Random-phase-approximation-type vertex corrections to the axial-vector current

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We study the isovector weak response of the core to an external perturbation using linear response theory in nuclear matter. We extend previous results to the kinematic region away from the long wavelength regime where the effects have been proven small. For the muon capture process we found a sizable quenching of the contribution of the induced pseudoscalar term. Applied to actual nuclei, however, the strong density dependence of the weak isovector response washes out the screening found at nuclear matter densities, slightly modifying muon capture rates.

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I. INTRODUCTION

It is well known [1] that the axial-vector current is considerably modified by hadronic processes when embedded in the nuclear medium. In particular, transitions driven by the zeroth component of the axial current appear to be strongly enhanced with respect to free space estimates (up to a factor of 2 for heavy nuclei [2]) when in the presence of other nucleons. In Ref. [3], as part of an investigation related to this enhancement in heavy nuclei, we calculated the isovector weak response of the core to the presence of a valence particle, and found that it was negligible for the small momentum transfers involved in nuclear β decay (in [3] we concentrated on first forbidden β transitions). (The contribution of meson-exchange currents within the same framework was studied in [4].)

Axial charge transitions, however, do not always take place at zero momentum transfer. In ordinary muon capture processes the four-momentum transfer q^2 is large ($\approx -0.9m_\mu^2$), and therefore it becomes of interest to investigate whether the weak response of the core generates any sizable contribution in this kinematic region. As it turns out, the response of the core, which we calculate in linear response theory within the context of relativistic hadrodynamics, differs appreciably from the zero momentum result. Modifications arise in the induced pseudoscalar term of the hadronic axial current which produce moderate to strong effects (with some model dependence) in the calculation of capture rates. In this paper we investigate the *isovector* weak response of the core to the external perturbation in the approximation of nuclear matter and in linear response theory summing the ring diagrams to all orders in the random phase approximation (RPA) —the so-called *backflow* correction [5,6]. We present results for the effects of this core polarization (of positive and negative energy nucleons) as a function of the momentum transfer and illustrate these results by calculating the effects expected in the transition

$$\mu^- + {}^{16}\mathrm{O}(0^+, \mathrm{g.s.}) \rightarrow {}^{16}\mathrm{N}(0^-, 120\ \mathrm{keV}) + \nu_\mu$$

using the local-density approximation (LDA). Earlier work on this same transition by Price and Walker [7] and by Nedjadi and Rook [8] utilized relativistic wave functions in finite nuclei to calculate the transition rate, but left unanswered the effect of the backflow corrections.

II. ISOVECTOR WEAK RESPONSE IN NUCLEAR MATTER

Our starting point will be the most general form of the hadronic axial current consistent with G invariance (absence of second-class currents) and pseudovector πNN coupling for the pion-pole dominance term,

$$\begin{aligned} J_{\mathrm{PV}}^{\mu 5\pm} &= \bar{\psi}(x)\left[g_A\gamma^\mu\gamma^5 - g_P\frac{q_\nu\gamma^\nu}{q^2-m_\pi^2}\gamma^5 q^\mu\right]\tau^\pm\psi(x) \\ &= \bar{\psi}(x)\Gamma_A(q)\psi(x). \end{aligned} \quad (2.1)$$

With Γ_A we indicate a generic matrix in spin-isospin space. Making $g_A = g_P$ in (2.1) the partial conservation of the axial current (PCAC) is straightforwardly satisfied, namely, $\lim_{m_\pi\to 0}\partial_\nu J^{\nu 5}(x) = 0$. This choice of the form factor is known as the Goldberger-Treiman value and differs somewhat from what is customarily defined as g_P in the literature (see, for example, Gmitro and Truöl [9]); they are related by

$$g_P^{\mathrm{here}} = g_P^{\mathrm{GT}}(-0.88m_\mu^2)\frac{m_\pi^2+m_\mu^2}{(m_p+m_n)m_\mu} = g_P^{\mathrm{GT}}\frac{1}{6.78}. \quad (2.2)$$

In free space, pseudoscalar or pseudovector coupling can be used indistinctly in the pion-pole term both satisfying the PCAC constraint. In the nuclear medium, however, pseudoscalar coupling has the disadvantage that, besides not satisfying PCAC at the level of operators, yields muon capture rates which are anomalously large [8].

To study the response of the core to the presence of a valence particle, a valence hole, or a particle-hole pair we start by calculating the matrix elements of a generic current operator J_{Γ_A} in the form

$$\begin{aligned} j_{\Gamma_A}(x) &= \langle\phi_{\rm final}|J_{\Gamma_A}(x)|\phi_{\rm initial}\rangle \\ &= \langle\phi_{\rm final}|\bar{\psi}(x)\Gamma_A\psi(x)|\phi_{\rm initial}\rangle \end{aligned}$$

between initial and final states that have in common the same core (the filled Fermi sphere) and may differ only by the addition of a particle (hereafter we shall consider only the valence-particle case) to an unoccupied state with the quantum numbers $(\mathbf{k}_i, s_i, \tau_i)$ for the initial and $(\mathbf{k}_f, s_f, \tau_f)$ for the final state. The extra particle gives rise to additional meson fields which modify the core. In nuclear matter, the addition of a nucleon to the Fermi sphere introduces changes in the single-particle core wave functions and consequently in their matrix elements. Thus, in the single-particle approximation, we write

$$|\phi_{\rm initial}\rangle = a^{\dagger}_{\mathbf{k}_i,s_i,\tau_i}|\phi_{\rm core}\rangle,$$

$$\langle\phi_{\rm final}| = \langle\phi_{\rm core}|a_{\mathbf{k}_f,s_f,\tau_f}, \tag{2.3}$$

and, therefore,

$$j_{\Gamma_A}(x) = \langle\phi_{\rm core}|a_{\mathbf{k}_f,s_f,\tau_f}\bar{\psi}(x)\Gamma_A(q)\psi(x)a^{\dagger}_{\mathbf{k}_i,s_i,\tau_i}|\phi_{\rm core}\rangle. \tag{2.4}$$

To calculate this we first write the operators corresponding to the creation or the annihilation of the added particle in terms of the field operators $\psi(x)$ and $\bar{\psi}(x)$,

$$a_{\mathbf{k}_f,s_f,\tau_f} = \int \frac{dx_f^3}{\sqrt{\Omega}}\,\bar{U}(\mathbf{k}_f, s_f, \tau_f)\gamma^0\psi(x)\exp\{ik_f x\}, \tag{2.5a}$$

$$a^{\dagger}_{\mathbf{k}_i,s_i,\tau_i} = \int \frac{dx_f^3}{\sqrt{\Omega}}\,\bar{\psi}(x)\gamma^0 U(\mathbf{k}_i, s_i, \tau_i)\exp\{-ik_i x\}, \tag{2.5b}$$

where $U(\mathbf{k}_i, s_i, \tau_i)$ and $\bar{U}(\mathbf{k}_f, s_f, \tau_f)$ are four spinors corresponding to solutions of the Dirac equation for the interacting baryon fields normalized such that

$$U^{\dagger}(\mathbf{k}, s, \tau)U(\mathbf{k}, s', \tau') = \delta_{s,s'}\delta_{\tau,\tau'}. \tag{2.6}$$

We work with a uniform system of A baryons in a box of volume Ω in the mean-field approximation in nuclear matter which means that the meson fields are substituted by their classical expectation values. The symmetries of the problem simplify things since, on the one hand, the expectation value of the pion field vanishes due to parity conservation and, on the other, rotational invariance around the $\hat{z}$ axis in isospin space makes the neutral component of the ρ field the only one to survive.

Substituting Eqs. (2.5) back in (2.4) and arbitrarily choosing initial and final times so as to introduce a time ordering we obtain

$$\begin{aligned} j_{\Gamma_A}(x) = \frac{1}{\Omega}\int dx_f^3\, dx_i^3\, &\exp\{ik_f x\}\,\exp\{-ik_i x\}\,\bar{U}_{\omega}(\mathbf{k}_f, s_f, \tau_f)\gamma^0_{\omega\alpha}\Gamma_{A\,\gamma\delta} \\ &\times\langle\phi_{\rm core}|T\left[\psi_{\alpha}(x_f)\bar{\psi}_{\gamma}(x)\psi_{\delta}(x)\bar{\psi}_{\beta}(x_i)\right]|\phi_{\rm core}\rangle\gamma^0_{\beta\epsilon}U_{\epsilon}(\mathbf{k}_i, s_i, \tau_i). \end{aligned} \tag{2.7}$$

This matrix element can be calculated using standard perturbation theory in terms of the unperturbed (by the extra particle) field operators. In the relativistic quantum hadrodynamical model QHD-II of Ref. [10] the interaction Hamiltonian reads

$$\mathcal{H}_I(\psi) = -\mathcal{L}_I(\psi) = g_{\omega}\bar{\psi}\gamma_{\mu}\omega^{\mu}\psi - g_{\sigma}\bar{\psi}\sigma\psi + \frac{1}{2}g_{\rho}\bar{\psi}\gamma_{\mu}\boldsymbol{\tau}\cdot\boldsymbol{\rho}^{\mu}\psi + \frac{1}{2}e\bar{\psi}\gamma_{\mu}A^{\mu}(1+\tau_3)\psi + ig_{\pi}\bar{\psi}\gamma^5\boldsymbol{\tau}\cdot\boldsymbol{\pi}\psi - \frac{1}{2}g_{\sigma\pi}m_{\sigma}\boldsymbol{\pi}\cdot\boldsymbol{\pi}\sigma \tag{2.8}$$

or $(-\frac{g_{\pi}}{2M}\bar{\psi}\gamma^{\mu}\gamma^5\psi\boldsymbol{\tau}\cdot\partial_{\mu}\boldsymbol{\pi})$ for pseudovector πN coupling. Schematically, each meson-baryon vertex in (2.8) is of the form

$$H_I(\psi) = \bar{\psi}(x)\Gamma^a_B\Pi^a(x)\psi(x), \tag{2.9}$$

where Γ_B is a matrix in Dirac space and $\Pi(x)$ is a meson field operator. There are also meson-meson vertices that are represented by

$$\mathcal{H}_I(\psi) = \Gamma_M\boldsymbol{\pi}\boldsymbol{\pi}\sigma, \tag{2.10}$$

with Γ_M a c number. These generalized vertices carry implicitly a summation over the different couplings to the mesons. Given the interacting Hamiltonian the matrix element

$$\tau(x_f, x', x, x_i) = \langle\phi_{\rm core}|\,T\,\left[\psi(x_f)\bar{\psi}(x')\psi(x)\bar{\psi}(x_i)\right]|\phi_{\rm core}\rangle \tag{2.11}$$

is expanded in perturbation theory, and the series summed to all orders in the RPA approximation as represented, diagramatically, in Fig. 1.

Finally, substituting Eq. (2.4) in the expression for $j_{\Gamma_A}(x)$ and taking the appropriate limits for the initial and final times, after some tedious calculations we arrive

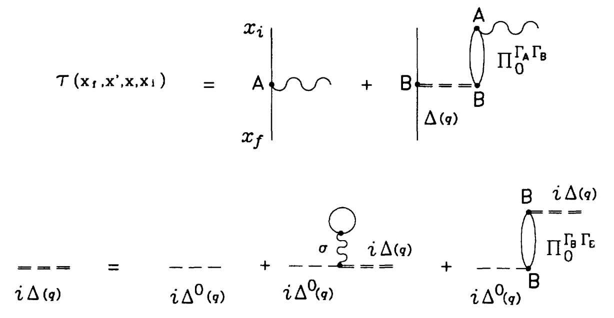

FIG. 1. Diagramatic representation of the perturbative calculation of the matrix element of Eq. (2.11).

at an expression for the current of the form

$$j_{\Gamma_A}(x) = \langle\phi_f|J^{sp}_{\Gamma_A}(x) + \delta J^{\rm RPA}_{\Gamma_A}(x)|\phi_i\rangle$$
$$= \left\{\langle u^0_f|\Gamma_A|u^0_i\rangle - \Delta(q)\Pi_0^{\Gamma_A\Gamma_B}(q)\langle u^0_f|\Gamma_B|u^0_i\rangle\right\}, \qquad (2.12)$$

where

$$\Pi_0^{\Gamma_A\Gamma_B}(q) = i\int \frac{dk^4}{(2\pi)^4}\exp ik^0\eta \;{\rm Tr}\left[\Gamma_A G(k+q)\Gamma_B G(k)\right] \qquad (2.13)$$

is the polarization insertion, $\Delta(q)$ is the dressed meson propagator, and $u^0_{i,f}$ are the initial and final valence single-particle wave functions

$$u^0_{i,f} = \frac{1}{\sqrt{\Omega}}e^{-ik_{i,f}x}\,U(\mathbf{k}_{i,f}, s_{i,f}, \tau_{i,f}). \qquad (2.14)$$

Equation (2.12) contains a single-particle current term plus an extra term originating on the polarization of the core —the *backflow* current. If we use this result to calculate the matrix element of the axial-vector current given in (2.1), we obtain

$$\langle\phi_f|J^{\mu5\pm}_{\rm PV}(x)|\phi_i\rangle \equiv \langle\phi_f|J^{\mu5\pm}_{sp}(x) + \delta J^{\mu5\pm}_{\rm RPA}(x)|\phi_i\rangle$$
$$= g_A\left[\langle u^0_f|\,\gamma^\mu\,\gamma^5\,\tau^\pm\,|u^0_i\rangle - \frac{q^\mu}{q^2-m^2_\pi}\langle u^0_f|\,q_\nu\,\gamma^\nu\,\gamma^5\,\tau^\pm\,|u^0_i\rangle\right]$$
$$-g_A\left[\left(\frac{g_\pi}{2M}\right)^2\Delta^\pi(q)\left(\Pi_0^{\mu5,5}(q) - \frac{q^\mu}{q^2-m^2_\pi}\Pi_0^{5,5}(q)\right)\langle u^0_f|\,q_\nu\,\gamma^\nu\,\gamma^5\,\tau^\pm\,|u^0_i\rangle\right]. \qquad (2.15)$$

Clearly, the only components of the polarization insertion tensor that contribute are $\Pi_0^{5,5}(q)$ and $\Pi_0^{\mu5,5}$. The other components vanish even at this value of the four-momentum transfer (as already shown in [3]). They are related in following fashion,

$$\Pi_0^{5,5}(q) = q_\mu\,\Pi_0^{\mu5,5}(q). \qquad (2.16)$$

Also, from the explicit calculation of the polarization insertion

$$\Pi_0^{\mu5,5}(q) = q^\mu\,f_{\rm PV}(q), \qquad (2.17)$$

and therefore

$$\Pi_0^{5,5}(q) = (q^\mu)^2\,f_{\rm PV}(q), \qquad (2.18)$$

with $f_{\rm PV}$ a complicated function of the four-momentum transfer, not singular at the origin and which depends on the density through M^* and k_F [11]. Substituting Eqs. (2.17) and (2.18) back into (2.15) and rearranging we arrive at an expression of the form

$$\langle\phi_f|J^{\mu5\pm}_{\rm PV}(x)|\phi_i\rangle = g_A\left[\langle u^0_f|\,\gamma^\mu\,\gamma^5 - \frac{q_\mu}{q^2-m^2_\pi}\,q_\nu\,\gamma^\nu\,\gamma^5\,\tau^\pm\,|u^0_i\rangle\right]$$
$$+g_A\left[\frac{q^\mu}{q^2-m^2_\pi}\,m^2_\pi\left(\frac{g_\pi}{2M}\right)^2\Delta^\pi(q)\,f_{\rm PV}(q)\,\langle u^0_f|\,q_\nu\,\gamma^\nu\,\gamma^5\,\tau^\pm\,|u^0_i\rangle\right]. \qquad (2.19)$$

Notice that the axial response of the core is proportional to the momentum transfer which indicates that it will vanish as the momentum decreases (as was the case for β-decay processes). This behavior is unlike that of the electromagnetic response which at zero momentum transfer makes a sizable contribution to the current (e.g., in the calculation of magnetic moments).

It pays to define the momentum-dependent function

$$C_{\rm PV}(q) \equiv 1 - m^2_\pi\left(\frac{g_\pi}{2M}\right)^2\Delta^\pi(q)\,f_{\rm PV}(q) \qquad (2.20)$$

since with this definition Eq. (2.19) may be recast in the form

$$\langle \phi_f | J^{\mu 5 \pm}_{\rm PV}(x) | \phi_i \rangle = g_A \left[\langle u^0_f | \gamma^\mu \gamma^5 \tau^\pm | u^0_i \rangle \, - \, C_{\rm PV}(q) \, \frac{q^\mu}{q^2 - m^2_\pi} \, \langle u^0_f | \, q_\nu \, \gamma^\nu \gamma^5 \tau^\pm \, | u^0_i \rangle \right] , \tag{2.21}$$

allowing for the definition of an *effective* axial–vector current of the form,

$$J^{\mu 5 \pm}_{\rm PVeff}(x) = \bar{\psi}_0(x) \, g_A \left[\gamma^\mu \gamma^5 \, - \, C_{\rm PV}(q) \, \frac{q^\mu}{q^2 - m^2_\pi} \, q_\nu \, \gamma^\nu \gamma^5 \right] \tau^\pm \, \psi_0(x). \tag{2.22}$$

Equation (2.22) is our main result. It shows that all effects originating on the isovector axial response of the core add up to a renormalization of the induced pseudoscalar term which can, alternatively, be accounted as a renormalization of the form factor $g_P(q) = g_A(q) C_{\rm PV}(q)$. This effective current satisfies PCAC since, from (2.20)

$$\lim_{m_\pi \to 0} C_{\rm PV} \, = \, 1. \tag{2.23}$$

One point to remark is that this polarization correction originates almost entirely in particle-hole excitations (the density-dependent terms in the Green's function) as opposed to particle-antiparticle excitations characteristic of the isoscalar electromagnetic response [5]. The relativistic dynamics leaks into the calculation of the polarization insertion through the effective mass M^*.

The function $C_{\rm PV}(q)$ depends on the momentum transfer as shown in Fig. 2. The value of q_0 has been fixed to that of muon capture, $q_0^2 = 0.01 m_\mu^2$. The dashed line corresponds to the mean-field approximation (MFA) where one-loop vacuum effects are ignored and the solid line to the relativistic Hartree approximation (RHA). The parameters for both calculations are given in Table I. The peaks in both curves relate to the resonance excitation, at finite density, of particle-hole pairs centered at $q = M^* q_0 / \kappa_F$ which gives $q = 20.6$ MeV/c in MFA ($M^*/M = 0.556$; $k_F = 1.42$ fm^{-1}) and $q = 29.6$ MeV/c in RHA ($M^*/M = 0.731$; $k_F = 1.3$ fm^{-1}). For the muon-capture kinematics $q = 99.5$ MeV/c which, from the figure, implies a quenching of the pseudoscalar induced term in the MFA case of $C^{\rm MFA}_{\rm PV}(-0.88 m^2_\mu) = 0.57$. In RHA the screening of the axial charge by vacuum excitations is not as strong, and one gets $C^{\rm RHA}_{\rm PV}(-0.88 m^2_\mu) = 0.82$. Since in both calculations the effect is substantial, it becomes important to investigate implications for finite nuclei where the density is not constant.

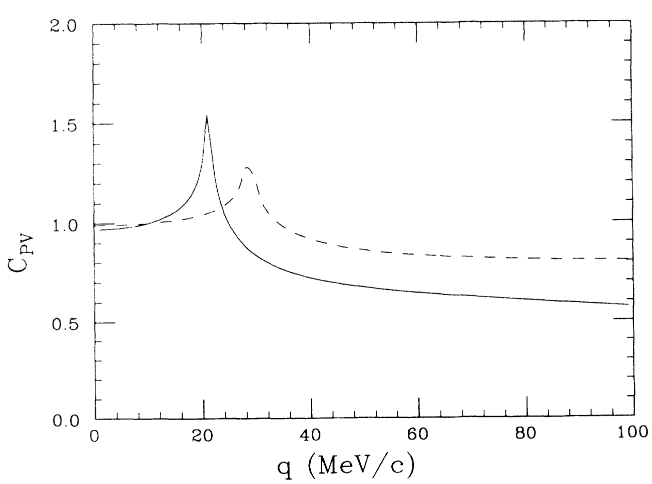

FIG. 2. The real part of the function $C_{\rm PV}$ as given by Eq. (2.20) in the text as a function of the three-momentum transfer $q = |\vec{q}|$. The solid curve is the mean-field-approximation (MFA) result. The dashed curve is the relativistic Hartree (RHA) calculation.

III. THE $\mu^- + {}^{16}\mathrm{O} \to {}^{16}\mathrm{N} + \nu_\mu$ REACTION REVISITED

To explore the consequences of the results in the previous section we analyzed the muon-capture process in ^{16}O using the local density approximation. This particular transition has been extensively studied in the literature, particularly employing relativistic dynamical models (Refs. [7,12–14]). The current understanding is that the capture rate is insensitive to relativistic dynamics with the proviso that pseudovector coupling be used at the πN vertex [8,14]. The claim is that the axial-charge density and the spacelike piece of the induced pseudoscalar term roughly cancel each other leaving a contribution for which the relativistic correction is not very large [14].

If, as shown in the previous section, the induced pseudoscalar term undergoes a sizable modification due to vacuum polarization processes then this cancellation may not take place and an enhancement appear. Unfortunately, a clean calculation is hindered by the nuclear structure of the problem which is not free from some ambiguity. Standard calculations use a closed-shell description of ^{16}O and a $(2s_{1/2} 1p^{-1}_{1/2})$ - $(1d_{3/2} 1p^{-1}_{3/2})$ mixing for the 0^-, $T = 1$ first-excited state in ^{16}N. The amount of the mixing, λ, is usually treated as a parameter used to fit simultaneously the muon-capture and the beta-decay rates.

TABLE I. Parameters used for the calculations presented in the text. The masses were fixed to their experimental values: $M = 939$ MeV, $m_\omega = 783$ MeV, $m_\rho = 770$ MeV, and $m_\pi = 139$ MeV. The coupling constant g^2_ρ corresponds to the value that fits the $\rho \to \pi\pi$ decay.

	$g^2_{\sigma N}$	m_σ	$g^2_{\omega N}$	g^2_π	g^2_ρ	κ_F (fm^{-1})	M^*/M
MFT	91.64	550	136.2	181	36.79	1.42	0.556
RHA	69.97	520	102.8	181	36.79	1.30	0.731

TABLE II. Muon-capture rates in $\sec^{-1}$ for the transition $\mu^- + ^{16}\mathrm{O} \rightarrow ^{16}\mathrm{N} + \nu_\mu$. The mixing parameter λ has been set to zero. $C_{\mathrm{PV}} = 1$ corresponds to no polarization effects, and C_{PV}=LDA to the results of the polarization calculation using the local density approximation. Pseudovector coupling is employed at the πN vertex.

MFA		RHA	
$C_{\mathrm{PV}} = 1$	C_{PV}=LDA	$C_{\mathrm{PV}} = 1$	C_{PV}=LDA
3193.9	3336.1	2762.1	2779.2

In Table II we show the results for the muon-capture rate obtained in MFA and RHA. For each case we quote the value without inclusion of the medium response (C_{PV}=1), and that using the function C_{PV} in (2.20) in LDA (C_{PV}=LDA). These results were obtained setting the mixing parameter $\lambda = 0$. A more realistic value (from our calculations $0.08 \lesssim \lambda \lesssim 0.12$) would not alter the conclusions appreciably. The important point to remark is that despite the strong quenching of the induced pseudoscalar term at nuclear matter densities in actual nuclei, its effect is small. The fact is that C_{PV} has a strong density dependence that can be appreciated by looking at Fig. 3 where the real part of C_{PV} is shown for the MFA calculation. At the center of the nucleus it has roughly the nuclear matter value and decreases as the density diminishes. Upon reaching the nuclear radius, it starts to increase and becomes larger than one, to finally fall asymptotically toward this value. This is the expected limit since, as mentioned, only contributions from the density-dependent terms in the Green's functions enter in the polarization. A similar analysis applies to the RHA case. The total effect on the capture rate is of the order of 4% in MFA and less than 1% in RHA.

For completeness we have included in Table III the results corresponding to calculations employing pseudoscalar coupling at the πN vertices. Clearly the effects are dramatic now, with the function C_{PS} producing an

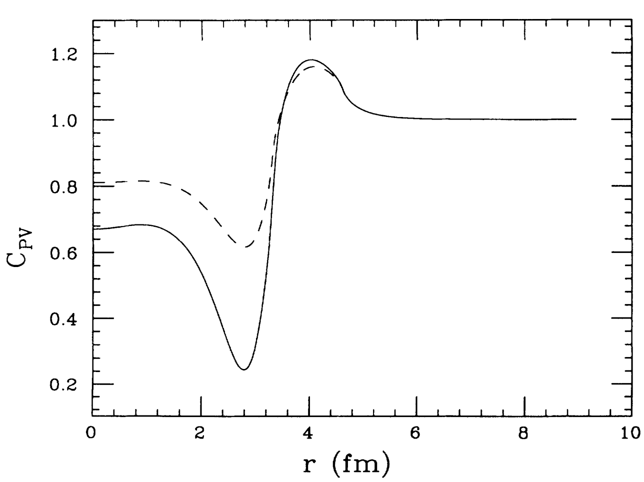

FIG. 3. The real part of the function C_{PV} calculated in local–density approximation for $^{16}\mathrm{O}$ as a function of the nuclear radius. The solid line corresponds to the MFA case and the dashed line to the RHA.

TABLE III. Same as Table II but for pseudoscalar coupling at the πN vertex.

MFA		RHA	
$C_{\mathrm{PS}} = 1$	C_{PS}=LDA	$C_{\mathrm{PS}} = 1$	C_{PS}=LDA
3807.0	10148.7	3288.3	10017.9

enhancement in the capture rates of 266% in MFA and 304% in RHA.

IV. SUMMARY AND CONCLUSIONS

Summarizing, we have investigated the isovector weak response of the core to an external perturbation driven by a particle, a hole or a particle-hole pair, using linear response theory in nuclear matter. The idea behind the work was to extend previous results to a kinematic region away from $q^2 \approx 0$ where we know that the effect is negligible. We concentrated on muon-capture processes. After some algebraic manipulation we noticed that the polarization contribution could be described as a factor multiplying the induced pseudoscalar term in the axial current. For the kinematics of muon capture we found a sizable quenching, due to this factor, of the contribution of the induced pseudoscalar term. When applied to actual nuclei, however, the strong density dependence of the weak isovector response washes out the screening found at nuclear matter densities, slightly modifying muon-capture rates.

The presence of the weak response of the core as calculated here must, however, be taken into account in relativistic calculations carried out assuming a uniform density distribution for the nucleus (see Ref. [15]). In such cases the induced pseudoscalar term is quenched as indicated, and relativistic effects become important. This backflow effect is substantial and may help one to better understand the experiment in these types of calculations.

On the other hand, the results of the previous section suggest that the uniform density approach to muon capture may be altogether flawed because of the strong dependence of the backflow current on the density. Fearing *et al.* [15] have already pointed out that muon-capture calculations are sensitive to a number of effects which, if not included, make results difficult to interpret.

Another important remark relates to the calculations of Chiang *et al.* [16] where it is concluded that medium renormalization effects are instrumental in bringing theory to accord with experiment specially for medium and heavy nuclei. We anticipate a larger effect from the weak response in these heavy nuclei. Work on this direction is currently in progress.

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[1] I.S. Towner, Comments Nucl. Part. Phys. **15**, 145 (1986).
[2] E.K. Warburton, Phys. Rev. Lett. **66**, 1823 (1991); Phys. Rev. C **44**, 233 (1991).
[3] A.O. Gattone, E.D. Izquierdo, and M. Chiapparini, Phys. Rev. C **46**, 788 (1992).
[4] G. Barenboim, A.O. Gattone, and E.D. Izquierdo, Phys. Rev. C **48**, 2537 (1993).
[5] R.J. Furnstahl and B.D. Serot, Nucl. Phys. **A468**, 539 (1987).
[6] J. Cohen, Phys. Rev. C **48**, 1346 (1993).
[7] M.W. Price and G.E. Walker, Phys. Rev. C **38**, 2860 (1988).
[8] Y. Nedjadi and J.R. Rook, Nucl. Phys. **A528**, 537 (1991).
[9] M. Gmitro and P. Truöl, in *Radiative Muon Capture and the Weak Pseudoscalar Coupling in Nuclei*, Advances in Nuclear Physics, Vol. 18, edited by J.W. Negele and E. Vogt (Plenum, New York 1987).
[10] B. D. Serot and J. D. Walecka, in *The Relativistic Nuclear Many-Body Problem*, Advances in Nuclear Physics, Vol. 16, edited by J.W. Negele and E. Vogt (Plenum, New York 1986).
[11] E.D. Izquierdo, Ph.D. thesis, Universidad de Buenos Aires, 1994.
[12] J.A. McNeil and J.R. Shepard, Phys. Rev C **31**, 686 (1985).
[13] G. Do Dang, M. Jaminon, and N. Van Giai, Phys. Lett. B **153**, 17 (1985).
[14] H.P.C. Rood, Phys. Rev. C **33**, 1104 (1986), comment on McNeil and Shepard, Phys. Rev C **31**, 686 (1985); J.A. McNeil and J.R. Shepard, Phys. Rev. C **33** 1105 (1986), reply to the comment.
[15] H.W. Fearing and M.S. Welsh, Phys. Rev. C **46**, 2077 (1992).
[16] H.C. Chiang, E. Oset, and P. Fernández de Córdoba, Nucl. Phys. **A510**, 591 (1990).