The coupling to doorway states and to compound states in hot nuclei

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The different role the doorway states and the compound states play (in the regime of chaotic intrinsic dynamics) in the damping of single-particle states and giant resonances at finite temperature, is discussed.

Introduction

The problem of the variety of mechanisms at work in the damping of collective motion of many-body systems is a fascinating one, and central in a variety of fields ranging from condensed matter to atomic nuclei.

According to the topics of this symposium and workshop, the main features of the damping in hot nuclei will reviewed, in particular for the Giant Dipole Resonance (GDR), following Ref. 1. See Ref. 2 and these proceedings[3)] for the most recent achivements in the experiments and theory.

The properties of the GDR are experimentally known at finite temperature T and angular momentum from the γ-decay of a compound nucleus produced in inelastic scattering and/or fusion reactions. In general terms, it is found that the centroid energy remains constant, while the total width increases at value of the order of 10–12 MeV at temperature of the order of 2.5–3 MeV in nuclei as ^{120}Sn or ^{208}Pb.

At the mean field level, the inhomogeneous damping mechanism is at work, reflecting the spatial anisotropy of the system. Indeed, we know that the energy centroid of the GDR is inversely proportional to the nuclear radius. In a spherical nucleus, the vibrations along any three perpendicular axis are equivalent ($\hbar\omega \approx 79/A^{1/3}$ MeV at $T = 0$), and the strength is equally distributed. In deformed nuclei (permanent deformation or deformation induced by temperature and/or angular momentum), the dipole strength splits. Each of the peak can undergo further splitting or eventually acquire a width, still in a pure mean field description of the nucleus, due to the following mechanisms:
(1) decay into single-particle motion, that is in uncorrelated particle-hole excitations, a phenomenon known as Landau damping in solid state physics.[4)] The relative contribution to the total width Γ of the giant resonances will depend on the particular mode and nucleus.
(2) Particle-decay: giant resonances lie above threshold for particle emission and may be damped through emission of nucleons. The corresponding width $\Gamma^{\uparrow}$ may not be negligible ($\Gamma^{\uparrow}/\Gamma \approx 10^{-1}$).

These two mechanisms are expected to be rather temperature independent, being strongly connected to mean-field properties,[5)] see however Ref. 6.

On the contrary, the contribution of inhomogeneous damping, is expected to depend on temperature. This is because at finite temperature, the system will display thermal, large amplitude fluctuations which will make the system probe larger deformations with larger probability. These probabilities are controlled by the Boltzmann factors, and the resulting spread in frequencies will result as a weighting average over all possible quadrupole distortions. Because the corresponding integral can be parametrized in most cases in terms of a Gaussian integral, a square root dependence with temperature will follow.[1,7)]

Beyond mean field (around each deformation), vibrational motion is damped by coupling to progressively more complicated states lying at the same excitation energy. The first in this hierarchy of states ("doorway states"[8)]) are 2particle-2hole excitations. The most effective ones among them are those containing (in the zero temperature limit) an uncorrelated particle-hole excitation and a collective surface vibration or two vibrations.[9)] The coupling to these doorway states leads to a breaking of the giant resonance strength. This is the phenomenon of collisional damping, and is essentially connected with the coupling of the giant modes to the quantal fluctuations of the nuclear surfaces. The associated spreading width $\Gamma^{\downarrow}$, of the order of few MeV, provides (at $T = 0$ and in heavy nuclei) an overall account of the experimental data for the width.[1,9)] In fact, coupling to more complicated states, and eventually to the totally mixed states describing the compound nucleus, gives rise to a smoothly varying strength function.[10)] The associated width does not differ substantially from the spreading produced by the doorway coupling, because the increase of the density of complicated states is compensated by the decrease of the coupling matrix elements.

The collisional damping is found to be essentially independent of temperature.[1,11,12)] This in spite of the fact that single-particle damping displays a linear dependence with T. Quantum interference between the different amplitudes involving the decay of a particle or of a hole, essentially cancels out this dependence.

In the next sections, we discuss in more detail the properties of the collisional damping.

The hierarchy of couplings: Compound coupling

In this Sect., we will discuss the "compound nucleus damping" mechanism, that is the contribution to the damping of the giant resonances due to the coupling to progressively more complicated states up to the fully mixed compound nucleus states, following Ref. 10.

We write the total Hamiltonian describing the system as

$$H = H_0 + V, \tag{1}$$

sum of a mean field Hamiltonian H_0, and a residual two-body interaction V. The eigenstates of the mean field Hamiltonian H_0 are the independent many-particle many-hole excitations of the system, that is the one-particle one-hole ($1p$-$1h$), $2p$-$2h$, $3p$-$3h$ etc. excitations. We diagonalize the total Hamiltonian H in a restricted basis of $1p$-$1h$ states, obtaining the giant resonances (GR) plus low-lying collective states, aside from a large number of uncharacteristic states being close in energy to the unperturbed $1p$-$1h$ excitations.

The diagonalized $1p$-$1h$ states, together with the set of $2p$-$2h$, $3p$-$3h$ etc. states form a complete set of basis states denoted by $|\mu\rangle$, with energies $E_\mu = \langle\mu|H|\mu\rangle$. The exact stationary states of the total Hamiltonian obtained by diagonalizing H in the complete basis of states $|\mu\rangle$ are denoted $|i\rangle$ with the energies E_i. The damping of the giant resonance state $|0\rangle$ is controlled by its coupling to the states $|\mu\rangle$. Among these states, in the $2p$-$2h$ class are the doorway states consisting of a surface vibration plus an incoherent $1p$-$1h$ state.

To procede, we introduce a model based on two main assumptions. First, the spreading of any state $|\mu\rangle$ into the compound nucleus eigenstates $|i\rangle$ associated with the intrinsic chaotic dynamics is characterized by random, uncorrelated amplitudes X^i_μ,

$$|\mu\rangle = \sum_i X^i_\mu |i\rangle. \tag{2}$$

The second assumption is that of the uniformity of the spectrum and of the structure of the states $|\mu\rangle$. In other words, the states $|\mu\rangle$ are all characterized by a single strength function P centered around each unperturbed energy E_μ,

$$P(E_i) = \rho(E_i)|\langle\mu|i\rangle|^2 = \rho(E_i)|X^i_\mu|^2, \tag{3}$$

where ρ denotes the total density of states with the same quantum numbers as those of the mode under consideration.

This in keeping with the fact that, for a single-particle damping width with a linear or quadratic dependence on the energy, as for the experimental data in Fig. 7.20 of Ref. 13, the damping width of a many-particle-many-hole state has only a rather weak dependence on the number of particle-hole excitations, cf. Ref. 14 appendix A.

We cannot apply this argument in discussing the strength function $P_0(E - E_0)$ of the giant resonance state $|0\rangle$, because of the coherent properties of this mode, as discussed in the next Sect. As a result we allow the strength function P_0 of the giant resonance to be different from the strength function P.

The uniformity assumption also requires that the coupling between the states $|\mu\rangle$ have a uniform distribution. For the associated average coupling strength per unit energy, we write

$$\begin{aligned} \rho_\mu(E_\nu)\overline{|H_{\mu\nu}|^2} &= \rho_\mu(E_\nu)\overline{|\langle\mu|H|\nu\rangle|^2} \\ &= f(E_\nu - E_\mu). \end{aligned} \tag{4}$$

Here, ρ_μ denotes the level density of those states which couple, through the interaction V, to a given state $|\mu\rangle$. Similarly, for the GR we write,

$$\begin{aligned} \rho_0(E_\mu)\overline{|H_{0\mu}|^2} &= \rho_0(E_\mu)\overline{|\langle 0|H|\mu\rangle|^2} \\ &= f_0(E_\mu - E_0), \end{aligned} \tag{5}$$

with ρ_0 being the density of doorway states associated with the state $|0\rangle$.

The diagonalization of the full Hamiltonian H proceeds in two steps. First H is diagonalized among all the states $|\mu\rangle$ except for the giant resonance state $|0\rangle$, thus producing a new set of eigenstates $|\alpha\rangle$. These states $|\alpha\rangle$ have no mutual interactions but couple only to the state $|0\rangle$ with typical matrix elements

$$\begin{aligned} \overline{|H_{0\alpha}|^2} &= \overline{|\sum_\mu H_{0\mu}\langle\mu|\alpha\rangle|^2} \\ &\approx \sum_\mu \overline{|H_{0\mu}|^2}\ \overline{|\langle\mu|\alpha\rangle|^2} \\ &\approx \sum_\mu \overline{|H_{0\mu}|^2}\ \frac{1}{\rho(E_\alpha)} P_\mu(E_\alpha), \end{aligned} \tag{6}$$

where the strength function P_μ is defined in Eq. (3). In writing the above equation, the main assumption have been used: the amplitudes are uncorrelated.

The coupling matrix elements $H_{0\alpha}$ control the last stages of the relaxation process leading to extremely complicated (chaotic) compound nucleus states $|i\rangle$, each carrying tiny fractions of the giant resonance strength. If N is the number of principal components in the compound nucleus states, then $\rho \sim N$ and, from Eq. (6), $H_{0\alpha} \sim N^{-1/2}$.[15] Due to this N-scaling, we may expect[16] the value of the spreading width to saturate, as the dependence on exponentially growing level density is eliminated.

This is seen in Ref. 10, where the second part of the diagonalization was carried out exactly, in the weak and strong coupling limit, with the strength function written as,[17]

$$\begin{aligned} P_0(E - E_0) &= \frac{1}{\pi} Im \left[E - E_0 - \Sigma_0(E - E_0)\right]^{-1}, \\ \Sigma_0(E - E_0) &= \sum_\alpha \frac{|H_{0\alpha}|^2}{E - E_\alpha - i0}\ . \end{aligned} \tag{7}$$

The self-energy part $\Sigma_0(E)$ of the giant resonance is expressed in terms of mixing matrix elements given in Eq. (6). Since we average over nearby shell model states we may write,

$$\begin{aligned} \Sigma_0(E - E_0) \approx \sum_\alpha \sum_\mu \frac{1}{E - E_\alpha - i0} \overline{|H_{0\mu}|^2} \frac{1}{\rho(E_\alpha)} \\ \cdot P(E_\alpha - E_\mu) \end{aligned}$$

$$\approx \int dE_\alpha \int dE_\mu \left[\rho_0(E_\mu)\overline{|H_{0\mu}|^2}\right] \cdot \frac{P(E_\alpha - E_\mu)}{E - E_\alpha - i0}. \tag{8}$$

Equations (7) and (8) relate the strength function P_0 of the Giant resonance to the strength function P of the states $|\mu\rangle$.

Using the assumption of uniformity, the infinite set of equations for P_μ is substituted by the integral equation for the common strength function $P(E)$,

$$P(z) = \frac{1}{\pi} Im\,[z - \Sigma(z)]^{-1},$$
$$\Sigma(z) \approx \int dx \int dy\ f(x-y)\ \frac{P(y)}{z - x - i0}. \tag{9}$$

The set of coupled Eqs. (7)–(9) can be solved by iteration and define the giant resonance strength function in terms of average coupling matrix elements introduced in Eqs. (4) and (5). The relations (7)–(9) implicitly contain the coupling of the collective excitation to a hierarchy of ever more complicated many-particle many-hole states $|\mu\rangle$. In ordinary shell model calculations, the set of equations is truncated at a maximum number of particle-hole excitations. Here instead, we have imposed uniformity in the equations so that strength functions for high seniority states, i.e. large number of particle-hole excitations, are not neglected but rather retained at a limiting form.

The strength functions obtained by solving the set of coupled Eqs. (7)–(9) depend both on the strength as well as the energy range of the residual interaction. Analytic solutions are found by parametrizing the couplings given in Eqs. (4) and (5) as

$$f_0(E) = \frac{n_0 v_0^2}{2\pi} \frac{W_0}{E^2 + (W_0/2)^2}, \tag{10}$$

and

$$f(E) = \frac{n v^2}{2\pi} \frac{W}{E^2 + (W/2)^2}. \tag{11}$$

The quantity W_0 measures the range in energies E_μ over which the giant resonance state $|0\rangle$ couples to the doorway states. The strength of the residual interaction $n_0 v_0^2$ has been written in terms of a typical r.m.s. matrix element v_0 and the number of doorway states n_0. If the collective state couples to an unbounded set of doorway states we may let both β_0 and W_0 go to infinity while keeping their ratio β_0/W_0 constant, the ratio being proportional to the density ρ_0 of doorway states. With this parametrization, the second moment of the strength function P_0 of the giant resonance $|0\rangle$ is given by

$$M_0^{(2)} = \langle 0|(H - \langle H\rangle)^2|0\rangle = \sum_\mu |\langle 0|H|\mu\rangle|^2$$
$$\approx \int dE_\mu\ f_0(E_\mu - E_0) = \sum_{\mu=1}^{n_0} v_0^2 = n_0 v_0^2, \tag{12}$$

and a similar expression holds for the second moment $M_\mu^{(2)}(= \sum_{\mu=1}^{n} v^2)$, of any state $|\mu\rangle$.

With the above parametrization, the self-energies acquire the form

$$\Sigma_0(z) = \int dx\ \frac{n_0 v_0^2}{z - x - iW_0/2}\ P(x),$$
$$\Sigma(z) = \int dx\ \frac{n v^2}{z - x - iW/2}\ P(x). \tag{13}$$

In the weak coupling limit, $\sqrt{n_0} v_0 \ll W_0$ and $\sqrt{n} v \ll W$, the self energy $\Sigma_0(z)$ is a slowly changing function of energy within the spreading width. Consequently, the collective strength function will not depend on the details of $P(x)$ and eventually attain a Breit-Wigner form (exponential decay),

$$P_0(E - E_0) = \frac{1}{2\pi} \frac{\Gamma_0}{(E - E_0)^2 + (\Gamma_0/2)^2}, \tag{14}$$

where the spreading width of the giant resonance is given by

$$\Gamma_0 = 2\ Im\Sigma_0(0) = \frac{4 n_0 v_0^2}{W_0}; \qquad \sqrt{n_0} v_0 \ll W_0. \tag{15}$$

This result also could have been obtained directly from Eqs. (5) and (10) using Fermi's Golden Rule, that is,

$$\Gamma^\downarrow \equiv \Gamma_0 = 2\pi v_0^2 \rho_0(E_\mu \approx E_0) = 2\pi f_0(0). \tag{16}$$

In keeping with the assumptions made above, this relation applies in those cases in which the damping width Γ_0 is much smaller than the range of the interaction W_0.

In the strong coupling limit for the intrinsic states, that is when $\sqrt{n} v \gg W$, the exact shape of the coupling parametrized in Eq. (10) is not important and the Hamiltonian can be modelled by a banded matrix with random matrix elements of magnitude v in a band of size n surrounding the diagonal. The resulting self-consistent strength function of generic background states is a semicircle of radius R, that is

$$P(z) = \frac{2}{\pi R^2} \sqrt{R^2 - z^2}, \quad R = 2\sqrt{n} v; \quad \sqrt{n} v \gg W. \tag{17}$$

The giant resonance strength function, solution of Eqs. (13) and (17) becomes

$$P_0(z) = \frac{2}{\pi} \frac{\sqrt{R^2 - z^2}}{R_0^2 + 4z^2(R^2/R_0^2 - 1)}, \tag{18}$$

where $R_0 = 2\sqrt{n_0} v_0$, that is, a hybride of the semicircle and Breit-Wigner shapes. The effective spreading width (FWHM) of the giant resonance varies from $\Gamma_0 = \sqrt{3} R_0$ at $R_0 = R$ to $\Gamma_0 = R_0^2/R$ at $R \gg R_0$. Again, the upper boundary for the width is determined solely by the coupling to the doorway states. Subsequent scattering into the the compound nucleus eigenstates may change the shape of the strength function but does not add to the damping width which may actually decrease. Note that the result given by Eq. (18) (strong coupling limit) could not have been obtained in perturbation theory.

The estimates of P_0 collected in Eqs. (14) and (18), give an upper limit for the damping width of the giant resonances. These results are expected to remain qualitatively valid also in the case where the compound nucleus has a very high excitation energy and the giant resonance is thermally excited, provided a thermal averaging of Eqs. (7)–(9) is carried out. This is because, the quantity $R_0 = 2\sqrt{n_0}v_0$, which is proportional to the square root of the expectation value of H^2 in the giant resonance state, cf. Eq. (12), does not contain any exponentially growing parameters. On the other hand, both n_0 and v_0 may change with temperature.[11)]

A simple example of this reasoning is offered by the case of the giant quadrupole resonance (GQR) of ^{208}Pb. One obtains from Eq. (15) ($n_0 \approx 60, v^2 \approx 0.02$ MeV^2, $W_0 \approx 2$ MeV from Fig. 6 of Ref. 18) $\Gamma_0 \approx 2.4$ MeV. This number is essentially equal to the value of $\Gamma^{\downarrow}_{GQR}$, obtained in Ref. 18 by the coupling to the doorway states. This is because to couple to the compound nucleus states, the giant resonance has first to couple to the doorway states. Then, the coupling to the compound nucleus provides a smooth averaging of the doorway states ($\Gamma^{\downarrow}_{CN} \approx 2.4$ MeV/60 ≈ 40 keV) and a physical interpretation of a fraction of the averaging parameter introduced in the coupling to the doorway states by many authors.[9)] In other words, both the centroid of the giant resonances as well as the width are determined by the detailed dynamics of the problem, namely shell structure and doorway coupling . On the other hand, the fluctuations in the energies and in intensities are expected to display Wigner and Porter-Thomas behaviour, respectively, as for a collective mode mixed with a background of chaotic stateswith statistical properties close to the random matrix theory limit.[15–17)]

This was the basis for the reanalysis of the electron scattering experiments carried out in Ref. 19, which removed the discrepancy between the electron and the hadron data about the strength of the GQR in ^{208}Pb. In fact, assuming a Porter-Thomas distribution for the strength of the different lines, part of what had been considered before background became giant quadrupole resonance. The resulting cross sections agreed within the experimental errors with the one obtained from hadron scattering.

In the next two sections, we will discuss the features of the coupling of the single-particle motion and of the giant vibrational motion to the corresponding doorway states.

The hierarchy of coupling: Doorway coupling

Assuming that single-particle and collective modes can be treated as independent nuclear degrees of freedom, the coupling to doorway states was studied[11)] in an unified approach within the framework of the thermal Matsubara Green's function theory.[20)]

As already mentioned, the doorway states most effective in the damping of the giant resonance motion contain an uncorrelated pair of fermions coupled to a collective (surface) vibration or two vibrations,[11)] see Ref. 21 for a different approach, but with many aspects in common.

In Fig. 1 of Ref. 22, three lowest-order processes contributing to the coupling of a giant resonance D (GDR) to a surface vibration λ are drawn.

One can expect, that the so-called "single-particle self-energy correction" (Fig. 1(a) and (b))[22)] reflects the same properties of the single-particle self-energy and therefore leads to a linear increase of the GR spreading width with T as described below. On the other side, the fact that the third graph, the so-called "vertex correction" (Fig. 1(c)),[22)] partially cancels the contribution of the first two, in the case of coupling to isoscalar, spin-independent collective modes, is a well known effect at zero temperature.[9,23)] It reflects basic properties as conservation and/or restauration of fundamental symmetries (number and current conservation, isospin...). Numerical calculations of giant resonance (in particular GDR) strength functions in several nuclei confirm this feature also at finite temperature.[11,12)] The reasoning is described below.

All the three graphs have a very similar analytical structure. One can cast them (in the model of Ref. 23) into the following single expression for the correction to the RPA propagator of a giant resonance D

$$Im\Pi_D(\omega,T) = \sum_{012} \phi_{02}(\omega)\left(\frac{F^2_{01,D}F^2_{12,\lambda}}{(\Delta\omega_{01})^2} + \frac{F^2_{12,D}F^2_{01,\lambda}}{(\Delta\omega_{12})^2}\right) - 2\phi_{02}(\omega)\frac{F_{01,D}F_{12,\lambda}F_{12,D}F_{01,\lambda}}{\Delta\omega_{01}\Delta\omega_{12}} \quad (19)$$

where

$$\phi_{02}(\omega) = \beta^2_D(n_F(\varepsilon_0) - n_F(\omega+\varepsilon_0)) \cdot (1 + n_B(\omega-\varepsilon_{20}) - n_F(\varepsilon_2))\ Im\Pi^{RPA}_\lambda(\omega-\varepsilon_{20}). \quad (20)$$

The quantity

$$\Pi^{RPA}_\lambda(\omega) = \sum_\alpha \beta^2_\alpha\left(\frac{1}{\omega-\omega_\alpha+i\eta} - \frac{1}{\omega+\omega_\alpha+i\eta}\right), \quad (21)$$

is the RPA propagator[24)] of the collective vibration in the doorway state $(2,1,\lambda)$. The quantities β_D and β_λ are related to the excitation probability of the corresponding vibrations with single-particle matrix elements F. The boson and fermion occupation factors n also appear.

The first two terms in Eq. (19) correspond to the single-particle self-energy corrections. The third, corresponding to the vertex renormalization, carries a -2 factor and therefore reduces the self-energy correction effect. To emphasize this self-consistency requirement, we discuss in some detail the properties of the single-particle self-energy $\Sigma^{ret}_i(\varepsilon,T)$.

Numerical results in finite nuclei recently suggested that $Im\Sigma^{ret}_i(\varepsilon,T)$ and the corresponding spreading width behaves as a linear function of energy and of the temperature.[11,12,25)] As an example, we report in detail the results for the nucleus ^{120}Sn from Ref. 12. They are well fitted by the relation

$$\Gamma^{\downarrow}_p(\varepsilon,T) = a + bT \quad (22)$$

with $a \approx 1.3$ MeV and $b = 0.9$ for $\varepsilon \approx \varepsilon_F$ and $a \approx 4.1$ MeV and $b = 0.2$ for $|\varepsilon - \varepsilon_F| \approx 8$ MeV (the energy range relevant to the giant resonance physics). The values of a are in very good agreement with the empirical ones.[13)]

The linear behaviour can be explained making use of simple arguments,[11,26] emphasizing the rôle of the surface in the dressing process, in contrast to the case of coupling of the state i to uncorrelated particle-hole states which leads, on the contrary, to the well known quadratic dependence on temperature and excitation energy.[27] The difference between these two cases is to be ascribed only to the different behaviour of the response function in the case of uncorrelated particle-hole excitations and of surface vibrations. In fact, the imaginary part of the single-particle self-energy has the same structure in both cases,

$$Im\Sigma_i^{ret}(\varepsilon,T) = -\sum_f F_{if}^2(1+n_B(\varepsilon-\varepsilon_f)-n_F(\varepsilon_f)) \cdot ImR(\varepsilon-\varepsilon_f)\,. \quad (23)$$

where $ImR(\varepsilon)$ is the imaginary part of the response function[24] of the state (particle-hole or surface vibration) the particle couples to.

In the case of coupling to uncorrelated particle-holes, $ImR^0(\varepsilon)$ has a typical linear behaviour (see, e.g., Ref. 28, p. 107ff, for the case of an infinite Fermi liquid). Inserting $ImR^0(\varepsilon) \sim \varepsilon$ in Eq. (23), substituting the sum over f with an integral and assuming constant matrix elements F_{ij}, one finds[26]

$$Im\Sigma_i^{ret}(\varepsilon,T) \sim (\pi T)^2 + \varepsilon^2, \quad (24)$$

which is the result obtained by Morel and Nozières as an estimation of the quasi-particle damping width in ^{3}He.[27]

The case of coupling to surface vibrations has been analytically studied in a model of semi-infinite nuclear matter (the slab model, presented in Ref. 29). It has been found that ImR^{RPA}, Eq. (21), can be considered as essentially constant as function of energy and temperature and this leads to the following result,[26]

$$Im\Sigma_i^{ret}(\varepsilon,T) \sim \sqrt{(2ln2\,T)^2+\varepsilon^2}, \quad (25)$$

displaying in particular a linear dependence on the temperature, at the Fermi energy.

In the case of surface dominated systems, as atomic nuclei are, the surface response is much larger than the uncorrelated one,[1,9,30] so that the dressing of single-particle states (in the energy range of interest here) has to be ascribed to their coupling to surface vibrations.

The real part of the single-particle self-energy $Re\Sigma_i^{ret}(\varepsilon,T)$ is obtained from $Im\Sigma_i^{ret}(\varepsilon,T)$ via dispersion relation.[31] Connected to it is the effective nucleon mass,[31] function in particular of the temperature,[26,32,33] and which determines the effective density of states in the spreading calculations.[34]

Coming back to the giant resonances, the reduction of the single- particle self-energy effects in Eq. (19) is found to cancel the weak temperature dependence of diagrams 1(a) and (b), and the spreading width of the giant resonances results to be quite constant as a function of temperature.[11,12,22]

This result, confirmed by the fully microscopic calculations in light nuclei of Ref. 35, is consistent with the experimental results reported above for the width of the GDR in hot nuclei. Indeed, the measured temperature and angular momentum dependence of the FWHM of the "experimental" strength functions (see the many contributions to these proceedings[3]) are reproduced by a refined model[7] which takes into account the large-amplitude thermal fluctuations, inhomogeneus damping and angular momentum, and the life time of the compound nucleus.[36] Recent examples are shown in Refs. 37 and 38, see also Ref. 39 for the highly related discussion of the restoration of the isospin symmetry in hot nuclei.

At this stage, a crucial question to be asked is, whether the eventual γ-decay of the GDR can be observed at a given T. This is because the time it takes for a giant resonance to exist as a thermal excitation in a compound nucleus after its formation, may be longer than the time it takes for the compound nucleus to cool through particle emission. Consequently, the γ-decay of the giant dipole resonance, may not be observable, within the limits of sensitivity of the available γ-arrays, at all temperatures. The highest temperature for the observation of the γ-decay of the giant dipole resonance is expected to correspond to that in which the damping width of the giant vibration and the width for particle emission from compound states have similar values.[40]

For the future and on the theory side, we have to work on a fully(?) microscopic theory that includes both thermal and quantal fluctuations.

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