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Rho-nucleon tensor coupling and charge-exchange resonances

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Abstract

The Gamow–Teller resonance in ^{208}Pb is discussed in the context of a self-consistent RPA, based on the relativistic mean field theory. We inquire on the possibility of substituting the phenomenological Landau–Migdal force by a microscopic nucleon–nucleon interaction, generated from the rho-nucleon tensor coupling. The effect of this coupling turns out to be very small when the short range correlations are not taken into account, but too large when these correlations are simulated by the simple extraction of the contact terms from the resulting nucleon–nucleon interaction.

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The quantum hadrodynamics (QHD) aims to describe the nuclear many-body system in terms of nucleons and mesons [1]. Proposed initially as a full-fledged renormalizable quantum field theory, nowadays it is seen as an *effective field theory*, derivable, in principle, from the quantum chromodynamics [2].

The relativistic mean field theory (RMFT), which can be thought as a mean field (Hartree) approximation to the QHD, has been applied with great success during the last few decades. For instance, it accounts for both (i) the nuclear matter saturation, and (ii) the ground state properties of finite nuclei along the whole periodic table [3]. More recently, the RMFT has also been exploited for the description of unstable nuclei all up to the nucleon drip lines [4].

Through a relativistic version of the random phase approximation (RRPA), various excited states and resonances have been studied in the context of the RMFT [5–9] as well. Quite recently we have also reported [10] the first calculation of this type for the Gamow–Teller (GT) and isobaric analogue (IA) resonances, excited from the ground states of ^{48}Ca, ^{90}Zr and ^{208}Pb nuclei.

Because of its pseudoscalar nature, the pion does not participate in the description of the the ground states in the RMFT. Thus, besides the nucleon and the Coulomb fields, only the σ, ω and ρ mesons are usually involved in the calculations. Yet, in dealing with isovector excitations it is essential to include, together with the ρ meson, the π meson as well. This has already been done in our previous work [10], with the pseudovector pion–nucleon coupling f_π fixed at its experimental value. For the remaining mesons, only the nonderivative couplings to the nucleon were in-

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cluded, as usually done in RMFT. With this prescription we were not able to reproduce the excitation energies of the just mentioned resonances. This has been possible only after introducing the repulsive Landau–Migdal (LM) delta force

$$V_{\mathrm{LM}}(1,2) = g'\left(\frac{f_\pi}{m_\pi}\right)^2 \boldsymbol{\tau}_1 \cdot \boldsymbol{\tau}_2\, \boldsymbol{\sigma}_1 \cdot \boldsymbol{\sigma}_2\, \delta(\boldsymbol{r}_1 - \boldsymbol{r}_2), \quad (1)$$

of the same magnitude ($g' = 0.7$) as the one used in the nonrelativistic calculations [11].

Here we wish to analyze whether the tensor (derivative) coupling of the ρ meson to the nucleon could generate a sufficiently repulsive nucleon–nucleon force in order to locate the GT resonance at the correct experimental energy and in this way substitute the phenomenological LM force. The IA resonance is practically not affected by this part of the ρ-meson-exchange potential and therefore it will not be discussed so exhaustively as we do with the GT resonance.

As mentioned above, it is not usual to include the tensor coupling of the vector mesons to the nucleon in RMFT. This is because its effect on the ground state is (rightly) thought to be small. On a more general perspective, however, there are two good reasons why one should do so. For one, according to the rules of effective field theory such terms should appear in the effective QHD Lagrangian [12]. For another, and perhaps more important reason for the phenomenological stand we are taking, it is well known that the tensor ρ-nucleon coupling gives a large contribution to the spin–isospin component of the nucleon–nucleon interaction [13], and as such it could have an important effect on the dynamics of the GT resonance.

Our Lagrangian density is now

$$\begin{aligned}
\mathcal{L} = {} & \bar\psi\left(i\gamma_\mu\partial^\mu - M\right)\psi \\
& + \frac{1}{2}\partial_\mu\sigma\partial^\mu\sigma - \frac{1}{2}{m_\sigma}^2\sigma^2 - \frac{1}{3}g_2\sigma^3 \\
& - \frac{1}{4}g_3\sigma^4 - g_\sigma\bar\psi\psi\sigma \\
& - \frac{1}{4}\Omega_{\mu\nu}\Omega^{\mu\nu} + \frac{1}{2}{m_\omega}^2\omega_\mu\omega^\mu - g_\omega\bar\psi\gamma_\mu\psi\omega^\mu \\
& + \frac{1}{2}\partial_\mu\boldsymbol{\pi}\cdot\partial^\mu\boldsymbol{\pi} - \frac{1}{2}{m_\pi}^2\boldsymbol{\pi}\cdot\boldsymbol{\pi} \\
& - \frac{f_\pi}{m_\pi}\bar\psi\gamma_5\gamma_\mu\boldsymbol{\tau}\psi\cdot\partial^\mu\boldsymbol{\pi} \\
& - \frac{1}{4}\boldsymbol{R}_{\mu\nu}\cdot\boldsymbol{R}^{\mu\nu} + \frac{1}{2}{m_\rho}^2\boldsymbol{\rho}_\mu\cdot\boldsymbol{\rho}^\mu \\
& - g_\rho\bar\psi\gamma_\mu\boldsymbol{\tau}\psi\cdot\boldsymbol{\rho}^\mu \\
& - \frac{f_\rho}{2M}\bar\psi\sigma_{\mu\nu}\boldsymbol{\tau}\psi\cdot\partial^\mu\boldsymbol{\rho}^\nu \\
& - \frac{1}{4}F_{\mu\nu}F^{\mu\nu} - e\bar\psi\gamma_\mu\frac{1+\tau_3}{2}\psi A^\mu, \qquad (2)
\end{aligned}$$

where

$$\begin{aligned}
\Omega^{\mu\nu} &= \partial^\mu\omega^\nu - \partial^\nu\omega^\mu, \\
\boldsymbol{R}^{\mu\nu} &= \partial^\mu\boldsymbol{\rho}^\nu - \partial^\nu\boldsymbol{\rho}^\mu - 2g_\rho\boldsymbol{\rho}^\mu\times\boldsymbol{\rho}^\nu, \\
F^{\mu\nu} &= \partial^\mu A^\nu - \partial^\nu A^\mu. \qquad (3)
\end{aligned}$$

This Lagrangian is identical to that of Ref. [10], except for the ρ-nucleon tensor coupling term (the one proportional to f_ρ).[1] Therefore, following the same route one arrives at identical equations for the mean boson fields, except for that of the ρ-meson (only the component ρ_3^0 survives for spherical, definite-charge nuclei), which now takes the form

$$\left(-\nabla^2 + {m_\rho}^2\right)\rho_3^0 = g_\rho\rho_3(r) + \frac{f_\rho}{2M}\nabla\cdot\boldsymbol{\rho}_{t3}(r), \qquad (4)$$

where the (vector) isovector density ρ_3 is as defined in [10] and we have introduced the tensor isovector density

$$\boldsymbol{\rho}_{t3} = \langle\bar\psi i\boldsymbol{\alpha}\tau_3\psi\rangle = \sum_{\alpha=1}^{A}\overline{\mathcal{U}}_\alpha i\boldsymbol{\alpha}\tau_3\mathcal{U}_\alpha. \qquad (5)$$

The summation is over all the occupied single-particle, positive-energy states $\mathcal{U}_\alpha$, which obey the mean-field Dirac equation. This is also modified to

$$\begin{aligned}
& \big\{-i\boldsymbol{\alpha}\cdot\nabla + \beta[M + V_s(r)] \\
& \quad + V_v(r) + (i\beta\boldsymbol{\alpha}\cdot\boldsymbol{r}/r)V_t(r)\big\}\mathcal{U}_\alpha = E_\alpha\mathcal{U}_\alpha. \qquad (6)
\end{aligned}$$

Again the scalar (V_s) and vector (V_v) potentials are as defined in [10], while the tensor potential,

$$V_t = -\frac{f_\rho}{2M}\frac{d\rho_3^0}{dr}\tau_3, \qquad (7)$$

is the contribution from the tensor-coupling term in (2).

[1] There is a minor correction to be made in [10]. One must replace g_ρ by $2g_\rho$ in Eqs. (1) and (2) of that reference for consistency with the remaining equations.

The general structure and derivation of the RRPA for charge-exchange excitations, in the discretized spectral version we use, has been delineated in [10]. An alternative, more detailed account can be found in [14]. The main ingredient is the residual interaction V. For a self-consistent calculation, this must be obtained from the same Lagrangian (1) used for the mean field. Also, since Fock terms are ignored in RMFT, we must consider only the *direct* matrix elements of V. Hence only the isovector mesons contribute, and we get $V = V_\pi + V_\rho$, with, in the instantaneous approximation,

$$
\begin{aligned}
&V_\pi(1,2)\\
&\quad = -\left(\frac{f_\pi}{m_\pi}\right)^2 \boldsymbol{\tau}_1\cdot\boldsymbol{\tau}_2\,(\boldsymbol{\sigma}_1\cdot\nabla_1\,\boldsymbol{\sigma}_2\cdot\nabla_2)\\
&\qquad \times Y(m_\pi, r_{12}),
\end{aligned}
\tag{8}
$$

$$
\begin{aligned}
V_\rho(1,2) = {}& \boldsymbol{\tau}_1\cdot\boldsymbol{\tau}_2\\
&\times\left[\left(g_\rho - \frac{f_\rho}{2M} i\beta\boldsymbol{\alpha}\cdot\nabla\right)_1 \left(g_\rho - \frac{f_\rho}{2M} i\beta\boldsymbol{\alpha}\cdot\nabla\right)_2\right.\\
&\quad - \left(g_\rho\boldsymbol{\alpha} + \frac{f_\rho}{2M}\beta\boldsymbol{\sigma}\times\nabla\right)_1\\
&\quad \left.\times\left(g_\rho\boldsymbol{\alpha} + \frac{f_\rho}{2M}\beta\boldsymbol{\sigma}\times\nabla\right)_2\right] Y(m_\rho, r_{12}),
\end{aligned}
\tag{9}
$$

where $r_{12} = |\boldsymbol{r}_1 - \boldsymbol{r}_2|$ and $Y(m,r) = \exp(-mr)/(4\pi r)$.

For the numerical values of the parameters we follow mostly the philosophy of [10], adopting the parameter set NL1 [3,15]. Yet, in view of the difficulties encountered by Ma et al. [9] in accounting for the E1 and E0 giant resonances with the NL1 parameters, a few results for the TM1 model, worked out by Sugahara and Toki [16], will be presented as well. [2] Taking experimental values for the pion, the only new parameter is the ρ-nucleon tensor coupling constant f_ρ. As mentioned in [17], the vector dominance model predicts for the ratio $f_\rho/g_\rho \equiv K_\rho$ a value equal to the isovector magnetic moment of the nucleon, i.e., $\mu_p - \mu_n - 1 = 3.7$. On the other hand, most meson-exchange models for the nuclear force use $K_\rho = 6.6$ [13]. The former choice was preferred in the description of the ground state properties in closed shell nuclei within the relativistic Hartree–Fock approximation [17]. Thus, the discussion that follows will mainly rely on the lower value for K_ρ, even though we are aware of the fact that the inclusion of Fock terms can considerably change the adjusted values of the QHD parameters [17].

Another point to consider is whether the inclusion of the tensor coupling term in the Lagrangian (2) does not sensitively affect the values of the remaining parameters. Fortunately, while the contribution of this term is not strictly zero in RMFT, its effects on the single particle energies as well as on the ground state properties are certainly very small. We therefore feel justified in keeping the remaining parameters fixed at their NL1 or TM1 values. With $K_\rho = 3.7$, for instance, the spin-orbit splitting is modified in less than 150 keV. [3] Similarly tiny effects on the energy per particle and the root-mean-square radii are displayed in Table 1. An interesting side remark can be made concerning the latter observables. It is well known that, at variance with the nonrelativistic calculations, it is a common feature of the relativistic models to

Table 1
RMFT results for the energy per particle and the root-mean-square radii of the neutron and proton point-particle distributions in ^{208}Pb computed with several values of the ρ-nucleon tensor coupling constant f_ρ. The remaining parameters are kept fixed at their NL1 and TM1 values

$K_\rho \equiv f_\rho/g_\rho$	$E/A - M$ [MeV]	$\sqrt{\langle r_n^2\rangle}$ [fm]	$\sqrt{\langle r_p^2\rangle}$ [fm]
NL1 parameter set:			
0	−7.884	5.795	5.474
3.7	−7.882	5.777	5.480
6.6	−7.883	5.763	5.485
TM1 parameter set:			
0	−7.874	5.755	5.485
3.7	−7.871	5.741	5.492
6.6	−7.871	5.730	5.497
Experiment	−7.868[a]	5.593[b]	5.452[c]

[a] Taken from Ref. [18].
[b] Taken from Ref. [19].
[c] Taken from Refs. [20] and [21, Eq. (6.1)].

[2] In the latter case, the ω-meson self-interaction term, not appearing in (2), was also included in the numerical calculations.

[3] For identical particles, the NL1 paramerization yields significantly larger spin-orbit splittings than the TM1 model, while the opposite happens for nonidentical particles.

overestimate the neutron skin thickness [16,23]. But, as seen from the results shown in Table 1, the tensor ρ-N coupling has the tendency to correct the RMFT for this handicap. This fact, in turn, could have very important consequences on the estimates of the atomic parity nonconservation [24,25].

The GT and IA resonances in ^{208}Pb were computed in RRPA, for both the NL1 and TM1 sets of parameters and within the same model space as that of [10], i.e., including only $0\hbar\Omega$ and $2\hbar\Omega$ excitations, and only those single-particle states that are bound at least for neutrons. For simplicity, we ignored the negative-energy states, although it has been shown [8] that they are required in principle even if the no-sea approximation is made for RMFT, since one needs a complete single-particle basis to develop a perfectly consistent RRPA. In fact, the transitions from Fermi- to Dirac-sea states are essential to ensure certain desirable features, such as current conservation and the removal of the spurious $J^\pi = 1^-$ translational state. However, such issues are not crucial for our present purposes and, furthermore, Ma et al. have shown in a recent calculation [26] that the contribution of the negative-energy states is of decisive importance only for the isoscalar modes. We therefore feel safe to leave their inclusion for a future, more sophisticated and detailed treatment of those isovector resonances.

In Fig. 1 are shown the NL1 results for the GT strength distribution, both in terms of the individual strengths,

$$s_\lambda = \left| \sum_{p\bar{n}} X^\lambda_{p\bar{n}} \langle p \| \boldsymbol{\sigma} \| \bar{n} \rangle + \sum_{n\bar{p}} Y^\lambda_{n\bar{p}} \langle \bar{p} \| \boldsymbol{\sigma} \| n \rangle \right|^2, \qquad (10)$$

and of a "strength function" obtained by replacing the spikes by Lorentzians of conveniently chosen widths Δ [10], i.e.,

$$S(E) = \frac{\Delta}{\pi} \sum_\lambda \frac{s_\lambda}{(E - E_\lambda)^2 + \Delta^2}, \qquad (11)$$

where $X^\lambda_{p\bar{n}}$ and $Y^\lambda_{n\bar{p}}$ are, respectively, the forward and backward going RPA amplitudes for the state at excitation-energy E_λ. The upper, middle and lower panels correspond, respectively, to: (a) $K_\rho = 0, g' = 0$; (b) $K_\rho = 3.7, g' = 0$ and (c) $K_\rho = 3.7, g' = 0.7$. From these results one is induced to conclude that the tensor ρ-N coupling has a very small effect on the GT resonance. That is, it seems as though this coupling could merely redistribute the GT strength in the energy region between 5 and 15 MeV, but was incapable of promoting it to the correct experimental energy. The latter is only achieved after introducing an LM force of the same magnitude that has been used in the previous calculation, where the just mentioned coupling has not been considered at all [10]. The issue of the NN-force generated by the ρ-N coupling is, however, not so

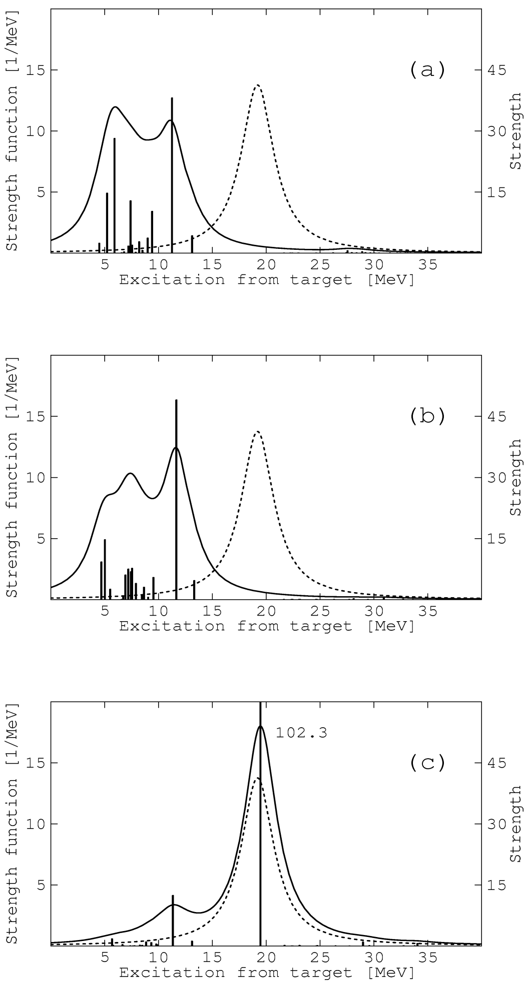

Fig. 1. Gamow–Teller strength distribution for the parent nucleus ^{208}Pb for the parametrization NL1. The upper, middle and lower panels correspond, respectively, to: (a) $K_\rho = 0$, $g' = 0$; (b) $K_\rho = 3.7, g' = 0$ and (c) $K_\rho = 3.7, g' = 0.7$. The spikes (r.h.s. scale) give the raw RRPA results and the continuous curve (l.h.s. scale), the strength function smoothed out by means of Lorentzians having widths of: (a) and (b) 3.0, and (c) 3.65 MeV. The strength function for the resonance peak extracted from experiment [22] is drawn in dotted line.

simple and deserves further discussion, which is presented below. Before proceeding, let us just mention that we have not noticed large differences between the NL1 and TM1 results for the IA and GT resonances. For instance, in the case (c) we get that these excitations are localized at: $E_{\mathrm{IA}}(\mathrm{NL1}) = 18.6$ MeV and $E_{\mathrm{GT}}(\mathrm{NL1}) = 19.5$ MeV and $E_{\mathrm{IA}}(\mathrm{TM1}) = 18.7$ MeV and $E_{\mathrm{GT}}(\mathrm{TM1}) = 20.3$ MeV, while the experimental results are: $E_{\mathrm{IA}}(\exp) = 18.8$ MeV and $E_{\mathrm{GT}}(\exp) = 19.2$ MeV. Thus, henceforth only the parametrization NL1 will be used.

In the upper panel of Fig. 2 are confronted several diagonal $J^\pi = 1^+$ proton-particle neutron-hole matrix elements for the V_{LM}, V_π, V_ρ^{VV}, V_ρ^{VT} and V_ρ^{TT} potentials. (The meaning of the upper indices is self-explanatory.) One can see, in particular, that the matrix elements of V_ρ^{TT} are very small in comparison with those coming from V_{LM}. However, when we rewrite V_ρ^{TT} in the form

$$
\begin{aligned}
V_\rho^{\mathrm{TT}}(1,2) = {} & \left(\frac{f_\rho}{2M}\right)^2 \boldsymbol{\tau}_1 \cdot \boldsymbol{\tau}_2 \beta_1 \beta_2 \\
& \times \biggl\{ - (\boldsymbol{\alpha} \cdot \nabla)_1 (\boldsymbol{\alpha} \cdot \nabla)_2 Y(m_\rho, r_{12}) \\
& \quad - \frac{1}{3} m_\rho^2 \left(\frac{3}{m_\rho^2 r_{12}^2} + \frac{3}{m_\rho r_{12}} + 1 \right) \\
& \quad \times Y(m_\rho, r_{12}) S_{12} \\
& \quad + \frac{2}{3} \left[m_\rho^2 Y(m_\rho, r_{12}) - \delta(\boldsymbol{r}_1 - \boldsymbol{r}_2) \right] \\
& \quad \times \boldsymbol{\sigma}_1 \cdot \boldsymbol{\sigma}_2 \biggr\}
\end{aligned}
\tag{12}
$$

and evaluate different parts separately, we find out that the Yukawa and contact pieces in the last term engender, each one, very large matrix elements. In fact, as shown in the lower panel, their individual values are larger than those of V_{LM}, but the overall contribution to V_ρ^{TT} is small, because they cancel each other very strongly. (A similar cancellation, though not so pronounced, also occurs in the case of V_π.)

It should be remembered that the contact terms in V_π and V_ρ^{TT} would be smeared over a finite region if finite-nucleon-size effects (FNSE) were introduced, and they would be totally killed by realistic short range

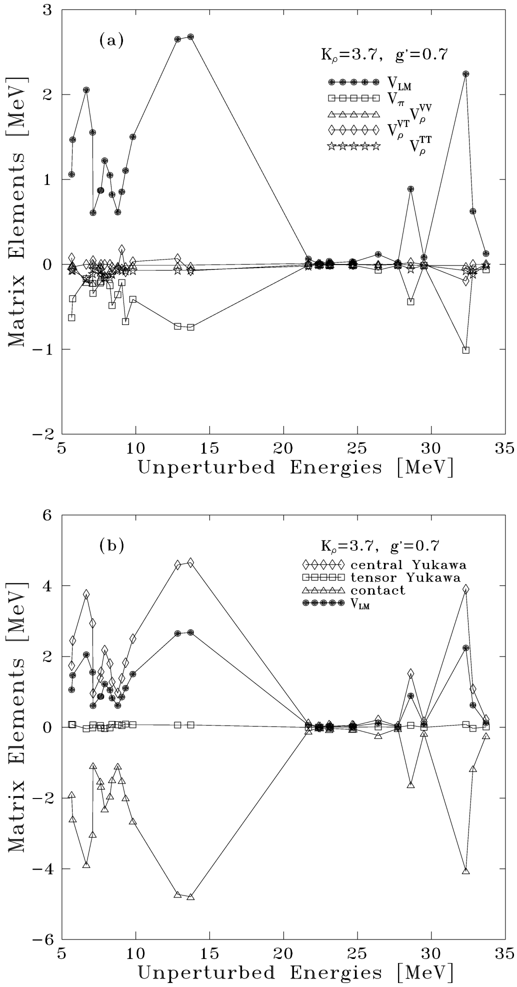

Fig. 2. Diagonal matrix elements of: (a) the several terms of the $\pi + \rho$ NN-interaction and (b) different pieces of V_ρ^{TT}, taken between proton-particle neutron-hole 1^+ states in ^{208}Pb. The matrix elements of the Landau–Migdal contact force are also shown in both panels for comparison. The states are positioned at their unperturbed energies.

correlations (SRC). [4] Yet, none of these two effects is considered in a mean field treatment, such as the present one. In return, it is common practice [5,17] to extract the contact parts from (8) and (9) by adding to the residual interaction the correction term $\delta V =$

[4] Note, however, that the contributions of the contact terms are nonzero when, both the FNSE, and the SRC are considered simultaneously [27].

$\delta V_\pi + \delta V_\rho$, with

$$\delta V_\pi(1,2) = \frac{1}{3}\left(\frac{f_\pi}{m_\pi}\right)^2 \boldsymbol{\tau}_1 \cdot \boldsymbol{\tau}_2\, \boldsymbol{\sigma}_1 \cdot \boldsymbol{\sigma}_2\, \delta(\boldsymbol{r}_1 - \boldsymbol{r}_2),$$

$$\delta V_\rho(1,2) = \frac{1}{3}\left(\frac{f_\rho}{2M}\right)^2 \boldsymbol{\tau}_1 \cdot \boldsymbol{\tau}_2\, \beta_1 \beta_2 \times (\boldsymbol{\alpha}_1 \cdot \boldsymbol{\alpha}_2 + 2\boldsymbol{\sigma}_1 \cdot \boldsymbol{\sigma}_2)\delta(\boldsymbol{r}_1 - \boldsymbol{r}_2). \quad (13)$$

For consistency, one must also perform such an extraction in the mean field part. Since, differently from the Hamiltonian formalism followed in [17], we are working within a Lagrangian formalism, we did this extraction in the baryon self-energy computed in the Hartree approximation (which is equivalent to RMFT). As a consequence the replacement $V_t \to V_t + \delta V_t$ has to be done in the Dirac equation (6), with

$$\delta V_t = \frac{1}{3}\left(\frac{f_\rho}{2M}\right)^2 \frac{\boldsymbol{\rho}_{t3} \cdot \boldsymbol{r}}{r} \tau_3 \quad (14)$$

being a correction that arises upon the extraction from the baryon self-energy of the contact part due to this derivative coupling in Eq. (2). But, when this recipe is implemented in the numerical calculation we get too much repulsion and the GT resonance is pushed up very high in energy. This comes from the fact that δV is basically a δ-force of the type (1), with

$$g'_{\pi+\rho} \cong \frac{1}{3} + \frac{2}{3}\left(\frac{f_\rho}{f_\pi}\right)^2 \left(\frac{m_\pi}{2M}\right)^2 = 1.6, \quad (15)$$

which is significantly larger than $g' = 0.7$.

Note that in the nonrelativistic approximation the contact term also appears in V_ρ^{VV} and V_ρ^{VT}, and instead of (15) one would have

$$g'_{\pi+\rho} \cong \frac{1}{3} + \frac{2}{3}\left(\frac{g_\rho + f_\rho}{f_\pi}\right)^2 \left(\frac{m_\pi}{2M}\right)^2 = 2.3. \quad (16)$$

There is no consensus on whether one should proceed in the same way in the relativistic case. Some authors exclude the contact terms only from V_π and V_ρ^{TT} [17], while others do that for the full $\pi + \rho$ interaction [5, 28]. That the potentials V_ρ^{VV} and V_ρ^{VT} also contain a contact term follows from the substitution [29]

$$\gamma_\mu \longleftrightarrow \frac{1}{2M}\left(2P_\mu + \sigma_{\mu\nu}\partial^\nu\right) \quad (17)$$

for the vector ρ-N coupling.

It is worth noting that Toki and Weise [30] have interpreted microscopically the LM force as arising from the $\pi + \rho$ meson-exchange model combined with the SRC and FNSE. In the static limit, which is used here, the result is [31]:

$$g'_{\mathrm{LM}}(\omega = q = 0) \cong \frac{1}{3}\left(\frac{\Lambda^2 - m_\pi^2}{\Lambda^2 + m_0^2}\right)^2 \frac{m_0^2}{m_0^2 + m_\pi^2} + \frac{2}{3}\left(\frac{g_\rho + f_\rho}{f_\pi}\right)^2 \left(\frac{m_\pi}{2M}\right)^2 \frac{m_0^2}{m_0^2 + m_\rho^2}, \quad (18)$$

where Λ is the cut-off mass for the pion–nucleon vertex and m_0^{-1} is the correlation length. For $\Lambda = 1$ GeV, $m_0 = m_\rho$ and $g_\rho + f_\rho = 17.2$ this leads to $g'_{\mathrm{LM}}(\omega = q = 0) = 0.67$ [31]. (In the present work $g_\rho + f_\rho = 23.4$.)

Our results can be summarized as follows:

(1) When the short range correlations are not considered, the tensor ρ-nucleon coupling plays only a minor role in the description of the GT resonances.
(2) If one tries to take these correlations into account by merely extracting the contact terms from the NN interaction, the GT resonance is pushed up too high in energy.

Thus, the simulation of the short range correlations by the simple-minded extraction of the contact terms alone is not a satisfactory procedure; at least not in the case of the heavier mesons. The explanation is that the contact terms in the $\pi + \rho$ NN-interaction are not the only ones to be strongly modified by the short range correlations. In particular, because of the large ρ-meson mass, also the Yukawa terms generated in (9) should be strongly reduced. We conclude hence that the implementation of, both realistic short range correlations, and finite-nucleon-size effects, in the context of the relativistic RPA, is required. Presently, we are working on this issue.

Finally, let us mention that the tensor ρ-nucleon coupling plays an important role in the transverse spin response, and that some progress in assessing this through a relativistic many-body calculation has been made quite recently by Yoshida and Toki [32].

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