

Double giant resonances in deformed nuclei

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(Received 11 December 1997)

We report on the first microscopic study of the properties of two-phonon giant resonances in deformed nuclei. The cross sections of the excitation of the giant dipole and the double giant dipole resonances in relativistic heavy ion collisions are calculated. We predict that the double giant dipole resonance has a one-bump structure with a centroid 0.8 MeV higher than twice the energy for the single giant dipole resonance in the reaction under consideration. The width of the double resonance equals 1.33 of that for the single resonance. [S0556-2813(98)07110-6]

PACS number(s): 24.30.Cz, 23.20.-g, 25.70.De, 25.75.Dw

One of the most exciting progresses in the field of giant resonances in atomic nuclei over the last few years was the experimental observation of two-phonon giant resonances [1]. Nowadays, we may speak about some systematics of their properties (the energy position, width, and excitation probability) in spherical nuclei although it is still sparse and some open questions in this field stimulate theoretical studies (see, e.g., Refs. [2–15]). An investigation on the properties of two-phonon giant resonances together with similar studies on low-lying two-phonon states [16] should give an answer about how far the harmonic picture of boson-type excitations holds in finite fermion systems such as atomic nuclei.

The possibility to observe two-phonon giant resonances in deformed nuclei with the present state of art experimental techniques is still questionable. This is mainly because of the fact that one has to expect a larger width of these resonances as compared to spherical nuclei. Also, the situation with the low-lying two-phonon states in deformed nuclei is much less clear than in spherical ones.

The first experiment with the aim to observe the double giant dipole resonance (DGDR) in ²³⁸U in relativistic heavy ion collisions (RHIC) was performed recently at the GSI/SIS facility by the LAND Collaboration [17]. It will take some time to analyze the experimental data and to present the first experimental evidence of the DGDR in deformed nuclei, if any. The first microscopic study of properties of the DGDR in deformed nuclei is the subject of the present paper. The main attention will be paid to the width of the DGDR and its shape.

In a phenomenological approach the GDR is considered as a collective vibration of protons against neutrons. In spherical nuclei this state is degenerate in energy for different values of the spin $J=1^-$ projection $M=0, \pm 1$. The same is true for the 2^+ component of the DGDR with projection $M=0, \pm 1, \pm 2$. In deformed nuclei with an axial symmetry like ²³⁸U, the GDR is split into two components $I^\pi(K) = 1^-(0)$ and $I^\pi(K) = 1^-(\pm 1)$ corresponding to vibrations against two different axes. In this approach one expects a

three-bump structure for the DGDR with the value $K=0$, $K=\pm 1$ and $K=0, \pm 2$, respectively. Actually, the GDR possesses a width and the main mechanism responsible for it in deformed nuclei is the Landau damping. Thus, the conclusion on how three bumps overlap and what is the real shape of the DGDR in these nuclei, i.e., either a three-bump or a flat broad structure, can be drawn out only from some consistent microscopic studies.

In the present paper we use the quasiparticle phonon model (QPM) [18] to investigate the properties of the GDR and the DGDR in ²³⁸U. The QPM, although with somewhat different technical details which reflect the difference between spherical and deformed nuclei, was used to investigate the same resonances in spherical nuclei in Refs. [6,8,10,15]. The model Hamiltonian includes an average field for protons and neutrons, monopole pairing and residual interaction in a separable form. We use in our calculations for ²³⁸U the parameters of Woods-Saxon potential for the average field and monopole pairing from the previous studies [19]. They were adjusted to reproduce the properties of the ground state and low-lying excited states. The average field has a static deformation with the deformation parameters $\beta_2=0.22$ and $\beta_4=0.08$. To construct the phonon basis for the $K=0$ and $K=\pm 1$ components of the GDR we use the dipole-dipole residual interaction (for more details, see, e.g., Ref. [18]). The strength parameters of this interaction are taken from Ref. [20] where they have been fitted to obtain the centroid of the $B[E1, 0_{g.s.}^+ \rightarrow 1^-(K=0, \pm 1)]$ strength distribution at the value known from experiment [21] and to exclude the center of mass motion. In this approach, the information on the phonon basis (i.e., the excitation energies of phonons and their internal fermion structure) is obtained by solving the RPA equations. For electromagnetic $E1$ -transitions we use the free values of the effective charges, $e_{\text{eff}}^{Z(N)} = eN(-Z)/A$.

The results of our calculation of the $B(E1)$ strength distribution over $|1_{K=0}^-(i)\rangle$ and $|1_{K=\pm 1}^-(i')\rangle$ GDR states are presented in Fig. 1, together with the experimental data. The index i in the wave function stands for the different random-phase approximation (RPA) states. All one-phonon states with the energy lower than 20 MeV and with the $B(E1)$ value larger than $10^{-4} e^2 \text{fm}^2$ are accounted for. Their total number equals to 447 and 835 for the $K=0$ and $K=\pm 1$

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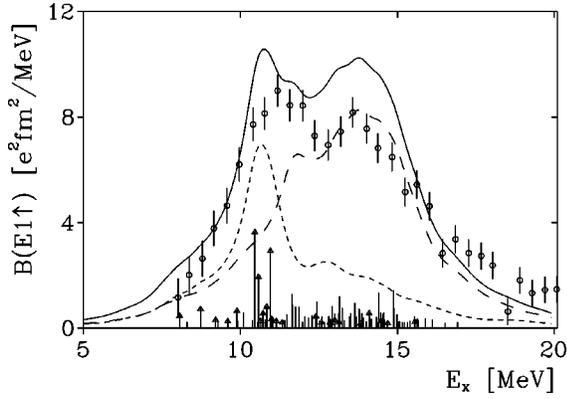


FIG. 1. The $B(E1)$ strength distribution over $K=0$ (short-dashed curve) and $K=\pm 1$ (long-dashed curve) 1^- states in ^{238}U . The solid curve is their sum. The strongest one-phonon 1^- states are shown by vertical lines, the ones with $K=0$ are marked by a triangle on top. Experimental data are from Ref. [21].

components, respectively. Only the strongest of them with $B(E1) \geq 0.2 e^2 \text{fm}^2$ are shown in the figure by vertical lines. Our phonon basis exhausts 32.6% and 76.3% of the energy weighted sum rules, $14.8 \times NZ/A e^2 \text{fm}^2 \text{MeV}$, by the $K=0$ and $K=\pm 1$ components, respectively. For a better visual appearance we also present in the same figure the strength functions averaged with a smearing parameter, which we take as 1 MeV. The short (long) dashed-curve represent the $K=0$ ($K=\pm 1$) components of the GDR. The solid curve is their sum. The calculation reproduces well the two-bump structure of the GDR and the larger width of its $K=\pm 1$ component. The last is consistent with the experiment [21] which is best fitted by two Lorentzians with widths equal to $\Gamma_1 = 2.99 \text{ MeV}$ and $\Gamma_2 = 5.10 \text{ MeV}$, respectively. The amplitudes of both maximums in the calculation are somewhat overestimated as compared to the experimental data. This happens because the coupling of one-phonon states to complex configurations is not taken into account which can be more relevant for the $K=\pm 1$ peak at higher energies. But in general the coupling matrix elements are much weaker in deformed nuclei as compared to spherical ones and the Landau damping describes the GDR width on a reasonable level.

The wave function of the 0^+ and 2^+ states belonging to the DGDR are constructed by the folding of two 1^- phonons from the previous calculation. When a two-phonon state is constructed as the product of two identical phonons its wave function gets an additional factor $1/\sqrt{2}$. The 1^+ component of the DGDR is not considered here since its excitation is quenched in RHIC for the same reasons as in spherical nuclei [10]. The anharmonicity effects that arise from interactions between different two-phonon states are also not included in the present study. It was shown that these effects have an $A^{-4/3}$ dependence on the mass number A [14] and that they are small for the DGDR [8,12,14], even for ^{136}Xe and ^{208}Pb .

The folding procedure yields three groups of the DGDR states:

$$[[1_{K=0}^-(i_1) \otimes 1_{K=0}^-(i_2)]_{0^+, 2^+_{K=0}}], \quad (1a)$$

$$[[1_{K=0}^-(i) \otimes 1_{K=\pm 1}^-(i')]_{2^+_{K=\pm 1}}], \quad (1b)$$

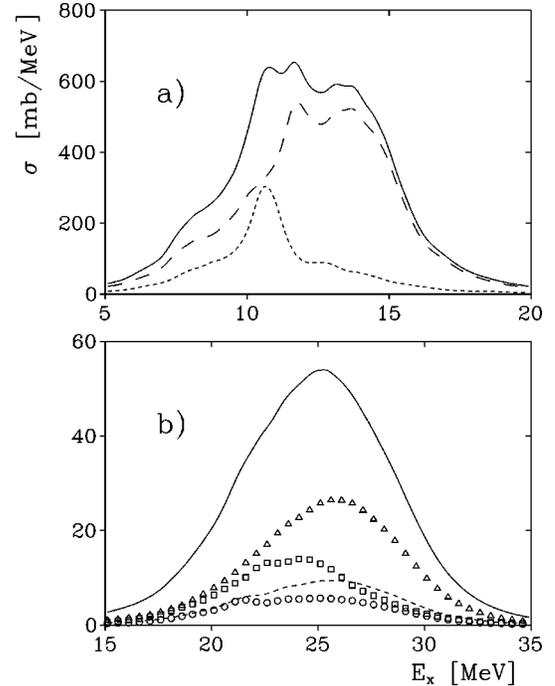


FIG. 2. The strength functions for the excitation: (a) of the GDR, and (b) of the DGDR in ^{238}U in the ^{238}U (0.5A GeV) + ^{208}Pb reaction. In (a), the short-dashed curve corresponds to the GDR ($K=0$) and the long-dashed curve to the GDR ($K=\pm 1$). In (b) the dashed curve corresponds to the DGDR $_{0^+}$ ($K=0$), the curve with circles to the DGDR $_{2^+}$ ($K=0$), the curve with squares to the DGDR $_{2^+}$ ($K=\pm 1$), and the curve with triangles to the DGDR $_{2^+}$ ($K=\pm 2$). The solid curve is the sum of all components. The strength functions are calculated with the smearing parameter equal to 1 MeV.

$$[[1_{K=\pm 1}^-(i'_1) \otimes 1_{K=\pm 1}^-(i'_2)]_{0^+, 2^+_{K=0, \pm 2}}]. \quad (1c)$$

The total number of nondegenerate two-phonon states equals to about 1.5×10^6 . The energy centroid of the first group is the lowest and of the last group is the highest among them. So, we also obtain the three-bump structure of the DGDR. But the total strength of each bump is fragmented over a wide energy region and they strongly overlap.

Making use of the nuclear structure elements discussed above, we have calculated the excitation of the DGDR in ^{238}U projectiles (0.5A GeV) incident on ^{120}Sn and ^{208}Pb targets, following the conditions of the experiment in Ref. [17]. These calculations have been performed in the second order perturbation theory [3], in which the DGDR states of Eqs. (1) are excited within a two-step process: g.s. \rightarrow GDR \rightarrow DGDR. As intermediate states, the full set of one-phonon $|1_{K=0}^-(i)\rangle$ and $|1_{K=\pm 1}^-(i')\rangle$ states was used. We have also calculated the GDR excitation to first order for the same systems. The minimal value of the impact parameter, which is very essential for the absolute values of excitation cross section has been taken according to $b_{\min} = 1.28 \times (A_t^{1/3} + A_p^{1/3})$.

The results of our calculations are summarized in Fig. 2 and Table I. In Fig. 2 we present the cross sections of the GDR (part a) and the DGDR (part b) excitation in the ^{238}U (0.5A GeV) + ^{208}Pb reaction. We plot only the

TABLE I. The properties of the different components of the GDR and the DGDR in ^{238}U . The energy centroid E_c , the second moment of the strength distribution m_2 in RHIC, and the cross sections σ for the excitation of the projectile are presented for: (a) ^{238}U (0.5A GeV) + ^{120}Sn , and (b) ^{238}U (0.5A GeV) + ^{208}Pb .

	E_c [MeV]	m_2 [MeV]	σ [mb]	
			(a)	(b)
GDR($K=0$)	11.0	2.1	431.2	1035.4
GDR($K=\pm 1$)	12.3	2.6	1560.2	3579.1
GDR(total)	12.0	2.6	1991.4	4614.5
DGDR $_0$ ($K=0$)	25.0	3.4	18.3	88.9
DGDR $_2$ ($K=0$)	24.4	3.5	11.8	58.7
DGDR $_2$ ($K=\pm 1$)	23.9	3.2	22.7	115.4
DGDR $_2$ ($K=\pm 2$)	25.3	3.4	49.7	238.3
DGDR(total)	24.8	3.4	102.5	501.3

smearred strength functions of the energy distributions because the number of two-phonon states involved is numerous. The results for ^{238}U (0.5A GeV) + ^{120}Sn reaction look very similar and differ only by the absolute value of cross sections. In Table I the properties of the GDR and the DGDR, and their different K components are given. The energy centroid E_c and the second moment, $m_2 = \sqrt{\sum_k \sigma_k \cdot (E_k - E_c)^2 / \sum_k \sigma_k}$, of the distributions are averaged values for the two reactions under consideration.

The two-bump structure can still be seen in the curve representing the cross section of the GDR excitation in ^{238}U in RHIC as a function of the excitation energy. But its shape differs appreciably from the $B(E1)$ strength distribution [see Fig. 2(a) in comparison with Fig. 1]. The reason for that is the role of the virtual photon spectra. First, for the given value of the excitation energy and impact parameter it is larger for the $K=\pm 1$ component than that for the $K=0$ one (see also the first two lines in Table I). Second, for both components it has a decreasing tendency with an increase of the excitation energy [3]. As a result, the energy centroid of the GDR excitation in RHIC shifts by the value 0.7 MeV to lower energies as compared to the same value for the $B(E1)$ strength distribution. The second moment m_2 increases by 0.2 MeV.

The curves representing the cross sections of the excitation of the $K=\pm 1$ and $K=\pm 2$ components of the DGDR in ^{238}U in RHIC have typically a one-bump structure [see the curves with squares and triangles in Fig. 2(b), respectively]. It is because they are made of two-phonon 2^+ states of one type: the states of Eq. (1b) and Eq. (1c), respectively. Their centroids should be separated by an energy approximately equal to the difference between the energy centroids of the $K=0$ and $K=\pm 1$ components of the GDR. They correspond to the second and the third bumps in a phenomenological treatment of the DGDR. The $K=0$ components of the DGDR include two groups of states: the states represented by Eq. (1a) and those of Eq. (1c). Its strength distribution has two bumps [see the curve with circles for the $2^+(K=0)$ and the dashed curve for the $0^+(K=0)$ components of the DGDR, respectively]. The excitation of the states given by Eq. (1a) in RHIC is enhanced because of their lower energies, while the enhancement of the excitation of the states

given by Eq. (1c) is related to the strongest response of the $K=\pm 1$ components to the external $E1$ Coulomb field in both stages of the two-step process.

Summing together all components of the DGDR yields a broad one-bump distribution for the cross section for the excitation of the DGDR in ^{238}U , as a function of excitation energy. It is presented by the solid curve in Fig. 2(b). Another interesting result of our calculations is related to the position of the DGDR energy centroid and to the second moment of the DGDR cross section. The centroid of the DGDR in RHIC is shifted to the higher energies by about 0.8 MeV from the expected value of two times the energy of the GDR centroid. The origin for this shift is in the energy dependence of the virtual photon spectrums and it has nothing to do with anharmonicities of the two-phonon DGDR states. In fact, the energy centroid of the $B(E1, \text{g.s.} \rightarrow 1_i^-) \times B(E1, 1_i^- \rightarrow \text{DGDR}_f)$ strength function appears exactly at twice the energy of the centroid of the $B(E1, \text{g.s.} \rightarrow \text{GDR})$ strength distribution because the coupling between different two-phonon DGDR states are not accounted for in the present calculation. The same shift of the DGDR from twice the energy position of the GDR in RHIC also takes place in spherical nuclei. But the value of the shift is smaller there because in spherical nuclei the GDR and the DGDR strength is less fragmented over their doorway states because of the Landau damping. For example, this shift equals 0.25 MeV in ^{208}Pb for the similar reaction. This effect is also seen when the DGDR position against the GDR is reported from experimental studies [22]. But the larger value of the shift under consideration in deformed nuclei should somehow simplify the separation of the DGDR from the total cross section in RHIC.

Another effect which also works in favor of the extraction of the DGDR from RHIC excitation studies with deformed nuclei is its smaller width than $\sqrt{2}$ times the width of the GDR, as observed with spherical nuclei. Our calculation yields the value 1.33 for the ratio $\Gamma_{\text{DGDR}}/\Gamma_{\text{GDR}}$ in this reaction. The origin for this effect is in the different contributions of the GDR $K=0$ and $K=\pm 1$ components to the total cross section, due to the reaction mechanism. It should be remembered that only the Landau damping is accounted for the width of both the GDR and the DGDR. But since the effect of narrowing of the DGDR width is due to the selectivity of the reaction mechanism it will still hold if the coupling to complex configurations is included in the calculation.

It may be argued that the procedure of independent excitations of two RPA phonons applied in this paper is not sufficient for a consistent description of the properties of the two-phonon giant resonances. This is true for the case of spherical nuclei where only the coupling of two GDR phonons to more complex, $3p3h$, configurations allows one to describe the DGDR width [8]. But the typical matrix element of this coupling in deformed nuclei does not exceed the value of 200 keV [23] while in spherical nuclei it is an order of magnitude larger. It means that because of the coupling, the strength of each GDR RPA phonon will fragment within the energy interval of 100–200 keV in deformed nuclei. The last value should be compared to the second moment m_2 presented in Table I which is the result of the Landau damping accounted for in our calculation. Taking into account that the reaction amplitude has a very weak energy dependence

and that mixing of different RPA phonons in the GDR wave function does not change the total strength [24], the total cross sections of the GDR and DGDR excitation in RHIC will be also conserved.

To conclude, we present the first theoretical studies based on microscopic calculations of the properties of the two-phonon giant dipole resonance in deformed nuclei in relativistic heavy ion collisions. We predict that the excitation function has a one-bump shape and that there are at least two

effects that work in favor of its experimental observation, namely, the energy shift to higher energies, and the narrowing of its width.

V.Yu. P. thanks the Instituto de Física of the Universidade Federal do Rio de Janeiro for the hospitality, and the CNPq for financial support. We thank Dr. L.A. Malov for fruitful discussions. This work was partially supported by the RFBR Grant No. 96-15-96729, by the FUJB/UFRJ, and by the MCT/FINEP/CNPq(PRONEX) Contract No. 41.96.0886.00.

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