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Coherent Contributions of Nuclear Mesons to Electroproduction and the HERMES Effect*

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Abstract

We show that nuclear σ, ω, and π mesons 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. We are able to describe recent HERMES inelastic lepton-nucleus scattering data at low Q^2 and small x using photon-meson and meson-nucleus couplings which are consistent with (but not determined by) existing constraints from meson decay widths, nuclear structure, deep inelastic scattering, and lepton pair production data. We find that while nuclear-coherent pion currents are not important for the present data, they could be observed at different kinematics. Our model for coherent meson electroproduction requires the assumption of mesonic currents and couplings which can be verified in separate experiments. The observation of nuclear-coherent mesons in the final state would verify our theory and allow the identification of a specific dynamical mechanism for higher-twist processes.

Nuclear targets provide a unique way to adiabatically modify the hadronic environment when testing QCD. The shadowing and antishadowing effects of the nuclear medium on the electroproduction cross section and nuclear structure functions in the Bjorken scaling region are typically less than 20%, in qualitative agreement with theoretical expectations. However, recent measurements by the HERMES collaboration of the inelastic lepton-nucleus cross section at low $Q^2 < 1.5$ GeV^2 and small $x < 0.06$ for a 27.5 GeV positron beam interacting on gas jet targets at HERA display an extraordinarily strong nuclear and virtual photon polarization dependence [1]. The HERMES data for deuterium, ^{3}He and ^{14}N targets show an anomalously strong nuclear dependence of the ratio $R^A = \sigma_L^A(x, Q^2)/\sigma_T^A(x, Q^2)$ at low momentum transfer Q^2 and small x. For example, R^N/R^D (nitrogen vs. deuterium) is $\simeq 5$ at $Q^2 \simeq 0.5$ GeV^2, for $x \simeq 0.01$. This ratio of five results from a nuclear enhancement of the longitudinally polarized virtual photoabsorption by about a factor of 2 and a reduction of the transverse cross section $\sigma_T(x, Q^2)$ by about a factor of 2.5. These nuclear effects are very much larger than the typical shadowing effects mentioned above, and very much larger than previous estimates [2] of nuclear enhancements of R^A. The nuclear target experiments in this kinematic regime are very challenging because of the large size of the radiative corrections.

It has long been recognized that a small value of R is the signature of spin 1/2 partons, and, conversely, a large value would be a signature of bosonic constituents [3].

Furthermore, the rapid decrease of the nuclear enhancement of σ_L with increasing Q^2 seen by HERMES is compatible with elastic scattering of the positron on a composite bosonic system: the power-law decrease of the square of the mesonic form factors $|F_M(Q^2)|^2 \propto 1/Q^4$ could account for the rapid fall-off of the R^A/R^D enhancement with momentum transfer. It is thus natural to interpret the unusual nuclear enhancement of R^A reported by HERMES in terms of leptons scattering on the mesonic fields of nuclei.

The fundamental microscopic description of the deep inelastic electroproduction cross section on nuclei in QCD is based on quark and gluon degrees of freedom, but at small Q^2 there exist intricate and non-local correlations which are more readily described in terms of mesonic currents. A familiar example is the emergence of pions as approximate Goldstone bosons, which is difficult to see directly in the quark language, but is immediate in the language of effective Lagrangians. Thus meson fields are natural degrees of freedom for the physics of leptoproduction at low values of Q^2 and x.

The nuclear enhancement of $\sigma_L^A(x, Q^2)$ for $x < 0.06$ observed by HERMES suggests constructive interference of amplitudes from mesons emitted by different nucleons. The minimum laboratory momentum transfer to the nucleus in diffractive production of a meson of mass m_M is $\Delta p_L \simeq (Q^2 + m_M^2)/2\nu \simeq 2x_{bj}M_N \approx 60$ MeV, which is comparable with the inverse nuclear size of nitrogen. Thus electroproduction can occur from higher-twist subprocesses in which the lepton scatters elastically on mesons emitted coherently throughout the entire nuclear volume.

Now let us consider the specific processes which could contribute. The most natural effect to consider is positron quasi-elastic scattering from a nuclear pion. The emission of a pion leads to a set of low-lying nuclear states, and one could find significant effects in the sum over states. However, the probability of finding a pion at small values of x in the nuclear medium vanishes relative to vector mesons, reflecting the connection between the Regge behavior of deep inelastic structure functions and the spin of the exchanged constituents [4]. Thus, even though the pion couples strongly to nucleons, we do not expect the pion to contribute significantly to the HERMES effect at small x. However, as we discuss below, pionic currents can yield significant effects for $x \approx 0.25$ and $Q^2 \sim 1$ GeV^2.

Other nuclear mesons besides pions are known to play an important in nuclear

physics. In simple models [5], the exchange of scalar mesons between nucleons leads to attractive forces which bind the nucleus. On the other hand, the exchange of vector mesons supplies the repulsive potential which prevents the collapse of the nucleus. The strengths of such fields in the nucleus are quite large [6], corresponding to a significant number of effective bosonic partons at very small values of x [7]. Furthermore, in high energy processes, the effects of vector mesons are enhanced because of the presence of factors of momentum in the interaction. This is an essential feature of the explanation provided here.

Nuclear deep inelastic lepton-scattering and Drell-Yan experiments place important constraints which limit the nuclear enhancements of mesons relative to the nucleon; see the summary [8]. For example, for infinite nuclear matter the nucleons must carry 90% of the light cone plus momentum [9], which implies mesons can carry no more than 10%. Our calculations shall use models and parameters which respect these constraints.

We begin a quantitative analysis by examining the consequences of nuclear-coherent vector and scalar mesons. We find that the process of Fig. 1a in which the interaction between the virtual photon γ^* of momentum q with a nuclear ω meson which produces a σ meson in the final state can give a significant contribution to $\sigma_L(A)$. To understand the shadowing of σ_T, we need a process which interferes destructively with the dominant process (at low x) of Fig. 1b. We find that the process of Fig. 1c in which the virtual photon converts a nuclear σ meson into a vector meson can supply the necessary destructive interference. If this amplitude is 1/3 of the dominant process it leads to a reduction of the cross section of a factor of 4/9, which is needed to account for the shadowing observed at low x. In order to evaluate these diagrams, we will adopt a procedure of postulating photon-meson interactions, consistent with gauge invariance, and then verifying that there is no conflict with available information.

Consider the longitudinal cross section. We denote the contribution (per nucleon) to the nuclear hadronic tensor caused by excess nuclear mesons as $\delta W^{\mu\nu}$. This is in addition to the gluonic effects which give a non-zero value for R for a free nucleon. Then, using standard kinematic relations [10], we take the ratio of nuclear to nucleon

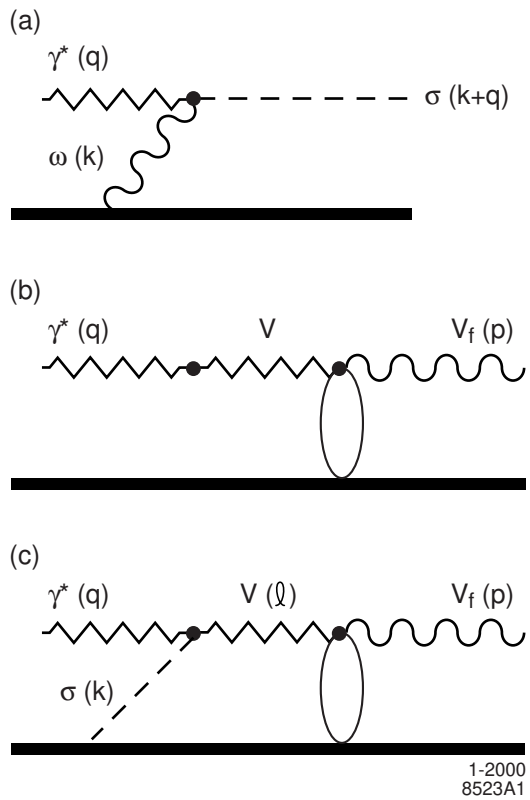

Figure 1: Diagrams for the high energy, low Q^2 γ^* nuclear scattering. For each mechanism, the nuclear target is represented by a heavy line. (a) Dominant contribution to $\sigma_L(A)$; (b) Dominant contribution to the transverse cross section; (c) Mesonic term that interferes with (b).

values of the longitudinal cross section to find

$$\frac{\sigma_L(A)}{\sigma_L(D)} = 1 + \frac{Q^4}{\nu(\nu^2+Q^2)} \frac{\delta W^{00}}{F_2^D R_D}(1+R_D). \tag{1}$$

The notation D represents the nucleonic value as represented by the deuteron cross section, with F_2^D and R_D taken from standard parameterizations of the data[11, 12]. The implication of Eq. (1) is that δW^{00} must vary as ν^3 to obtain a significant effect at small x.

In order to evaluate the effects of Fig. 1a, we need to determine the $\gamma\omega\sigma$ interaction. We postulate a gauge-invariant form

$$\mathcal{L}_I = \frac{g\ e}{2m_\omega} F^{\mu\nu}(\omega_\nu \partial_\mu \sigma - \omega_\mu \partial_\nu \sigma) \tag{2}$$

where $F^{\mu\nu}$ is the photon field strength tensor. In momentum space one can use

$\mathcal{M} = \frac{g}{m_\omega}(p \cdot q\, \epsilon^\gamma \cdot \epsilon^\omega - p \cdot \epsilon^\gamma\, q \cdot \epsilon^\omega)\, F_V(Q^2)$ in which p_μ is the momentum of the σ, and 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 $\omega \to \pi^+\pi^-\gamma < 3.6 \times 10^{-3}$ [13], which we assume to come from the process $\omega \to \gamma\sigma$ followed by the two pion decay of the σ meson. This process supplies a contribution to the ω width, $\delta\Gamma = (e^2/12\pi)(q_0^3 g^2/m_\omega^2)$, where q_0 is the photon energy in the cm frame, $q_0 = (m_\omega^2 - m_\sigma^2)/2m_\omega$ and $F_V(Q^2 = 0) = 1$. The predicted value of $\delta\Gamma$ and the extracted value of g depend strongly on the mass of the σ meson, which is given [13] as the $f(400 - 1200)$. If we use the average value of 800 MeV, the decay would not occur, and the width of the ω would provide no constraint on the value of g. In these first calculations we choose $m_\sigma = 600$ MeV, and determine an upper limit for g: $g_{UL}^2\alpha = .013 \approx 2\alpha$. We shall take $g_{UL}^2 = 2$ as a nominal value.

The standard formula for a contribution to the hadronic tensor per nucleon is

$$\delta W^{\mu\nu} = \frac{1}{4\pi M_A}\frac{1}{A}\int d^4\xi\; e^{iq\cdot\xi}\; \langle P|J^\mu(\xi)J^\nu(0)|P\rangle. \tag{3}$$

The current J^μ is obtained from our interaction (2) using $J^\mu = \delta\mathcal{L}_I/\delta A_\mu$. The state $|P\rangle$ is the nuclear ground state, normalized as $\langle P'|P\rangle = 2E(P)(2\pi)^3\delta^{(3)}(P-P')$. The only terms of J_ν we need to keep are those which are proportional to the large momentum of the outgoing σ meson. Keeping these and evaluating Eq. (3), gives the result:

$$\delta W^{\mu\nu} = g^2 F_V^2(Q^2)\,\frac{p\;\nu^2}{2A}\,\frac{p^\mu p^\nu}{m_\omega^2}\int d\Omega_{\mathbf{p}}\;\omega^0(\mathbf{p}-\mathbf{q})\omega^{\mathbf{0}}(\mathbf{p}-\mathbf{q}), \tag{4}$$

in which negligible retardation effects in the ω propagator are ignored, p^ν is the momentum of the outgoing σ meson, and p is the magnitude of its three-momentum, $\sqrt{\nu^2 - m_\sigma^2}$. Note that only the time ($\mu = 0$) component of the field ω^μ has a significant value. The term $\omega^0(\mathbf{p}-\mathbf{q})$ is the Fourier transform of the nuclear vector potential,

$$\omega^0(\mathbf{p}-\mathbf{q}) = \int \frac{d^3r}{(2\pi)^{3/2}} e^{i(\mathbf{p}-\mathbf{q})\cdot\mathbf{r}}\omega^0(\mathbf{r}). \tag{5}$$

which contains the nuclear form factor. The momentum transfer **p-q** will be huge on a scale of nuclear momenta $\sim 1/R_A$, and the nuclear form factor nearly vanishes unless **p** is parallel to the direction of the virtual photon. As a result, the angular integration gives a factor $\sim 1/(R_A^2\nu^2)$, which, when multiplied by a large factor $\sim R_A^6$ arising from the product of two volume integrals over the entire size of the nucleus, leads to $\delta W^{00} \sim R_A\;\nu^3$. This term increases roughly as $\mathrm{A}^{1/3}$, representing the net coherent effect of the nuclear ω field, and it has the necessary dependence on ν.

The essential input is $\omega^0(\mathbf{p}-\mathbf{q})$. For the ^{14}N target, we use a Fermi form: $\omega_F^0(r) = V^0/[1+e^{(r-R)/a}]$. The resulting δW_F^{00} is computed numerically using Eqs. (4,5). The term V^0 is the central value of the field. In mean field theory, V^0 is the vector potential for the nucleon divided by the ω-nucleon coupling constant. The value is $V^0 = 28$ MeV in nuclear matter (along with the nucleon effective mass, $M^* = 0.56M$) using QHD1 [5]. However, the very small nucleon effective mass causes the value of light cone plus momentum carried by nucleons to be far too small for consistency with lepton-nuclear deep inelastic scattering data. However, there are many versions of the theory. One could include meson-meson interactions in the Lagrangian, as well as RPA and Brueckner correlations. For example, in a mean field theory which includes σ^3 and σ^4 terms, the value $V^0 = 13$ MeV ($M^* = .84M$) gives a reasonable description of many nuclear properties [14], and is also consistent with nuclear deep inelastic scattering [15]. The quark-meson coupling model [16] also gives $M^* \simeq 0.8M$. Thus the expected range of V^0 is $13 < V^0 < 28$ MeV. In our present calculations we adopt the average value of $V^0 = 20$ MeV as a nominal value.

In order to complete the description of our calculation we need to specify the form factor. The expression: $F_V(Q^2) = (1+Q^2/m_\rho^2)^{-1.5}$ is an analytic representation of the quark-loop diagram calculation of Ito and Gross [17]. The results obtained with this form factor (labeled IG) are shown in Fig. 2, using $R = 2.89$ fm, $(g^2/g_{UL}^2)(V^0/20\text{MeV})^2 = 1$, a value which is consistent with available information. To display the sensitivity to parameters, we also use a dipole form factor (results labeled dipole in Fig. 2), along with the values $(g^2/g_{UL}^2)(V^0/20\text{MeV})^2 = 1.2$. In either case, the values of $\sigma_L(A)$ are large, and a qualitative reproduction of the HERMES data is obtained. Clearly, the values of the couplings could be much smaller; for example if $(g^2/g_{UL}^2) = 1/4$, then the computed ratio of $\sigma_L(A)/\sigma_L(D)$ would be ~ 1.2 for $x \sim 0.01$.

We now return to the effects of the pion. The $\gamma^* - \pi$ interaction is represented by the current: $(\partial^\mu \phi^*)\phi + \phi^* \partial^\mu \phi$. Its use in Eq. (1) gives the pionic contribution as:

$$\frac{\sigma_L^\pi(A)}{\sigma_L(D)} = \frac{Q^4}{2\nu^3}\frac{1+R_D}{F_2^D R_D}\left[\frac{\nu^2}{Q^2 M_N}\right] f_{\pi/A}(x) F_\pi^2(Q^2) = x f_{\pi/A}(x)\frac{1+R_D}{F_2^D R_D} F_\pi^2(Q^2). \quad (6)$$

The effects of the pion form factor are included via the term $F_\pi^2(Q^2)$. The pion distribution function $f_{\pi/A}(x)$ gives the probability that a nuclear pion has a plus momentum of xM_N. As noted above, $f_{\pi/A}(x)$ vanishes as x approaches zero. The

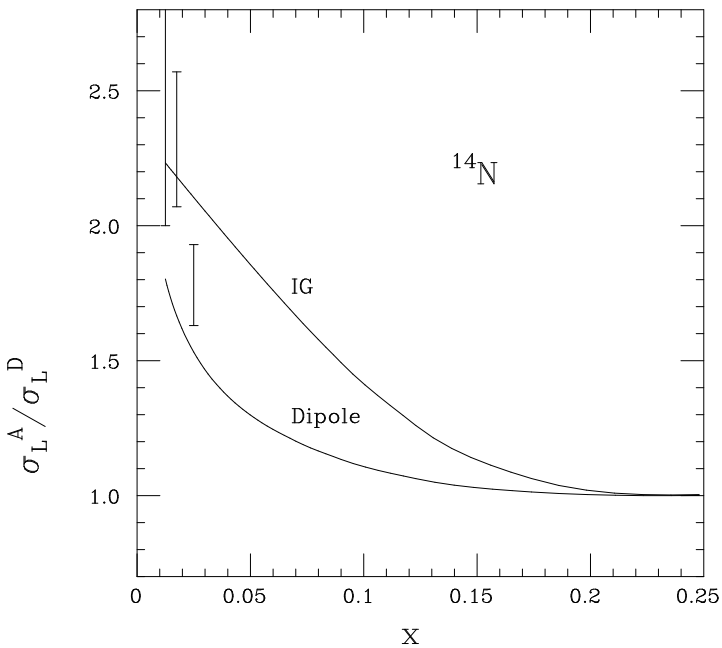

Figure 2: $\frac{\sigma_L(A)}{\sigma_L(D)}$, A=14, data from Ref. [1]. The labels IG, dipole refer to form factors, see text.

value of $\sigma_L^\pi(A)/\sigma_L(D)$ for the data point at $x = 0.0125$ is estimated using [11, 12] $F_2^D(x) = 0.22$, $R_D = 0.36$, and $f_{\pi/A}(x) \approx 0.01$ from the calculation of Ref. [18] for infinite nuclear matter. This calculation uses a version of the Ericson-Thomas [19] model in which the parameters are chosen to be consistent with nuclear deep inelastic and Drell-Yan data. The result is $\sigma_L^\pi(A)/\sigma_L(D) \sim 0.01$ which is negligible. The function $f_{\pi/A}(x)$ peaks at $x = 0.25$ with a value (for the charged pions) of 0.24. For this x, $F_2^D \approx 0.23$, $R_D = 0.3$. These values give $\sigma_L^\pi(A)/\sigma_L(D) \approx F_\pi^2(Q^2)$. For the kinematics in the HERMES experiment, $Q^2 \approx 3.1$ GeV^2 and the square of the form factor is ≈ 0.03. Thus one obtains a contribution of about 0.03. This is small compared to the vector meson current contributions, but it is not entirely negligible. However, in an experiment at $x = 0.25$ with $Q^2 \sim 1$ GeV^2, nuclear-coherent pions would contribute significantly to σ_L.

The presence of nuclear pions opens the door to a variety of new experiments. One could study the exclusive final state of pionic reactions such as $\gamma^*\ {}^3\mathrm{He} \to \pi^+\ {}^3\mathrm{H}$, $\gamma^*\ {}^{14}\mathrm{N} \to \pi^-\ {}^{14}\mathrm{O}$, and $\gamma^*\ {}^{14}\mathrm{N} \to \pi^+\ {}^{14}\mathrm{C}$, where the final state nuclides can be formed in their ground or excited states. These nuclear-diffractive electroproduction

processes can be used to test predictions from QCD for the Q^2 and t dependence of the off-shell mesonic form factors $F_M(Q^2,t)$. Conversely, by tagging the nuclear final state in an inelastic reaction, one could identify the structure functions of off-shell mesons, as in the Sullivan process. Such processes are analogous to the measurements of the pomeron structure function in diffractive deep inelastic scattering.

Next we examine the nuclear value of the transverse cross section σ_T. Our explanation of the data requires a significant destructive interference effect at low $Q^2 \approx 0.5 - 2$ GeV^2, 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 [20]. Thus consistency with all available data demands an amplitude for $\gamma^*\sigma \to V$ which vanishes, or is small, as the Q^2 of the virtual photon γ^* approaches 0. This means that measuring the real photon decays of the vector mesons provides no constraints on the coupling constant. We postulate a sum of two gauge-invariant forms:

$$\mathcal{L}_{\rho\gamma\sigma} = \frac{g_{\rho\gamma\sigma}}{2m_\sigma}\left[F_{\mu\nu}\rho^{\mu\nu}\sigma + \lambda F_{\mu\nu}(\rho^\mu\partial^\nu\sigma - \rho^\nu\partial^\mu\sigma)\right] \tag{7}$$

$$\to \frac{-g_{\rho\gamma\sigma}}{m_\sigma}\left[(q\cdot(k+q)\epsilon^\gamma\cdot\epsilon^\rho - q\cdot\epsilon^\rho k\cdot\epsilon^\gamma) + \lambda(q\cdot\epsilon^\rho k\cdot\epsilon^\gamma - q\cdot k\epsilon^\gamma\cdot\epsilon^\rho)\right]. \tag{8}$$

The choice $\lambda = 1$ leads to an effective Lagrangian $\frac{g_{\rho\gamma\sigma}}{m_\sigma}\left(Q^2\epsilon^\gamma\cdot\epsilon^\rho + q\cdot\epsilon^\gamma\; q\cdot\epsilon^\rho\right)F_V(Q^2)$, which is the form that we adopt. We choose a gauge such that $q\cdot\epsilon^\gamma = 0$. Note also that the difference in the two terms of Eq. (7) is due to the source of the electromagnetic field, so that $\mathcal{L}_{\rho\gamma\sigma}$ depends on the virtuality of the photon. A form factor $F_V(Q^2)$ is introduced. We simplify the calculation by considering only one intermediate vector meson state, the ρ meson. Then the usually dominant term $\mathcal{M}_{\rm dom}$ of Fig. 1b takes the form:

$$\mathcal{M}_{\rm dom} = e\frac{m_\rho^2}{f_\rho}\frac{-1}{Q^2+m_\rho^2}\,\epsilon^\gamma\cdot\epsilon^\rho\int d^4x'\; T_A(x')e^{i(p-q)\cdot x'}, \tag{9}$$

in which p is the momentum of the final vector meson, and $T_A(x)$ represents the purely imaginary final-state interaction with the target nucleus which converts the intermediate vector meson to the final vector meson.

The nuclear-coherent term of Fig. 1c takes the form

$$\mathcal{M}_\sigma = Q^2F_V(Q^2)\epsilon^\gamma\cdot\epsilon^\rho\frac{g_{\rho\gamma\sigma}}{m_\sigma}\int d^4x\int\; d^4x'\int\frac{d^4l}{(2\pi)^4}e^{ip\cdot x'}T_A(x')\frac{e^{-il\cdot(x'-x)}}{l^2-m_\rho^2+i\epsilon}e^{-iq\cdot x}\sigma(x). \tag{10}$$

The evaluation is straightforward. The integration over d^4x gives a delta function setting l equal to $k+q$. The propagator contains a factor $(k+q)^2 - m_\rho^2 + i\epsilon \approx -2\nu k^3 - Q^2 - m_\rho^2 + i\epsilon$ in which the direction of the photon momentum is taken as the positive $z(3)$ axis. The approximation of neglecting $k^2 (\ll m_\rho^2)$ yields a propagator of the eikonal form, so that after integration

$$\mathcal{M}_\sigma = \frac{-iQ^2 F_V(Q^2)}{2\nu} \epsilon^\gamma \cdot \epsilon^\rho \frac{g_{\rho\gamma\sigma}}{m_\sigma} \int d^4x' e^{i(p-q)\cdot x'} T_A(x') \int_{-\infty}^{z'} dz e^{-i\frac{Q^2+m_\rho^2}{2\nu}(z'-z)} \sigma(\mathbf{x}'_\perp, z) \quad (11)$$

The overall factor of i arises from the eikonal propagator, and another phase appears in the exponential. This phase factor vanishes for extremely large values of ν, but it is important here because the large nuclear radius enters ($z'-z \sim R_A$). Thus $\mathcal{M}_\sigma$ can interfere with $\mathcal{M}_{\rm dom}$. Note that for large enough values of ν, $\mathcal{M}_\sigma \propto (Q^2 + m_\rho^2/2\nu) R_A$. This gives a rough guide to the dependence on Q^2 and A. The amplitude $\mathcal{M}_\sigma$ is proportional to the central value of the nuclear σ field $\sigma(0)$ which takes on a value of -43 MeV in QHD1. If the non-linear model [14] is used, $\sigma(0) = -29$ MeV. In our present calculations we choose the average value $\sigma(0) = -36$ MeV. The results shown in Fig. 3 are obtained by numerical integration and assuming that the time dependence of $T_A(x')$ is the same for all positions within the nucleus. The coupling constants and form factors are taken as $g_{\rho\gamma\sigma}$ =5, 6, and $F_V(Q^2) = \exp\left[-(Q^2 - m_\rho^2) R_V^2/6\right]$, with $R_V = 0.99$ fm. An exponential, rather than power-law fall-off, is essential in reproducing the data. This form factor implies a strong Q^2 dependence at fixed values of x. The parameter R_V takes on the value of a typical hadronic size.

How large should $g_{\rho\gamma\sigma}$ be? We can compare the strength of the $\gamma \to \rho\sigma$ transition with that of $\gamma \to \rho$ at $Q^2 = m_\rho^2$: $g_{\rho\gamma\sigma} s\sigma/m_\sigma$ *vs.* e/f_ρ. The factor σ is expected from the σ model to be of the order of $M/g_{\pi NN} \approx 70$ MeV. Thus the couplings are comparable if $g_{\rho\gamma\sigma} \approx 0.5$. This is about ten times smaller than the coupling needed to describe the strong shadowing observed by HERMES. However, $g_{\rho\gamma\sigma}$ =5 is not ruled out by any existing data. On the other hand, if $g_{\rho\gamma\sigma} = 0.5$, $\sigma_T(A)/\sigma_T(D) \approx 0.95$ for $x \sim 0.01$. The nuclear enhancement of R is obtained from computing the ratio of the results of Eqs. (4,11). This is shown in Fig. 4, where it is seen one has a reasonably good description of the data.

Data also exist for the ^{3}He target. To address these data properly, one should perform a three-body calculation. We do not attempt this here. To understand if our theory has a reasonable dependence on A, we simply rescale the radius parameter R

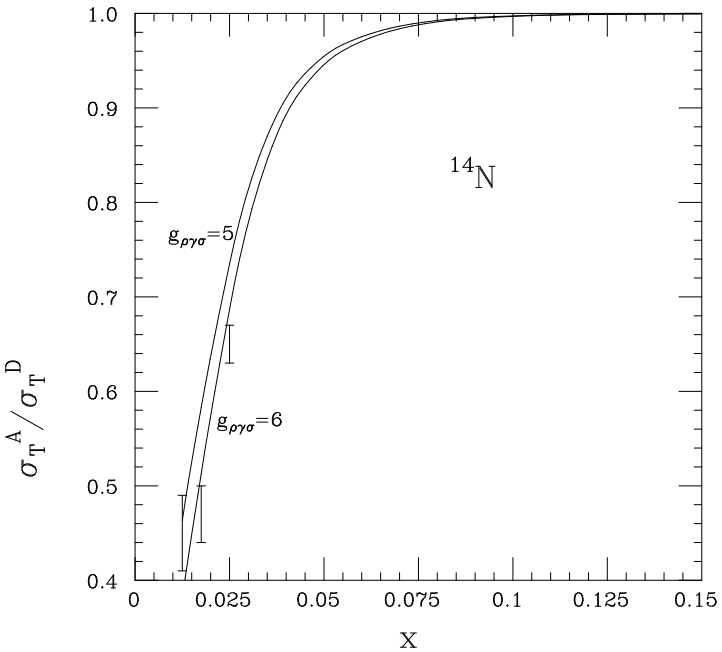

Figure 3: $\sigma_T(A)/\sigma_T(D)$, A=14, data of Ref. [1]

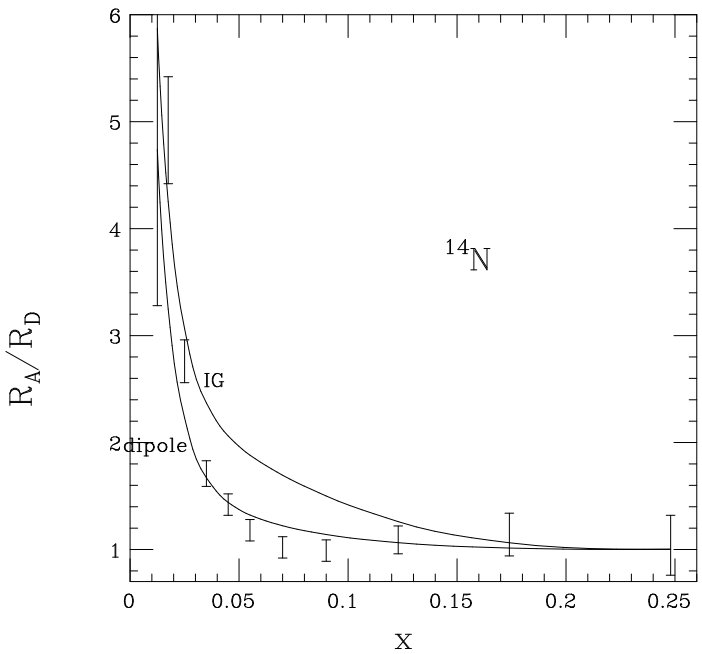

Figure 4: $R(A)/R(D)$, A=14

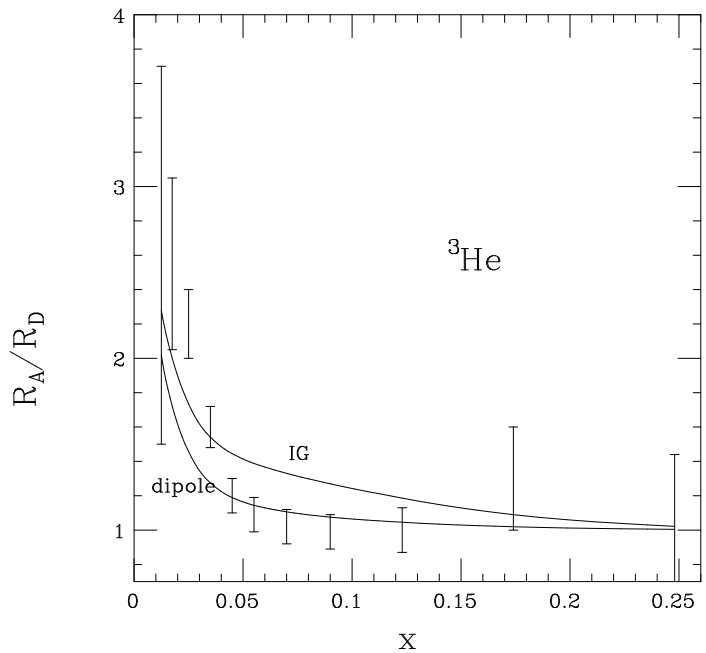

Figure 5: $R(A)/R(D)$, A=3, data from Ref. [1]

by a factor $(3/14)^{1/3}$ and take the number of nucleons to be 3. The result, as shown in Fig. 5, is in qualitative agreement with the HERMES ^{3}He data.

Our analysis shows that it is possible, with reasonable coupling strengths, to reproduce the salient features of the HERMES data. The calculations employ a particular choice of couplings, optimized to reproduce the HERMES, without contradicting other experimental constraints. For example, the strengths of the meson fields in nuclear medium which we use are at least roughly consistent with measured nuclear binding energies, nuclear densities, and nuclear deep inelastic and Drell-Yan data. However, the constants $g, g_{\rho\gamma\sigma}$ have never been measured, and their magnitudes could turn out to be small. The increase in the longitudinal cross section is readily explained by the exclusive process $eA \to e'\sigma A$ via ω exchange. The nuclear enhancement follows from the coherence of the ω field. The strong shadowing of σ_T at small Q^2 requires a special choice of the effective Lagrangian which is theoretically and empirically allowed, but does not seem to follow from any general principle. Thus, more conservatively, we cannot rule out the HERMES data using our theory.

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 σ 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 $Q^2 \sim 1$ 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.

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