The width of the giant dipole resonance at finite temperature

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Abstract

We have proposed a method to evaluate the effect of the change of the Fermi sea on the width of the giant dipole resonance at finite temperature. In a schematic model we have found that indeed in ^{208}Pb the width increases very sharply up to T=4 MeV about but shows a much weaker variation for higher temperature.

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Experimental studies of giant dipole resonance (GDR) in highly excited nuclei show that the energy of the GDR is to a good precision independent of the excitation energy of the thermal equilibrated state on which it is built while its width increases strongly [1,2] at moderate excitation energy(up to roughly 130 MeV in Sn isotopes) with a saturation for higher excitation energies[3,4,5]. All theoretical studies of the GDR built on excited states of the nucleus agree to explain the energy independence of its energy but give no clear explanation of the observed increase of its width. The damping of the GDR built on ground state is known to be largely due to the coupling of particle and hole with surface collective modes[6,7,8]. However this contribution to the damping of the GDR in excited nuclei has been shown to be nearly unaffected by temperature and even to slightly decrease when the excitation energy of the nucleus increases[9]. Thermal fluctuations of the nuclear shape have also been invoked but they give a width increasing approximately as $T^{\frac{1}{2}}$ which is not enough to explain the data. Moreover it has been shown that when thermal fluctuations are treated dynamically a motional narrowing process arises which inhibits the previous result above T=1MeV about and strongly weakens the temperature dependence of the width[10]. These features are confirmed by a recent microscopic time dependent model calculation[11]. Another effect on the width which seems important comes from the angular momentum transferred to the compound nucleus. Indeed if the transferred angular momentum is high, the rapid rotation of the nucleus induces a deformation of the nucleus which can be responsible for a broadening of the resonance. This is supported by the measurements in Sn isotopes where the maximum spin that the compound nucleus can sustain before fissioning and the width saturate at similar energies[3]. Such effects have been calculated [12,13] and for very high spin ($60\hbar$ at T=2MeV in Sn) give a large dipole resonance. However all possible values of spin are transferred to the nucleus and what is observed is more likely an average over angular momenta that probably weakens the effect on the observed width.Even though this spin effect certainly exists it does not seem to be sufficient to explain the data, mainly at low temperatures. One can see this on fig.1 where we have plotted the widths measured in ^{63}Cu [14] in four different reactions, ^{4}He+^{59}Co, ^{6}Li+^{57}Fe , ^{12}C+^{51}V and ^{18}O+^{45}Sc. All the values of energies, widths, average temperatures and average spins are summarised in the table 1 of ref.[5] and we have plotted in fig.1 the width as a function of temperature (curve a) and angular momentum (curve b) . In each plot the values of $< J >$ for curve a and $< T >$ for curve b are indicated for each experimental point. These plots seem to indicate that indeed high angular momenta play an important role on the width but that one needs further effects, directly related to temperature, to explain the variation of the width over the all domain.

In the present work we look for a different phenomenon which could also be responsible for the broadening of the resonance. Namely , when the temperature is different from zero , new particule-hole configurations due to the changing of occupation numbers arise. These new configurations have a quite broad energy spectrum what may increase the Landau damping of the RPA resonance. This was already suggested by the results of ref.[15] where RPA equations were solved for ^{40}Ca in a particle-hole subspace changing with temperature and including all allowed $1\hbar\omega$ particle-hole states. At $T >$ 1-2 MeV, instead of one major peak, we got several smaller peaks apart from the main one. It was checked that the spreading over several peaks was due to the spreading of particle-hole energies by assuming all these unperturbed energies equal in which case we recovered one peak only whatever was the temperature. This result, however, was not confirmed by the work of Sagawa and Bertsh[16]. In a self-consistent RPA calculation including continuum states their strength function for the GDR of ^{40}Ca does not change with increasing temperature. It is not easy to understand why the results of these two works are so different. In the litterature it has been attributed to self-consistency but, even though self-consistency plays an important role, there is no reason why it will wash out peaks. On the contrary, at zero temperature, self consistent RPA gives in general less collectivity and a larger Landau damping due to the spreading of the strength over particle-hole states than non self consistent RPA. In ref.[16] the dipole resonance is already very broad at $T = 0$ and one could say that new peaks are hidden in this large resonance. This is not a very satisfying argument because the strength functions given at $T = 0, 3, 6$ MeV are very similar except for a larger strength appearing at low energy for $T = 6$ MeV. This difference could also come from the continuum states inclusion. Indeed the new particle-hole configurations are in ref.[15] (1f-2p)$^{-1}$ (1g-2d-3s)^{1} while in ref.[16] they correspond to a hole in the 1f-2p shell and a particle in the continuum which could not perhaps generate new peaks but could affect the distribution of the strength. Therefore we think that the apparent difference between these two works does not allow to make a definitive statement about the important or non important effect of the increasing of particle-hole subspace on RPA results and in the following we concentrate our attention on the effect on the GDR energy and width of these new configurations which have energies distributed over a quite large domain.

We proceed as follows. We first define $\mathcal{E}_0$, the particle-hole subspace at $T = 0$ and $\Delta\mathcal{E}_T$ the subspace formed of the new particle-hole states allowed at a given finite temperature T. $\mathcal{E}_0$ and $\Delta\mathcal{E}_T$ are defined in Fig.2. At $T \neq 0$ the RPA particle -hole subspace is $\mathcal{E}_0 + \Delta\mathcal{E}_T$.

Let's assume that the RPA equations have been solved for any temperature in the subspace $\mathcal{E}_0$ defining a set of independent phonons q of energy ω_q at each temperature. Then we will calculate the effect on the phonons q of their interactions with the particle-hole pairs of $\Delta\mathcal{E}_T$. This problem is analogous to treating the interaction between phonons and electrons in statistical physics and we now follow the books of Abrikosov et al[17] and Mahan[18]. We use the finite temperature Green's functions introduced by Matsubara[19] for imaginary times and frequencies. They can be calculated by the same method of Feynman diagram technics as established for zero temperature Green's functions. If $D_q^{(0)}$ is the propagator of our free phonon q, we calculate the lowest order correction to $D_q^{(0)}$ as represented in Fig.3 leading to the perturbed phonon propagator $D_q^{(1)}$ defined by the equation:

$$D_q^{(1)}(i\omega_n) = D_q^{(0)}(i\omega_n) + D_q^{(0)}(i\omega_n)\,\Pi_q^{(1)}(i\omega_n)\,D_q^{(0)}(i\omega_n) \qquad (1)$$

where $\Pi_q^{(1)}$ is the lowest order contribution to the phonon mass operator and is given by:

$$\Pi_q^{(1)}(i\omega_n) = \frac{1}{\beta}\sum_{\alpha,\beta\in\Delta\mathcal{E}_T}\sum_{ip_n}\mathcal{G}_\alpha^{(0)}(i\omega_n + ip_n)\,\mathcal{G}_\beta^{(0)}(ip_n)\,|V(\alpha\beta;q)|^2 \qquad (2)$$

β is the inverse of the temperature expressed in MeV. The frequencies are for the phonon $\omega_n = 2n\pi/\beta$ and for the fermions $p_n = (2n+1)\pi/\beta$, the $\mathcal{G}^{(0)}$'s are the Hartree Fock one particle Green's function and $V(\alpha\beta;q)$ the interaction between a particle-hole pair $(\alpha\beta)$ and the free phonon q. The summation over ip_n in eq.(2) is easily performed [18] and leads to:

$$\Pi_q^1(i\omega_n) = \sum_{\alpha,\beta\in\Delta\mathcal{E}_T}|V(\alpha\beta;q)|^2\,\frac{n_\beta - n_\alpha}{i\omega_n - \epsilon_\alpha + \epsilon_\beta} \qquad (3)$$

n_i is the occupation number of a particle of energy ϵ_i and is given by:

$$n_i = \frac{1}{1 + exp(\beta(\epsilon_i - \mu))}$$

More generally the phonon Green's function obeys the Dyson's equation:

$$D(i\omega_n) = D^{(0)}(i\omega_n) + D^{(0)}(i\omega_n)\,\Pi(i\omega_n)D(i\omega_n) \qquad (4)$$

where Π is the exact phonon mass operator.Eq.(4) writes with short notations:

$$D = \frac{1}{D^{(0)-1} - \Pi} \tag{5}$$

where $D^{(0)}$, the unperturbed phonon propagator is:

$$D_q^{(0)}(i\omega_n) = \frac{2\omega_q}{(i\omega_n)^2 - \omega_q^2} \tag{6}$$

ω_q is the unperturbed energy of the phonon q. Substituting eq.(6) into eq.(5) we get:

$$D(i\omega_n) = \frac{2\omega_q}{(i\omega_n)^2 - \omega_q^2 - 2\omega_q\,\Pi(i\omega_n)} \tag{7}$$

Let's assume that the only effect of the phonon self energy is to change the unperturbed frequencies ω_q into a new set of renormalised frequencies Ω_q so that we may write:

$$D_q(i\omega_n) = \frac{2\omega_q}{(i\omega_n)^2 - \Omega_q^2} \tag{8}$$

Ω_q is then calculated as the pole of the retarded phonon Green's function obtained from $D_q(i\omega_n)$ by analytic continuation on the real axis, that is to say by replacing in eq.(8) $i\omega_n$ by $\omega + i\eta(\eta \to +0)$. This gives a general equation for the renormalised phonon energy:

$$\Omega_q^2 = \omega_q^2 + 2\omega_q\,\Pi(\Omega_q) \tag{9}$$

To go from this general expression to our problem, therefore from eq.(4) to our eq.(1) we first replace the self energy Π by its lowest order contribution $\Pi^{(1)}$ of eq.(3), then we replace Ω_q in the right hand side of eq.(9) by ω_q what is equivalent to replace D by $D^{(0)}$ in the right hand side of eq.(4) to get our eq.(1). Therefore the energy of the phonon q modified by the contribution of the diagram of Fig.3 is determined by the following equation:

$$\Omega_q^2 = \omega_q^2 + 2\omega_q \lim_{\eta\to+0} \sum_{\alpha\beta\in\Delta\mathcal{E}_T} |V(\alpha\beta;q)|^2 \frac{n_\beta - n_\alpha}{\omega_q - \epsilon_\alpha + \epsilon_\beta + i\eta} \tag{10}$$

This is the general equation defining the new energies of the phonon q when the diagram of Fig.3 is taken into account. The purpose of this paper is to show that the dispersion of energies $\epsilon_\alpha - \epsilon_\beta$ gives rise to an imaginary contribution to Ω_q which is equivalent to say that the phonon strength function exhibits a damping width equal to twice the imaginary part of Ω_q with a Lorentzian distribution. To do this we will use a schematic model in which the RPA equations can be solved exactly and analytically in both subspaces $\mathcal{E}_0$ and $\mathcal{E}_0 + \Delta\mathcal{E}_T$. This well known schematic model relies on the assumptions that within a major shell all states have the same energy, that shells are equidistant with equal degeneracy and that particle-hole matrix element are separable and independent of the states $(\alpha\beta)$. For dipole phonons we assume a dipole-dipole interaction and write:

$$V(\alpha_i\beta_i;\alpha_j\beta_j) \approx \lambda D(\alpha_i\beta_i)\, D(\alpha_j\beta_j) \approx \lambda D^2 \tag{11}$$

We further assume that the Hartree Fock energies are independent of temperature.

In this model the eigenvalues of the RPA equations at finite temperature when they are solved in the $\mathcal{E}_0$ subspace (constructed with the α_0 and β_0 states of Fig.2) are simply given by[20]:

$$\omega_q^2 = \epsilon^2(1 + C\delta n_0) \tag{12}$$

where

$$\epsilon = \epsilon_\alpha - \epsilon_\beta \ , \ \delta n_0 = n_{\beta_0} - n_{\alpha_0} = \tanh \beta\epsilon/4 \ , \ C = \frac{2\lambda D^2}{\epsilon}$$

We note that ω_q, the energy of our unperturbed phonon of eq.(1) , has a strong energy dependence.

When the RPA equations are solved in the total subspace $\mathcal{E}_0 + \Delta\mathcal{E}_T$ one gets[15] :

$$\omega'^2_q = \epsilon^2(1 + C) \tag{13}$$

which is independent of temperature in agreement with more realistic RPA calculations and experiments. The comparison of eqs.(12) and (13) shows the importance of working with the complete particle-hole subspace.

The interaction $V(\alpha\beta;q)$ of eq.(10) where the phonon q is the collective solution of RPA solved in $\mathcal{E}_0$ and where the particle-hole state $\alpha\beta$ belong to $\Delta\mathcal{E}_T$ is easily calculated with eq.(11) and found to be:

$$V(\alpha\beta;q) = \lambda D^2 \sqrt{\frac{\epsilon\, \delta n_0}{\omega_q}}$$

The renormalised phonon energy of eq.(10) is then simply given by:

$$\begin{aligned} \Omega_q^2 &= \omega_q^2 + 2\epsilon\lambda^2 D^4\, \delta n_0 \left\{ \frac{\sum_{\alpha>\beta} n_{\alpha\beta}}{\omega_q - \epsilon + i\eta} + \frac{\sum_{\alpha>\beta} n_{\beta\alpha}}{\omega_q + \epsilon + i\eta} \right\} \\ n_{\alpha\beta} &= n_\beta - n_\alpha \end{aligned} \tag{14}$$

The summation over $\alpha\beta$ has been restricted to states such that $\epsilon_\alpha > \epsilon_\beta$, consequently the second term in the bracket corresponds to the backward going diagram. $\sum_{\alpha>\beta} n_{\alpha\beta}$ is easily calculated: the summation starts with the highest "occupied" state with $n_\beta = 1$, then all intermediate occupation numbers cancel up to n_{β_0} since the $(\alpha_0\beta_0)$ states should be eliminated from the summation; it starts again with n_{α_0}, all further occupation numbers cancelling with the summation ending with the first "unoccupied" state with $n = 0$. This gives:

$$\sum_{\alpha>\beta} n_{\alpha\beta} = -\sum_{\alpha>\beta} n_{\beta\alpha} = 1 - n_{\beta_0} + n_{\alpha_0} = 1 - \delta n_0 \tag{15}$$

Ω_q^2 of eq.(14) has an imaginary part obtained by writing:

$$\begin{aligned} \frac{1}{\omega_q \pm \epsilon + i\eta} &= \mathcal{P}\frac{1}{\omega_q \pm \epsilon} - i\pi\delta(\omega_q \pm \epsilon) \\ &= \mathcal{P}\frac{1}{\omega_q \pm \epsilon} - i\pi lim_{\gamma\to+0}\frac{\gamma}{(\omega_q \pm \epsilon)^2 + \gamma^2} \end{aligned} \tag{16}$$

where $\mathcal{P}$ means the principal value. Taking into account that the energies of the $(\alpha\beta)$ configurations do not have a well defined energy ϵ but energies distributed around ϵ can be achieved by replacing the small γ in eq.(16) by a finite quantity which means that the $(\alpha\beta)$ states have a Breit Wigner distribution centered at $\epsilon_\alpha - \epsilon_\beta = \epsilon$ with a width 2γ. Let's note that the same result could have been obtained by replacing $\pm\epsilon$ by $\pm\epsilon - i\gamma$ in eq.(14).

With eqs.(14-16) we may now write our renormalised phonon energy as $\Omega_q - i\Gamma_q$. Ω_q is the energy of the phonon and $\Gamma = 2\Gamma_q$ is the width of the Lorentzian phonon strength function. They are given by:

$$\begin{aligned} \Omega_q^2 = & \ \omega_q^2 + \frac{1}{2} C^2 \epsilon^3 \delta n_0 \left[\frac{1}{\omega_q - \epsilon} - \frac{1}{\omega_q + \epsilon} \right] = \omega'^2_q \\ \Gamma = & \ 2\Gamma_q \\ = & \ \frac{1}{2\Omega_q} C^2 \epsilon^3 \delta n_0 (1 - \delta n_0) \left[\frac{\gamma}{(\omega_q - \epsilon)^2 + \gamma^2} - \frac{\gamma}{\omega_q + \epsilon)^2 + \gamma^2} \right] \end{aligned} \tag{17}$$

We see that the renormalised phonon energy is, in our schematic model, equal to the exact RPA phonon energy calculated in the total subspace and given by eq.(13). This is an interesting feature of our model of eq.(1) and we shall go back to this property later on.

These simple formulas are now applied to the giant dipole resonance of ^{208}Pb. The parameters appearing in eq.(17) are $\epsilon = \hbar\omega \simeq 7$MeV and C$=2\lambda D^2/\epsilon = 3$ such that the GDR has an energy of 14MeV at $T = 0$. The width γ of the particle-hole states of subspace $\Delta\mathcal{E}_T$ has to be evaluated. In ^{208}Pb the number of such states is very high and we have looked at ^{40}Ca nucleus. In Fig.4 are plotted the values $|D_{\alpha\beta}|^2$ of the dipole strength for each configuration constructed from a hole in the (1f-2p) shell and a particle in the (3s-2d-1g) shell as a function of their energy. From the corresponding points we have constructed an histogram by summing the $|D_{\alpha\beta}|^2$ within energy domains of 2 MeV. We see that the dispersion of energies is quite large and if we qualitatively adjust a Breit Wigner distribution centered at $\epsilon_\alpha - \epsilon_\beta = \epsilon \simeq 13$ MeV, we get $\gamma \simeq 3$MeV. This value is adopted for ^{208}Pb but we shall see that the results are not very sensitive to small changes of

γ.

The results for Ω_q and Γ are given in table 1 for temperatures between 1 and 6 MeV. We have plotted the values of Ω_q to remind that they are equal to the RPA energies. The values of Γ given in the table are drawn as a function of T in our Fig.5. We see that it increases quite strongly at lower temperature with a shallower behaviour above T=4 MeV about. If we remember that at T=0 the total width of the GDR in ^{208}Pb is 4 MeV about, this correction is quite large, even though our calculation is schematic and absolute values of Γ should be taken with some care. Indeed the schematic model put all the collectivity in one state only and then always overestimates correlations effects. However it has always been very successful in explaining general properties like RPA energies[15,21] or temperature dependence of level density parameter[22]. In the Fig.5 we have also shown the values of Γ obtained for γ = 2.5 MeV. We see that there is not much sensitivity to γ.

For $^{108-110}$Sn , the measured widths at E=130 and 240 MeV are similar. If we adopt the value a=A/8 for the level density parameter, it corresponds to T=3 and 4.2 MeV respectively but if we follow experimental measurement of a [23] which says that $a \simeq A/13$ for T >3-4 MeV we get T=3.6 and 5.1 MeV and this is compatible with the curve of Fig.5

As already mentionned our values of Ω_q which are calculated in second order perturbation theory applied to the phonon Green's function are equal to the exact solution of the RPA which sums up all particle-hole bubble diagrams to all orders. This result which has been obtained in a simple model and is a consequence of relation(15) is expected to be nearly preserved in a more realistic calculation of eq.(10). It can be understood by looking at the diagrams involved in both models. Up to second order the diagram expansions of both, our phonon Green's function and the RPA one, are identical. The third order diagrams corresponding to our $D_q^{(1)}$ are shown in Fig.6 where $(\alpha'\beta')$, the particle-hole intermediate states, belong to $\mathcal{E}_0 + \Delta\mathcal{E}_T$, therefore can be either the $(\alpha_0\beta_0)$ or the $(\alpha\beta)$ states of our model. For simplicity we have represented upward going diagrams only but the comparison is the same for all "time" orderings. We see that our Green's function differs from the RPA's one (diagram a) by a diagram (diagram b) proportionnal to $(n_{\alpha\beta})^2 = (n_\beta - n_\alpha)^2$ which is always a small quantity. One can verify that in n^{th} order, the two Green's functions differ by diagrams proportionnal to $(n_{\alpha\beta})^2$,$(n_{\alpha\beta})^3$....$(n_{\alpha\beta})^{n-1}$. Therefore $D_q^{(1)}$ should be a good approximation to the exact RPA propagator. This can be an interesting feature of our model of eqs(1-10) which could perhaps be used to calculate properties of hot nuclei since $D^{(1)}$ is much easier to calculate than G_{RPA}.

We have found that at finite temperature the new particle-hole configurations induce a large broadening of the Landau type width of the GDR with a behaviour in terms of temperature similar to what is observed. However one may ask a further question: in such a model what would be the Landau damping due to the dispersion of the particle-hole energies in the $\mathcal{E}_0$ subspace. This question can be qualitatively answered by a simple evaluation, even though perhaps not rigorous. Indeed we may assume that in eq.(12) the energy ϵ is replaced by $\epsilon - i\gamma_0$ where γ_0 is the half-width of a Breit Wigner distribution of the dipole strength in $\mathcal{E}_0$ as shown in Fig.7 for ^{40}Ca. The corresponding width of the phonon ω_q is for a temperature T(zero or finite):

$$\Gamma_{q0} = (1 + \frac{1}{2} C\,\delta n_0)((1 + C\,\delta n_0)^{-\frac{1}{2}}\,\gamma_0 = f(T)\,\gamma_0 \tag{18}$$

In table 2 we give the coefficient f (T) in ^{208}Pb for T between 0 and 6 MeV. We see that f has a weak temperature dependence and decreases slightly with increasing temperature what, if we calculate the total GDR width due to Landau damping, will just have as consequence to attenuate slightly the increase of the GDR width at large T. Note that, assuming $\gamma_0 \simeq 0.5$ MeV accordingly to Fig.7, gives Γ_0, the Landau damping width of the GDR at T=0, equal to 1.25 MeV about what is compatible with what is usually admitted[7]. This result could perhaps be interpreted as a further test of the reliability of such a derivation.

Our result is not in contradiction with the fact that angular momentum plays an important role in the broadening of the GDR. The effects of angular momentum are certainly important for heavy systems and high temperatures for which the transferred angular momenta are large. The two effects seem to us complementary and we are convinced that both contribute to the width . In fact, as shown on the curves of Fig.1 where we have summarised all measurements on ^{63}Cu, both are necessary to explain experimental observations. A further remark concerns the possible overlap between what we have calculated in this paper and a contribution to the width due to a deformation of the nucleus. Indeed the temperature, by changing the occupation numbers, creates "quasi particles" so that the highly excited nuclear states are mixtures of multiparticle-multihole states when referred to the ground state and one knows that the excitation of a certain number of nucleons induces a deformation of the nucleus[24,25]. This is supported by the results of finite temperature RPA. Indeed, as soon as the temperature is finite, the RPA spectrum of ^{40}Ca exhibits a low lying 2^+ state[16,26] which, following the results of refs.[24,25], is a deformed state with very large 4p-4h components.

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T (MeV)	1	2	3	4	5	6
ω_q(MeV)	13.7	12.2	11.2	10.5	9.9	9.5
Ω_q(MeV)	14	14	14	14	14	14
Γ(MeV)	0.30	1.66	2.82	3.55	3.95	4.1

Table 1: The energies ω_q of our free dipole phonon and Ω_q of the renormalised phonon and the width Γ of the dipole resonance for temperatures between 1 and 6 MeV in ^{208}Pb.

T (MeV)	0	1	2	3	4	5	6
$f(T)$	1.25	1.23	1.16	1.11	1.08	1.06	1.05

Table 2: The ratio between the width of the RPA dipole resonance calculated in $\mathcal{E}_0$ and the width of the particle-hole dipole strength function in the same subspace $\mathcal{E}_0$ in ^{208}Pb.

Figure caption

Fig.1 The measured GDR width in ^{63}Cu: as a function of $< T >$ (curve a) and $< J >$ (curve b). To each point are associated the corresponding values of $< J >$ in curve a and of $< T >$ in curve b.

Fig.2 Definition of the subspaces $\mathcal{E}_0$ and $\Delta\mathcal{E}_T$.

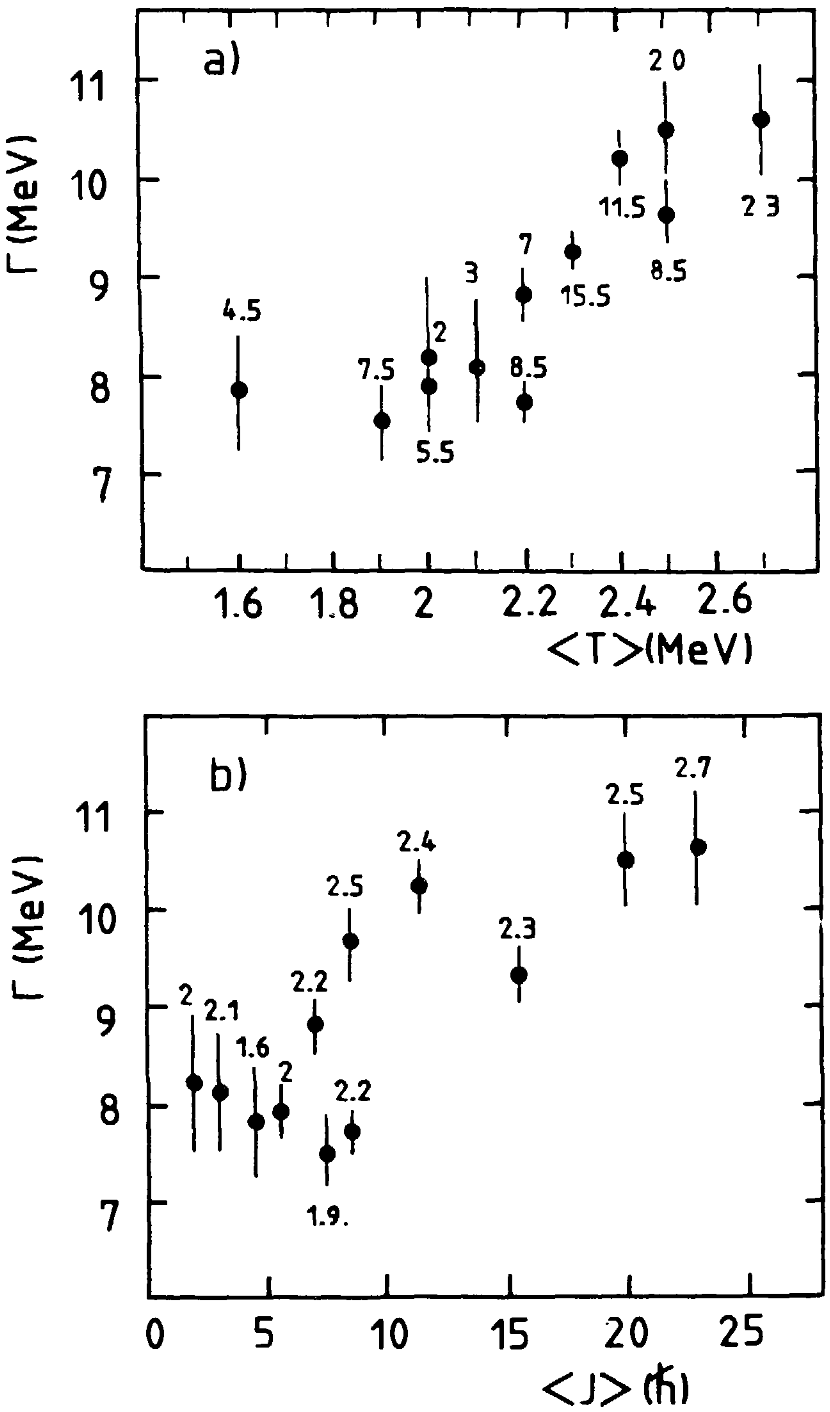
Fig.3 The second order diagram corresponding to our eq.(1). The points represent the interaction $V(\alpha\beta; q)$

Fig.4 The dipole strength function $|D_{\alpha\beta}|^2$ as a function of $\epsilon_\alpha - \epsilon_\beta$, the particle-hole energies in $\Delta\mathcal{E}_T$ for ^{40}Ca.

Fig.5 The GDR width of eq.(17) as a function of temperature in ^{208}Pb. The points correspond to $\gamma = 3$ MeV and the crosses to $\gamma = 2.5$ MeV

Fig.6 Third order diagrams included in our phonon Green's function of eq.(1) in terms of particle and hole propagators.

Fig.7 The dipole strength function $|D_{\alpha_0\beta_0}|^2$ as a function of $\epsilon_{\alpha_0} - \epsilon_{\beta_0}$, the particle-hole energies in $\mathcal{E}_0$ for ^{40}Ca.



a)
Γ (MeV)
11
10
9
8
7
4.5
7.5
2
5.5
3
7
8.5
15.5
11.5
2 0
8.5
2 3
1.6
1.8
2
2.2
2.4
2.6
<T>(MeV)
b)
Γ (MeV)
2
2.1
1.6
2
2.2
2.5
2.4
2.2
1.9.
2.3
2.5
2.7
0
5
10
15
20
25
<J> (ħ)

Fig. 1

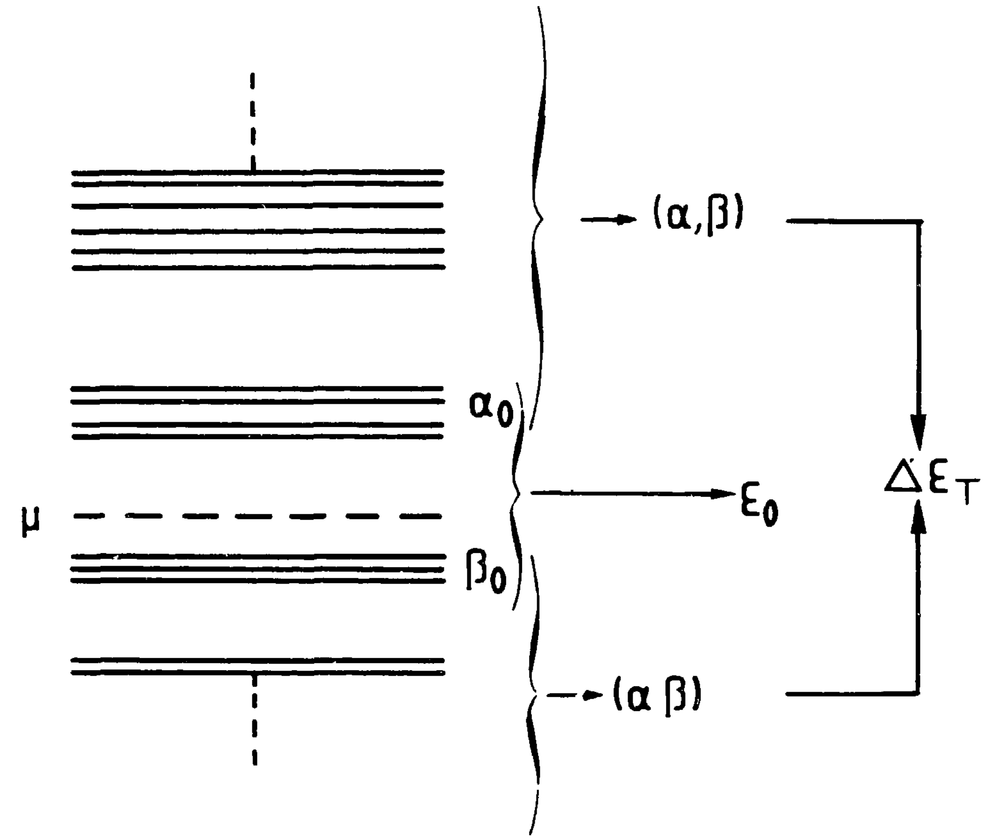


Fig. 2

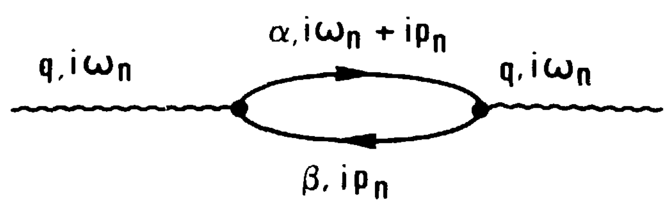


Fig. 3

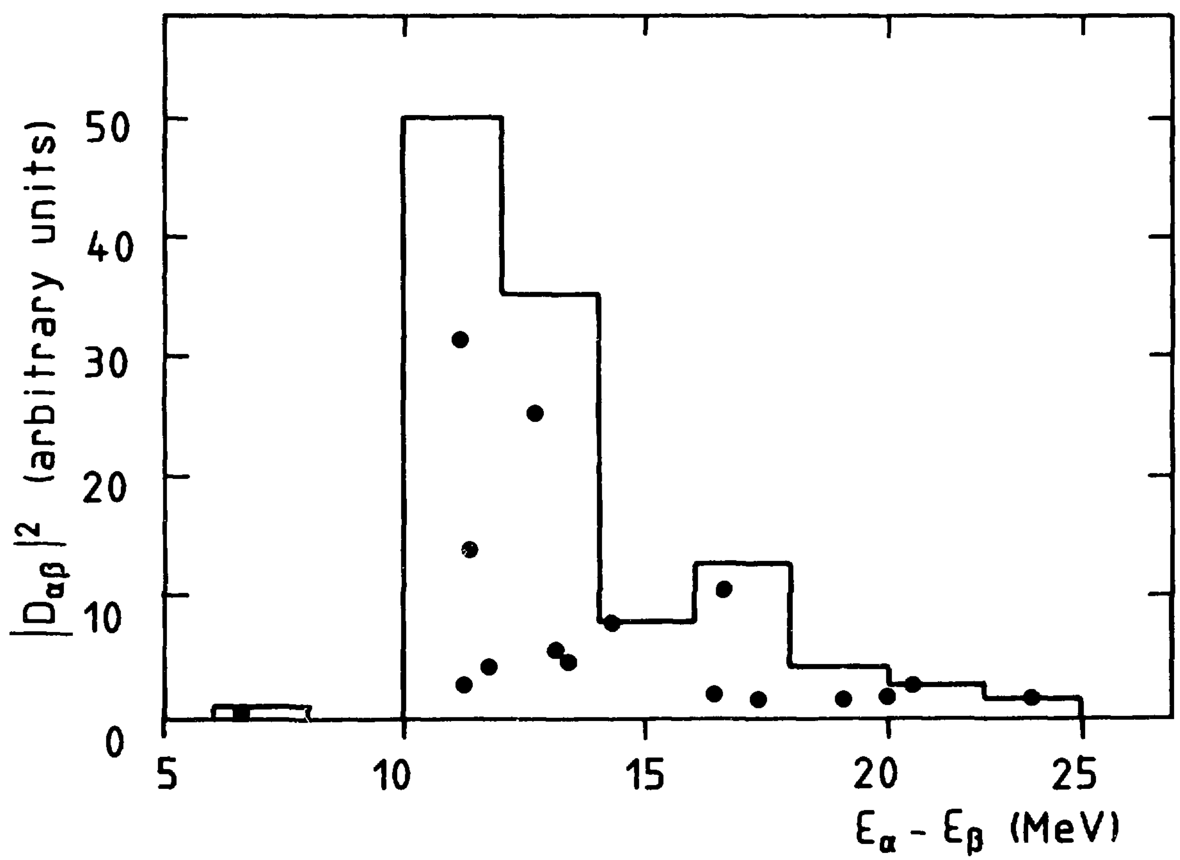

$|D_{\alpha\beta}|^2$ (arbitrary units)
0
10
20
30
40
50
5
10
15
20
25
$E_\alpha - E_\beta$ (MeV)

Fig. 4

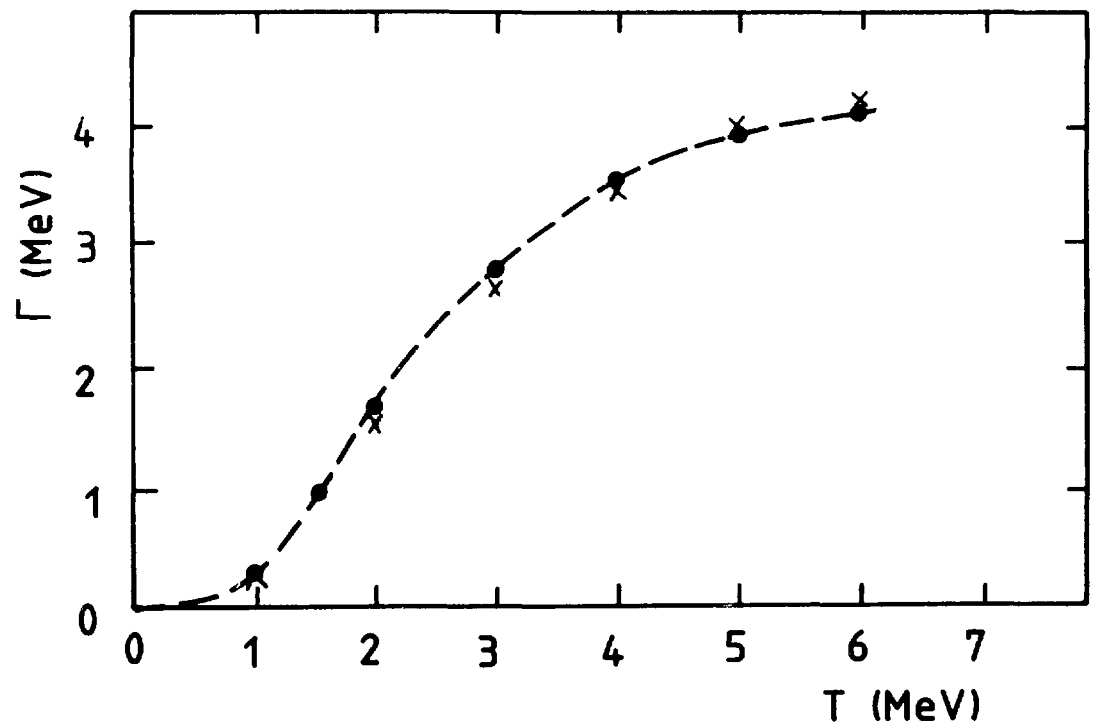
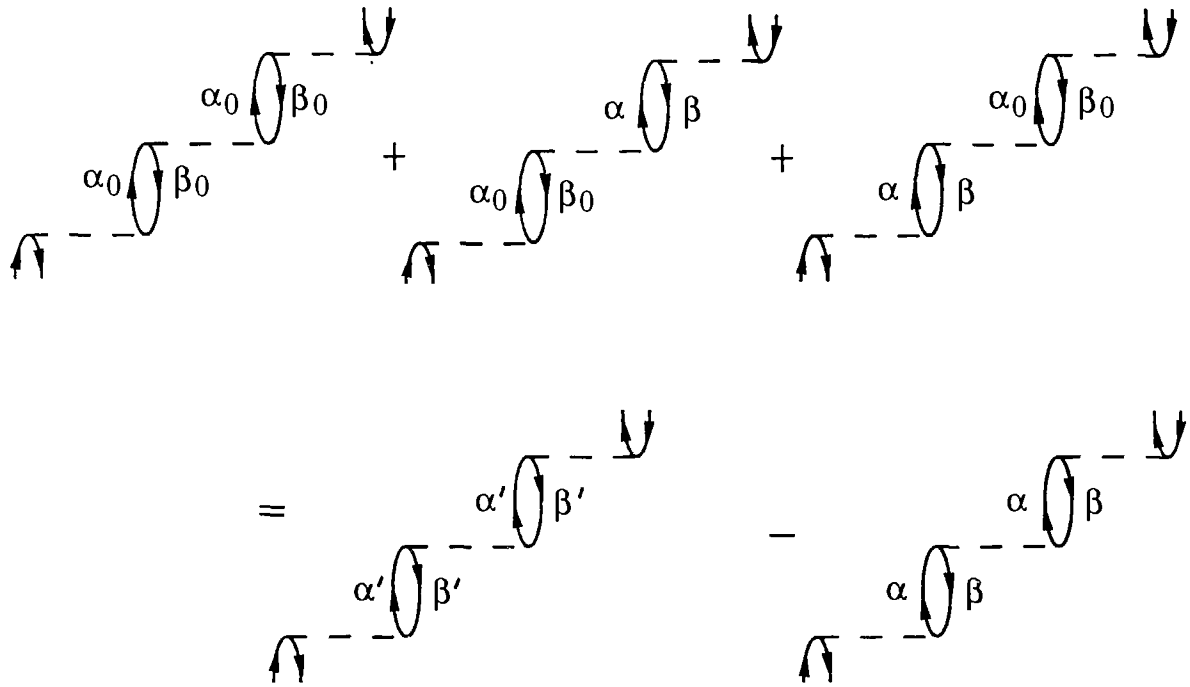


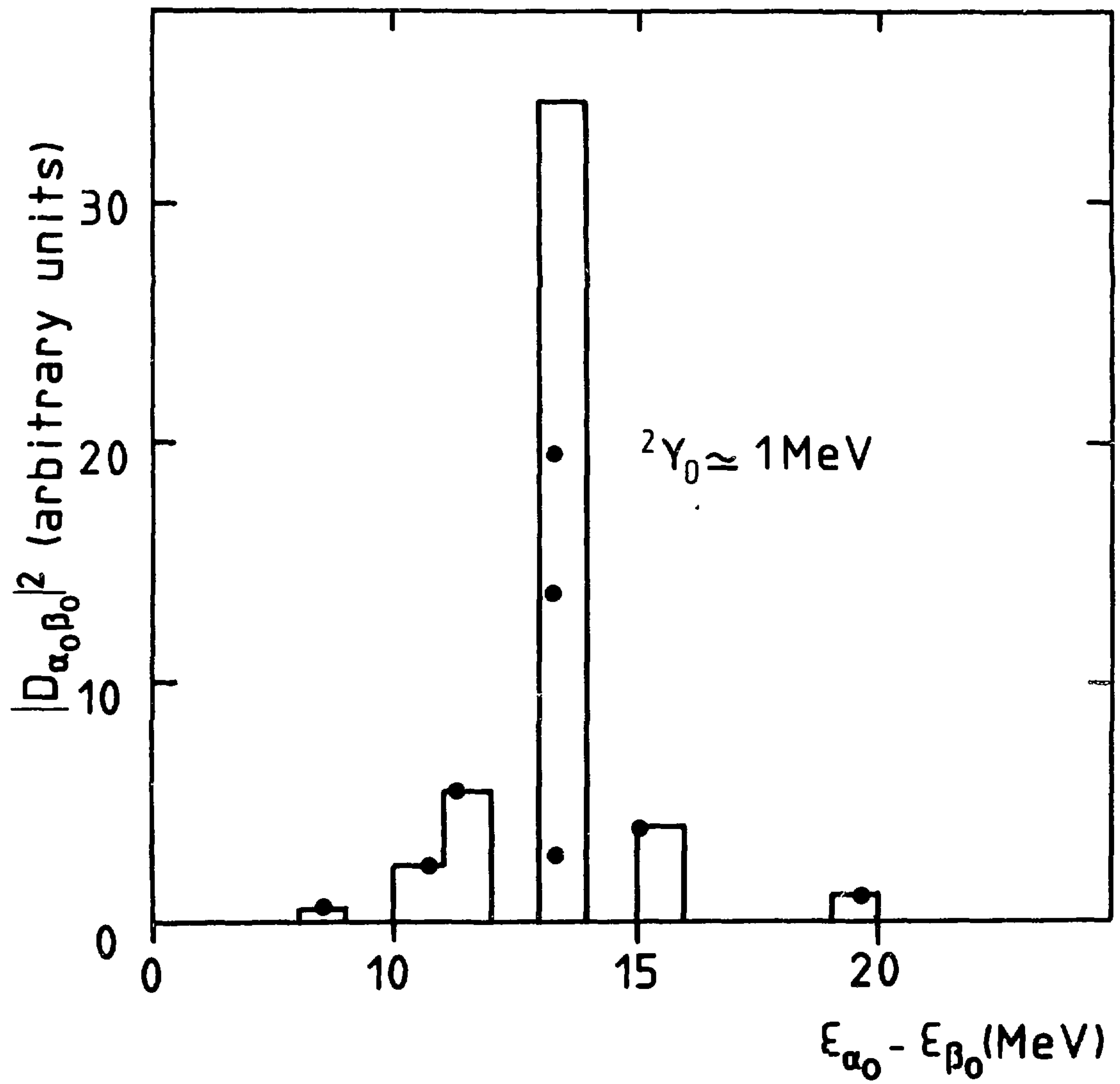
Fig. 5



α_0 β_0
α_0 β_0
+
α β
α_0 β_0
+
α_0 β_0
α β
=
α' β'
α' β'
−
α β
α β

(a) (b)

Fig. 6



$|D_{\alpha_0\beta_0}|^2$ (arbitrary units)
0
10
20
30
0
10
15
20
$^2Y_0 \simeq 1$MeV
$\varepsilon_{\alpha_0} - \varepsilon_{\beta_0}$(MeV)

Fig. 7