Evidence for Different Time Scales Controlling Thermal Fluctuations in Hot Nuclei

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A comparison between the predictions of a theoretical model describing the giant dipole resonance in hot nuclei, which includes the coupling to time-dependent thermal fluctuations of the nuclear surface, and experimental data on the cross section and the angular distribution associated with the dipole decay of ^{92}Mo* is presented. Effects of different time scales for fluctuations in the deformation and orientation degrees of freedom are observed.

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The study of nuclear structure at finite temperature is one of the central issues lying at the forefront of nuclear research. In it, the interplay between single-particle and collective degrees of freedom typical of nuclear structure studies [1] acquires a new dimension because of the ubiquitous role played in this coupling by the compound nucleus.

The modern view of nuclear research at finite temperature is based on the Axel-Brink hypothesis, which states that one can build a giant dipole resonance (GDR) upon any excited state. Recent work [2–5] has shown the need to generalize this hypothesis by taking into account the fact that the excited states do not merely act as an undifferentiated base upon which one builds a collective excitation, but can also couple to it and modify its properties in an important way. The study of this coupling is proving to be important in the understanding of the stability of collective motion in nuclei (cf., e.g., Ref. [6]) and the role played by the time dependence of fluctuations of the nuclear surface in the damping of the GDR, in particular, the relevant time scales related to thermal fluctuations of the deformation and orientation degrees of freedom [4,5].

In the present paper, we analyze the strength function and angular distribution associated with the recently measured decay of ^{92}Mo* [7]. We primarily investigate the influence of time-dependent thermal fluctuations of the nuclear surface on the properties of GDR, focusing the attention on the possibility that fluctuations in the deformation and orientation degrees of freedom may exhibit different time scales.

The coupling between the dipole and the quadrupole shape leads to three fundamental vibrations with frequency

$$\omega_k = \omega_D \exp[-\sqrt{5/4\pi}\beta\cos(\gamma+2\pi k/3)] , \qquad (1)$$

where $k=1,2,3$ denote the principal axes in the intrinsic frame. The quadrupole deformation is characterized by the parameters β and γ which measure the ratio between axes of the ellipsoid and its departure from axial symmetry [1].

In the calculations presented below, the GDR is simulated by a rotating, three-dimensional harmonic oscillator, whose equation of motion in the rotating frame is [5]

$$\ddot{d}_k + \Gamma_k^i \dot{d}_k + \omega_k^2 d_k = -[\boldsymbol{\omega}_R \times (\boldsymbol{\omega}_R \times \mathbf{d})]_k - 2(\boldsymbol{\omega}_R \times \mathbf{d})_k , \qquad (2)$$

where Γ_k^i is the intrinsic dipole width, ω_k is given by Eq. (1), and $\boldsymbol{\omega}_R$ is the angular velocity of the rotating nucleus. The rotational frequency was assumed to be along the fixed laboratory z axis, and was determined from the average angular momentum for the compound nuclear reaction including thermal fluctuations, i.e.,

$$\langle J\rangle = Z^{-1}\int d\tau\, \mathcal{J}(\beta,\gamma,\theta,\psi)\omega_R \exp\{-[F_0(\beta,\gamma,T) - \tfrac{1}{2}\mathcal{J}(\beta,\gamma,\theta,\psi)\omega_R^2]/T\} , \qquad (3)$$

where $d\tau = \beta^4 d\beta \sin 3\gamma\, d\gamma\, d\Omega$, $Z=\int d\tau\, e^{-F/T}$, F_0 is the nuclear free energy at $\omega_R=0$, and $\mathcal{J}(\beta,\gamma,\theta,\psi)$ is given by

$$\mathcal{J}(\beta,\gamma,\theta,\psi) = \mathcal{J}_1\cos^2\psi\sin^2\theta + \mathcal{J}_2\sin^2\psi\sin^2\theta + \mathcal{J}_3\cos^2\theta ,$$

where ψ and θ are two of the Euler angles defining the orientation of the system, and the $\mathcal{J}_k$ are the nuclear moments of inertia and are functions of β and γ. Here, rigid-body values were used, assuming the radius to be $R=1.2A^{1/3}$ fm. Also, the free energies were evaluated using the Nilsson-Strutinsky method and the Nilsson and liquid-drop parameters of Refs. [8] and [9], respectively. Including the effects of thermal fluctuations, we find for ^{92}Mo at $T=2.0$ MeV and $\langle J\rangle = 33\hbar$, $\hbar\omega_R = 1.11$ MeV.

The intrinsic dipole width Γ_k^i appearing in Eq. (2) is due to the coupling between the GDR and low-lying two-particle–two-hole states. Making use of the parametrization of Ref. [10], the intrinsic width for each component k is taken to be $\Gamma_k^i = \Gamma[\omega_k/\omega_D]^\delta$, where Γ, ω_D, and δ are parameters generally determined from the ground-state data of the GDR.

The GDR photoabsorption cross sections can be obtained by applying a harmonic, external force $\mathbf{f}(t)$ to Eq. (2) [5]. If the magnitude of the applied force is unity, then

$$\sigma_\mu(E) = \frac{4\pi^2\hbar}{M}\frac{ZN}{A}\frac{2}{3}\langle P_\mu(E)\rangle , \tag{4}$$

where $\mathbf{f}(t)$ is composed of spherical tensor (μ) components in the laboratory frame with frequency $\omega = E/\hbar$, and $\langle P_\mu(E)\rangle = 1/T\int_0^T \mathbf{f}_\mu(t)\cdot\dot{\mathbf{d}}(t)dt$ is the average power delivered by the external force. The total cross section is given by the sum $\sigma_T = \sum\sigma_\mu$, while the angular-distribution quadrupole a_2 coefficient, defined in the relation $\sigma(E,\theta) = \sigma_T(E)[1 + a_2(E)P_2(\cos\theta)]$, is given by

$$a_2 = \frac{1}{\sigma_T}[0.5\sigma_0 - 0.25(\sigma_1 + \sigma_{-1})] , \tag{5}$$

where for the cases studied here θ is the angle between the emitted photon and the incident beam direction.

As the excitation energy of the nucleus increases, large-amplitude, thermal fluctuations of the nuclear shape become important [2–4,11]. If the time scale of these fluctuations is long compared to the shift in the GDR frequency induced by the shape changes, the adiabatic approximation is valid, and the GDR cross section is given by the superposition

$$\sigma_\mu(E) = Z^{-1}\int d\tau\, \sigma_\mu(E,\beta,\gamma,\Omega)e^{-F/T} . \tag{6}$$

However, as the time scale for the fluctuations decreases, Eq. (6) is no longer valid, and we have accounted for time-dependent fluctuations using a model based on the Kubo-Anderson process [12]. In this model, jumps are made between the various deformations and orientations independent of the initial starting point, and the conditional probability $P(\boldsymbol{\alpha},t,\boldsymbol{\alpha}_0,t_0)$ of having deformation $\boldsymbol{\alpha} \equiv (\beta,\gamma,\Omega)$ at time t, after having been at the point $\boldsymbol{\alpha}_0$ at time t_0, is

$$P(\boldsymbol{\alpha},t,\boldsymbol{\alpha}_0,t_0) = e^{-\lambda(t-t_0)}\delta(\boldsymbol{\alpha}-\boldsymbol{\alpha}_0) + (1-e^{-\lambda(t-t_0)})p(\boldsymbol{\alpha}) , \tag{7}$$

where λ is the mean jumping rate and $p(\boldsymbol{\alpha})$ is the stationary probability distribution $Z^{-1}e^{-F/T}$. A general feature of time-dependent fluctuations on the GDR is that as λ increases, the effects of thermal fluctuations on the line shape are mitigated: a phenomenon known as motional narrowing.

We evaluate σ_μ via Eq. (4) by selecting an ensemble of 500 (1000 for the adiabatic calculations) initial points distributed according to the stationary probability distribution, and then integrating the equations of motion for the time period $t_{\max} \approx 100\lambda$ (≈ 100 jumps) using a fifth-order Runge-Kutta algorithm with time step $\Delta t \approx 2\pi/40\omega_D$. The time step was chosen so as to conserve the energy in the undamped and adiabatic limit to within 1 part in 10^6 for approximately 200 oscillator periods. At each integration time step the probability of making a jump to a new deformation and orientation is $\lambda\Delta t$. Finally, σ_μ is then taken as the ensemble average. We note that in the adiabatic limit ($\lambda = 0$), this procedure is identical to integrating Eq. (6) using Monte Carlo techniques. The theoretical uncertainties were obtained from the ensemble variance and were largest in the adiabatic limit, but were typically less than 0.5 mb.

The quantity $1/\lambda$ may be viewed as the relaxation time for the quadrupole degrees of freedom. We note that it is possible that the relaxation times for deformation and orientation coordinates may be different, as they are related to the coupling of vibrations and rotations in a hot nucleus with states of the compound nucleus. Here, we account for different time scales by treating deformation and orientation jumps separately. Assuming two relaxation times $\lambda_{\beta\gamma}$ and λ_Ω, with $\lambda_{\beta\gamma} \geq \lambda_\Omega$, the probability of making a jump in the (β,γ) coordinates in the time interval Δt is $\lambda_{\beta\gamma}\Delta t$. If this condition was satisfied at any integration step, the conditional probability of $\lambda_\Omega/\lambda_{\beta\gamma}$ was used to determine if a jump also occurred in the orientation coordinates. In either case, the new deformations and/or orientations were chosen according to the stationary probability distribution.

Very little information regarding $\lambda_{\beta\gamma}$ and λ_Ω is known, and, therefore, we have evaluated σ_T and a_2 with different combinations of $\lambda_{\beta\gamma}$ and λ_Ω in order to best reproduce the experimental data by minimizing the error-weighted χ^2 between the theoretical and experimental points for both the cross section and a_2 coefficient for energies greater than 10 meV. We define χ^2/ν as

$$\chi^2/\nu = (\chi^2_{\sigma_T} + \chi^2_{a_2})/(N - n_p) ,$$

where N is the number data points, n_p is the number of fitted parameters, and $\chi^2_{\sigma_T}$ and $\chi^2_{a_2}$ are the error-weighted chi square between experiment and theory for the cross section and a_2 coefficient, respectively. The cutoff in the experimental data was introduced because below this energy the gamma-ray spectra are dominated by statistical, nondipole $E1$ transitions, and the quoted experimental uncertainties do not reflect any uncertainties in the unfolding of the statistical $E1$ spectrum.

The two remaining parameters that enter our calculations are the total strength of the resonance and the intrinsic width Γ. For the total strength of the resonance, we multiply Eq. (4) by the factor S and minimize the χ^2/ν with respect to S. Generally, we find $S \approx 0.9$–1.1. As for the intrinsic width, the zero-temperature value of Γ is ~ 5 MeV [13], but the cross section is not described by a single Lorentzian with this width very well. Also, we note that although theoretical studies indicate that the in-

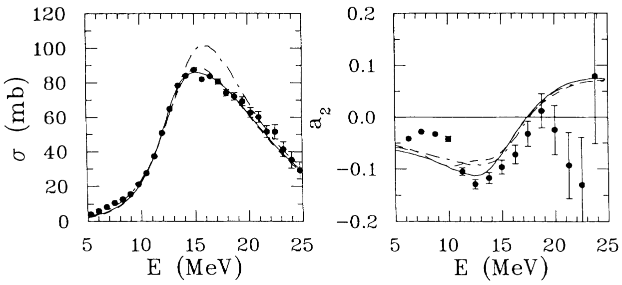

FIG. 1. Comparison between experimental data [6] and the present model for ^{92}Mo at $T=2.0$ MeV, $\langle J\rangle=33\hbar$. The dashed line represents the adiabatic calculation ($\Gamma=5.6$ MeV, $S=0.904$), while the solid and dash-dotted lines represent the "restricted" ($\Gamma=6.75$ MeV, $\lambda_{\beta\gamma}=7.5$ MeV, $\lambda_\Omega/\lambda_{\beta\gamma}=0.1$, $S=0.97$) and "full" ($\Gamma=6.75$ MeV, $\lambda_{\beta\gamma}=\lambda_\Omega=3$ MeV, $S=1.05$) motionally narrowed regimes, respectively.

trinsic width is expected to be essentially temperature independent [14,15], changes of the order 20%–30% are not excluded. Further, as in Ref. [5], we find that the experimental data are better described by a larger width. For this reason, we have also included Γ as a free parameter in our fits to experimental data. Last, for ^{92}Mo, we take $\hbar\omega_D=16.82$ MeV and $\delta=1.8$.

In the adiabatic limit ($\lambda_{\beta\gamma}=\lambda_\Omega=0$), the minimum χ^2/ν occurs for $\Gamma=5.6$ and $s=0.904$ with a value of 4.76 ($n_p=2$). A comparison with experimental data is exhibited by the dashed lines in Fig. 1, where it is seen that the a_2 underestimates the data by approximately 30% at $E\approx 12$ MeV. On the other hand, when time-dependent fluctuations are included, the minimum χ^2/ν reduces to ≈ 1.7 ($n_p=4$) at $\Gamma=6.75$, $\lambda_{\beta\gamma}=7.5$, $\lambda_\Omega=0.75$, and $S=0.97$. This "best" fit and "restricted" motionally narrowed situation is represented by the solid lines in Fig. 1, where it is seen that the a_2 is significantly enhanced over the adiabatic limit. Also, for illustrative purposes, the "full" motionally narrowed limit with $\Gamma=6.75$ MeV, $\lambda_{\beta\gamma}=\lambda_\Omega=3$, and $S=1.05$ is displayed in Fig. 1 by the dash-dotted line.

To illustrate the sensitivity of χ^2/ν, contour plots of χ^2/ν as a function of $\lambda_{\beta\gamma}$ and λ_Ω are shown in Fig. 2 for $\Gamma=6.25$ MeV (top), 6.75 MeV (middle), and 7.25 MeV (bottom). In the $\Gamma=6.75$ MeV case, we see that χ^2/ν is quite flat and that the data are essentially equally well described with $5.5\le\lambda_{\beta\gamma}\le 10$ and $\lambda_\Omega/\lambda_{\beta\gamma}\approx 0.8$. Finally, we note that even in the restricted motionally narrowed, best fit to the data, the a_2 still underestimates the data by approximately 15%.

We conclude that ^{92}Mo experimental data for the GDR at high excitation energy reflect the time dependence of thermal fluctuations of the nuclear surface. Furthermore, the data provide evidence for the fact that the relaxation times associated with the shape and orientation degrees of freedom may differ considerably. In fact, the "hopping time" associated with the surface degrees of freedom is found to be much shorter than that associated with the orientation degrees of freedom.

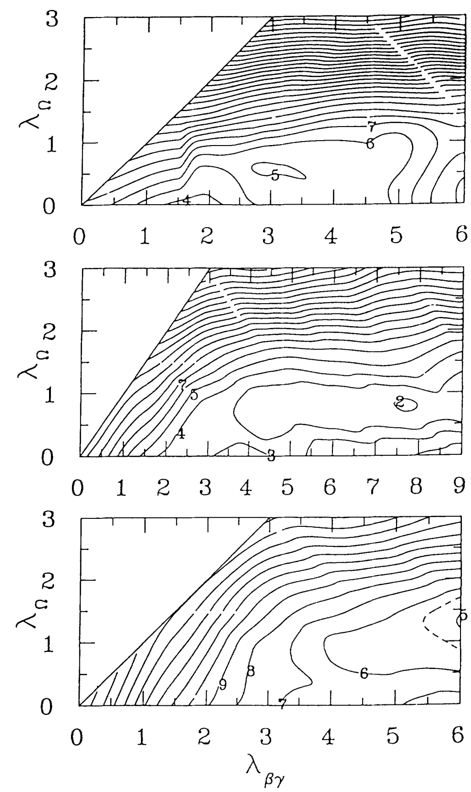

FIG. 2. Contour plot of the χ^2/ν as a function of $\lambda_{\beta\gamma}$ and λ_Ω with $\Gamma=6.25$, 6.75, and 7.25 MeV.

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