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VECTOR-MESON EFFECTS IN THE ELECTRIC FORM-FACTOR OF ^{3}He *

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ABSTRACT

Assuming vector meson dominance of ^{3}He elastic electric scattering, we calculate double impulse "shadowing" corrections, similar to those discussed for deuterium by Gunion and Blankenbecler. The corrections are too small to account for the high momentum transfer structure of the form factor. We then discuss the possibility of double-counting of the Born terms for these processes and find that even using this possibility, the corrections cannot simultaneously account for the form factor dip at 11.8 f^{-2} and the behavior at higher momentum transfer. The tentative conclusion reached is that the high momentum transfer behavior ($10\ f^{-2} < \Delta^2 < 28\ f^{-2}$) of the ^{3}He electric form factor is caused by high momentum components of the wave function generated directly by the strong interactions, rather than being of electromagnetic origin.

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I. INTRODUCTION

As of the summer of 1970 the experimentally measured electric form factor of ^{3}He, as a function of momentum transfer, was in reasonable qualitative agreement with the predictions of most of the competing theories of the nucleon-nucleon interaction[2] and could be fit to assuming rather simple phenomenological forms for the ^{3}He wave function.[1] At that time, McCarthy et al.[3] reported the results of their measurement of this form factor in the region $\Delta^2 = 8\ f^{-2}$ to $\Delta^2 = 20\ f^{-2}$, where Δ is the four-momentum of the incident virtual photon. This region had not been included in the earlier experiment of Collard et al.[4] They found a diffraction dip with a minimum near $\Delta^2 = 11.8\ f^{-2}$, then a form factor, the absolute value of which then rises rapidly with Δ^2 just above $\Delta^2 = 11.8\ f^{-2}$ becoming relatively flat near a value $|F^{e\ell}| = 5 \times 10^{-3}$ from $\Delta^2 = 14\ f^{-2}$ to the maximum Δ^2 of the experiment; $\Delta^2 = 20\ f^{-2}$ (Fig. 1).

Several researchers[5,6] have been unable to fit this data with a realistic two-nucleon potential, with or without a hard core. These attempted fits have the undesirable features that:

1. for the correct position of the diffraction dip, the binding energy is too small.
2. Attempts to adjust the binding energy move the diffraction dip out by about $3\ f^{-2}$.
3. In any case, the amplitude of the high momentum transfer tail of the electric form factor is an order of magnitude too small.

On the other hand, they all give reasonable fits to the form factor for $\Delta^2 \leq 8\ f^{-2}$.

II. THE SHADOWING OR STRONG INTERACTION CORRECTION

Gunion and Blankenbecler[7] have noted that if we assume vector dominance for the interaction of a deuteron with a virtual photon, then the vector meson, which couples to the photon, can scatter strongly from one of the nucleons and be absorbed by the other. With the assumption that for the deuteron to remain bound each nucleon receives approximately equal momentum transfer, and assuming a simple phenomenological form for the momentum transfer dependence of the vector meson nucleon scattering amplitude, they are able to approximate the integrals involved in the calculation of this contribution to the electric and magnetic form factors of deuterium.

With a further assumption on the relative strength of the $t=0$ [8] ρ and ω photoproduction amplitudes (see their paper for details, by this we mean their condition $f_1=f_2$), they obtain the normalization of the vector meson-nucleon scattering amplitude by requiring that this shadowing process account for the correction required to the value of the deuteron magnetic moment calculated using a Partovi wave function. They then note that the Partovi or other "good" wave functions may predict a deuteron electric form factor which is a bit too low in the region of momentum transfer; $\Delta^2 \geq 24\ f^{-2}$ and which falls off somewhat too rapidly in this region. The suggestion is then made that a slight relaxation of their requirement $f_1=f_2$ can give a less rapidly falling form factor which could be in better agreement with experiment than is that calculated by the single impulse approximation.

They finally suggest that similar double and triple impulse diagrams in ^3He (Fig. 2) could be an explanation for the dip and tail in the ^3He elastic form factor which was discussed above.

In Section IV we will undertake to give an estimate of the size of such effects. Let us first made a few general remarks. First, we expect a substantial contribution to the ^{3}He magnetic form factor from mesonic exchange currents and the like. (See, for example, the review article of Delves and Philips.[2]) Therefore, we shall have no more to say about the magnetic form factor in this paper. Next, the single impulse approximation to the electric form factor; $F_1^{e\ell}(\Delta^2)$, can be expressed as a product of a sum of single nucleon form factors and a body form factor $F_B(\Delta^2)$, where Δ is the momentum transfer and $F_B(\Delta^2)$ is the Fourier transform of the square of the ^{3}He wave function.[7] Now in the double impulse contribution; $F_2^{e\ell}(\Delta^2)$, in order to maintain a bound state each impulsed nucleon should receive a momentum transfer of about $\Delta/2$. This is equivalent to keeping the two struck nucleons fixed and giving the third an impulse of $-\Delta/2$. Therefore the body form factor that appears in $F_2^{e\ell}(\Delta^2)$ will be centered on $F_B(\Delta^2/4)$ (likewise the triple impulse contribution will depend on $F_B(2\Delta^2/9)$). As a consequence of this, in the tail region from Δ^2=12 f^{-2} to Δ^2=20 f^{-2}, where $\Delta^2/4 \leq 5$ f^{-2} the wave function factor in $F_2^{e\ell}(\Delta^2)$ is from a region where several phenomenological models and most "realistic" potentials give a good fit to the form factor. We thus expect $F_2^{e\ell}(\Delta^2)$ to be almost model independent and, in particular, to be independent of any high momentum transfer structure that might appear in the model we use.

Our computation will be done with a Gaussian wave function $\psi(r_{12}, r_{23}, r_{13}) = A \exp\left(-\frac{1}{2}\alpha^2\left(r_{12}^2+r_{23}^2+r_{23}^2\right)\right)$[9, 10] where the r_{ij} are the inter-nucleon distances and A and α are constants. A being determined by the normalization condition that the integral of $|\psi^2|$ over the two independent particle position vectors in a given coordinate system be unity. Since we are interested in the magnitude of a correction, we will not include the spacially antisymmetric

S' state and the D state contributions to the ^{3}He wave function as the sum of these contributors perhaps 10% of the total wave function normalization.[2]

In the next section possible criticisms of the Gunion-Blankenbecler contributions and further correction to $F_1^{e\ell}(\Delta^2)$ are discussed.

III. FURTHER CORRECTIONS

Consider the Born terms of Fig. 2 (Fig. 3). If we had a wave function that was fully relativistic and included exactly many-body effects then, for example, that part of Fig. 3a to the left of the broken line would already be included in the Bethe-Salpeter iteration of the wave function. Thus the Born terms would be included in the single impulse approximation and should thus be subtracted off from the full vector scattering amplitude of Fig. 2 to avoid double counting. Indeed, the difference between the full vector meson scattering amplitude and its Born terms includes such relativistic effects as production of resonances and multiparticle states along one of the nucleon lines, followed by their vector decay.

However, the wave functions, that we use, approximate solutions of the Schroedinger equation with static nonrelativistic potentials that describe low energy nucleon-nucleon scattering. In this view, we look at the Born term for vector meson-nucleon scattering and the other diagrams that contribution to the full scattering amplitude as additional pieces in the form factor, which are not included in the potential generated wave function. Instead they are due to many-body and other relativistic effects (i.e., distortion of the ^{3}He wave function by the vector meson field). This means that, depending on the details of the potential, there is at most some double-counting. To the extent that the potential contains nonrelativistic reductions of such effects as intermediate particle creation and absorption we may still have a degree of double-counting which, possibly, varies with energy.

In the latter part of the next section we calculate the contribution of the Born terms (Fig. 3) $F_2^{e\ell B}(\Delta^2)$ to the electric form factor in the momentum transfer region $\Delta^2 = 8\ f^{-2}$ to $\Delta^2 = 20\ f^{-2}$, again using a Gaussian wave function. We then multiply $F_2^{e\ell B}(\Delta^2)$ by a factor $-f$ representing the double-counting. This factor is taken to be constant throughout the momentum transfer range and is expected to lie in the range zero to one. An upper limit for f is estimated from the data. Calculations of pion photoproduction from nucleons indicate that we should not be surprised if $|F_2^{e\ell B}(\Delta^2)| \gg |F_2^{e\ell}(\Delta^2)|$.[11]

IV. DETAILS OF THE CALCULATION

We will now explicitly compute the contributions to the elastic form factor of ^{3}He of the diagrams of Fig. 2. The energy transfer for which the helium nucleus remains bound is given by $\Delta_0 = \Delta^2/6M$; thus for $\Delta^2 \ll 27\ f^{-2}$ we can take $\Delta^2 = \vec{\Delta}^2$. Since we expect approximately half the incoming momentum Δ to be transferred to each impacted nucleon[7] we represent the momentum of the exchanged vector meson as $\vec{\Delta}/2 - \vec{\delta}$ where we integrate over all $\vec{\delta}$ but anticipate important contributions only when $|\vec{\delta}|$ is small.

Next we give our approximation for the vertices (ignoring magnetic effects). The ρ-nucleon vertex is given by $G_p(q^2)\vec{\tau}_N$ where $\vec{\tau}_N$ is the nucleon isospin and the isovector form factor is

$$F_V(q^2) = g_\rho \frac{1}{m_\rho^2 + q^2} G_\rho(q^2) \tag{4.1a}$$

Here g_ρ is the photon-vector meson vertex constant. Likewise for the ω-nucleon vertex, $G_\omega(q^2)$ is defined by the isoscalar form factor

$$F_S(q^2) = g_\omega \frac{1}{m_\omega^2 + q^2} G_\omega(q^2) \tag{4.1b}$$

The vector meson-nucleon scattering amplitudes are assumed to be spin and energy independent and of the form[12]

$$A_{\omega\omega}(t) = 4\pi a\, e^{+4.3\,t} \tag{4.2a}$$

$$A_{\rho\omega}(t) = 4\pi b\, e^{+4.3\,t} \tag{4.2b}$$

$$A_{\rho\rho}(t) = 4\pi c\, e^{+4.3\,t} \tag{4.2c}$$

The results of Gunion and Blankenbecler are consistent with SU(3) D-coupling for the vector meson-nucleon scattering amplitudes. This gives $b = \frac{1}{\sqrt{3}} a$ and $c = \frac{1}{3} a$. We combine this with their result $a = -.12$ f to get our vertices. Finally, the photon-vector meson vertices are given by $g_\omega = m_\omega^2/2\gamma_\omega$ and $g_\rho = m_\rho^2/2\gamma_\rho$. From SU(3) $\gamma_\omega = \sqrt{3}\gamma_\rho$ and we take $\gamma_\rho^2/4\pi = \frac{1}{2}$.[13] In all our computations we will ignore the ρ-ω mass difference and take $m_\omega^2 = m_\rho^2 = m_V^2 = 14.9\ f^{-2}$.

Evaluating the isospin factors of the vector exchange matrix elements using the fully antisymmetric spin-isospin wave function of Schiff[4] we obtain the contribution of the double impulse diagram to the ^{3}He electric form factor:

$$F_2^{e\ell}(\Delta^2) = 4\pi a \int \frac{d^3\delta}{(2\pi)^3} \left(\frac{1}{\Delta^2 + m_V^2}\right) \times \exp\left(-4.3\left(\frac{\vec{\Delta}}{2} + \vec{\delta}\right)^2\right) F_B(q_r, q_\rho)$$

$$\times \left\{\frac{2}{3} F_V\left(\left(\frac{\vec{\Delta}}{2} - \vec{\delta}\right)^2\right) - 4F_S\left(\left(\frac{\vec{\Delta}}{2} - \vec{\delta}\right)^2\right)\right\} \tag{4.3}$$

where the Jacobi momenta are $\vec{q}_r = -\frac{\vec{\Delta}}{4} + \frac{\vec{\delta}}{2}$, $\vec{q}_\rho = \frac{\vec{\Delta}}{6} + \vec{\delta}$ and $F_B(\vec{q}_r, \vec{q}_\rho)$ the body form factor is determined in terms of the wave function $\psi(\vec{r}, \vec{\rho})$ (expressed as a function of the Jacobi coordinates) as:

$$F_B(\vec{q}_r, \vec{q}_\rho) = \int d^3\vec{r} \int d^3\vec{\rho}\, e^{-i\vec{q}_r \cdot \vec{r}}\, e^{-i\vec{q}_\rho \cdot \vec{\rho}}\, |\psi(\vec{r}, \vec{\rho})|^2 \tag{4.4}$$

We have chosen a Gaussian wave function which gives a good fit to the data out to $\Delta^2 = 8\ f^{-2}$. In this case (4.4) gives $F_B(\vec{q}_r, \vec{q}_\rho) = \exp\ -(q_r^2/6\alpha^2) + (q_\rho^2/8\alpha^2)$ with $\alpha^2 = .14 f^{-2}$. F_S and F_V are taken from Ref. 14. The integral (4.3) can be evaluated if we approximate $F_{S,V}\left(\left(\frac{\Delta}{2} - \delta\right)^2\right) = F_{S,V}(\Delta^2/4)$. This yields

$$F_2^{e\ell} = +\ 4.1 \times 10^{-3}\ (1/(1+\Delta^2/14.9)) \cdot \Big(-1.665 + 5.42/(1+\Delta^2/59.2)$$
$$- .55/(1 + \Delta^2/150.4)\Big) \cdot \exp\left(\frac{-7\Delta^2}{20}\right) \tag{4.5}$$

We note that since $F_2^{e\ell}(0) = +\ 4.1 \times 10^{-3}$, in practice we need not renormalize the form factor to five $F_1^{e\ell}(0) + F_2^{e\ell}(0) = 1$. At $\Delta^2 = 5 f^{-2}$, $F_2^{e\ell}(5) = +\ 4 \times 10^{-3}$ compared with the Gaussian single impulse form factor of 10^{-1}. By $\Delta^2 = 10 f^{-2}$ we have $F_2^{e\ell} = +\ 2.3 \times 10^{-4}$ which is about 3% of the measured value of $F^{e\ell} = 6.5 \times 10^{-3}$. $|F_2^{e\ell}(\Delta^2)|$ then continues to fall off exponentially. It, therefore, cannot begin to explain the relatively flat tail in the data from $\Delta^2 = 14 f^{-2}$ out to $\Delta^2 = 20\ f^{-2}$ where $|F^{e\ell}| \simeq 5 \times 10^{-3}$. In fact, $F_1^{e\ell}(\Delta^2)$; the single impulse contribution for the Gaussian is given by (see Table 1 for numerical values):

$$F_1^{e\ell}(\Delta^2) = \left(\frac{3}{2}\, F_S(\Delta^2) + \frac{1}{2}\, F_V(\Delta^2)\right) \exp\ (-8\ \Delta^2/20) \tag{4.6}$$

So $F_1^{e\ell}(\Delta^2)$ falls off only slightly faster than $F_2^{e\ell}(\Delta^2)$ (see Eq. (4.5)). Rough estimates show that the triple impulse contribution; $F_3^{e\ell}(\Delta^2)$ is about 10% of $F_2^{e\ell}(\Delta^2)$ and due to the rapid fall off of the two vector meson-nucleon scatterings present will fall off quite rapidly. Thus the corrections to the single impulse diagrams suggested by Gunion and Blankenbecler[7] can explain neither the tail nor the dip in the observed ^{3}He electric form factor.

Next we estimate the effects due to the possible partial double-counting of the Born terms for intermediate vector meson exchange (Fig. 3). By taking a nonrelativistic limit for the nucleon spinors and the intermediate propagators we can derive the nucleon vector-meson scattering vertex in Born approximation. The kinematics are chosen so as to agree with the computation of the first part of this section. Consider the diagram of Fig. 3a, ignoring magnetic effects, the amplitude in the static limit for the absorption of an isoscalar photon and the omission of a is given by

$$A_{\gamma\rho}^{B_1}(\Delta^2) = \mu^+(p_1') \frac{(-i(\not{p}_1+\not{\Delta}) + M_N)}{(p_1+\Delta)^2 + M_N^2} \mu(p_1)\ F_S(\Delta^2) G_V \left(\frac{\vec{\Delta}}{2} + \vec{\delta}\right) \vec{\tau}_N \tag{4.7}$$

where p_1 is the four-momentum of the impacted nucleon, Δ is that of the photon, q_ρ if that of the ρ and $p_1' = p_1 + \Delta - q_\rho$. With the approximation $p_{10} = M_N$ and using $\Delta_0 = \Delta^2/6M_N$ and $p_1^2 = -M_N^2$ we obtain

$$A_{\gamma\rho}^{B_1}(\Delta^2) = X_1^+ \frac{2M_N}{2\Delta^2/3} X_1\ F_S(\Delta^2) G_V \left(\frac{\vec{\Delta}}{2} + \vec{\delta}\right) \vec{\tau}_N \tag{4.8}$$

with X_1 the spinor part of the ^{3}He wave function for the first nucleon. Similar expressions hold for the diagrams of Fig. 3b and for the scatterings involving the other combinations of vector mesons.

Equation (4.8) no longer has any p_1 dependence therefore the Born term vector meson-nucleon scattering amplitude is effectively local. Also, the exponential form of the full nucleon vector meson scattering amplitude is replaced by a rational function of Δ^2. Therefore, we expect this term $F_2^{e\ell B}(\Delta^2)$ to fall off like $F_1^{e\ell}(\Delta^2/4)$ which is more slowly than $F_2^{e\ell}(\Delta^2)$. The term (4.8) diverges at Δ^2=0. However, as in the case of photon bremsstrahlung,[16] the divergence is cancelled by vertex and self interactions (here the binding energy of the ^{3}He nucleus plays the same role as the finite energy resolution of the detector in bremsstrahlung in providing a cutoff which makes each term finite). The terms that cancel the divergence at $\Delta^2 = 0$ involve interactions with single nucleon line and thus fall off like $F_1^{e\ell}(\Delta^2)$. So in the Δ^2 region of interest they should be only a few percent of the Born terms and we will ignore them.[14] The expression for the Born term contribution to the electric form factor is then of the form

$$F_2^{e\ell B}(\Delta^2) = \sum_{\substack{\text{all vector}\\ \text{mesons}}} \int \frac{d^3\delta}{(2\pi)^3} \left\{ \left(\left(\frac{\vec{\Delta}}{2} + \vec{\delta} \right)^2 + M_N^2 \right) \cdot F_{S,V}(\Delta^2) \cdot F_{S,V}^2 \left(\left(\frac{\vec{\Delta}}{2} - \vec{\delta} \right)^2 \right) \cdot 2M_N \right.$$

$$\left. \times \exp\left(-\Delta^2/72\,\alpha^2 - \delta^2/6\,\alpha^2\right) \times 1/\left(g_{\rho,\omega}^2 \cdot \frac{5\Delta^2}{12} \right) \right\} \qquad (4.9)$$

The handling of the isospin is as in the case of the full scattering amplitude. As above we ignore the δ dependence of F^V and F^S. This gives an approximation for $F_2^{e\ell B}(\Delta^2)$

$$F_2^{e\ell B}(\Delta^2) = \frac{16}{\sqrt{\pi}} \; \frac{M_N}{M_V^4} \left(6\alpha^2\right)^{3/2} \left(\frac{1}{4} + \frac{\left(M_V^2 + 9\alpha^2\right)}{A^2} \right) \times$$

$$\times \left(\left(F_V(\Delta^2) - 3\,F_S(\Delta^2)\right) \cdot F_V^2\left(\frac{\Delta^2}{4}\right) + 3\left(F_V(\Delta^2) + 3\,F_S(\Delta^2)\right) \cdot F_S\left(\frac{\Delta^2}{4}\right) \right) \times$$

$$\times \exp\left(-\Delta^2/72\,\alpha^2\right) \qquad (4.10)$$

First we note that $F_1^{e\ell} + F_2^{e\ell} - F_2^{e\ell B}$ fits the data at $\Delta^2 = 14\ f^{-2}$. However, in this case $F_1^{e\ell} + F_2^{e\ell} - F_2^{e\ell B}$ shows no dip near $\Delta^2 = 12$. In addition at $\Delta^2 = 20$, $\left(F_1^{e\ell} + F_2^{e\ell} - F_2^{e\ell B}\right)^2 = 10^{-6}$ which is much too small. Therefore, corrections of the form $F_2^{e\ell B}$ if adjusted to fit the magnitude of the tail of the electric form factor, cannot fit its shape and, in addition, move the dip much too far in. Next let us try to adjust f so as to fit the dip, which for computational convenience we take at $\Delta^2 = 12\ f^{-2}$. If we assume $F^{e\ell}(12) = 0$, we get $f F_2^{e\ell B} = F_1^{e\ell}(12) + F_2^{e\ell}(12)$. This gives f = .338. We summarize our numerical results for $F_1^{e\ell}$, $F_2^{e\ell}$ and $F_2^{e\ell B}$ in Table 1. With this crude attempt to fit the dip $|F_1^{e\ell} + F_2^{e\ell} - f F_2^{e\ell B}|$ is no greater than a few percent of $|F^{e\ell}|$ in the tail region from $\Delta^2 = 14\ f^{-2}$ to $\Delta^2 = 20\ f^{-2}$. Thus to avoid too drastic a dip in the form factor we have, at most, about one-third double-counting of the Born term. On the other hand, with, say, one third double-counting the calculated tail of the electric form factor is at least an order of magnitude too small. Our results are given in Fig. 1 for f=1 and for f=0.338.

V. CONCLUSIONS AND SUMMARY

The corrections to the ^{3}He electric form factor, calculated from phenomenological nonrelativistic nucleon pair potentials can be divided into two categories. The first is due to corrections to the internal dynamics of the three-body system. This includes virtual elementary particle production resonance excitation etc.; also included here are any explicit three-body effects that are present. All these effects occur in the isolated three nuclear system regardless of how we probe it. The other category we shall call electromagnetic corrections. These come about as the incident virtual photon can induce processes in the target nucleus which are not present in the isolated three-nucleon system. In fact, to the extent that the (virtual) photon in hadronic electromagnetic interactions is dominated by vector meson poles, the wave function for ^{3}He will be distorted by the presence of a fourth particle the interactions of which with the constituent nucleons are as strong as those among the nucleons themselves. Let us discuss the meaning of the results of our computation of the simplest electromagnetic effects, namely the shadowing effect discussed in Section II.

In the region of interest $\Delta^2 = 10\ f^{-2}$ to $\Delta^2 = 20\ f^{-2}$ the electromagnetic correction is expected to have little dependence on the internal dynamics correction. The reason for this is that both the full scattering corrections and any Born term subtractions at momentum transfer Δ^2 show a dependence on the square of the ^3He wave function which is strongest in a region near $\Delta^2/4$, which is less than $5\ f^{-2}$ here. In this region the ^3He electric form factor is well described by a wave function calculated from phenomenological two nucleon potentials. Therefore we should be able to calculate the electromagnetic corrections using a simple uncorrected ^3He wave function such as the Gaussian we have used.

The contribution of processes similar to those described by Gunion and Blankenbecler for deuterium seem of insufficient magnitude to describe the structure of the ^3He electric form factor in the region $\Delta^2 = 10\ f^{-2}$ to $\Delta^2 = 20\ f^{-2}$. In addition, $F_2^{e\ell}(\Delta^2)$ shows a too rapid fall off in Δ^2 to fit the tail of the form factor. This fall off of at least one quarter as fast as the single impulse contribution (on a semi-log plot) can be seen from the details of our calculation to apply to double impulse corrections in general with reasonable assumptions about the behavior of the vertex functions included, therefore, it is unlikely that the inclusion of other processes similar in form to those generating $F_2^{e\ell}(\Delta^2)$ can explain the high momentum transfer behavior of the form factor which is the most salient feature of the results of McCarthy et al.[3]

The possibility of double-counting the Born term of $F_2^{e\ell}$ led us to consider $F_2^{e\ell B}$. Indeed, it is much larger than $F_2^{e\ell}$.[11] But if we attempt to fit the tail of the electric form factor by it, we find that first, the shape is not right. $F_2^{e\ell B}$

falls off too rapidly and second, the correction term is so large in the region $10\ f^{-2} \leq \Delta^2 \leq 14\ f^{-2}$ as to render the intermediate momentum transfer to behavior of the form factor completely incorrect. On the other hand, if we have about 34 percent double-counting, the subtraction of this percentage of $F_2^{e\ell B}$ gives a qualitatively reasonable fit to the diffraction minimum but yields much too small a form factor in the tail region with the wrong shape (even if we assume $F_2^{e\ell B}$ is the only contribution to the electric form factor here). This result is qualitatively similar to that of Ref. 6, which is not surprising as vector meson exchange can be used to generate a core. Therefore the most we can say for the Born term subtraction and electromagnetic corrections, in general, is that they may be important near the diffraction minimum.

Our main conclusion is then that the large amplitude of the high momentum transfer (small distance) part of the ^3He electric form factor and its constancy of shape will have to be explained by those features of the internal dynamics of the two and three nucleon systems which generate the large momentum components in the ^3He wave function and not in the nature of the interaction with the virtual photon. In addition at very small distances it may not even be reasonable to assume the ^3He is simply a bound state of three nucleons and we may have to take into account those parts of the wave function which are due to the binding, of say, a Δ and two nucleons and the like. We cannot now solve the full strong interaction problem so the task of explaining the results of McCarthy et al.[3] will probably be carried out by either (semi) phenomenologically treating the short distance part of the ^3He wave function or by devising two and three nucleon potentials to do this.

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FOOTNOTES AND REFERENCES

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17. For this reason we limit our discussion of $F_2^{e\ell B}(\Delta^2)$ to the region $\Delta^2 \geq 8\,f^{-2}$.

FIGURE CAPTIONS

1. The square of the ^{3}He elastic form factor from McCarthy et al. With our calculation of $\left(F_1^{e\ell} + F_2^{e\ell} - f\, F_2^{e\ell B}\right)^2$ for f=1 and f=0.338.
2. The vector meson scattering correction of Gunion and Blankenbecler in ^{3}He.
3. The Born terms for the vector meson scattering correction in ^{3}He.

TABLE 1

The Results of Our Calculation of $F_1^{e\ell}$, $F_2^{e\ell}$, and $F_2^{e\ell B}$

Δ^2	$F_1^{e\ell}$	$F_2^{e\ell}$	$F_2^{e\ell B}$
8	1.86×10^{-2}	$+ 5.24 \times 10^{-4}$	2.92×10^{-2}
9	1.15×10^{-2}	$+ 3.45 \times 10^{-4}$	2.09×10^{-2}
10	7.09×10^{-3}	$+ 2.27 \times 10^{-4}$	1.52×10^{-2}
11	4.39×10^{-3}	$+ 1.50 \times 10^{-4}$	1.12×10^{-2}
12	2.73×10^{-3}	$+ 9.94 \times 10^{-5}$	8.36×10^{-3}
13	1.70×10^{-3}	$+ 6.59 \times 10^{-5}$	6.30×10^{-3}
14	1.06×10^{-3}	$+ 4.37 \times 10^{-5}$	4.79×10^{-3}
15	6.61×10^{-4}	$+ 2.90 \times 10^{-5}$	3.68×10^{-3}
16	4.14×10^{-4}	$+ 1.93 \times 10^{-5}$	2.84×10^{-3}
17	2.59×10^{-4}	$+ 1.29 \times 10^{-5}$	2.21×10^{-3}
18	1.63×10^{-4}	$+ 8.58 \times 10^{-6}$	1.73×10^{-3}
19	1.03×10^{-4}	$+ 5.73 \times 10^{-6}$	1.38×10^{-3}
20	6.47×10^{-5}	$+ 3.83 \times 10^{-6}$	1.08×10^{-3}

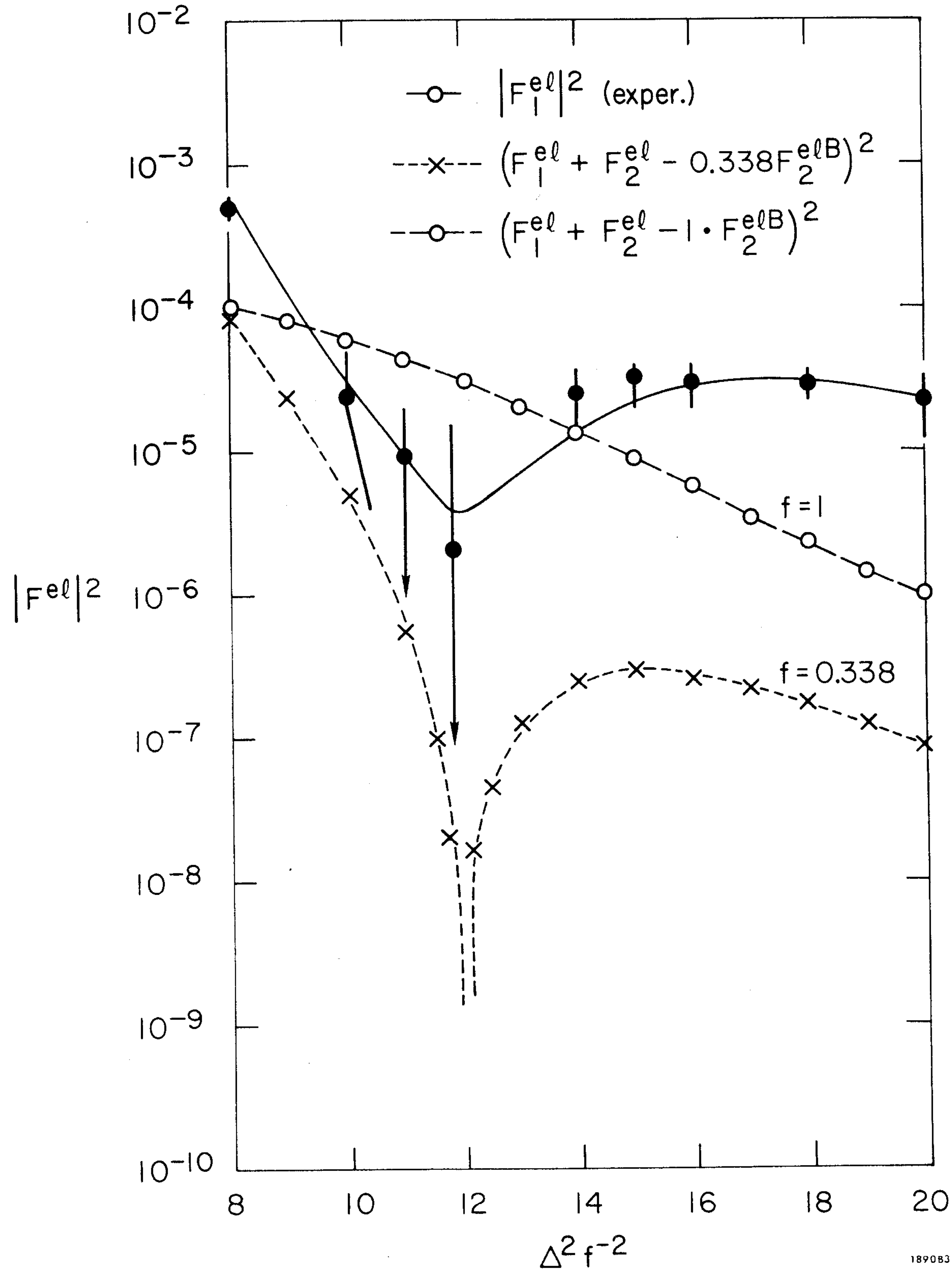
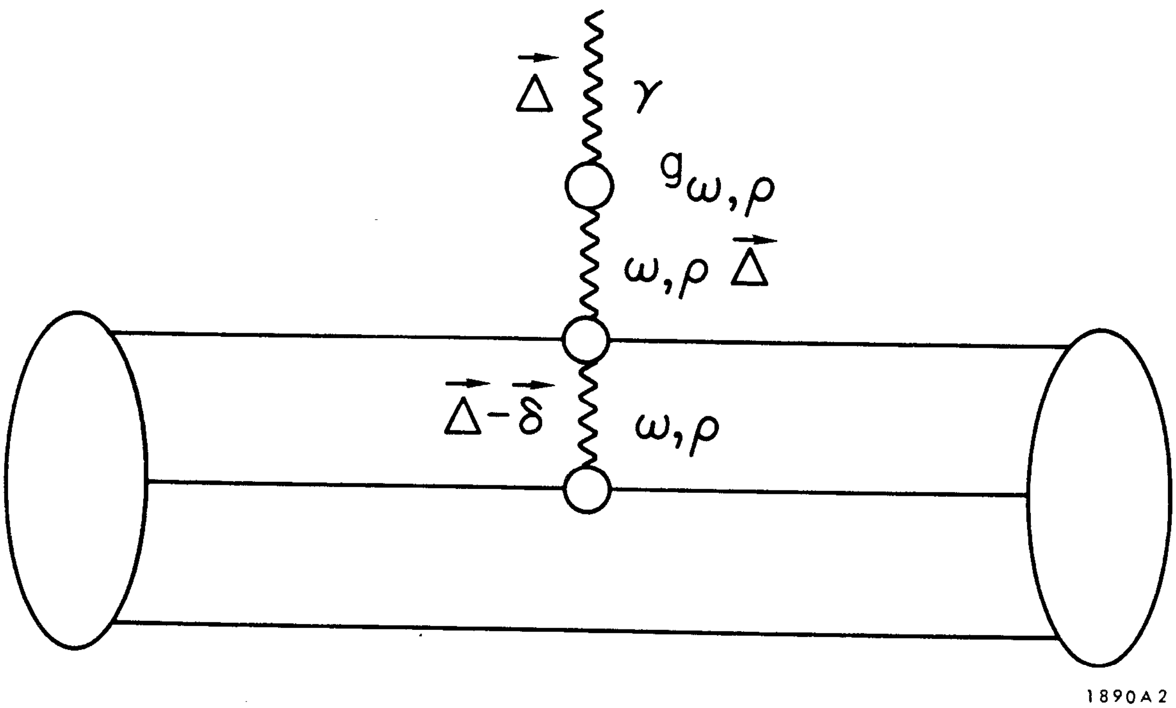
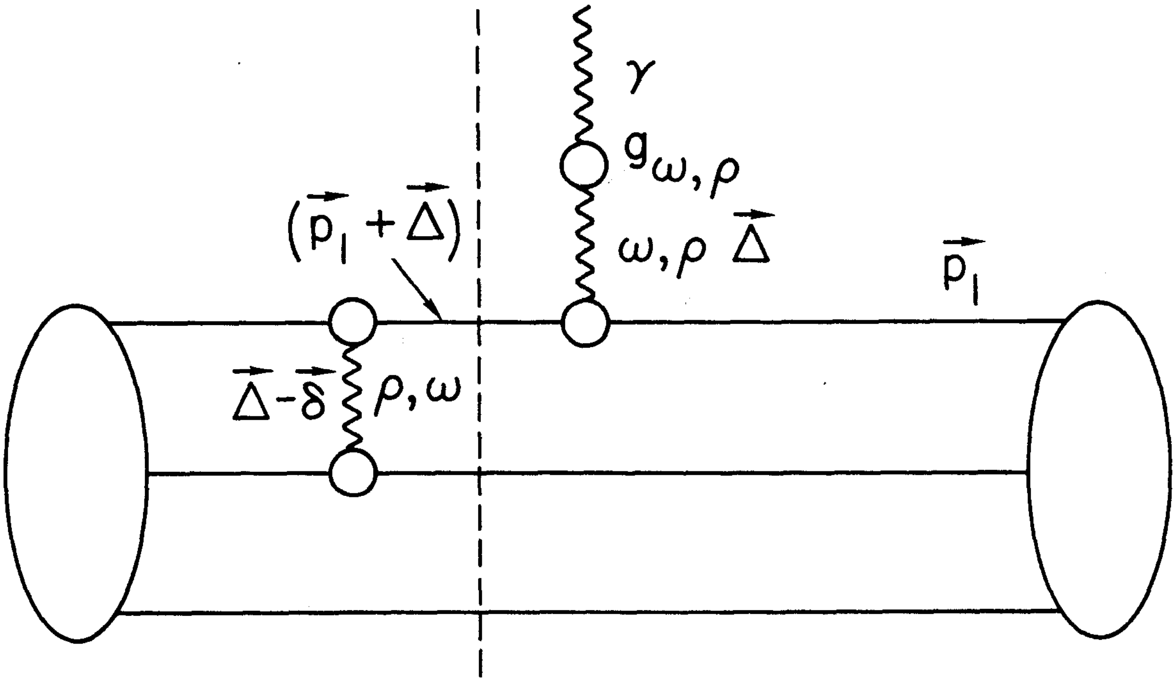

$|F_1^{el}|^2$ (exper.)
$(F_1^{el} + F_2^{el} - 0.338F_2^{elB})^2$
$(F_1^{el} + F_2^{el} - 1 \cdot F_2^{elB})^2$
$|F^{el}|^2$
10^{-2}
10^{-3}
10^{-4}
10^{-5}
10^{-6}
10^{-7}
10^{-8}
10^{-9}
10^{-10}
f = 1
f = 0.338
8
10
12
14
16
18
20
$\Delta^2 f^{-2}$
189083

Fig. 1

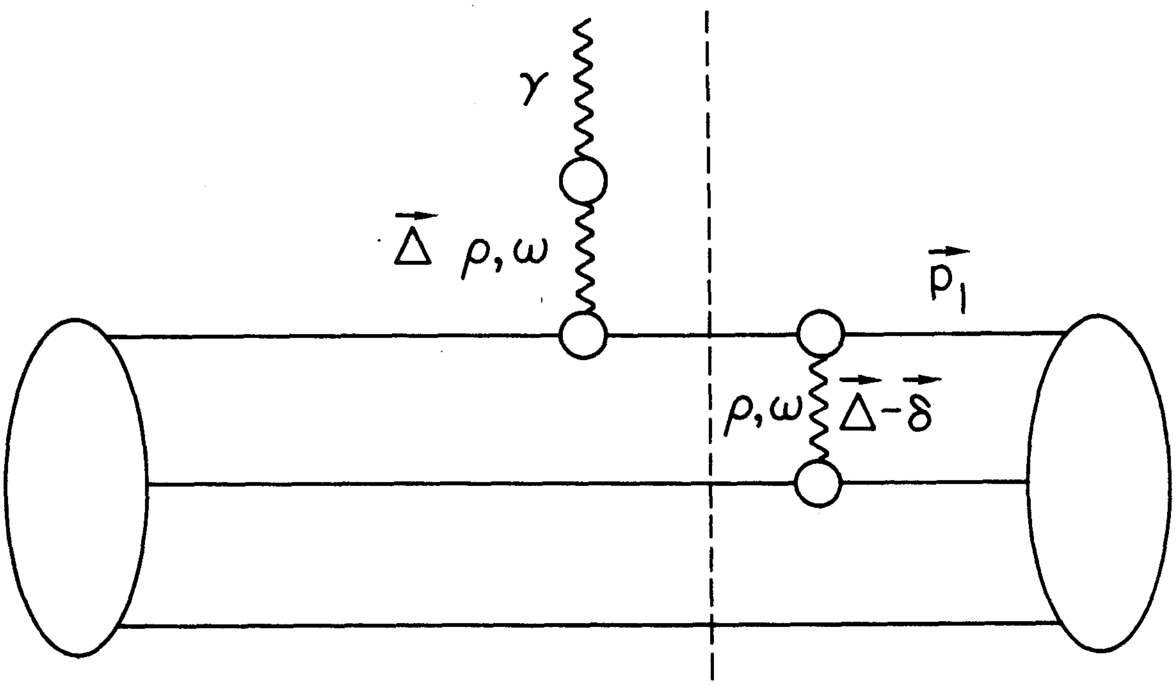


$\vec{\Delta}$
γ
$g_{\omega,\rho}$
ω,ρ $\vec{\Delta}$
$\vec{\Delta}-\vec{\delta}$
ω,ρ
1890A2

Fig. 2



(a)



(b)

1890A1

Fig. 3