediction for the ratio of the nucleon structure function in a nucleus to that in a deuterium (a) carbon, (b) iron-copper-calcium, and (c) tin-silver. The data are taken from Ref. 24.

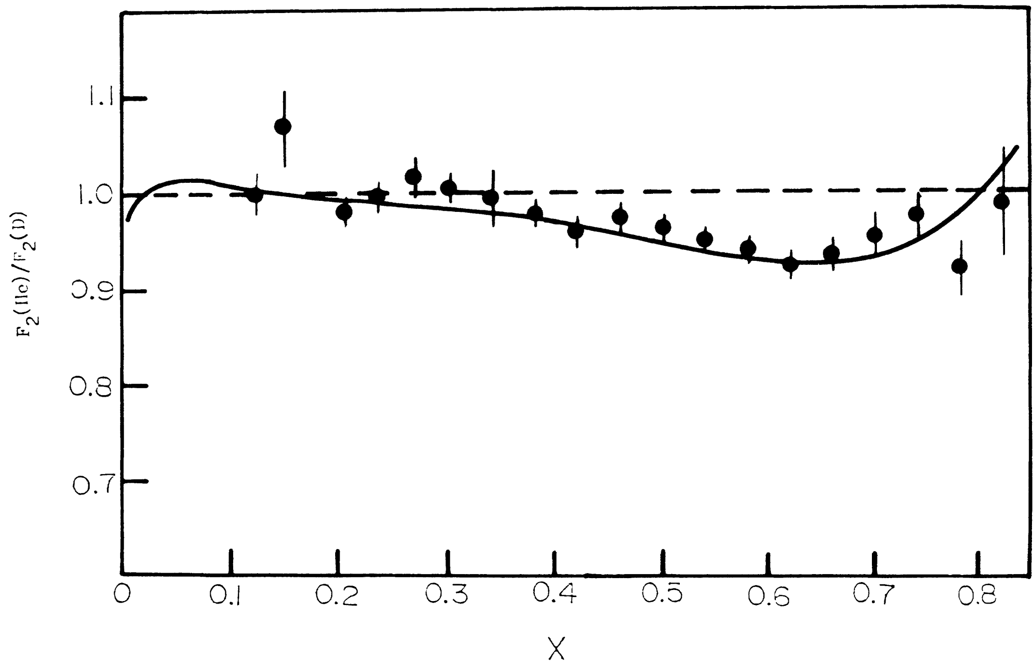

FIG. 4. The prediction for the ratio of the structure function for helium to that for a free nucleon. The data are taken from Ref. 20.

shadowing and antishadowing into account. Based on this consideration the behaviors of the ratios for $x \lesssim 0.3$ are evidently dependent on the mass number A. Furthermore the A dependence of the ratios R_μ for $x > 0.3$ are caused by the nuclear surface effect. Figure 4 shows the ratio of a light nucleus He to deuterium. A better approach reducing the systematic errors is that in which the data on two different nuclear targets are simultaneously

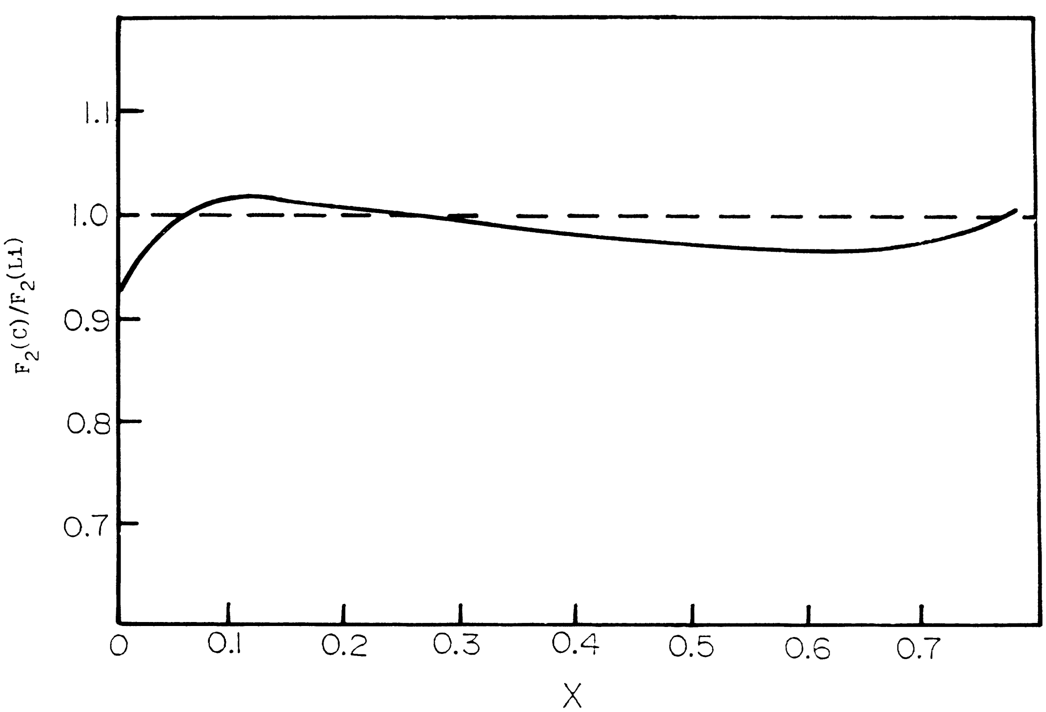

FIG. 5. The prediction for the ratio of the structure function for carbon to that for lithium.

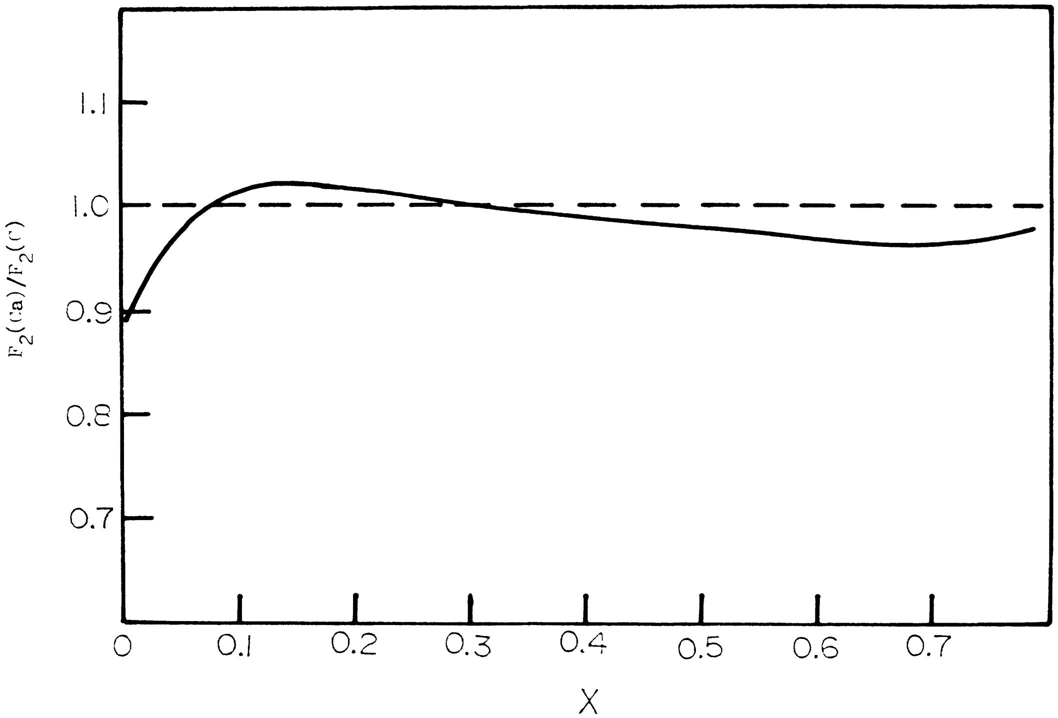

FIG. 6. The same as Fig. 5, but for the ratio of that for calcium to that for carbon.

measured in an experiment. The ratios F_c^C/F_2^{Li} and F_2^{Ca}/F_2^C are calculated and given in Figs. 5 and 6, respectively. The A dependence of the ratios at a fixed x is drawn in Fig. 7.

Recent experiments[23,24] show that the first crossover point x_{c1} of the ratio $R_\mu = F_2^A/F_2^D$ with the horizontal axis $R_\mu = 1$ runs evidently towards the x-increasing direction with A rising. The phenomenon is contrary to the estimation given by Nikolaev and Zakharov.[36] Qiu conceived[38] of the independence of the point x_n on the mass number A. Here x_n is approximately corresponding to the crossover point x_{c1}. Later Berger and Qiu[40] further considered the nuclear surface effect. They found that the point x_n varies in its value from 0.112 to 0.100 for a nuclear target that varies from tin to carbon. The A dependence of the point x_n is so small that it is insufficient to account for the phenomenon observed in the experiments: The value of the crossover point x_{c1} distinctly decreases when a nuclear target varies from tin to carbon. Close and Roberts[41] suggested that the point x_A, corresponding to the supposed saturation of shadowing, is not strongly dependent on the mass number A; at least, such is the case for $A \gtrsim 12$ such that the point x_n mentioned above evidently depends on A. Another consistent interpretation about the dependence of the crossover point x_{c1} on A is shown in our result. It is just the coexistence of the antishadowing in part in the shadowing region ($x \leq x_n$) so that the cross-section ratio $R_\mu(x)$ rises to unity for $x < x_n$. As a consequence, the lighter a nuclear target is, the weaker its shadowing, and the more rapidly the cross-section ratio $R_\mu(x)$ rises: that is, the lighter a nuclear target is, the larger Qiu's shadowing factor $R_s(x=0, A)$ is and the smaller the value of the crossover point x_{c1}. We find that the crossover point varies in the value from 0.018 to 0.079 as the mass number A varies from 4 to 200. This is quite consistent with the experimental observation. The reason is primarily due to the observation that after considering the modification of the shadowing and antishadowing, the distribution function of sea quarks evidently becomes harder and harder with A increasing, see the fourth row of Table I.

Let us discuss the neutrino- (antineutrino-) induced deep-inelastic scattering on nuclear targets. Because of the parity violation the process is used to extract separating terms from its differential cross section. The antineutrino cross section is

TABLE II. The Fermi momentum k_F and the binding energy for a series of nuclei with $4 \leq A \leq 200$ used in the calculations.

A	4	6	12	20	27	40	56	100	120	200
k_F (fm)$^{-1}$	0.93	0.86	1.12	1.18	1.20	1.27	1.28	1.30	1.32	1.34
B_A (MeV)	7.0	5.3	7.6	8.0	8.6	8.6	8.6	8.6	8.5	8.0

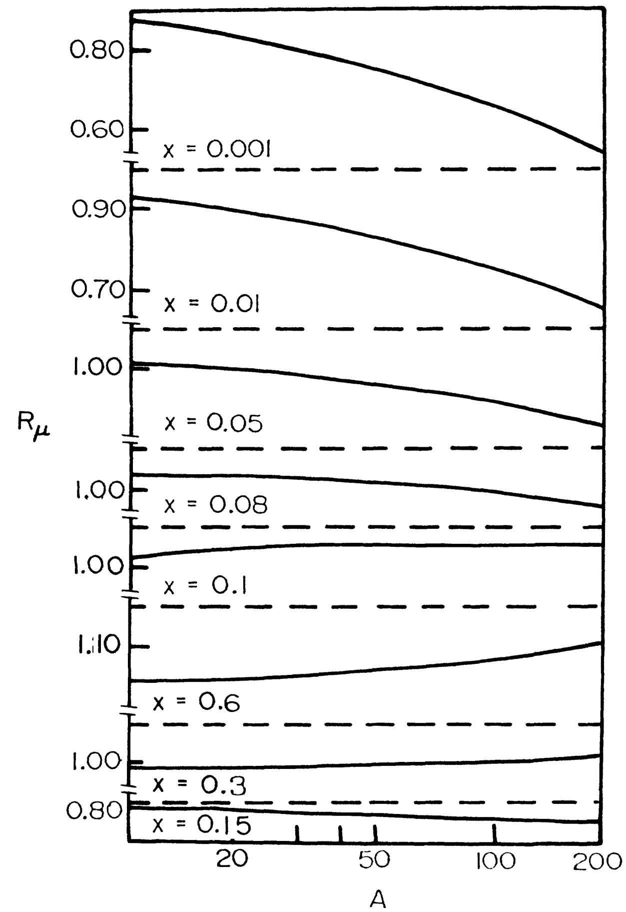

FIG. 7. The A dependence of the cross section at a fixed x in the CQM.

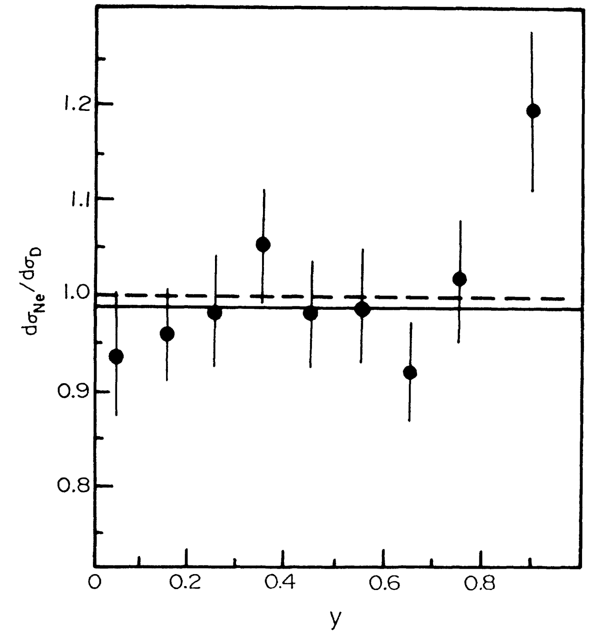

FIG. 9. The same as Fig. 8, but vs y.

$$d\sigma^{\bar{\nu}}/dx\,dy = (G^2 M E_{\bar{\nu}}/\pi)x\,[q(x)(1-y)^2+\bar{q}(x)] \ . \quad (5.2)$$

The predictions of the CQM for the ratio of the antineutrino cross section on a neon target to that on a deuterium target are shown in Fig. 8 with the data coming from Ref. 42. For the sake of comparison, the curves with and without the shadowing and antishadowing predicted by the CQM are simultaneously drawn in Fig. 8.

The content of antiquarks in a bound nucleon can also be seen from the y dependence of the neutrino or antineutrino cross section extracted from experiments:

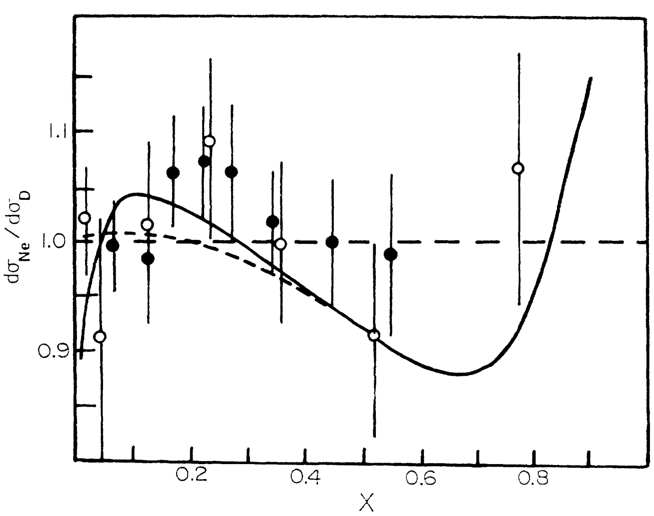

FIG. 8. The antineutrino-induced cross-section ratio (neon/deuterium) as a function of x in the CQM. The data are taken from Ref. 42.

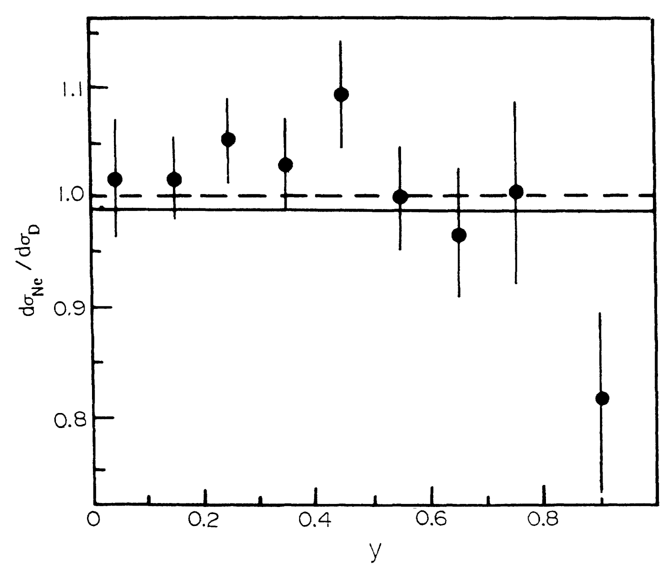

FIG. 10. The same as Fig. 9, but for the neutrino-induced cross-section ratio.

$$\sigma_y(\bar{\nu}N)=\sigma_0[\langle q\rangle_2(1-y)^2+\langle\bar{q}\rangle_2]\ ,$$
$$\sigma_y(\nu N)=\sigma_0[\langle q\rangle_2+(1-y)^2\langle\bar{q}\rangle_2] \tag{5.3}$$

in which $\langle q\rangle_2$ and $\langle\bar{q}\rangle_2$ are the average nucleon momenta carried by a quark and an antiquark, respectively. Taking the nuclear binding effect into account and neglecting a small contribution from the annihilation of overlapping sea quarks, the average momentum carried by each species in a bound nucleon will be reduced by a factor of η_A in which $\eta_{\rm He}=0.991$. Hence, the ratios

$$R_y^{\bar{\nu}}=\sigma_y(\bar{\nu}{\rm Ne})/\sigma_y(\bar{\nu}D)$$

and

$$R_y^{\nu}=\sigma_y(\nu{\rm Ne})/\sigma_y(\nu D)$$

are attained and are drawn in Figs. 9 and 10. The results are compatible with the existing data.[42]

VI. CONCLUSION

Within the framework of the CQM we have studied in detail the distortion behaviors of the nucleon structure function in the nuclear medium. As far as its physics essence is concerned, this is a concise model. It is supposed that each bound nucleon consists of three constituent quarks and each constituent quark is composed of partons.

There are two crucial points that give rise to the difference between the structure function for a bound nucleon and that for a free nucleon. The nuclear medium makes the oscillator constant of the three consistent quarks inside a nucleon weaken so that the distribution of constituent quarks is softened. At the same time, under the DIS condition, the gluons of the three constituent quarks recombine with those in other bound nucleons so that the sea quark distribution is hardened.

The first point, i.e., the weakening of the oscillator constants in the nuclear medium, results in the experimental observation that the cross-section ratio $R_\mu(x)$ in the region $0.3\lesssim x\lesssim 0.7$ is lower than unity and that this depression is enhanced with the mass number A increasing. The second point, i.e., Mueller-Qiu's gluon recombination mechanism which is improved by taking the coexistence of nuclear shadowing and antishadowing into account and by obeying the momentum conservation of the parton-quark system, results in another experimental observation that the ratio $R_\mu(x)$ in the region $0\lesssim x\lesssim 0.3$ gradually rises from the value below unity to that above unity as x increases. Especially, the improved gluon recombination predicts that there exist two cross over points x_{c1} and x_{c2} of the cross-section ratio $R_\mu(x)$ with the horizontal axis $R_\mu(x)=1$, which run towards the x-increasing direction as the mass number A rises. Here it should be stressed to be the universality about the structure of the constituent quarks themselves. In contrast with other models, it is just the universality that leads to the low-x behavior of the cross-section ratio $R_\mu(x)$ without considering nuclear shadowing, i.e., it tends to unity as x approaches zero.[10,11] Except for those two points, of course, the behavior of the ratio $R_\mu(x)$ at high x ($x\gtrsim 0.7$) can be ascribed to the smearing over the distribution of the constituent quark, which is caused by the Fermi motion among the bound nucleons.

It is meaningful at last to remark on the momentum conservation of the parton-quark system. The distribution of partons in a nucleon is a convolution of the distribution of the partons inside a constituent quark and the distribution of the constituent quarks inside a nucleon in the CQM. Therefore the average nucleon momentum fraction carried by the partons is the product of the average nucleon momentum fraction carried by the constituent quarks and average momentum fraction of the constituent quarks carried by the partons. As in the case of a free nucleon, one-third of the average nucleon momentum fraction is carried by each constituent quark in a bound nucleon in which the binding energy is neglected. As just mentioned, if there is no shadowing, the parton distribution in a constituent quark is universal. As a result the average momentum carried by each parton species in a nucleon is a conserved quantity. According to Mueller-Qiu's mechanism, the gluon evolution governs the shadowing of gluons and that of sea quarks. The gluon-gluon interaction $GG\to G$ does not change the average momentum of gluons and also does not result in the momentum transfer between gluons and sea quarks or gluons and valence quarks. In this approximation, the average nucleon momentum fraction carried by each parton species (valence quarks, sea quark-antiquark pairs, and gluons) in a nucleon is preserved respectively. This is one of the characteristics of the CQM. Considering nuclear shadowing and omitting the recombination of sea quarks,[37,38] the momentum conservation of the CQM still remains as before. Only on account of a weak nuclear binding effect do we have

$$I=\tfrac{18}{5}\int_0^1 dx[F_2^{\rm Fe}(x)-F_2^{D}(x)]=-(1-\eta)/2=-0.005$$

if it is assumed that the quark momentum is one-half the nucleon momentum. Of course, the I value would be smaller than that mentioned above if the contribution of the annihilation of the overlapping sea quarks is considered. In this case, we have $I=-0.0086$.

The CQM successfully predicts the lepton-induced DIS cross-section ratio for a variety of nuclear targets with $4\le A\le 200$ within the kinetic region $0.01<x<0.9$ and the predictions are in fair agreement with the existing data.

ACKNOWLEDGMENTS

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APPENDIX

We attempt to analyze the reason for the weakening of spring tension between the constituent quarks in nuclear surrounding. The harmonic oscillator model[27,28] is a phenomenological description for the interactions among the constituent quarks. One knows that this kind of interaction originates from the fundamental interactions of QCD. So one should look for the answer for the change of spring tension from QCD. The QCD interactions for the bound state involve in nonperturbative QCD. A way of avoiding that difficult is to ascribe the nonperturbative effect to the QCD vacuum.[43–46] A constituent quark in the CQM is a cluster consisting of a valence quark with its surrounding gluons and sea quarks (even mesonic cloud). Such a cluster has a space scale. If the space for the constituent quarks might be considered to be a nontrivial QCD vacuum, the QCD vacuum for the three constituent quarks constructing a color singlet forms the so-called baglike state [Fig. 11(a)]. As shown in Ref. 43, such a bag vacuum could be considered to be a superconductor consisting of the color-magnetic field and expelling the color-electric field. Relative to a normal vacuum, the QCD vacuum possesses the so-called vacuum energy density C. Its contribution to the nucleon Hamiltonian, i.e., the volume energy, is an effective statement for the nonperturbative QCD interaction among the constituent quarks. It is remarked here that the baglike state composed of the constituent quarks differs from the MIT bag. For the former, the constituent quarks extending to a space is introduced, and the configuration space occupied by the three constituent quarks is equivalent to the bag space. Hence the oscillation of the constituent quarks around their equilibrium positions could be considered as the bag oscillation. Let us estimate the restoring force for the bag oscillation. For simplicity, one takes into account two limit cases. One of them is the oscillation of a sphere-symmetric bag [Fig. 11(b)]. Omitting a small perturbative quark interaction the restoring force for that oscillation is $F \sim \partial E/\partial R \sim CR^2$ with R the bag scale. Another of them is the case in which a constituent quark is far from a pair of constituent quarks. In that case the bag has a larger deformation to form a stringlike bag [Fig. 11(c)]. The stringlike bag was discussed by a number of authors. It turns out to be the bag restoring force $F \sim \partial E/\partial L \sim C^{1/2}L^0$ with L the string scale. It is presumed that the bag oscillation such as Fig. 11(a), is an intermediary between Figs. 11(b) and (c), and in that case the bag restoring force becomes $F \sim C^{3/4}R$ where the power of C comes from the dimensional requirement. From this the spring coefficient is $K \sim C^{3/4}$. It is shown, as analyzed previously, that the deviation G_N^c for low y originates from the large scale correlation of quarks in the small-y region, i.e., a large deformation of the nucleon. The harmonic oscillator model in that case is no longer a reasonable phenomenological description.

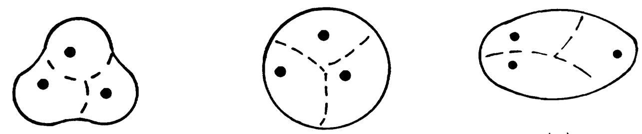

FIG. 11. The nucleon baglike state (a) and its two limit cases: (b) the sphere-symmetric state, (c) the stringlike state.

Now we return to the issue of what happens as two nucleons in a nucleus are close to each other. It is interesting, as proposed by Noble[47] that there exists a possibility that the overlapping of two nucleons is responsible for the weakening of the vacuum constant C in the overlapping volume and $C \rightarrow C' \equiv C - C\delta\Omega/\delta V_N$ where $\delta\Omega$ is the overlapping volume and V_N the volume of the QCD vacuum occupied by the nucleons. In an infinite nuclear matter without taking into account the edge effect, as estimated by Ref. 47, $\delta\Omega/\delta V_N = (R_N/R_0)^3/2$ where the R_N is the nucleonic "glue" radius and R_0 is the Wigner-Seitz radius around 1.2 fm. $C'/C \sim 0.74-0.50$ if R_N takes 0.7 fm–1.12 fm and in that case the spring coefficient is correspondingly weakened, $K_A/K_N = 0.8-0.6$, i.e., $dk_0 = 0.2k_N - 0.4k_N$. The parameter value $dk_0 = 0.32k_N$ we take is acceptable.

[1]Europoean Muon Collaboration, J. Aubert *et al.*, Phys. Lett. **123B**, 275 (1983).
[2]C. H. Lewellyn Smith, Phys. Lett. **128B**, 107 (1983).
[3]M. Ericson and A. M. Thomas, Phys. Lett. **128B**, 112 (1983).
[4]E. L. Berger, F. Coester, and R. B. Wiringa, Phys. Rev. D **29**, 398 (1983); E. L. Berger and F. Coester, Phys. Rev. D **32**, 1071 (1985).
[5]S. V. Akulinichev, S. A. Kulaging, and G. M. Vagradov, Phys. Lett. **158B**, 485 (1985).
[6]F. E. Close, R. G. Roberts, and G. G. Ross, Phys. Lett. **129B**, 346 (1983).
[7]C. E. Carlson and T. J. Havens, Phys. Rev. Lett. **51**, 261 (1983).
[8]R. L. Jaffe, Phys. Rev. Lett. **50**, 228 (1983).
[9]R. L. Jaffe, F. E. Close, R. G. Roberts, and G. G. Ross, Phys. Lett. **134B**, 449 (1984).
[10]C. H. Chang and W. Zhu, Phys. Lett. B **187**, 405 (1987).
[11]W. Zhu and J. G. Shen, Phys. Lett. B **219**, 107 (1989).
[12]R. C. Hwa, Phys. Rev. D **22**, 1593 (1980).
[13]R. C. Hwa and M. S. Zahir, Phys. Rev. D **23**, 2593 (1981).
[14]E. V. Shuryak, CERN Report CERN 83-01, 1983.
[15]C. G. Callan, R. Dashen, and D. J. Gross, Phys. Rev. D **17**, 2717 (1978).
[16]E. V. Shuryak, Nucl. Phys. **B203**, 93 (1982); **B203**, 116 (1982); **B203**, 140 (1982).
[17]J. D. Sullivan, Phys. Rev. D **5**, 1732 (1972).
[18]G. Altarelli and G. Parisi, Nucl. Phys. **B126**, 298 (1977).
[19]A. Bodek *et al.*, Phys. Rev. Lett. **50**, 1431 (1983).
[20]A. Bodek *et al.*, Phys. Rev. Lett. **51**, 534 (1983).
[21]R. G. Arnold *et al.*, Phys. Rev. Lett. **52**, 727 (1984).
[22]Bologna-CERN-Dubna-Munich-Sacley (BCDMS) Collaboration, A. C. Benvenuti *et al.*, Phys. Lett. B **189**, 483 (1987).
[23]European Muon Collaboration, J. Ashman *et al.*, Phys. Lett. B **202**, 603 (1988).
[24]European Muon Collaboration, J. Ashman *et al.*, Phys. Lett. B **211**, 493 (1988).
[25]R. C. Hwa and C. S. Lam, Phys. Rev. D **26**, 2338 (1982).
[26]P. E. Hussar, Phys. Rev. D **23**, 2781 (1981).
[27]R. P. Feynman, M. Kisslinger, and F. Ravndal, Phys. Rev. D **3**, 2705 (1971).
[28]Y. S. Kim and M. E. Noz, *Theory and Application of the Poin-*

caré Group (Reidel, Dordrecht, 1986).
[29]F. Martin, Phys. Rev. D **18**, 1382 (1979).
[30]W. Zhu and J. G. Shen, Chin. Phys. **3**, 968 (1983).
[31]D. W. Duke and J. F. Owens, Phys. Rev. D **30**, 49 (1984).
[32]D. von Harrach, Nucl. Phys. **A478**, 290 (1988).
[33]A. Bohr and B. R. Mottelson, *Nuclear Structure* (Benjamin, New York, 1969), Vol. II. p. 160.
[34]C. H. Llewelly Smith, Phys. Lett. **128B**, 107 (1983); M. Ericson and A. W. Thomas, Phys. Lett. **128B**, 112 (1983).
[35]R. P. Feynman, *Photon-Hadron Interactions* (Benjamin, New York, 1972).
[36]N. N. Nikolaev and V. I. Zakharov, Phys. Lett. **55B**, 397 (1975).
[37]A. H. Mueller and J. Qiu, Nucl. Phys. **B268**, 427 (1986).
[38]J. Qiu, Nucl. Phys. **B291**, 746 (1987).
[39]W. Zhu and J. G. Shen, Phys. Lett. B **235**, 170 (1990).
[40]E. L. Berger and J. Qiu, Phys. Lett. B **206**, 141 (1988).
[41]F. E. Close and R. G. Roberts, Phys. Lett. B **213**, 91 (1988).
[42]J. Guy *et al.*, CERN Report CERN/EP 86-217 (1986).
[43]S. Mandelstam, Phys. Rep. **23**, 245 (1976).
[44]C. G. Callasn, R. Dashen, and D. J. Gross, Phys. Rev. D **19**, 1826 (1979).
[45]T. D. Lee, Phys. Rev. D **19**, 1802 (1979).
[46]E. V. Sharyak, Nucl. Phys. **B203**, 93 (1982); **B203**, 116 (1982); **B203**, 140 (1982).
[47]J. V. Nobel, Phys. Lett. B **178**, 285 (1986).