Light Front Nuclear Theory and the HERMES Effect

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Abstract

I discuss applications of relativistic light front dynamics (the use of light cone variables) to computing the nucleonic and mesonic components of nuclear wave functions. Our method is to use a Lagrangian and its associated energy-momentum tensor $T^{\mu\nu}$ to define the total momentum operators $P^{\mu}$, with $P^+$ as the plus-momentum and $P^-$ the light cone time development operator. The aim of this unusual approach to nuclear physics is the desire to use wave functions, expressed in terms of plus-momentum variables, which are used to analyze high energy experiments such as deep inelastic scattering, Drell-Yan production, $({e,e'})$ and $(p,p')$ reactions. We discuss or mention the topics: infinite nuclear matter within the mean field approximation; finite nuclei using the mean field approximation; nucleon-nucleon scattering, within the one boson exchange approximation; and, infinite nuclear matter including the effects of two-nucleon correlations. Standard good results for nuclear saturation properties are obtained, with a possible improvement in the lowered value, 180 MeV, of the computed nuclear compressibility. In our approach, manifest rotational invariance can’t be used to simplify calculations. But for each of the examples reviewed here, manifest rotational invariance emerges at the end of the calculation. Thus nuclear physics can be done in a manner in which modern nuclear dynamics is respected, boost invariance in the $z$-direction is preserved, and in which rotational invariance is maintained. A salient feature is that $\omega, \sigma$ and $\pi$ mesons are obtained from the nuclear structure theory as important constituents of nuclei. I then argue that these constituents can contribute coherently to enhance the electroproduction cross section on nuclei for longitudinal virtual photons at low $Q^2$ while depleting the cross section for transverse photons. Thus the recent HERMES inelastic lepton-nucleus scattering data at low $Q^2$ and small $x$ can be described using photon-meson and meson-nucleus couplings which are consistent with (but not determined by) constraints obtained from meson decay widths, nuclear structure, deep inelastic scattering, and lepton pair production data. Our model makes a variety of predictions testable in future experiments.
I. INTRODUCTION

This talk consists of two parts. The first is concerned with the subject of light front nuclear theory \[1\] - \[4\], and the second with the application \[5\] of that theory to an effect discovered in high energy lepton-nucleus scattering by the HERMES collaboration \[6\].

In this work, we will consider the regime in which the nucleus treated as being made of nucleons and mesons. Our goal is to get the ground state wave function at zero temperature in terms of light front coordinates. Using these coordinates to evaluate the consequences of a given Lagrangian is also called light front dynamics. We shall begin by giving a few more details concerned with answering the following questions. What is light front nuclear theory? Why do it?

We shall present examples and results of three studies: infinite nuclear matter treated using the mean field approximation; finite nuclei also treated with that approximation; and, infinite nuclear matter including correlations between two nucleons (light front Brueckner Theory). In each application we find that vector and scalar mesons are prominent components of nuclear wave functions.

The second part is concerned with searching for experimental consequences of these components. Thus we discuss the HERMES effect \[6\] as a signature of these nuclear mesons. This effect is observed in the interaction of 27.5 GeV positrons with nuclei, and arises because virtual photons have longitudinal \(L\) polarization as well as the usual transverse \(T\) polarization. The HERMES collaboration measures a ratio of cross sections for the scattering of photons with different polarizations. They find \[6\]:

\[
\frac{(\sigma_{L}(A))}{(\sigma_{L}(D))} \div \frac{(\sigma_{T}(A))}{(\sigma_{T}(D))} \approx 5, \quad (1.1)
\]

for \(x = 0.01, \text{ and } Q^2 = 0.5 \text{ GeV}^2\). This is truly a remarkable result. It has long been known that for deep inelastic scattering from a free nucleon one measures \((\sigma_{L}\sigma_{T}) \ll 1\). The vanishing of this ratio is known (in other notation) as the Callan-Gross relation \[7\], and its verification provided the evidence for the finding that the partons observed in deep inelastic scattering are fermions (quarks). Now we have an experiment that seems to indicate that in nuclei the partons are bosons.

II. WHAT IS LIGHT FRONT DYNAMICS?

I try to mention only the most essential features. This dynamics is a relativistic many-body dynamics in which the fields are quantized at a “time”, \(t + z = x^0 + x^3 \equiv x^+\) and the canonically conjugate “energy” is given by \(p^0 - p^3 \equiv p^-\). Indeed, \(p^-\) is the \(x^+\) evolution operator, just as the Hamiltonian, \(p^0\) is the \(x^0\) time evolution operator. If \(x^+\) is “time”, then for “space” we must have \(x^- \equiv t - z\), with the canonically conjugate “momentum” as \(p^+ = p^0 + p^3 \equiv p^+\). The other coordinates are treated as usual \(\vec{x}_\perp, \vec{p}_\perp\). The use of the \(p^+\) as the momentum variable is, for me, the reason behind the use of light front dynamics. This is because for a particle moving with large speed such that \(\vec{v} \approx c\hat{e}_3\), \(p^+\) is BIG, and for this reason is the experimentalists variable. Many observables are best expressed using
this notation. In deep inelastic scattering the famous Bjorken variable \( x_{Bj} = Q^2/2M\nu \) is actually a ratio of plus momenta of the struck quark to that of the entire struck nucleon. With our choice of variables: \( A^\pm = A^0 \pm A^3 \), and the dot product of two four vectors is given by

\[
A \cdot B = A^\mu B_\mu = \frac{1}{2} (A^+ B^- + A^- B^+) - \vec{A}_\perp \cdot \vec{B}_\perp.
\]  

(2.1)

The most important application of Eq. (2.1) is the energy-momentum relation for a free particle:

\[
p^\mu p_\mu = m^2 = p^+ p^- - p^2 \perp,
\]

(2.2)

so that

\[
p^- = \frac{1}{p^+} (p^2 \perp + m^2).
\]

(2.3)

One has a relativistic expression for the energy without a square root operator. This is an enormous simplification when one wants to separate the coordinates of the center-of-mass from the rest of the wave function. We may provide an approximate summary of light front dynamics: Do ordinary quantum mechanics with energy denominators expressed in terms of \( p^- \).

Another feature is that, when one uses the Lagrangians of nuclear physics the usual lore about light front dynamics should be true. That is the vacuum really is empty. This is because nucleons are heavy enough so that nucleon pairs do not form vacuum condensates. Thus we should not ask what the light front dynamics can do for nuclear physics. Instead we should ask what nuclear physics can do for light front dynamics. This is to provide solutions of realistic, four-dimensional problems with relevance to observables.

### III. MOTIVATION FOR LIGHT FRONT NUCLEAR PHYSICS

Since much of this work is motivated by the desire to understand nuclear deep inelastic scattering and related experiments, it is worthwhile to review some of the features of the EMC effect [8][10]. One key experimental result is the suppression of the structure function for \( x \sim 0.5 \). This means that the valence quarks of bound nucleons carry less plus-momentum than those of free nucleons. Some other degrees of freedom must carry the plus-momentum, and some authors therefore postulate that mesons carry a larger fraction of the plus-momentum in the nucleus than in free space [11][12]. While such a model explains the shift in the valence distribution, one consequently obtains a meson (i.e. anti-quark) distribution in the nucleus, which is enhanced compared to free nucleons, and which should be observable in Drell-Yan experiments [13]. However, no such enhancement has been observed experimentally [14], and this has been termed as a severe crisis for nuclear theory in Ref. [15].

The EMC effect is rather small, so that one may begin by regarding the nucleus as being made of nucleons. In this case, we say that deep inelastic scattering proceeds when
a virtual photon is absorbed by a quark carrying plus-momentum \( p^+ \), which came from a nucleon carrying a plus-momentum \( k^+ \). In the parton model, the kinematic variable \( x_{Bj} = Q^2/2M_N\nu \) is given by

\[
x_{Bj} = \frac{p^+}{k^+}.
\] (3.1)

Thus one needs to know the probability \( f_N(k^+) \) that a nucleon has \( k^+ \). One also wants to know the related probability for a meson, for example: \( f_\pi(k^+) \).

Light front dynamics applies to nucleons within the nucleus as well as to partons of the nucleons, and this is a useful approach whenever the momentum of initial or final state nucleons is large compared to their mass [16]. For example, this technique can be used for \((e, e'p)\) and \((p, 2p)\) reactions at sufficiently high energies.

The essential technical advantage of using light cone variables is that the light cone energy \( P^- \) of a given final state does not appear in the delta function which expresses the conservation of energy and momentum. Thus one may use closure to perform the sum over final states which appears in the calculation of an exclusive nuclear cross section. The result is that the cross sections may be expressed in terms of the probabilities:

\[
\sigma \propto f_N(k^+) \sim \int d^2k_\perp \cdots |\Psi_{A,i}(k^+, k_\perp, \cdots)|^2,
\] (3.2)

where \( \Psi \) represents the ground state wave function.

For all these reasons we are concerned with calculating \( f_N(k^+) \). One standard approach to the calculation, based on using the shell model equal time formulation is that: \( E_\alpha + k^3 \rightarrow k^+ \). But this can not be correct because \( k^+ \) is a continuous kinematic variable which is not related to any discrete eigenvalues. Thus we need realistic calculations, with real dynamics and symmetries. This brings me to the conclusion that it is necessary to redo nuclear physics on the light front.

\[\text{IV. LIGHT FRONT QUANTIZATION LITE}\]

Our attitude towards this topic is that we need \( \mathcal{L} \) no matter how bad! This is because, in contrast with approaches based on symmetries, we try to obtain all of the necessary operators from a given Lagrangian \( \mathcal{L} \). The basic idea is to use the standard procedure to go from \( \mathcal{L} \) to \( T^{\mu\nu} \), with the essential difference from the usual procedure being that

\[
P^\mu = \frac{1}{2} \int d^2x_\perp dx^- T^{\mu+}.
\] (4.1)

A technical challenge arises because we need to express \( T^{\mu+} \) in terms of independent variables. For example, the spin 1/2 nucleon is always represented by 4 component spinor. Thus it has only 2 independent degrees of freedom. One needs to express the two dependent variables in terms of the two independent variables. This procedure is discussed in the references.

We use two Lagrangians. The first is that of the Walecka model [17]: \( \mathcal{L}(\phi, V^\mu, N) \) which contains the fields: nucleon \( N \), neutral vector meson \( V^\mu \), neutral scalar meson \( \phi \). This is the
simplest model which provides a reasonable caricature of the nucleus. The binding is caused by the attractive effects occurring at relatively long range when nucleons exchange scalar mesons. The nucleus is prevented from collapsing by the short distance repulsion arising from the exchange of vector mesons.

We also shall show results obtained using a more complicated chiral Lagrangian: in which the fields are $N, \pi, \sigma, \omega, \rho, \eta, \delta$. Our plan is to first use the Walecka model in the mean field approximation, and then to include $NN$ correlations using the chiral Lagrangian.

V. INFINITE NUCLEAR MATTER IN MEAN FIELD APPROXIMATION–WALECKA MODEL

The basic idea behind the solution is very well known. One assumes that the sources are strong, and produce sufficiently many mesons to justify a classical treatment. In infinite nuclear matter, one works in a limit in which the nuclear volume $\Omega$ is considered to be infinite, so that all positions, and directions equivalent in the nuclear rest frame. In this limit the fields $\phi$ and $V^\pm$ become constant, with $V_\perp = 0$. These features simplify the solution of the field equations. One easily obtains the operators $T^\pm$, and the light front "momentum" and "energy" are given by

\[
\frac{P^\pm}{\Omega} = \langle T^\pm \rangle, \tag{5.1}
\]

in which the expectation value is over the nuclear ground state.

The nuclear momentum content is the essential feature we wish to understand here. The results are that

\[
\frac{P^-}{\Omega} = m_s^2 \phi^2 + \frac{4}{(2\pi)^3} \int_F d^2 k_\perp d k^+ \frac{k^2_\perp + (M + g_s \phi)^2}{k^+} \tag{5.2}
\]

\[
\frac{P^+}{\Omega} = m_v^2 (V^-)^2 + \frac{4}{(2\pi)^3} \int_F d^2 k_\perp d k^+ k^+ \tag{5.3}
\]

The first term of $P^+$ is the plus momentum carried by vector mesons, and the second term is the plus momenta carried by the nucleons. Here $g_s$ is the scalar-meson-nucleon coupling constant, and the vector meson-nucleon coupling constant $g_v$ enters in the expression for $V^-$. The interpretation of these results is aided by a change of variables:

\[
k^+ \equiv \sqrt{(M + g_s \phi)^2 + k^2 + k^3}, \tag{5.4}
\]

which defines the variable $k^3$. Using this variable one can show that rotational invariance is respected and obtain a spherical Fermi surface. Furthermore, one may show that $E \equiv \frac{1}{2} (P^- + P^+)$ is the same as the usual expression obtained in the Walecka model.

For nuclear matter in its rest frame we need to obtain $P^+ = P^- = M_A$. This is the light front expression of the statement that the pressure on the system must vanish \[3\]. Indeed the minimization

\[
\left( \frac{\partial(E/A)}{\partial k_F} \right)_\Omega = 0, \tag{5.5}
\]
determines the value of the Fermi momentum and is an expression that gives \( P^+ = P^- = M_A \).

We can quickly obtain the relevant numerical results using the 1974 parameters of Chin & Walecka. These are

\[
g_s^2 M_N^2 / m_s^2 = 195.9 \quad g_v^2 M_N^2 / m_v^2 = 267.1.
\] (5.6)

With these parameters \( M_N + g_s \phi = 0.56 M_N \) and \( g_v V^- = 270 \text{ MeV} \). These are the HUGE scalar and vector potentials which are characteristic of the Walecka model. The interesting variables are those associated with the total nuclear plus momentum \( P^+ \). With the above parameters, the vector meson contribution to this quantity: \( \int F d^2 k_{\perp} dk^+ k^+ \) is a monumental 0.35 \( \frac{P^+}{M} \), while the nucleon contribution \( \frac{4}{(2\pi)^3} \int_F d^2 k_{\perp} dk^+ k^+ \) is only 0.65 \( \frac{P^+}{M} \). Only 65% of \( P^+ \) carried by nucleons, but 90% is needed to understand the EMC effect in infinite nuclear matter \([18]\). This difference is huge. One can’t plot the results of the theory in comparison with the experimental results using pages of ordinary size.

So the large plus momentum carried by the vector mesons is a feature of the Walecka model, which seems very odd. One needs to find means to reduce this percentage from 35%, but one also expects that within any model the vector mesons would carry some plus-momentum and it therefore becomes desirable to search for an experimental verification of this feature.

VI. BEYOND INFINITE NUCLEAR MATTER– MEAN FIELD THEORY

One possibility is that the large plus momentum carried by vector mesons arises as an artifact of using infinite nuclear matter. We (Blunden, Burkardt and I, \([2]\)) therefore investigated the subject of light front Hartree Fock theory by performing nucleus mean field theory calculations.

The motivation is the need to find a way to reduce the plus-momentum carried by vector mesons, but one encounters a significant difficulty in the theory because the use of the variables \( (x^- , \vec{x}_{\perp}) \) entails the loss of manifest rotational invariance. We found that the procedure of minimizing \( \langle P^- \rangle \) subject to the constraint \( \langle P^+ \rangle = \langle P^- \rangle \), led to the recovery of rotational invariance. This was seen by counting the number of degenerate levels \( (2j + 1) \) that emerges from a numerical calculation in which only rotational invariance about the \( z \)-axis could be used to simplify the calculation. This is important for understanding the existence of magic numbers in nuclei. The result of doing the lengthy calculations was that nucleons were found (using the Walecka model) to carry about 70% of total \( P^+ \). This was only a modest improvement over the nuclear matter result of 65%, and does not resolve the problem.

VII. BEYOND MEAN FIELD THEORY

The interactions between nucleons are strong, and the mean field approximation is unlikely to provide a description of nuclear properties which involve high momentum observables. We developed a version of light front theory in which the correlations between two nucleons are included. The theory was applied to infinite nuclear matter.
The calculation required three principal steps. (1) Light front quantization of chiral $\mathcal{L}$. (2) Derive Light Front version of the $NN$ one boson exchange potential. This could be done exploiting the relationship between the Weinberg equation and the Blankenbecler-Sugar equation. A sample of results for the phase shift is shown in Fig. 1. The third step is to develop the many body theory. This turns out to be a long story [3]. The net result is that the light front theory looks like the usual relativistic Brueckner theory except that the Blankenbecler-Sugar equation is used, and the effects of retardation are kept. The resulting nuclear matter saturation curve is shown in Fig. 2. Standard good results for nuclear saturation properties are obtained, with a possible improvement in the lowered value, 180 MeV, of the computed nuclear compressibility.

The results for deep inelastic scattering and the related Drell-Yan process seem very promising. The nucleons carry at least 84% of the nuclear plus momentum. This is calculated using only the Fermi gas part of the wave function. The percentage might increase to about 90% (the target value for deep inelastic scattering) if the effects of the two-particle two-hole states are included. In these calculations the nucleus does have a pionic component, which arises as a result of going beyond the mean field approximation. The number of excess pions per nucleon is about 5%, and although the distribution function of nuclear pions has not yet been computed, this seems small enough to avoid a contradiction with the Drell-Yan data.
FIG. 2. Nuclear matter binding energy per particle. Solid: our theory, dashed: non-relativistic theory, dotted ignoring the effects of retardation.

It seems that nuclear physics can be done in a manner in which modern nuclear dynamics is respected, boost invariance in the $z$-direction is preserved, and in which the rotational invariance so necessary to understanding the basic features of nuclei is maintained. A salient feature is that $\omega$, $\sigma$ and $\pi$ mesons are important constituents of nuclei. Another point, not much discussed here, is that it seems possible to find Lagrangians that yield reasonable descriptions of nuclear deep inelastic scattering and Drell-Yan reactions.

In the remainder of this talk I pursue the idea that the HERMES effect provides a signature for the presence of nuclear $\omega$ and $\sigma$ mesons.

VIII. COHERENT CONTRIBUTIONS OF NUCLEAR MESONS TO ELECTROPRODUCTION– HERMES EFFECT

Let us discuss the nature of the HERMES effect [1]. Everyone is very familiar with the idea that lepton-nucleon or lepton nucleus scattering proceeds by the exchange of a virtual photon of four momentum $q = (\nu, \vec{q})$. The important structure functions depend mainly on $x_{Bj}$, with dependence on the logarithm of $Q^2 = -q^2$. However, the cross sections depend on a third variable: $y = \frac{\nu}{E}$, where $E$ is the energy of the incident lepton. The energy at which HERMES runs, 27.5 GeV, is small enough so that the experiment covers a wide range of values of $y$. This feature is what allows a new effect to be observed. In particular, the cross section depends on the scattering of transversely $T$ and longitudinally $L$ polarized photons [20]:

$$\sigma \propto \sigma_T + \epsilon \sigma_L$$  \hspace{1cm} (8.1)

$$\epsilon \approx \frac{4(1 - y)}{4(1 - y) + 2y^2}.$$  \hspace{1cm} (8.2)
It is conventional to make the definition
\[ R \equiv \frac{\sigma_L}{\sigma_T}. \] (8.3)
In that case, one may manipulate the standard relations between the cross section and structure function to find
\[ \frac{\sigma_A}{\sigma_D} = \frac{F_A^2}{F_D^2} \frac{1 + \epsilon R_A}{1 + \epsilon R_D}, \] (8.4)
A linear dependence of the ratio on \( \epsilon \) provides a signature of a large value of \( R_A \). Indeed, the HERMES collaboration extracted \( \frac{F_A^2}{F_D^2} \) and \( R_A \) from \( x, Q^2 \), and \( \epsilon \) dependence of \( \frac{\sigma_A}{\sigma_D} \). The results are that
\[ \frac{\sigma_L(A)}{\sigma_L(D)} > 1, \quad \frac{\sigma_T(A)}{\sigma_T(D)} < 1, \] (8.5)
with the largest effects
\[ \frac{R_A}{R_D} \approx 5 \] (8.6)
obtained for \( x \approx 0.01, \quad Q^2 = 0.5 \text{ GeV}^2 \). As noted in the Introduction, this represents a huge violation of the Callan-Gross relation \( [7] \), a violation large enough to indicate that, in nuclei, bosons are the partons of deep inelastic scattering! In the following, I describe the work of Ref. \( [5] \).

\section*{IX. NUCLEAR ENHANCEMENT OF \( \sigma_L \)}

We wish to describe the nuclear enhancement of \( \sigma_L \) and the nuclear suppression of \( \sigma_T \) using a single input theory. We start with \( \sigma_L \). We found that a process in which a virtual photon of four momentum \( q \) is converted in its interaction with a nuclear \( \omega \) meson into a scalar meson of four momentum \( p \), with \( p^0 \approx \nu \), produces the desired enhancement.

To evaluate the effects of this process, we need to determine the \( \gamma \omega \sigma \) interaction. We postulate a gauge-invariant form
\[ \mathcal{L}_I = \frac{g e}{2m_\omega} F^{\mu \nu}(\omega_\mu \partial_\nu \sigma - \omega_\nu \partial_\mu \sigma) \] (9.1)
where \( F^{\mu \nu} \) is the photon field strength tensor. In momentum space one can use
\[ \mathcal{M} = \frac{g}{m_\omega} (p \cdot q \, \epsilon^\gamma \cdot \epsilon^\omega - p \cdot \epsilon^\gamma \, q \cdot \epsilon^\omega) F_V(Q^2) \] (9.2)
in which we include a form factor \( F_V \). We seek a constraint on the value of \( g \) from the decay: \( \omega \to \sigma \gamma \). The branching ratio for \( \omega \to \pi^+ \pi^- \gamma < 3.6 \times 10^{-3} \) \( [21] \), which we assume to come from the process \( \omega \to \gamma \sigma \) followed by the two pion decay of the \( \sigma \) meson. We may determine an upper limit for \( g \): \( g_{UL}^2 \alpha = .013 \approx 2\alpha \).
Using this coupling constant, we obtained

$$\delta W^{00} \sim (V^-)^2 A^{1/3} \nu^3 F_V^2(Q^2),$$

(9.3)

in which $V^-$ is the value of the vector meson field at the center of the nucleus. We use values which are consistent with nuclear saturation, and DIS, DY reactions. Notice the presence of the $\nu^3$ term which arises from the factors of momentum appearing in Eq. (9.2). This dependence is essential because standard kinematics gives the result

$$\frac{\sigma_L(A)}{\sigma_L(D)} = 1 + \frac{Q^4}{\nu(\nu^2 + Q^2)} \frac{\delta W^{00}}{F_V^2 R^2}(1 + R_D).$$

(9.4)

The form factor $F_V(Q^2)$ is obtained from Ref. [22], we also use a dipole form factor. The results are shown in Fig. (3). One is able to account for the large enhancement.

**X. NUCLEAR SUPPRESSION OF $\sigma_T(A)$**

Our explanation of the nuclear transverse cross section $\sigma_T$ data requires a significant destructive interference effect at low $Q^2 \approx 0.5 - 2 \text{ GeV}^2$. For small values of $x$, in the target rest frame, the interaction proceeds by $\gamma^*$ decaying into a $q\bar{q}$ pair, which then interacts with a target nucleon and emerges as a vector meson $V_f$. For a nuclear target, we suppose that the virtual photon interacts with a nuclear $\sigma$ meson and is converted to an intermediate vector boson $V$, which is converted into the final vector meson $V_f$ by a final state interaction.

We need to find the necessary $\gamma^* \sigma \rightarrow V$ interaction. For consistency with data taken at larger value of $Q^2$, we need an interaction which decreases rapidly as $Q^2$ increases. Furthermore, the shadowing of the real photon ($Q^2 = 0$) is not very strong, and it is well explained...
by conventional vector meson dominance models [23]. Thus consistency with all available data demands an amplitude for $\gamma^* \sigma \rightarrow V$ which vanishes, or is small, as the $Q^2$ of the virtual photon $\gamma^*$ approaches 0. This means that measuring the real photon decays of the vector mesons provides no constraints on the coupling constant.

We postulate the gauge-invariant interaction

$$\delta \mathcal{L}_I = \sum_V \frac{g_{\gamma V \sigma} e}{2 m_\sigma} F_{\mu \nu} [V^{\mu \nu} \sigma + V^\mu \partial^\nu \sigma - V^\nu \partial_\mu \sigma] F_V^V,$$

which in momentum space is $\propto Q^2$. The details of the application are given in the published work. The results are shown in Fig. (4).

The nuclear enhancement of $R$ is obtained from computing the ratio of the previous results. This is shown in Fig. (5), where it is seen one has a reasonably good description of the data.

We may summarize the results of our studies of the HERMES effect. We find that $\sigma_L(A)$ is enhanced by the presence of nuclear vector mesons, and that $\sigma_T(A)$ is depleted by the presence of nuclear scalar mesons. Both types of mesons are needed to obtain nuclei with correct binding energies and densities. The values of the strong coupling constants used are also roughly consistent with data on nuclear deep inelastic scattering data taken at larger values of $Q^2$.

Much verification of the present model is needed. Further tests of our model are possible. An immediate consequence would be the observation of exclusive mesonic states in the current fragmentation region. In particular, our description of $\sigma_L(A)$ implies significant nuclear-coherent production of $\sigma$ mesons along the virtual photon direction. Our model for the strong shadowing of coherent meson effects in $\sigma_T(A)$ can be tested by measurements performed at the same value of $x$ but different values of $Q^2$ than HERMES used.
The prospect that the mesonic fields which are responsible for nuclear binding can be directly confirmed as effective fundamental constituents of nuclei at small $x$ and $Q^2 \sim 1\text{ GeV}^2$ is an exciting development at the interface of traditional nuclear physics and QCD. The empirical confirmation of nuclear-coherent meson contributions in the final state would allow the identification of a specific dynamical mechanism for higher-twist processes in electroproduction. Clearly, these concepts should be explored further, both experimentally and theoretically.

ACKNOWLEDGMENTS

This work was supported in part by the United States Department of Energy under contract and DE-FG03-97ER4104. I thank all of the collaborators whose work I discuss here: P. Blunden, S. Brodsky, M. Burkardt, M. Karliner, and R. Machleidt.
REFERENCES

[1] G.A. Miller, Phys. Rev. C 56 (1997) R8; 56 (1997) 2789.
[2] P.G. Blunden, M. Burkardt, and G.A. Miller, Phys. Rev. C59 (1999) R2998; Phys. Rev. C 60 (1999) 55211.
[3] G.A. Miller and R. Machleidt, Phys. Lett. B455 (1999) 19; Phys. Rev. C60 (1999) 035202-1.
[4] G. A. Miller, Prog. Part. Nucl. Phys. 45, 83 (2000).
[5] G. A. Miller, S. J. Brodsky and M. Karliner, Phys. Lett. B481, 245 (2000).
[6] K. Ackerstaff et al. [HERMES Collaboration], hep-ex/9910071.
[7] C.G. Callan and D.J. Gross, Phys. Rev. Lett. 22 (1969) 156.
[8] J. Aubert et al., Phys. Lett. 123B (1982) 275; R.G. Arnold et al., Phys. Rev. Lett. 52 (1984) 727; A. Bodek et al., Phys. Rev. Lett. 51 (1983) 534.
[9] Piller and Weise L.L. Frankfurt and M.I. Strikman, Phys. Rep. 160(1988) 235; M. Arneodo, Phys. Rep. 240 (1994) 301; D.F. Geesaman, K. Saito, A.W. Thomas, Ann. Rev. Nucl. Part. Sci. 45 (1995) 337.
[10] L.L. Frankfurt and M.I. Strikman, Phys. Rep. 160 (1988) 235.
[11] C.H. Llewellyn Smith B128 (1983) 107.
[12] M. Ericson and A.W. Thomas, Phys. Lett. B128 (1983) 112.
[13] R.P. Bickerstaff, M.C. Birse, and G.A. Miller, Phys. Rev. Lett. 53, (1984) 2532; M. Ericson and A.W. Thomas, Phys. Lett. 148B (1984) 191.
[14] D.M. Alde et al., Phys. Rev. Lett. 64 (1990) 2479.
[15] G.F. Bertsch, L. Frankfurt, and M. Strikman, Science 259 (1993) 773.
[16] L.L. Frankfurt and M.I. Strikman, Phys. Rep. 76, (1981) 215.
[17] B.D. Serot and J.D. Walecka, Adv. Nucl. Phys. 16 (1986)1; R. J. Furnstahl and B. D. Serot, nucl-th/9912048.
[18] I. Sick and D. Day, Phys. Lett. B274 (1992) 16.
[19] H. Jung and G.A. Miller, Phys. Rev. C41(1990) 659.
[20] R. G. Roberts, Cambridge, UK: Univ. Pr. (1990) 182 p. (Cambridge monographs on mathematical physics).
[21] C. Caso et al., Eur. Phys. J. C3, 1 (1998).
[22] H. Ito and F. Gross, Phys. Rev. Phys. Rev. Lett. 71, 2555 (1993).
[23] T. H. Bauer, R. D. Spital, D. R. Yennie and F. M. Pipkin, Rev. Mod. Phys. 50 (1978) 261.