Description of Decay Mechanisms of the Giant Dipole Resonances with a $RPA+FKK$ Approach

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The decay mechanisms of giant resonances have been revisited to investigate the isoscalar dipole resonance in ^{208}Pb nucleus.

I. INTRODUCTION

The purpose of this work is to present a feasible scheme to calculate the giant resonances (*GR*'s) widths in order to investigate the decay mechanisms of these collective excitation modes. This scheme is based on a calculation model where the escape and the spreading widths are directly connected. The method is applied for isoscalar giant dipole resonances (*ISGDR*) in ^{208}Pb, since new results are indicating the necessity of improvements in the structure calculations around the *ISGDR* [1–5]. The escape width was calculated using a previous version of the continuum *RPA* approach [6], which was modified to take into account the differences among the neutron and proton radii in nuclei with neutron excess [7, 8]. To calculate the spreading width we present a new formulation of an approximated semi-microscopic method, based in the statistical Multi-Step Compound Theory (*MSC*) of Feshbach, Kerman and Koonin (*FKK*) [9] in connection with the *RPA* calculations.

In the Sect. II of this work, we describe the theoretical approach used in the spreading width calculation. The results are discussed in the Sect. III.

II. $RPA+FKK$ APPROACH

A recent version [7] of the continuum Random Phase Approximation (*RPA*) that accounts for a careful treatment on the differences between the neutron and proton densities was used to calculate the escape width. The details of this calculation may be seen in the contribution for this proceeding: "Partial Escape Width for Nuclei with Neutron Excess" [8].

In order to perform the calculation of the spreading width, we have used the Statistical Multi-Step Compound Theory (*MSC*) of Feshbach, Kerman and Koonin (*FKK*) [9]. In this approach, the excitation of the *GR* takes place in a number of stages, with the particle emission being allowed in any of them. Each stage is represented by a level of complexity, characterized by the number of particle-hole pairs that are excited by an external field [10]. Assuming an energy and angular momentum factorization of the state density, the spreading width in this formalism is written as: $\langle\Gamma_{nJ}^{\downarrow n+1}(E)\rangle = X_{nJ}^{\downarrow n+1}Y_n^{\downarrow n+1}(E)$. The angular momentum structure included in the particle-particle two-body interaction and the spin distribution of the single-particle levels are contained in the X function. The Y function contains all the dependence on excitation energy originated from the final state density, the particle-hole distinguishability, and describes the available phase space for the transition. Thus, the X and Y functions are calculated separately, resulting in a complete uncoupling between the angular momentum of the excited particle-hole pairs and the energy considered in the calculation of the width. This way, all $1p-1h$ pairs are treated on equal footing in any energy, making the results strongly dependent on the particle-hole basis (for more details of the FKK approach see Refs. [9] and [10]). On the other hand, in the *RPA* approach, the excitation probability of a specific particle-hole pair depends on the excitation energy and angular momentum coupling. Therefore, since we have accounted for a sufficiently complete $1p-1h$ basis in the *RPA* calculation, the results do not change considerably when another more internal (or more external) single-particle level is included in the configuration basis. Thus, intending to minimize the dependence on the $1p-1h$ basis and the number of possible intermediate coupling effects, we have implemented some modifications in the original form [9] of the X function to include the microscopic information calculated by *RPA*. In Ref. [9] the density of single-particle(hole) states is derived from the equidistant single-particle model, resulting in a n-stage $Y_n(E)$ function with a direct dependence on E^2. Therefore, for resonances at energies far away from the nucleon binding energy (as in isoscalar *GDR*) this outcome is not a good option because the spreading width involves only intermediate bound states. This way, we use the $Y_n(E)$ function proposed by Oblozinsky [11] where the level density is obtained restricting the nucleons to bound states, limiting the energy dependence.

Thus, the main improvement proposed in this work consists of taking into account the excitation probability of each $1p-1h$ pair that is accessed in the energy in which the calculations are performed (this proposal is hereafter referred to as *RPA*$+FKK$ approach). The proposed new X function is given by:

$$
\begin{aligned}
&X_{nJ}^{\downarrow(n+1)}(\varepsilon_m) = \\
&2\pi \sum_{jQj_4j_3} \mathcal{P}_{Qj_4}^{J} \frac{R_1(j)R_1(Q)R_{N-1}(j_4)}{R_N(J)} \frac{(2j_3+1)}{(2Q+1)} \Delta(QJj_4) \\
&\sum_{j_1j_2} \mathcal{P}_{j_1j_2}^{j_3} R_1(j_1)R_1(j_2)\left(\langle\{j_1j_2\}_{j_3}|V|\{Qj\}_{j_3}\rangle\right)^2 L(\varepsilon_m,E_2)
\end{aligned}
\tag{1}
$$

TABLE I: The escape [8] and spreading width of *GDR* in ^{208}Pb.

	$E(MeV)$	$\Gamma^{\uparrow}(MeV)$	$\Gamma^{\downarrow}(MeV)$	$\Gamma(MeV)$
This work				
IVGDR	10.8	0.0	4.4	4.4
ISGDR	24.4	2.7	3.3	6.0
Experimental values				
IVGDR	13.5			4.0[12]
ISGDR	$20-23$			$2.5-10$[1–4]

where

$$\mathcal{P}^J_{j_p j_h} = \frac{N_P |x^J_{j_p j_h}|^2}{\sum\limits_{j_p j_h(\Gamma_p \approx 0)} |x^J_{j_p j_h}|^2}\ , \tag{2}$$

and N_p is the number of bound (or with single-particle width too smaller than *GR* width) $1p-1h$ configurations coupling to $\overrightarrow{J}$. The functions $R_N(J)$ and $\Delta(abc)$ are given, respectively, by Eqs. (5.7) and (5.26) from Ref. [9]. The amplitudes $x^J_{Qj_4}$ and $x^{j_3}_{j_1 j_2}$, used in the evaluation of $P^J_{Qj_4}$ and $P^{j_3}_{j_1 j_2}$, are calculated in the excitation energy ε_m and intermediary phonon energy ε_{j_3}, respectively. All energies ε_j, as well as the amplitudes $x^J_{jj'}$, are obtained with *RPA* calculation. The quantity $L(\varepsilon_m, E_{2p-2h})$ is a lorentzian type function:

$$L(\varepsilon_m, E_{2p-2h}) = \frac{\eta^2}{(\varepsilon_m - E_{2p-2h})^2 + \eta^2}, \tag{3}$$

with $E_{2p-2h} = \varepsilon_j + \varepsilon_{j_1} - \varepsilon_{j_4} - \varepsilon_{j_2}$. We would like to emphasize that the calculation of the X function is performed only in the excitation energy (E) which correspond to the energy solution (ε_m) of the *RPA* formalism. The modifications in the original X function [9] are given by the $P^J_{j_p j_h}$ and $L(\varepsilon_m, E_{2p-2h})$ factors.

Another point that deserves a special attention is the part that deals with the residual interaction used in the calculation. The coherence of the $RPA+FKK$ approach is guaranteed by using the same two-body residual interaction, as the Landau-Migdal interaction given in Ref. [7], to diagonalize the *RPA* equations and to compute the X function in Eq.(1). Proceeding this way, the parameters of the residual interaction were adjusted at level of the *RPA* calculation in a standard way, in order to reproduce the excited states with the lowest energies and to eliminate the 1^- spurious solution at zero energy. However, in the *FKK* calculations [9, 10], the two-body interaction was assumed to be the simplest zero-range form where the strength of the interaction is a free parameter.

III. RESULTS AND DISCUSSIONS

The first application test refers to the spreading width calculation of the isovector *GDR* in ^{208}Pb. The general statement found in the literature is that the decay of this resonance in heavy nucleus is largely dominated by the statistical mechanism, corresponding to a more complex structure than $1p-1h$, which is in agreement with our results showed in TABLE I.

The calculated spreading width is approximately equal to the total width, showing the largest domain of the statistical decay. On the other hand, the isoscalar *GDR* in ^{208}Pb is located far away from the isovector one, and it presents a large direct decay. The strength of this excitation (calculated by *RPA* in Ref. [8]) is separated in two main components, in agreement with all experimental data. The lower-energy component is due to the remaining isovector contribution with a spurious state mixing character. The highest energy and broadest peak is identified as the *ISGDR*. Our calculations predict a considerable strength in the energy region above 20 MeV, which is composed by the presence of various narrow peaks superposing to exhaust about 82% of the *EWSR* between $20-30$ MeV. These peaks are mainly composed by $3\hbar\omega$ transitions involving the neutrons and protons of the external shells. Using the results provided by *RPA* calculation, we obtained the spreading width $\langle\Gamma^{\downarrow}\rangle = 3.3$ MeV (see TABLE I), resulting a total width $\langle\Gamma\rangle = 6.0$ MeV, in agreement whith the fits from Ref.[4], and also, with the value of 3.2 MeV used in the analysis of the Ref.[5]. The results show that the spreading and escape widths have compatible values. The large direct decay reflects the fact that this resonance is located in higher energy than the isovector *GDR* and far beyond the neutron threshold. Thus, the direct decay channel is quite favored by this high energy and it competes equally with the statistical mechanism, unlike of what happens with the isovector one.

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