Nuclear Physics A404 (1983) 15–28

THERMAL ASPECTS OF THE PAIRING CORRELATIONS IN FINITE NUCLEI

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Received 20 November 1982
(Revised 14 February 1983)

Abstract: The influence of temperature on the pairing correlations is investigated. The pairing collective motion is described within the temperature-dependent RPA formalism. The specific two-particle transfer mechanism associated with the pairing degree of freedom is discussed as a function of the temperature. Applications are shown for the case of the Sn isotopes in the mass region $106 \leq A \leq 124$.

1. Introduction

Evidence concerning the description of nuclear properties at finite temperatures is growing fast. Recently [1-8]) a great deal of attention has been paid to the study of the thermal properties of nuclear systems. The thermal Hartree–Fock approximation [1,3,4]), the Hartree–Fock-Bogoliubov approximation [1-4]), and a cranked pairing model [2,3]), were studied in connection with the thermodynamic properties of finite systems. One should also mention in this line the study of the persistence of shell effects for finite temperatures which has been reported in refs. [4,9]). From the experimental point of view, the study of the thermal aspects of nuclear structure is motivated by its impact upon the current development in the field of the heavy ion reactions [10]).

In general, the available literature shows that many of the assumptions concerning the behaviour of finite nuclear systems at non-zero temperature seem to be valid, at least, in the so-called asymptotic (Fermi gas) limit [11]). A very rich domain of physical situations remains however still unexplored, as for instance the thermal aspects of the nuclear collective motion at low temperatures. This problem requires the simultaneous consideration of the well-known concepts involving the microscopic description of the collective motion and the analysis of the various assump-

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tions which are usually claimed in connection with the statistical features of the nuclear motion at non-zero temperature. Early studies of this problem focussed on the nuclear level densities at low and intermediate energies [12-14]), where the effects of the pairing interaction cannot be neglected. In this respect, the analysis of the superfluid to normal phase transition and the collapse of the pairing gap constitute relevant questions [13]).

In this work we concentrate on the study of the collective motion associated with the pairing interaction at finite temperatures, within the RPA formalism, for realistic nuclear systems. Consequently, we focus our attention on the dependence of physical observables upon the temperature, and in particular we study the two-particle transfer mechanism.

The paper is organized as follows. Sect. 2 is devoted to a brief revision of the theoretical background (subsect. 2.1) and a derivation of temperature-dependent RPA equations is presented (subsect. 2.2). Sect. 3 contains the numerical results for the tin region and the conclusions are presented in sect. 4.

2. Formalism

2.1. TEMPERATURE-DEPENDENT BCS FORMALISM

The inclusion of the thermal effects on a fermion system interacting via the monopole pairing force can be found in detail in refs. [1-3]). We only give here a brief summary.

The pairing hamiltonian is written [16])

$$H_p = -\tfrac{1}{4} G \sum_{\nu\nu'} a^+_\nu a^+_{\bar\nu} a_{\bar\nu'} a_{\nu'}\,, \tag{1}$$

where G is the pairing constant and a^+_ν (a_ν) is the creation (annihilation) operator of an independent nucleon in shell-model states characterized by the quantum numbers $\nu = \{N, l, j, m\}$, $\bar\nu$ being the time-reversed state.

The standard BCS procedure allows for the transformation of (1) to the quasiparticle basis. For $T \neq 0$, the ordinary BCS vacuum is replaced by an (average) reference state characterized by that temperature[†]. We thus define the following expectation values for the quasiparticle operators (α^+_ν, α_ν) on the reference state

$$\begin{aligned} \langle \alpha^+_\nu \alpha_{\nu'} \rangle &= f_\nu \delta_{\nu\nu'}\,, \\ \langle \alpha_\nu \alpha^+_{\nu'} \rangle &= (1 - f_\nu)\delta_{\nu\nu'}\,, \end{aligned} \tag{2}$$

f_ν being the (thermal) occupation numbers. Without the explicit consideration of any "*a priori*" guess for the occupation numbers [since they are explicitly determined through the minimization of the free energy [1-4]], one can perform the standard

[†] Equivalently one may think of a statistical ensemble defined by the partition function $Z(T, N)$ with which expectation values are evaluated [12,13,20]).

normal ordering procedure (Wick's theorem at $T \neq 0$)[19]) leading to the following expression for H_p:

$$H_p = -(\Delta^2/G) + 2\Delta \sum_\nu u_\nu v_\nu \alpha_\nu^+ \alpha_\nu - \tfrac{1}{2}\Delta \sum_\nu (u_\nu^2 - v_\nu^2)(\alpha_\nu^+ \alpha_{\bar\nu}^+ + \alpha_{\bar\nu}\alpha_\nu)$$
$$-\tfrac{1}{2}G \sum_{\nu\nu'} (\Omega_\nu \Omega_{\nu'})^{1/2} (u_\nu^2 u_{\nu'}^2 + v_\nu^2 v_{\nu'}^2)(A_\nu^+ A_{\nu'} + A_{\nu'}^+ A_\nu)$$
$$+\tfrac{1}{2}G \sum_{\nu\nu} (\Omega_\nu \Omega_{\nu'})^{1/2} (u_\nu v_{\nu'} + u_{\nu'} v_\nu)(A_\nu^+ A_{\nu'}^+ + A_\nu A_{\nu'}) , \tag{3}$$

where

$$\Delta = G \sum_\nu \Omega_\nu u_\nu v_\nu (1 - 2f_\nu) ,$$
$$\Omega_\nu = (j_\nu + \tfrac{1}{2}) . \tag{4}$$

The pair operators $A_\nu^+ (A_\nu)$ are given by

$$A_\nu^+ = \frac{1}{\sqrt{\Omega_\nu}} \sum_{m_\nu > 0} \alpha_\nu^+ \alpha_{\bar\nu}^+ . \tag{5}$$

The single-particle hamiltonian $H_{s.p.}$ can now be added to (3) to obtain

$$H' = H_{s.p.} + H_p - \lambda N , \tag{6}$$

where

$$H' = H_{00} + H_{11} + H_{20} + H_{22} + H_{40} , \tag{7}$$

and

$$H_{00} = \sum_\nu 2\Omega_\nu v_\nu^2 (\varepsilon_\nu - \lambda) - \Delta^2/G ,$$
$$H_{11} = \sum_\nu [(u_\nu^2 - v_\nu^2)(\varepsilon_\nu - \lambda) + 2\Delta u_\nu v_\nu] \alpha_\nu^+ \alpha_\nu ,$$
$$H_{20} = \sum_\nu [u_\nu v_\nu (\varepsilon_\nu - \lambda) - \tfrac{1}{2}\Delta (u_\nu^2 - v_\nu^2)](A_\nu^+ + A_\nu) ,$$
$$H_{22} = -\tfrac{1}{2}G \sum_{\nu\nu'} (\Omega_\nu \Omega_{\nu'})^{1/2} (u_\nu^2 u_{\nu'}^2 + v_\nu^2 v_{\nu'}^2)(A_\nu^+ A_{\nu'} + A_{\nu'}^+ A_\nu) ,$$
$$H_{40} = \tfrac{1}{2}G \sum_{\nu\nu'} (\Omega_\nu \Omega_{\nu'})^{1/2} (u_\nu^2 v_{\nu'}^2 + u_{\nu'}^2 v_\nu^2)(A_\nu^+ A_{\nu'}^+ + A_\nu A_{\nu'}) , \tag{8}$$

ε_ν being the single particle energies.

One should notice that Δ in eq. (8) is an explicit function of T. The functional dependence given in eq. (4) exhibits a blocking factor $(1 - 2f_\nu)$, which is responsible for the decreasing trend of $\Delta(T)$. This is a well-known feature which has been pointed out long ago [12–14]).

At this point, one can again follow the standard BCS procedure, and with the condition $H_{20}=0$ one finally obtains, together with the gap equation (4), the equations

$$E_\nu=(\Delta^2+(\varepsilon_\nu-\lambda)^2)^{1/2},$$
$$v_\nu^2=\tfrac{1}{2}(1-(\varepsilon_\nu-\lambda)/E_\nu), \qquad u_\nu^2=\tfrac{1}{2}(1+(\varepsilon_\nu-\lambda)/E_\nu),$$
$$N=\sum_\nu 2\Omega_\nu[v_\nu^2(1-f_\nu)+u_\nu^2 f_\nu], \tag{9}$$

where

$$f_\nu=(1+\exp(E_\nu/T))^{-1}. \tag{10}$$

The set of eq. (10) completely determines the properties of the quasiparticle motion at $T\neq 0$. The well-known features of the thermally induced blocking effects [12–14]) are naturally contained therein.

2.2. COLLECTIVE MOTION AT $T\neq 0$

The thermal effects (represented by the temperature T which is assumed to be externally fixed) are numerically evaluated through the introduction of the occupation numbers f. This approach relies upon the validity of the mean field (Hartree–Fock) approximation at finite temperatures. It has been shown that this assumption is still valid up to values of T of the order of the single-particle energy-spacing within a shell [4,9,11]). As it has been shown in the previous subsection, the gap parameter Δ depends upon T via the factor f_ν. The detailed feature of this dependence is discussed later in sect. 3 but at first glance, it becomes evident that for a certain T the gap collapses leading to a "superfluid to normal" phase transition [12–14]).

Thus, for certain T_c the open-shell system would display again similar characteristics than a normal one, but with a diffused Fermi surface.

The collapse of the pairing gap not only modifies strongly the properties of the reference state but also produces major changes in the structure of the excited states that can be constructed on top of it.

To describe the collective states originated by the pairing interaction the standard procedure is to diagonalize the two-quasiparticle contributions contained in H_{22} and in H_{40} within the RPA formalism [16]). Since in the present case we deal with a (thermal) reference state, the two-quasiparticle coherent contributions must be collected accounting also for the thermal effects; this leads to the appearance of blocking effects that prevent the pairs to coherently contribute to the vibration. This situation can be assimilated to the one which is found when the pairing vibrations are studied as a function of the pairing constant [18]).

In order to set up the formalism, we start with the treatment of the collective vibrations. We follow the RPA procedure [16]) of linearizing the $H_{\mathrm{RPA}}=H_{22}+H_{40}$ hamiltonian with the inclusion of the thermal occupations f_ν.

We thus define the bosons

$$\Gamma_n^+ = \sum_\nu X_{\nu,n} A_\nu^+ - Y_{\nu,n} A_\nu , \tag{11}$$

where the forward and backward amplitudes $X_{\nu,n}$ and $Y_{\nu,n}$ are given by

$$X_{\nu,n} = \Lambda_n \frac{\sqrt{\Omega_\nu}}{2E_\nu - W_n} \left(a_n(u_\nu^2 - v_\nu^2) - b_n\right),$$

$$Y_{\nu,n} = \Lambda_n \frac{\sqrt{\Omega_\nu}}{2E_\nu + W_n} \left(a_n(u_\nu^2 - v_\nu^2) + b_n\right). \tag{12}$$

The energies W_n are the collective energies which are the solutions of the determinantal equation

$$\det \begin{vmatrix} -1 + G \sum_\nu \dfrac{2\Omega_\nu E_\nu (u_\nu^2 - v_\nu^2)^2 (1-2f_\nu)}{4E_\nu^2 - W_n^2} & GW_n \sum_\nu \dfrac{\Omega_\nu (u_\nu^2 - v_\nu^2)(1-2f_\nu)}{4E_\nu^2 - W_n^2} \\ GW_n \sum_\nu \dfrac{\Omega_\nu (u_\nu^2 - v_\nu^2)(1-2f_\nu)}{4E_\nu^2 - W_n^2} & -1 + G \sum_\nu \dfrac{2E_\nu \Omega_\nu (1-2f_\nu)}{4E_\nu^2 - W_n^2} \end{vmatrix} = 0 . \tag{13}$$

In eq. (12), a_n and b_n are defined by

$$a_n = GW_n \sum_\nu \frac{(u_\nu^2 - v_\nu^2)\Omega_\nu (1-2f_\nu)}{4E_\nu^2 - W_n^2},$$

$$b_n = -1 + G \sum_\nu \frac{(u_\nu^2 - v_\nu^2)^2 2E_\nu \Omega_\nu (1-2f_\nu)}{4E_\nu^2 - W_n^2}. \tag{14}$$

The strength functions Λ_n are thus given by

$$\Lambda_n = \left\{ \sum_\nu (1-2f_\nu)\Omega_\nu \left[\frac{(a_n(u_\nu^2 - v_\nu^2) - b_n)^2}{(2E_\nu - W_n)^2} - \frac{(a_n(u_\nu^2 - v_\nu^2) + b_n)^2}{(2E_\nu + W_n)^2} \right] \right\}^{-1/2}, \tag{15}$$

with the normalization condition for the amplitudes $X_{\nu,n}$ and $Y_{\nu,n}$ written as

$$\sum_\nu (1-2f_\nu)(X_{\nu,n}^2 - Y_{\nu,n}^2) = 1 . \tag{16}$$

The inversion formulae for the pair operators A_ν^+ read

$$A_\nu^+ = (1-2f_\nu) \sum_n (X_{\nu,n} \Gamma_n^+ + Y_{\nu,n} \Gamma_n) . \tag{17}$$

These, together with the normalization condition eq. (16), obviously satisfy the commutation relation

$$\langle [A_\nu, A_{\nu'}^+] \rangle = (1-2f_\nu)\delta_{\nu\nu'} \tag{18}$$

that has been used throughout.

In order to estimate the thermal effects upon the physical observables, we look at the pair transfer operator

$$\mathcal{T}_\nu = \sqrt{2} \sum_{m_\nu > 0} a^+_\nu a^+_{\bar{\nu}} \tag{19}$$

as the specific tool [17]). When $\mathcal{T}_\nu$ is written in terms of the quasiparticle operators, the resulting expression reads

$$\mathcal{T}_\nu = \mathcal{T}^{\text{pair}}_\nu + \mathcal{T}^{\text{q.p.}}_\nu + \mathcal{T}^0_\nu , \tag{20}$$

where

$$\begin{aligned} \mathcal{T}^{\text{pair}}_\nu &= \sqrt{\Omega_\nu}(u^2_\nu A^+_\nu - v^2_\nu A_\nu) , \\ \mathcal{T}^{\text{q.p.}}_\nu &= -u_\nu v_\nu N_\nu , \\ \mathcal{T}^0_\nu &= u_\nu v_\nu \Omega_\nu , \end{aligned} \tag{21}$$

with $N_\nu = \sum_{m_\nu} \alpha^+_\nu \alpha_\nu$.

We write now $\mathcal{T}^{\text{pair}}$ in terms of the phonon operators Γ^+_n, thus obtaining

$$\mathcal{T}^{\text{pair}}_\nu = \sqrt{\Omega_\nu}(1-2f_\nu) \sum_n \{(u^2_\nu X_{\nu,n} - v^2_\nu Y_{\nu,n})\Gamma^+_n + (u^2_\nu Y_{\nu,n} - v^2_\nu X_{\nu,n})\Gamma_n\} . \tag{22}$$

As it becomes evident from eq. (22), $\mathcal{T}^{\text{pair}}$ can only connect states which differ in one pairing phonon. The matrix element of $\mathcal{T}_\nu$ connecting the reference state with one-phonon states reads

$$\langle n=1|\mathcal{T}_\nu|n=0\rangle = \langle n=1|\mathcal{T}^{\text{pair}}_\nu|n=0\rangle = \sqrt{\Omega_\nu}(1-2f_\nu)(u^2_\nu X_{\nu,n} - v^2_\nu Y_{\nu,n}) , \tag{23}$$

which naturally contains the $T=0$ limit for which $f_\nu = 0$ [refs. [17,18])].

The blocking effect is noticeable since the factor $(1-2f_\nu)$ inhibits the transfer to states nearby the Fermi surface. As it is also evident from the set of eq. (21) the expectation value of $\mathcal{T}_\nu$ on the reference state is given by

$$\langle n=0|\mathcal{T}_\nu|n=0\rangle = u_\nu v_\nu \Omega_\nu (1-2f_\nu) , \tag{24}$$

which fulfils the condition

$$\sum_\nu \langle n=0|\mathcal{T}_\nu|n=0\rangle = \sum_\nu u_\nu v_\nu \Omega_\nu (1-2f_\nu) = \Delta/G . \tag{25}$$

To summarize the results of sect. 2 the following features should be pointed out:

(i) Reference state properties: the pairing gap is obtained as an explicit function of T (eq. (4)). For a critical value T_c of the temperature the pairing gap eventually collapses allowing for a superfluid to normal phase transition.

(ii) Collective properties: The blocking effects which result from the thermal average leads to the disappearance of the low-lying collective vibrations for increasing temperatures. The same blocking effect strongly inhibits the two-particle transfer process to the reference state as well as to the excited (vibrational) states. Both

effects, the collapse of the pairing gap and the hindrance of the two-particle transfer lead to the assignment of a characteristic temperature T_c, for which the pairing correlations vanish.

TABLE 1

Single-particle energies for the neutron levels [similar to the ones of ref. [23])]

Level	Energy (MeV)
$0g_{7/2}$	0.80
$1d_{5/2}$	0.00
$1d_{3/2}$	2.80
$2s_{1/2}$	1.40
$0h_{11/2}$	2.50

3. Results and discussion

In this section, we perform the numerical applications of the formalism previously presented for the case of the tin isotopes, in the mass region $106 \leq A \leq 124$.

The single-particle orbits for the neutrons in the $N = 4$ shell are given in table 1. The effective pairing constant is fixed at the value $G = 25/A$ MeV. Fig. 1 shows the behaviour of the pairing gap Δ, as a function of the T, for the neutron numbers $N = 56$, 64 and 74 which correspond to the isotopes ^{108}Sn, ^{114}Sn and ^{124}Sn respectively. The general feature of these curves indicates the decreasing value of the pairing gap for increasing temperatures within the range $0 \leq T \leq T_c$. The T_c values for each isotope ($T_c \simeq 0.75$, 0.78 and 0.65 MeV, respectively) are approximately

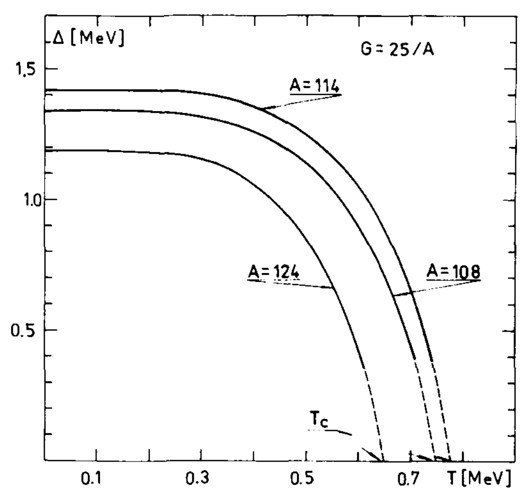

Fig. 1. Thermal variation of the gap exemplified for ^{108,114,124}Sn. The values T_c indicated in the plot are obtained by extrapolating the results of the calculation.

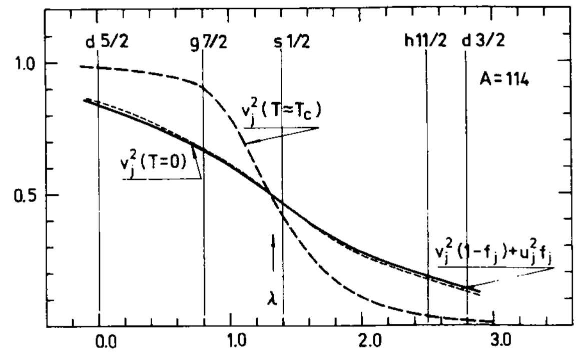

Fig. 2. Occupation numbers (v_j^2) for $T=0$ in the ^{114}Sn are displayed in full line. For $T=0.74$ MeV the occupation number $[v_j^2(1-f_j)+u_j^2f_j]$ is shown in dotted line while v_j^2 is displayed in dashed line.

given by the estimate

$$T_c \cong 0.55\Delta(T=0)\,, \tag{26}$$

which is in good agreement with the value predicted for extended systems [15]),

$$T_c = 0.567\Delta(T=0)\,. \tag{27}$$

The decreasing trend (eq. (4)) can be understood in terms of the blocking factor $(1-2f_\nu(T))$ which characterizes the effect of the thermal excitations upon the system. For relatively low values of the temperature the gap collapses, and a superfluid to normal transition appears [12,13]).

The changes in the occupation numbers with the temperature, that are another manifestation of the above-mentioned phase transition are shown in fig. 2. As expected, the major changes occur for the single-particle orbits which are nearer to the Fermi level. The noticeable aspect displayed by the results is the tendency to populate the single-particle orbits with occupation numbers, that for $T \simeq T_c$, are given by the Fermi distribution

$$n_\nu^F(T) = (1+\exp(\varepsilon_\nu - \lambda)/T)^{-1}\,. \tag{28}$$

To see this we note that

$$\lim_{\Delta\to 0}\,[v_\nu^2(1-f_\nu)+u_\nu^2 f_\nu] = n_\nu^F(T)\,.$$

Alternatively, one may think that v_ν^2 and u_ν^2 present the occupation numbers due to the presence of the pairing field while $v_\nu^2(1-f_\nu)$ and $u_\nu^2 f$ are particle and hole occupation numbers that correspond to a thermally diffused Fermi surface. The pair-correlated distribution, gives therefore rise to an open system of partially occupied levels with occupation numbers $n_\nu^F(T)$.

Another characteristic shown in fig. 2 is that the particle occupations at $T=0$ (i.e. v_ν^2) and at $T \simeq T_c$ (i.e. $v_\nu^2(1-f_\nu)+u_\nu^2 f_\nu$) are similar. For $T=0$ the slope of the particle occupations at the Fermi energy is $-(2\Delta)^{-1}$ ($T=0$) while for $T \simeq T_c$ its value is $-(4T_c)^{-1}$. Due to (27) they are very close and therefore the particle dispersions created by the pairing correlation will be of the same order of magnitude as the dispersions caused by the thermal excitations.

The occurrence of a phase transition for $T \simeq T_c$ can also be displayed by evaluating for each temperature T, the maximum value of G for which the pairing gap is zero. The gap equation (eq. (4)), if written in terms of the quasiparticle energies, results in the dispersion relation (for finite temperature)

$$\frac{1}{G} = \sum_\nu \frac{\Omega_\nu(1-2f_\nu)}{2E_\nu}. \tag{29}$$

The abovementioned critical value of the coupling constant G_c can be determined by taking the limit (for each value of T):

$$\frac{1}{G_c} = \lim_{\Delta\to 0}\left(\sum_\nu \frac{\Omega_\nu(1-2f_\nu)}{2E_\nu}\right) = \sum_\nu \frac{\Omega_\nu}{2|\varepsilon_\nu-\lambda|}\left[1-2\left(1+\exp\left(\frac{|\varepsilon_\nu-\lambda|}{T}\right)\right)^{-1}\right]. \tag{30}$$

The results of eq. (30), for the various Sn isotopes considered in this work, are shown in fig. 3. The departure from a constant value, around T_c can be attributed to the predominant role of the single-particle levels near the Fermi surface. The value of G_c increases for increasing temperature thus showing that a stronger pairing interaction is needed to overcome thermal effects. For $T \simeq T_c$, G_c approaches the value $25/A$ MeV used in the calculation thus leading to the occurrence of a (open shell) normal system.

The results presented hitherto stress the importance of the blocking "antipairing" effects played by the temperature and reflected in the factors f_ν.

The temperature has been considered up to now as an external parameter. In order to relate it to a scale of excitation energies we calculate the difference between the expectation values of H (eq. (8)) at a given T (energy of the reference state)

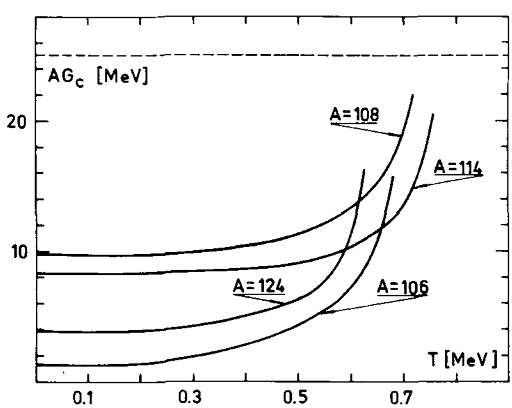

Fig. 3. The critical value of G as a function of the temperature as defined in the text (cf. eq. (30)).

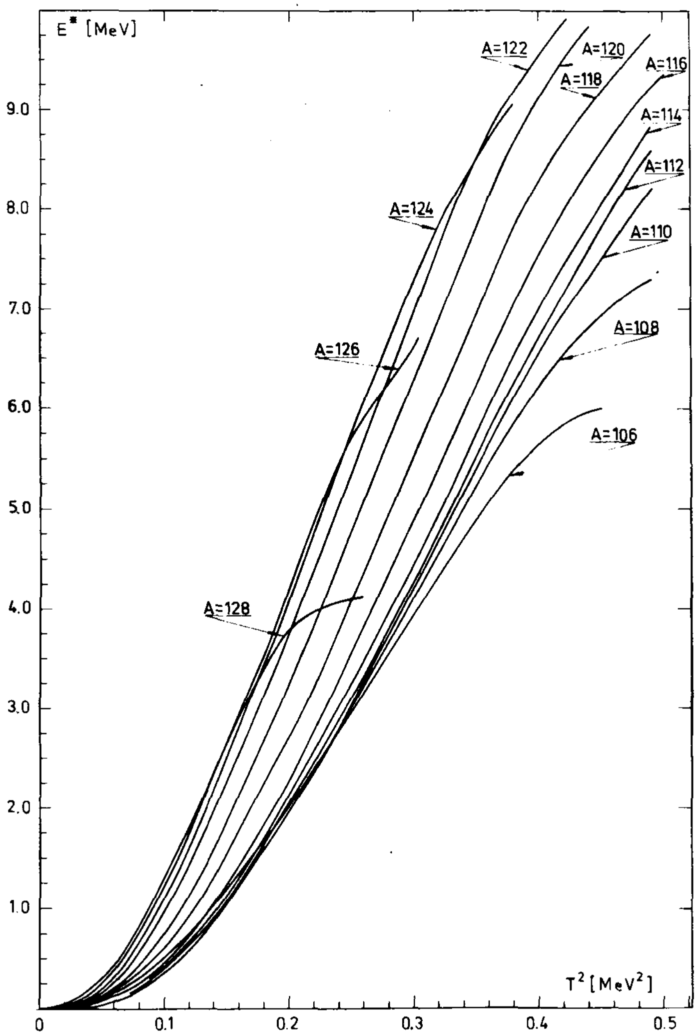

Fig. 4. Excitation energy of the reference state E^*_{RS} as a function of the temperature for even Sn isotopes.

and that at $T=0$ (ground state energy). The result is shown in fig. 4 for several Sn isotopes. This can be taken to be the scale to convert T into excitation energy E^*. Thus for $A=106$–124 the collapse of the gap is expected to occur at $E^*=6$ to 10 MeV. For values of T greater than T_c, E^* will approach the aT^2 dependence for a non-interacting Fermi gas [11]). This region can not be described within the BCS formalism that is only applicable in the superfluid phase. To that end one should instead describe the system in terms of the Fermi occupation numbers $n^F_\nu(T)$.

Figs. 1 to 4 reflect the main effect of T upon the reference state and single-quasiparticle degrees of freedom. We now turn to discuss the collective motion induced by the pairing interaction. The loss of correlations induced by the temperature also modifies the structure of the coherent two-quasiparticle states.

It should be noted that the treatment of these effects within the RPA scheme has not been explored previously. We here focus our attention on the solutions of the dispersion relation for the pairing phonons (eq. (13))†.

The variation of the collective energies W_n as a function of T is shown in fig. 5a and fig. 5b for the $a = 106$ and $A = 124$ Sn isotopes. The more drastic changes can be observed in the lowest excited states that retain the strongest collectivity. The disappearance of the first collective solution as T approaches T_c is also understood in terms of the pair-breaking mechanism previously discussed to explain the collapse of Δ.

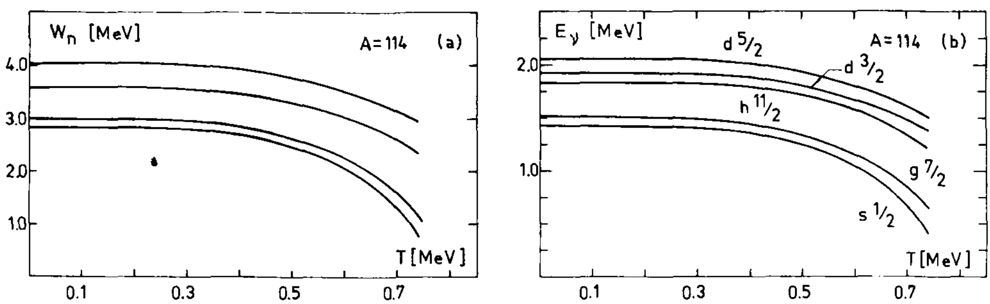

Fig. 5. Thermal variation of the excitation energies (measured with respect to the RS) of the coherent two-quasiparticle states (a) and of the one-quasiparticle states (b) for ^{114}Sn.

The values of the two-particle transfer populating the zero- and one-phonon states (eqs. (22), (23)) are shown in fig. 6. The absolute value for the transfer between the reference states of A and $A+2$ drops as $T \to T_c$ following the same behaviour as the gap (fig. (1)). The population of the coherent two-quasiparticle states is also hindered in the same fashion. The resemblance between T and G/G_c discussed in connection with fig. 3 can also be traced in the ratios

$$R_n = \frac{\sum_\nu |\langle n(A+2)|\mathcal{T}_\nu|\mathrm{RS}\,(A)\rangle|^2}{\sum_\nu |\langle \mathrm{RS}\,(A+2)|\mathcal{T}_\nu|\mathrm{RS}\,(A)\rangle|^2} \tag{31}$$

that diverge at the phase transition. This effect appears enhanced due to the limitations of the BCS approximation that is unable to reproduce accurately the region $G/G_c \approx 1$ (or equivalently $T/T_c \approx 1$). An exact treatment in which the finiteness of the system is properly accounted for would give finite values for these ratios of the order of unity for the collective state (two-pairing-phonon state) in the limit $G/G_c = 0$ ($T \to \infty$). The behaviour of the weakest state is not relevant to

† The extension of the temperature dependent RPA equation of motions for ph excitations will be reported elsewhere [21]).

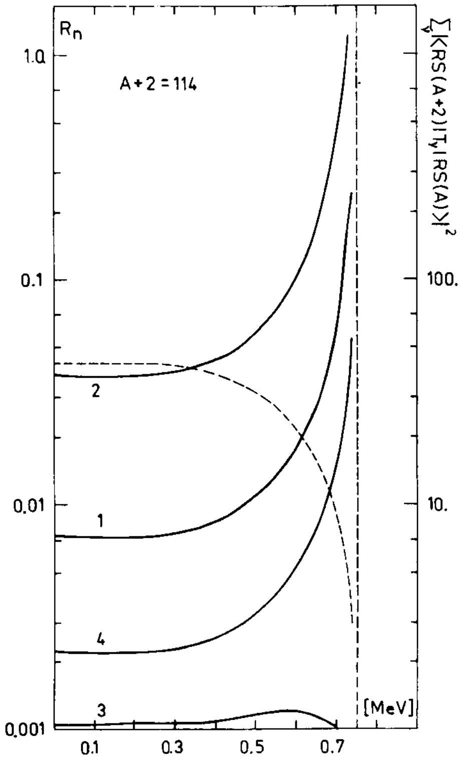

Fig. 6. Thermal variation of the two-particle transfer. The transfer between the reference states of A and $A+2$ (dashed line) is measured with the scale at the right. The ratios R_n (full lines) defined in the text (cf. eq. (31)) are measured with the scale at the left. The numbers 1, 2, etc corresponds to the label n individualizing each of the collective RPA roots. The value T_c for ^{114}Sn is indicated with a vertical dashed line.

this discussion. Its intensity drops to zero faster than the gap, as T approaches T_c, only because it retains the lowest collectivity.

4. Conclusions

We have discussed the temperature dependent BCS treatment for the Sn isotopes, in the mass region $106 \leq A \leq 124$ in what refers to both the quasiparticle and collective degrees of freedom.

The thermal treatment of pairing correlations leads, in agreement with previous estimates [1-3,12-14]), to the appearance of a phase transition from a superfluid to a normal phase. This is manifested through the collapse of the pairing gap, as a consequence of the "antipairing" character of the thermal rearrangement of particles in the s.p. levels. This destruction of the pairing correlations is also seen in the collapse of the collective two-quasiparticle excitations.

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