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Polarized EMC effect in the QMC model

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ABSTRACT

The ratios of the in-medium to free nucleon structure functions for the unpolarized and polarized cases are obtained using the MIT bag model for the free case, along with the QMC model to incorporate the in-medium modifications of the structure functions. As discussed in earlier work, the observed nuclear EMC effect is reasonably well described. This gives us confidence to investigate the predictions of the model for the polarized EMC effect, the ratio $g_1^*(x)/g_1(x)$ for a bound proton to that of a free proton. This ratio is found to be substantially different from unity and very similar in shape and size to that found in the unpolarized case. This prediction of such a fundamental change in the valence structure of a bound nucleon needs to be tested experimentally at the earliest opportunity.

1. Introduction

Our understanding of nuclear structure was severely challenged by the unexpected experimental results released by the European Muon Collaboration (EMC) in 1983. Roughly speaking, EMC compared the structure function of a free nucleon to that of a bound nucleon and found there was a significant difference between them [1,2], especially in the valence region. The ratio of the in-medium to free structure function of the nucleon was found to drop below one in that region, indicating a suppression of the bound nucleon's structure function, which became known as the EMC effect [3–5].

The discovery that the nucleon structure functions differ substantially in-medium compared to the free case surprised the nuclear physics community, pointing to a potential new approach to nucleon structure. One such approach is the Quark Meson Coupling (QMC) model [6–8], which explicitly allows the quark degrees of freedom to respond self-consistently to the nuclear mean fields and leads naturally to changes in the internal structure of the bound nucleons. Initial work within the QMC model did in fact reproduce the main features of the EMC effect [9,10] but this work required some finesse to overcome technical difficulties with momentum conservation in the bag model [11].

More recently Cloët, Bentz and Thomas [12,13] overcame the limitations of the bag model by using the covariant NJL model [14] for the structure of the nucleon within a QMC-like framework in which, once again, the internal structure of the nucleon was self-consistently solved in the mean fields generated in nuclear matter [15]. The calculations made in this model were successful in reproducing the unpolarized EMC data across the periodic table. These authors also extended their calculations to make predictions for the spin dependent structure function of odd-A nuclei. There they found a considerably larger change, referred to as the "polarized EMC effect", than that found in the unpolarized case.

Here we aim to investigate the model dependence of the results reported by Cloët et al. by calculating the spin dependent EMC effect in the QMC model, using the MIT bag model [16–19] for the structure of the nucleon. This is especially important as there is yet no consensus on the origin of the EMC effect. The work done within the QMC model and its generalizations to the NJL and chiral quark soliton [20] models, is based upon mean-field theory, with the structure of all of the bound nucleons modified in proportion to the local mean fields.

On the other hand, there has recently been much interest in the suggestion that the EMC effect might only be associated with those nucleons involved in short-range correlations [21–23]. As the contribution of correlated nucleons to the spin dependent structure function of a nucleus is expected to be suppressed, a measurement of this effect promises to be a valuable source of information in separating these two explanations. In parallel with such a mea-

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surement, we need the very best estimates of the polarized EMC effect within any particular model.

2. Calculation of the structure functions

In this work we employ the MIT bag model [16–19] to describe the structure of the free nucleon. It is a simple, yet successful phenomenological model for quark confinement, within which three non-interacting quarks are bound to a spherical region of space, with the boundary condition that the quark vector current normal to the surface vanishes. The quarks being treated as non-interacting is justified by appealing to the idea of asymptotic freedom, and the hard boundary condition is a crude implementation of quark confinement. An attractive advantage of the model is that many calculations can be carried through analytically, giving valuable insight into the origin of the effects of the medium on the quark distributions.

As we already remarked, while the usual formulation of the bag model is convenient for calculating static properties, it is not ideal for the study of deep inelastic scattering where the conservation of energy and momentum is essential to ensuring the correct support of the parton distribution functions (PDFs). For this reason Schreiber et al. [11] chose to perform the calculations of the PDFs in a framework where energy and momentum conservation was imposed [24] before any approximations were made. The formulation of Schreiber et al. builds in these conservation laws and, in addition, uses the Peierls–Yoccoz method [25] to construct approximate momentum eigenstates.

The PDF of a longitudinally polarized quark of flavor f inside of a longitudinally polarized nucleon of mass M can be calculated by evaluating:

$$q_f^{\uparrow\downarrow}(x) = \sum_n \delta\left(1 - x - \frac{p_n^+}{P^+}\right) \left|\langle n|\Psi_{+,f}|P,s\rangle\right|^2, \tag{1}$$

where x is the light-cone $(+)$[1] momentum fraction of the nucleon carried by the quark, which is described by wave function Ψ_f, while P and s describe the momentum and the spin of the nucleon. The notation $\uparrow\downarrow$ indicates aligned or anti-aligned helicities of the quark and the nucleon and the polarized PDF can be expressed (e.g., [26]) in terms of the conventional unpolarized, $f(x)$, and the helicity dependent, $\Delta f(x)$, distributions as $q^{\uparrow\downarrow}(x) = 1/2(f(x) \pm \Delta f(x))$. The sum is over all possible intermediate states n with momenta p_n when probing the quark in a deep inelastic scattering process. We consider only diquark intermediate states, that provide the dominant contribution to the quark PDF [11], which yields the following expression

$$q_f^{\uparrow\downarrow}(x) = \frac{M}{(2\pi)^2} \sum_m \langle\mu|P_{f,m}|\mu\rangle \times \int_{\frac{\left|M^2(1-x)^2 - M_n^2\right|}{2M(1-x)}}^{\infty} p_n dp_n \frac{|\phi_2(\mathbf{p}_n)|^2}{|\phi_3(\mathbf{0})|^2} |\tilde{\Psi}_m^{\uparrow\downarrow}(\mathbf{p}_n)|^2. \tag{2}$$

The spin-flavor wave function of the initial nucleon at rest is denoted by $|\mu\rangle$, and $P_{f,m}$ is a projector operator onto quark of flavor f and spin m. The functions ϕ_3 and ϕ_2 are the normalizations of the 3-quark and diquark wave functions, and M_n is the mass of the intermediate state.

In the zero mass quark case the MIT bag wave function for a spatial coordinate $\boldsymbol{r}$ takes the form [17,19]

$$\Psi_m(\boldsymbol{r}) = N \begin{pmatrix} j_0\left(\frac{\Omega|\boldsymbol{r}|}{R}\right)\chi_m \\ i\boldsymbol{\sigma}\cdot\hat{\boldsymbol{r}}\, j_1\left(\frac{\Omega|\boldsymbol{r}|}{R}\right)\chi_m \end{pmatrix} \Theta(R - |\boldsymbol{r}|), \tag{3}$$

with lowest energy eigenfrequency solution $\Omega \simeq 2.04$. Here R is the bag radius, j_0 and j_1 are the spherical Bessel functions of the first kind, χ_m are spinors, and $\boldsymbol{\sigma}$ are the Pauli spin matrices. The normalization of the wave function is given by

$$N^2 = \frac{1}{4\pi} \frac{\Omega^3}{2R^3(\Omega - 1)\sin^2(\Omega)}. \tag{4}$$

Using the Peierls–Yoccoz method to obtain the approximate eigenstates of momentum, we find

$$|\phi_2(\mathbf{p}_n)|^2 = \frac{4\pi R}{u} \left(\frac{2\pi N^2 R^4}{\Omega^4}\right)^2 \int \frac{dv}{v} \sin\left(\frac{2vu}{\Omega}\right) T^2(v) \tag{5}$$

and

$$|\phi_3(\mathbf{0})|^2 = 4\pi \left(\frac{2\pi N^2 R^4}{\Omega^4}\right)^3 \int \frac{dv}{v} T^3(v), \tag{6}$$

where $T(v)$ is the overlap function for quarks in displaced bags, and the following substitutions have also been made

$$v = \frac{|\boldsymbol{r}|\Omega}{2R}, \quad u = |\mathbf{p}_n|R. \tag{7}$$

The overlap function can be evaluated using the bag wave function, yielding

$$T(v) = T_t(v) + T_b(v), \tag{8}$$

$$T_t(v) = \Big[(\Omega - v)\sin(2v) + (1 - \sin^2(\Omega)) - \cos(\Omega)\cos(\Omega - 2v)\Big], \tag{9}$$

$$T_b(v) = \Bigg[\left(1 - \frac{4v^2}{2\Omega^2}\right)\sin^2(\Omega) - \frac{2}{\Omega}\sin(\Omega)\cos(\Omega - 2v) + \frac{2}{\Omega}\sin(\Omega)\cos(\Omega) + \Omega\sin(2v) - \sin(\Omega)\sin(\Omega - 2v) - v\sin(2v)\Bigg], \tag{10}$$

where $T_t(v)$ corresponds to the overlap integral of the upper component of the bag wave function and $T_b(v)$ to the lower component.

The Fourier transform of Ψ is given as

$$|\tilde{\Psi}_m^{\uparrow\downarrow}(\mathbf{p}_n)|^2 = \frac{1}{2}\Big[f(\mathbf{p}_n) \pm (-1)^{m+3/2} g(\mathbf{p}_n)\Big], \tag{11}$$

where

$$f(\mathbf{p}_n) = \frac{\pi R^3}{2} \frac{\Omega^3}{\left(\Omega^2 - \sin^2(\Omega)\right)} \times \left[s_1^2(u) + 2\frac{p_n^z}{|\mathbf{p}_n|} s_1(u)s_2(u) + s_2^2(u)\right], \tag{12}$$

[1] The + component of a 4-vector a here is defined as $a^+ = a^0 + a^3$.