Gamow-Teller transitions at finite temperatures in the extended quasiparticle random phase approximation

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Gamow-Teller single β decay transitions ($\Delta J=1$, $\Delta\pi=0$, $\Delta T_z=\pm 1$), at finite temperatures and between states of even-even and odd-odd mass nuclei, are calculated in the framework of the extended quasiparticle random phase approximation. Finite temperature blocking effects and particle-particle interactions are discussed in connection with the fragmentation of strength. It is shown that the inclusion of scattering terms, both in the phonons and in the transition operator, is needed to achieve the complete conservation of the transition strength. The calculations were performed in the mass region $A=106$ and for temperatures lower than the critical value T_c associated to the collapse of the pairing gaps ($T_c \cong 0.7$ MeV).

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I. INTRODUCTION

Large scale calculations of single β decay rates are a necessary element in astrophysical estimates of electroweak stellar processes [1]. The quasiparticle random phase approximation (QRPA) [2,3] has been used extensively in the calculation of β^- and β^+/EC transition rates for a large number of electron (positron) emitters [1]. From the nuclear structure point of view, the use of the QRPA to calculate charge-exchange transition rates in stellar conditions deserves some attention. To start with, since the electroweak processes in stellar media, like a supernova, take place at temperatures of the order of a few hundreds keV, one should treat particle and particle-hole states with nontrivial statistical occupancies. Thermal blocking effects can modify the pattern of the ground state occupation numbers and at relatively low temperatures particle-particle and hole-hole configurations are activated and multiple scattering across the Fermi surface can occur. In consequence, the QRPA formalism should be extended to accommodate scattering terms at finite temperatures. In this work we have adapted the formalism of [4] to account for charge dependent interactions. The suitability of the method was demonstrated in a previous work [5] in dealing with Fermi β decay transitions. Another aspect, concerning the calculation of Gamow-Teller (GT) transitions in the QRPA formalism, is the fragmentation of the strength produced by renormalized particle-particle and hole-hole interactions [6]. The problem discussed in [5] was addressed, recently, in the context of a schematic model space [7]. In spite of the fact that the model assumptions of Refs. [5,7] are different the main conclusion of both works is the same, that is the transition strength is strictly conserved if one and two quasiparticle terms are included in the equation of motion. In order to continue with the discussion advanced in [5,7], and in view of potential applications to astrophysics, we have calculated GT transitions at finite temperatures. As shown in [8,9] the calculation of strength distributions of GT transitions becomes relevant if information about energy thresholds of neutrino and electron capture processes is needed, for instance, in the context of the physics of supernovae [10]. As an example about the use of the formalism, we have taken the case of $0^+\rightarrow 1^+$ transitions between even-even and odd-odd mass nuclei in the $A=106$ mass region. We have used a Hamiltonian consisting of a separable force between protons and neutrons and monopole pairing interactions. The essentials of the formalism that are needed to describe GT transitions are presented in Sec. II. The results of the calculations are discussed in Sec. III. Finally, some conclusions are drawn in Sec. IV.

II. FORMALISM

The Hamiltonian, which we have adopted for the description of $J^\pi=1^+$ excitations in double-odd mass nuclei, is the separable one introduced by K'zmin and Soloviev [11], namely,

$$H=\sum_{pj} e_{pj}N_{pj}+\sum_{j} e_{nj}N_{nj}-G_pS_p^\dagger S_p-G_nS_n^\dagger S_n +2\chi\sum_{\mu} \beta^-_{1\mu}(-)^\mu\beta^+_{1-\mu}-2\kappa\sum_{\mu} P^-_{1\mu}(-)^\mu P^+_{1-\mu}, \tag{1}$$

and it containts single-particle energies, monopole pairing terms, and charge-exchange particle-hole and particle-particle (hole-hole) terms. The corresponding definitions are the following:

$$N_{qj}=\sum_m a^\dagger_{qjm}a_{qjm},$$

$$S^\dagger_q=\sum_{jm} a^\dagger_{qjm}a^\dagger_{q\overline{jm}},\quad S_q=(S^\dagger_q)^\dagger,\quad q=p,n,$$

$$\beta^-_{1\mu}=\sum_{jmj'm'} \langle pjm|\sigma_\mu|nj'm'\rangle a^\dagger_{pjm}a_{nj'm'},$$

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$$\beta^{+}_{1\mu}=(-)^{\mu}(\beta^{-}_{1-\mu})^{\dagger},$$

$$P^{-}_{1\mu}=\sum_{jmj'm'}\langle pjm|\sigma_{\mu}|nj'm'\rangle a^{\dagger}_{pjm}a^{\dagger}_{\overline{nj'm'}},$$

$$P^{+}_{1\mu}=(-)^{\mu}(P^{-}_{1-\mu})^{\dagger}, \tag{2}$$

and these operators are the number operator, the monopole pair operator, and the charge-exchanging particle-hole and particle-particle operators, respectively. Proton and neutron single-particle orbits, of angular momentum j and projection m, are denoted by the index q ($q=p$ for protons and $q=n$ for neutrons) and $a^{\dagger}_{qjm}$ is a particle creation operator and $a^{\dagger}_{q\overline{jm}}=(-1)^{j-m}a^{\dagger}_{qj-m}$ its time reversal. In the BCS representation [12] we can write H, of Eq. (1) as

$$\begin{aligned}H=&\sum_{\nu}E_{\nu}\alpha^{\dagger}{}_{\nu}\alpha_{\nu}+\sum_{pn,p'n'}r(pn,p'n')\\&\times[A^{\dagger}(pn,1\mu)A(p'n'1\mu)+A^{\dagger}(p'n',1\mu)A(pn,1\mu)]\\&+\sum_{jj'}s(pn,p'n')[A^{\dagger}(pn,1\mu)A^{\dagger}(p'n',\overline{1\mu})\\&+A(p'n'\overline{1\mu})A(pn,1\mu)]+\sum_{jj'}u(pn,p'n')\\&\times[B^{\dagger}(pn,1\mu)B(p'n',1\mu)+B^{\dagger}(p'n',1\mu)B(pn,1\mu)]\\&+\sum_{jj'}v(pn,p'n')[B^{\dagger}(pn,1\mu)B^{\dagger}(p'n',\overline{1\mu})\\&+B(p'n',\overline{1\mu})B(pn,1\mu)]+\sum_{jj'}t(pn,p'n')\\&\times[A^{\dagger}(pn,1\mu)B(p'n',1\mu)+B^{\dagger}(p'n',1\mu)A(pn,1\mu)]\\&+\sum_{jj'}w(pn,p'n')[A^{\dagger}(pn,1\mu)B^{\dagger}(p'n',\overline{1\mu})\\&+B(p'n',\overline{1\mu})A(pn,1\mu)],\end{aligned}\tag{3}$$

where E_{ν} are the quasiparticle energies and the index ν indicates single-particle states. The transformed operators and coefficients of Eq. (3) are given by

$$A^{\dagger}(pn,1\mu)=[\alpha^{\dagger}_{p}\alpha^{\dagger}_{n}]_{1\mu},$$

$$A^{\dagger}(pn,\overline{1\mu})=(-1)^{1-\mu}A^{\dagger}(pn,1-\mu),$$

$$B^{\dagger}(pn,1\mu)=[\alpha^{\dagger}_{p}\alpha_{\bar{n}}]_{1\mu},$$

$$B^{\dagger}(pn,\overline{1\mu})=(-1)^{1-\mu}B^{\dagger}(pn,1-\mu),$$

$$\begin{aligned}r(pn,p'n')=&\chi(\sigma_{pn}\sigma_{p'n'}+\bar{\sigma}_{pn}\bar{\sigma}_{p'n'})\\&-\kappa(\sigma^{U}_{pn}\sigma^{U}_{p'n'}+\sigma^{V}_{pn}\sigma^{V}_{p'n'}),\end{aligned}$$

$$\begin{aligned}s(pn,p'n')=&\chi(\sigma_{pn}\bar{\sigma}_{p'n'}+\bar{\sigma}_{pn}\sigma_{p'n'})\\&+\kappa(\sigma^{U}_{pn}\sigma^{V}_{p'n'}+\sigma^{V}_{pn}\sigma^{U}_{p'n'}),\\u(pn,p'n')=&\chi(\sigma^{U}_{pn}\sigma^{U}_{p'n'}+\sigma^{V}_{pn}\sigma^{V}_{p'n'})\\&-\kappa(\sigma_{pn}\sigma_{p'n'}+\bar{\sigma}_{pn}\bar{\sigma}_{p'n'}),\\v(pn,p'n')=&-\chi(\sigma^{U}_{pn}\sigma^{V}_{p'n'}+\sigma^{V}_{pn}\sigma^{U}_{p'n'})\\&-\kappa(\sigma_{pn}\bar{\sigma}_{p'n'}+\bar{\sigma}_{pn}\sigma_{p'n'}),\\t(pn,p'n')=&2\chi(\bar{\sigma}_{pn}\sigma^{V}_{p'n'}-\sigma_{pn}\sigma^{U}_{p'n'})\\&+2\kappa(\sigma^{V}_{pn}\bar{\sigma}_{p'n'}-\sigma^{U}_{pn}\sigma_{p'n'}),\\w(pn,p'n')=&2\chi(\sigma_{pn}\sigma^{V}_{p'n'}-\bar{\sigma}_{pn}\sigma^{U}_{p'n'})\\&+2\kappa(\sigma^{V}_{pn}\sigma_{p'n'}-\sigma^{U}_{pn}\bar{\sigma}_{p'n'}),\end{aligned}\tag{4}$$

and the factors

$$\begin{aligned}&\sigma_{pn}=\sigma(pn)u_{p}v_{n}, \quad \bar{\sigma}_{pn}=\sigma(pn)u_{n}v_{p},\\&\sigma^{U}_{pn}=\sigma(pn)u_{p}u_{n}, \quad \sigma^{V}_{pn}=\sigma(pn)v_{n}v_{p},\end{aligned}\tag{5}$$

are the reduced matrix elements of the spin operator

$$\sigma(pn)=\frac{\langle p||\sigma||n\rangle}{\sqrt{3}}, \tag{6}$$

multiplied by the proper proton and neutron occupation factors. The creation (annihilation) of quasiparticles is represented by the operator $\alpha^{\dagger}_{\nu}$ (α_{ν}) while u_{ν} and v_{ν} are BCS occupations factors.

In the standard form of the QRPA method [2,3] the Hamiltonian H is diagonalized in the phonon basis and only pair creation and pair annihilation operators $A^{\dagger}$ and A are included in the definition of the QRPA phonons. This form can be generalized to include the operators $B^{\dagger}$ and B in the definition of the phonons, as done in [4]. In the present case of proton-neutron excitations the extended QRPA phonon is written

$$\begin{aligned}\Gamma^{\dagger}(k,1\mu)=&\sum_{pn}[X(k,pn)A^{\dagger}(pn,1\mu)-Y(k,pn)A(pn,\overline{1\mu})\\&+Z(k,pn)B^{\dagger}(pn,1\mu)-\bar{Z}(k,pn)B(pn,\overline{1\mu})],\end{aligned}\tag{7}$$

where the extra terms

$$Z(k,pn)B^{\dagger}(pn,1\mu)-\bar{Z}(k,pn)B(pn,\overline{1\mu}) \tag{8}$$

are added to the conventional definition of the phonon operator. As shown in Ref. [4] vacuum expectation values can be replaced by thermal averages to account for temperature dependent effects. The thermal QRPA matrix equation can be written

$$\begin{pmatrix} \tilde{A} & \tilde{B} \\ \tilde{B}^* & \tilde{A}^* \end{pmatrix} = \omega \begin{pmatrix} \tilde{S} & 0 \\ 0 & -\tilde{S} \end{pmatrix} \begin{pmatrix} \tilde{X} \\ \tilde{Y} \end{pmatrix}, \tag{9}$$

and it doubles the number of dimensions of the standard QRPA problem. The forward ($\tilde{A}$) and backward ($\tilde{B}$) matrices, the metric matrix ($\tilde{S}$), and the amplitudes ($\tilde{X}$ and $\tilde{Y}$) are defined by

$$\tilde{A} = \begin{pmatrix} A & C \\ E & G \end{pmatrix},$$

$$\tilde{B} = \begin{pmatrix} B & D \\ F & H \end{pmatrix},$$

$$\tilde{S} = \begin{pmatrix} S & 0 \\ 0 & T \end{pmatrix},$$

$$\tilde{X} = \begin{pmatrix} X \\ Z \end{pmatrix},$$

$$\tilde{Y} = \begin{pmatrix} Y \\ \bar{Z} \end{pmatrix}. \tag{10}$$

The above matrices are written

$$A_{pn,p'n'} = \langle [A(pn,1\mu),[H,A^\dagger(p'n',1\mu)]] \rangle,$$
$$B_{pn,p'n'} = -\langle [A(pn,1\mu),[H,A(p'n',\overline{1\mu})]] \rangle,$$
$$C_{pn,p'n'} = \langle [A(pn,1\mu),[H,B^\dagger(p'n',1\mu)]] \rangle,$$
$$D_{pn,p'n'} = -\langle [A(pn,1\mu),[H,B(p'n',\overline{1\mu})]] \rangle,$$
$$E_{pn,p'n'} = \langle [B(pn,1\mu),[H,A^\dagger(p'n',1\mu)]] \rangle,$$
$$F_{pn,p'n'} = -\langle [B(pn,1\mu),[H,A(p'n',\overline{1\mu})]] \rangle,$$
$$G_{pn,p'n'} = \langle [B(pn,1\mu),[H,B^\dagger(p'n',1\mu)]] \rangle,$$
$$H_{pn,p'n'} = -\langle [B(pn,1\mu),[H,B(p'n',\overline{1\mu})]] \rangle,$$
$$S_{pn,p'n'} = \langle [A(pn,1\mu),A^\dagger(p'n',1\mu)] \rangle,$$
$$T_{pn,p'n'} = \langle [B(pn,1\mu),B^\dagger(p'n',1\mu)] \rangle, \tag{11}$$

and these matrix elements are obtained after evaluation of the corresponding commutators and double commutators, leading to the expressions

$$A_{pn,p'n'} = \delta_{p,p'}\delta_{n,n'}(1-f_n-f_p)(E_n+E_p) + r(pn,p'n')(1-f_{n'}-f_{p'})(1-f_n-f_p),$$

$$B_{pn,p'n'} = s(pn,p'n')(1-f_{n'}-f_{p'})(1-f_n-f_p),$$

$$C_{pn,p'n'} = t(pn,p'n')(f_{n'}-f_{p'})(1-f_n-f_p),$$

$$D_{pn,p'n'} = w(pn,p'n')(f_{n'}-f_{p'})(1-f_n-f_p),$$

$$E_{pn,p'n'} = C_{pn,p'n'},$$

$$F_{pn,p'n'} = D_{pn,p'n'},$$

$$G_{pn,p'n'} = \delta_{p,p'}\delta_{n,n'}(f_n-f_p)(E_p-E_n) + u(pn,p'n')(f_{n'}-f_{p'})(f_n-f_p),$$

$$H_{pn,p'n'} = v(pn,p'n')(f_{n'}-f_{p'})(f_n-f_p),$$

$$S_{pn,p'n'} = \delta_{p,p'}\delta_{n,n'}(1-f_n-f_p),$$

$$T_{pn,p'n'} = \delta_{p,p'}\delta_{n,n'}(f_n-f_p), \tag{12}$$

where f_ν are thermal occupation factors for single quasiparticle states

$$f_\nu = [1+\exp E_\nu/T]^{-1}. \tag{13}$$

The expectation values that appear in Eq. (11) have been calculated at finite temperature and the quantity T, of the quasiparticle occupation factor f_ν of Eq. (13), represents the nuclear temperature in units of energy.

The normalization condition for the phonons is

$$\langle [\Gamma(k,1\mu),\Gamma^\dagger(k',1\mu')] \rangle = \delta_{kk'}\delta_{\mu\mu'}\sum_{pn} \{(f_n-f_p)[Z(k,pn)^2-\bar{Z}(k,pn)^2] + (1-f_n-f_p)[X(k,pn)^2-Y(k,pn)^2]\}, \tag{14}$$

where the sum runs over proton-neutron two quasiparticle configurations. Next, we shall write the transition operators $\beta^\pm$, which are the GT operators, in the quasiparticle basis. The explicit expressions are the following:

$$\beta^-_{1\mu} = \sum_{pn} [\sigma_{pn}A^\dagger(pn,1\mu) + \bar{\sigma}_{pn}A(pn,\overline{1\mu}) - \sigma^U_{pn}B^\dagger(pn,1\mu) + \sigma^V_{pn}B(pn,\overline{1\mu})],$$

$$\beta^+_{1\mu} = (-1)^\mu(\beta^-_{1-\mu})^\dagger. \tag{15}$$

By using inversion formulas one can express these transition operators in the QRPA phonon basis. They are written

$$\beta^-_{1\mu} = \sum_k [a_k\Gamma^\dagger(k,1\mu) + b_k\Gamma(k,\overline{1\mu})],$$

$$\beta^+_{1\mu} = (-1)^\mu(\beta^-_{1-\mu})^\dagger, \tag{16}$$

where the amplitudes a_k and b_k,

$$a_k = \langle [\Gamma(k,1\mu),\beta^-_{1\mu}] \rangle,$$

$$b_k = \langle [\beta^-_{1\mu},\Gamma^\dagger(k,\overline{1\mu})] \rangle, \tag{17}$$

represent the thermal expectation values of the commutator of the transition operators with the QRPA phonons. They can also be expressed in terms of the quasiparticle pair and scattering amplitudes of the QRPA phonons,

$$a_k = a_k^{pair} + a_k^{scatt},$$

$$b_k = b_k^{pair} + b_k^{scatt}, \tag{18}$$

where

$$a_k^{pair} = \sum_{pn} [\sigma_{pn} X(k,pn) + \overline{\sigma}_{pn} Y(k,pn)](1 - f_n - f_p),$$

$$a_k^{scatt} = \sum_{pn} [\sigma_{pn}^V \overline{Z}(k,pn) - \sigma_{pn}^U Z(k,pn)](f_n - f_p),$$

$$b_k^{pair} = \sum_{pn} [\sigma_{pn} Y(k,pn) + \overline{\sigma}_{pn} X(k,pn)](1 - f_n - f_p),$$

$$b_k^{scatt} = \sum_{pn} [\sigma_{pn}^V Z(k,pn) - \sigma_{pn}^U \overline{Z}(k,pn)](f_n - f_p). \tag{19}$$

The transition strength is defined by

$$S^{\pm} = \sum_{k,\mu} |\langle \Gamma(k,1\mu)\beta_{1\mu}^{\pm}\rangle|^2, \tag{20}$$

and the Ikeda sum rule is given by the difference

$$\begin{aligned} S^- - S^+ &= \sum_{\mu} \langle [\beta_{1\mu}^+, (-)^{\mu}\beta_{1-\mu}^-]\rangle \\ &= \sum_{pn} [(\sigma_{pn}^2 - \overline{\sigma}_{pn}^2)(1 - f_n - f_p) \\ &\quad + (\sigma_{pn}^{U\,2} - \sigma_{pn}^{V\,2})(f_n - f_p)] \\ &= 3(N - Z). \end{aligned} \tag{21}$$

This sum rule can also be written in terms of quasiparticle pair and scattering amplitudes [see Eq. (17)] and the result is

$$\begin{aligned} S^- - S^+ &= 3\sum_k (a_k^2 - b_k^2) \\ &= 3\sum_k [(a_k^{pair})^2 - (b_k^{pair})^2 + (a_k^{scatt})^2 - (b_k^{scatt})^2 \\ &\quad + 2(a_k^{pair}a_k^{scatt} - b_k^{pair}b_k^{scatt})], \end{aligned} \tag{22}$$

and it should be compared with the conventional result that contains only pair contributions. As in the case of Fermi transitions [5] the cancellation of the interference between scattering and pair terms [the last term of Eq. (22)] is guaranteed by the orthonormalization of the QRPA phonons.

So far, the expressions listed above are general and the dependence with the Hamiltonian is given by the coefficients of Eq. (4). Concerning the use of the Hamiltonian of Eq. (1), as shown in a number of publications (see [13]) it describes a good amount of the correlations that are specific of charge-exchange channels with $J^{\pi} = 1^+$. It illustrates the main mechanism leading to the hindrance of low-energy charge-exchange transitions, namely the repulsion due to particle-hole interactions and the attraction induced by pairing and particle-particle interactions. Concerning the use of thermal averages and the extension of the standard equations of the QRPA, the procedure seems to be adequate to describe excited final states built upon a thermal reference state. Both conditions, i.e., thermal excitations and the competition between particle-hole, particle-particle, hole-hole-terms, and multiple scattering terms as they are accounted for by the extended QRPA, are the relevant ones in dealing with GT transitions in excited systems.

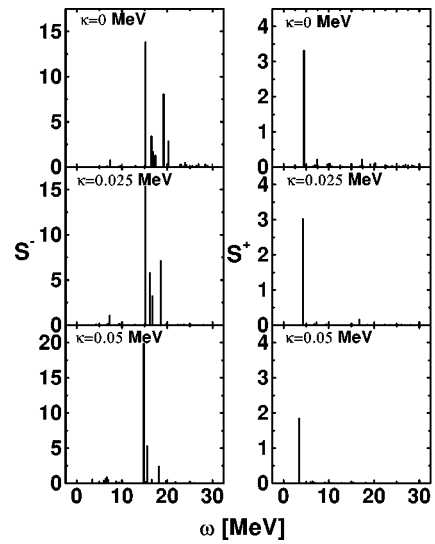

FIG. 1. Distribution of strength for β^- (S^-) and β^+ (S^+) transitions from the ground state of the nucleus ^{106}Cd. The quantity ω represents the energy of the QRPA phonons, which are excited to the 1^+ states of ^{106}In and ^{106}Ag, respectively. The value of the coupling constant κ is given in each case and the results were obtained with $T=0$ and $\chi=0.3$ MeV.

III. RESULTS AND DISCUSSION

In this section we are going to discuss the results of the above presented formalism, for the case of GT transitions from the ground state of ^{106}Cd. We have chosen this nucleus as an example of an open shell system. The single particle levels used in the calculations are harmonic oscillator levels with $N=3,4$, and 5, both for protons and neutrons. The calculations were performed for values of the temperature $0 \leqslant T \leqslant 0.5$ MeV. The pairing coupling constants G_n and G_p, which were used to solve the BCS equations, were fixed at the values $17/A$ MeV and $20/A$ MeV, respectively [12]. The position of the GT giant resonance was reproduced by fixing the coupling constant χ at zero temperature and at zero strength for the particle-particle channels. The corresponding value of χ is equal to 0.3 MeV. The renormalized coupling to particle-particle channels κ was taken as a free parameter in the range $0 \leqslant \kappa \leqslant 0.05$ MeV. By choosing values of κ in this interval a significant hindrance of the β^- and β^+ was obtained without producing the collapse of the QRPA. We have solved the QRPA equations with and without the extra terms of Eq. (8). With the help of Eqs. (18) and (19) we have calculated the transition strength of Eq. (20)

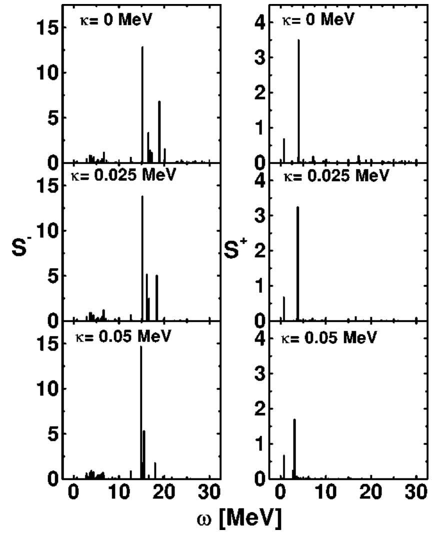

FIG. 2. Same as Fig. 1 for the temperature $T=0.5$ MeV.

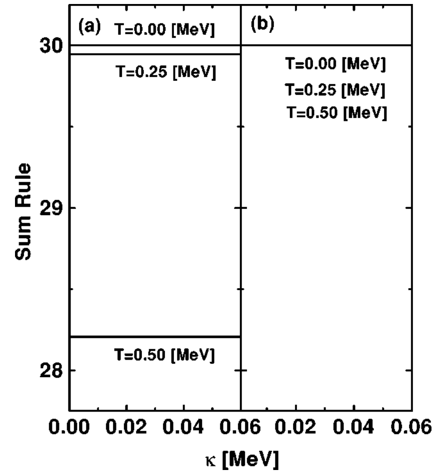

FIG. 3. Results corresponding to the Ikeda sum rule $S^- - S^+$, of Eq. (22), for different values of the temperature T and for different values of the coupling constant κ [see Eq. (1)]. The values indicated by (a) are the results of the conventional QRPA calculation. The results obtained with the present formalism are indicated by (b).

and the results are shown in Figs. 1 and 2. Figure 1 displays the distribution of strength corresponding to β^-, left-hand side of the figure and β^+, right-hand side of the figure, transitions as a function of the phonon energy ω. The calculations were performed at temperature $T=0$ MeV and the coupling constant κ takes the value 0 MeV, 0.025 MeV, and 0.05 MeV, as indicated in the figure. As is known from already published results [14] one obtains a fragmentation of the β^+ strength for increasing values of κ. The effect of the scattering terms are shown in Fig. 2, which displays the results obtained with $T=0.5$ MeV. This temperature is still lower than the critical value $T_c=0.7$ MeV, associated to the collapse of the pairing gaps. In addition to the fragmentation of the strength induced by particle-particle correlations, the inclusion of scattering terms adds up to the redistribution of the β^- and β^+ strengths at low energies. In spite of the fact that the number of proton-neutron two-quasiparticle configurations increases substantially, in going from $T=0$ MeV to $T=0.5$ MeV, the QRPA equations are still nonsingular and numerically solvable. At this point we shall address the problem of the conservation of the strength in GT transitions. The results of the present calculations are shown in Fig. 3. The left-hand side of Fig. 3 shows the temperature dependence of the standard QRPA results, also as a function of the renormalized coupling κ. It is observed that a sizable amount of the strength is lost in going from $T=0$ MeV to $T=0.5$ MeV. This is a thermal blocking mechanism that removes strength from transitions across the Fermi level. In the same figure the results corresponding to the calculations including scattering terms are shown, as a function of the temperature and for different values of the renormalized coupling constant κ. It is seen that the total strength is conserved. This is the main result of the present calculation.

IV. CONCLUSIONS

In this work we have considered the case of GT transitions at finite temperatures. We have applied an extended version of the QRPA formalism, which includes scattering terms in the definition of the phonons. It is found that thermal blocking effects, which are partly responsable for the loss of intensity, are compensated by the transitions induced by scattering terms. Like in the case of Fermi transitions the total intensity of GT transitions is conserved only if the scattering terms are included. Also, a certain amount of fragmentation leading to low energy transitions has been obtained as a consequence of thermal activation of particle-like configurations. This trend is particularly evident for the case of β^+ transitions. We think that these results may be of some significance in astrophysical calculations, where the centroids of Fermi and Gamow-Teller strength distributions are taken as inputs for leptonic capture and emission processes.

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