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Possible enhancement of the $E1$ decay rates of double giant dipole states

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The $E1$ transitional rates between the double giant dipole resonance (DGDR) states and giant dipole resonance (GDR) ones are studied using the wave functions, which include one-, two-, and many-phonon components. It is shown that there is a possibility for the enhancement of the $E1$ transition rates from the DGDR states to the GDR ones as compared to the harmonic limit. This enhancement arises from the matrix elements for the $E1$ transitions between many-phonon terms of the wave functions, which differ by one $E1$ phonon. [S0556-2813(97)50408-0]

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The electromagnetic excitation of the double giant dipole (DGDR) in relativistic heavy-ion reactions has been studied in a wide range of nuclei [1–6]. The observed excitation energy of the DGDR is about twice as much compared to the energy of the giant dipole resonance (GDR). The observed width of the DGDR is larger than the GDR width by 1.6 times [3–6]. The most striking point is that the strength of the DGDR was found to be strongly enhanced as compared to any theoretical estimations available so far. Among theoretical studies, which have been devoted to the description of the DGDR states [7–13], we notice Ref. [12], where the general properties of the DGDR were treated making use of a sum rule approach. One of the major results obtained in Ref. [12] is, if the mean energy of the DGDR is twice as much compared to the GDR energy, the DGDR strength is two times square of that of the GDR. It means there is a discrepancy between the strength of the DGDR determined by the sum rules and the experimental systematics.

The aim of the present work is to study a possibility for the enhancement of the $E1$ decay as well as excitation of DGDR in a two-step process.

As has been shown in Ref. [14], the wave function of an excited state in the mean-field representation can be expressed in terms of one-, two-, ..., many phonon operators, acting on the wave function $|\Psi_0\rangle$ of the ground state of an even-even nucleus. This general form of the wave function was used recently in Ref. [15] to study the order-to-chaos transition in nuclei. The calculations in Ref. [16] have shown that there are fast $E1$ and $M1$ transitions between large components of the wave functions differing by the octupole and quadrupole phonons in deformed nuclei. The experimental observation of the fast $E1$ or $M1$ transitions between excited states indicates the existence of large many-phonon components in the wave function, which is very important for understanding the nuclear structure and the order-to-chaos transition.

The wave function of the $I^\pi=1^-$ state in the GDR region can be written in terms of many-phonon components as

$$|\Psi_n(1^-M)\rangle=\Big\{\sum_k R^n_{1k}Q^\dagger_{1Mk}+\sum_{\lambda_2 i_2,\lambda_3 i_3} P^n_{\lambda_2 i_2,\lambda_3 i_3}(1)\times[Q^\dagger_{\lambda_2 i_2}\otimes Q^\dagger_{\lambda_3 i_3}]_{1M}+\cdots\Big\}|\Psi_0\rangle, \quad (1)$$

where $Q^\dagger_{1Mk}$ and $Q^\dagger_{\lambda i}$ are the creation operators of a dipole phonon and a phonon with the multipolarity λ, respectively. The dipole phonons ($\lambda=1$) responsible for the $E1$ transitions are numerated by the index k, while i stands for the numeration of the sequence of phonons-spectators with the multipolarity λ, including $\lambda=1$. The index n denotes the state with $I^\pi=1^-$. The wave function in Eq. (1) has three-, four-phonon etc., terms. The condition of orthogonality and normalization of the wave function in Eq. (1) has the form

$$\sum_k R^n_{1k}R^{n'}_{1k}+\sum_{\lambda i\geqslant\lambda' i'} P^n_{\lambda i,\lambda' i'}P^{n'}_{\lambda i,\lambda' i'}+\mathcal{L}^{nn'}=\delta_{nn'}, \quad (2)$$

where the function $\mathcal{L}^{nn'}$ is responsible for the inclusion of three-, four-phonon, etc., terms in the normalization. It is worth noticing that, when the energy of the 2^+_1 state is not too large, the energy centroid of the two-phonon states, each of which consists of one dipole-phonon state from the GDR region and the first quadrupole-phonon 2^+_1 one, will be in the energy region of the GDR. The wave function in Eq. (1), limited to one- and two-phonon terms, was used in Refs. [17] to describe the spreading width of giant resonances and in Ref. [18] to estimate the contribution of the hexadecapole one-phonon and double-γ-vibrational components of the $K^\pi=4^+$ wave function in ^{168}Er. The latter was also studied in Ref. [19] making use of a similar wave function with one- and two-boson terms.

The $B(E1)$ value for the $E1$ transition from the ground state $|\Psi_0\rangle$ to the state with $I^\pi=1^-$ has the form

$$B(E1;0^+_{\rm g.s.}\to 1^-_n)=\Big|\sum_k R^n_{1k}\mathcal{M}^k(E1)\Big|^2. \quad (3)$$

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The explicit form of the matrix element $\mathcal{M}^k(E1)$ is included in Eq. (2.124) in Ref. [20]. The $E1$ strength is obtained by summing up the $B(E1)$ values over all the 1^- states in the GDR region: $S_1=\Sigma_n B(E1;0^+_{\rm g.s.}\to 1^-_n)$. This strength S_1 is slightly different from the nonenergy-weighted sum of the GDR mentioned in Ref. [12].

The wave function of the $J^\pi=0^+$ or 2^+ states in the DGDR region can be chosen in the form

$$|\Psi_f(J^+M')\rangle=\Bigg\{\sum_{i_1} R^f_{Ji_1}Q^\dagger_{JM'i_1}+\sum_{k_1,k_2} P^f_{1k_1,1k_2}(J)[Q^\dagger_{1k_1}\otimes Q^\dagger_{1k_2}]_{JM'}+\sum_k\sum_{\lambda_2 i_2,\lambda_3 i_3} F^f_{1k,\lambda_2 i_2,\lambda_3 i_3}(J)\times[Q^\dagger_{1k}\otimes[Q^\dagger_{\lambda_2 i_2}\otimes Q^\dagger_{\lambda_3 i_3}]]_{JM'}+\cdots\Bigg\}|\Psi_0\rangle. \quad (4)$$

As a matter of fact the inclusion of the two-phonon terms in the wave function does not change the energy centroid of the GDR. The similar situation is expected for the energy centroid of the DGDR after including many-phonon terms in the DGDR wave function. In odd-mass nuclei all terms of the wave functions in Eqs. (1) and (4) must be multiplied by the relevant quasiparticle operator $\alpha^\dagger_{jm\sigma}$, which describes the odd nucleon.

In order to study the Coulomb excitation of the DGDR in a two-step dipole process, we calculate the matrix elements of the $E1$ transition from the ground state to the GDR states and from the GDR states to the DGDR ones making use of the wave functions in Eqs. (1) and (4). The total strength of the DGDR states is obtained by summing over all final states $|f\rangle$ and intermediate $|n\rangle$ ones. As a result one has

$$S_2=\sum_f\left|\sum_{n,k,k'}\mathcal{M}^k(E1)\mathcal{M}^{k'}(E1)R^n_{1k}\left[\sum_{k''}R^n_{1k''}P^f_{1k',1k''}(J)+\sum_{\lambda i,\lambda' i'}P^n_{\lambda i,\lambda' i'}(J)F^f_{1k',\lambda i,\lambda' i'}(J)+\cdots\right]\right|^2. \quad (5)$$

Equation (5) contains many terms, which correspond to the transitions between three-phonon components in Eq. (1) and four-phonon components in Eq. (4) and so on. The second term in Eq. (5) corresponds to the GDR constructed on the two-phonon components of the wave function in Eq. (1). In Eq. (5) there are also terms, which correspond to the GDR built on three-, four- and so on components of the wave function in Eq. (1).

If from Eq. (2) it followed that

$$\sum_n R^n_{1k}P^n_{\lambda i,\lambda' i'}(J)=0,$$

$$\sum_n R^n_{1k}F^n_{\lambda i,\lambda_2 i_2,\lambda_3 i_3}(J)=0,\ \ldots, \quad (6)$$

there would be no enhancement in the Coulomb excitation of DGDR states, which arises from the $E1$ transitions between many-phonon components differing by one-dipole-phonon operator. However a possibility cannot be excluded, that the sums at the l.h.s. of Eqs. (6) do not vanish because of the terms $\mathcal{L}^{nn'}$ at the l.h.s. of Eq. (2). Therefore, generally speaking, the Coulomb excitation of the DGDR states can be larger than the estimation within the limit of independent harmonic oscillators. It is also important to notice that the largest part of the $E1$ strength, but not all the total $E1$ strength, is shared by the states within the GDR region and is observed in the experiments. Indeed, only a part of the two-phonon strength is shared by the states, which are defined by the wave function in Eq. (1) and, therefore, are excited via the $E1$ transitions. The rest of the two-phonon strength is distributed over the states within and beyond the GDR region, which are not excited via $E1$ transitions. The same situation takes place with many-phonon components in Eq. (1).

Let us consider now the $E1$ transitions from DGDR states to the GDR ones. The $B(E1)$ value for $E1$ transition from $J^\pi=0^+$ and 2^+ states in the DGDR region to the $I^\pi=1^-$ states in the GDR region has the form

$$B(E1;J^+_f\to 1^-_n)=\left|\sum_k\mathcal{M}^k(E1)\Bigg\{\sum_{k'}C^1_J R^n_{1k'}P^f_{1k,1k'}(J)+\sum_{\lambda_2 i_2\lambda_3 i_3}C^1_{\lambda_2\lambda_3}(J)P^n_{\lambda_2 i_2,\lambda_3 i_3}(1)\times F^f_{1k,\lambda_2 i_2,\lambda_3 i_3}(J)+\cdots\Bigg\}\right|^2, \quad (7)$$

where C^1_J and $C^1_{\lambda_2\lambda_3}(J)$ denote the angular-momentum coupling coefficients. It should be without confusion that now indices f denote the initial states and n—the final states.

The total strength of the $E1$ transitions from the DGDR states to the GDR ones is obtained by summing over all initial states $|f\rangle$ and final $|n\rangle$ ones. If taking Eq. (6) into account, one ends up with

$$\sum_J \sum_{n,f} B(E1;J_f^+ \to 1_n^-) = \sum_J \sum_{n,f} \left[\left| \sum_{kk'} \mathcal{M}^k(E1) C_J^1 R_{1k'}^n P_{1k,1k'}^f(J) \right|^2 + \left| \sum_k \mathcal{M}^k(E1) \sum_{\lambda_2 i_2, \lambda_3 i_3} C_{\lambda_2\lambda_3}^1 P_{\lambda_2 i_2,\lambda_3 i_3}^n (J) F_{1k,\lambda_2 i_2,\lambda_3 i_3}^f (J) \right|^2 + \cdots \right]. \quad (8)$$

The first term at the r.h.s. of Eq. (8) corresponds to the transitions from the two-dipole-phonon to the one-dipole-phonon configurations. It becomes, therefore, clear from Eqs. (8) that the $E1$ strength between the DGDR and GDR can be enhanced because of the large matrix elements $\mathcal{M}^k(E1)$ for the transitions between three- and two-phonon components as well as between many-phonon components, which differ by one dipole-phonon operator. Thus, the $E1$ transitions from the relevant three-phonon components of the wave function in Eq. (4) to the two-phonon components $[1_i^- \otimes 2_1^+]_{1^-}$ of the wave function in Eq. (1) can be strongly enhanced. It may be the situation observed in ^{136}Xe and ^{197}Au. A similar large enhancement is unlikely to take place in ^{208}Pb as the energy of the first quadrupole phonon is large. It is a proposal of the present paper that this enhancement can be observed in the γ decay, which corresponds to the $E1$ transition from the DGDR to the GDR.

An enhancement is expected to occur also in the excitation of the DGDR states via the consequent $E1$ transitions from the ground state to the GDR states and from the latter to the DGDR states. As concerning the regions beyond the GDR and DGDR, the $E1$ transitions between many-phonon components, which differ by one dipole phonon, should lead to the background because there is no such large $E1$ matrix elements.

In conclusion, the present paper shows that:

(1) The $E1$ transition rates from the DGDR states to the GDR ones can be enhanced as compared to those between the two-dipole-phonon components in the wave functions of the DGDR states and the one-dipole-phonon components in the wave functions of the GDR states. This enhancement can be observed experimentally in the γ-decay spectra of the $E1$ transition from the DGDR to the GDR.

(2) It is possible to expect an enhancement of the DGDR states, which are formed in a two-step process of two consequent $E1$ transitions rather than in a Coulomb excitation.

The observation of the DGDR also may serve as a signature of orders, which take place at such high excitation energies in nuclei.

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