Structure and direct nucleon decay properties of isoscalar giant monopole and dipole resonances

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The strength function and partial widths for the direct nucleon decay of the isoscalar giant monopole and dipole resonances are analyzed within an extended continuum-random-phase-approximation approach. Calculations are performed for several medium and heavy mass nuclei with the use of a phenomenological nuclear mean field, the Landau-Migdal particle-hole interaction, and some partial self-consistency conditions. Calculation results are compared with available experimental data.

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I. INTRODUCTION

By now a number of experiments on the excitation of the isoscalar giant monopole and dipole resonances (ISGMR and ISGDR, respectively) have been performed mainly by means of the (α,α') reaction. Some experimental data on partial widths (branching ratios) for direct neutron decay of the ISGMR in ^{90}Zr [1], ^{124}Sn [2], and ^{208}Pb [3] were deduced from the $(\alpha,\alpha' n)$ reaction. Recently, the ISGDR has been identified in several nuclei: ^{90}Zr, ^{116}Sn, ^{144}Sm, and ^{208}Pb [4]. There is a plan to measure the partial branching ratios for the neutron decay of the ISGDR in ^{208}Pb [5].

Interest in the location of the excitation energies of isoscalar compression modes in nuclei is mainly connected with the hope of determining the nuclear compressibility from these data (see, e.g., Ref. [6]). However, study of the direct nucleon decay of various giant resonances (GRs) is of special interest, because such study allows one to obtain information on the microscopic structure of giant resonances. The peculiarities of the ISGDR are connected with the fact that this resonance is an overtone of the spurious state, associated with the center-of-mass motion. Moreover, because of the rather large excitation energy, the single-particle continuum makes an essential contribution to formation of the ISGDR.

This work is stimulated by the accumulation of the above-mentioned experimental data. The data are analyzed within an extended continuum-random-phase-approximation approach (CRPA), which also includes a phenomenological description of the doorway-state coupling to many-quasiparticle configurations. In the present work we adopt the approach formulated in Refs. [7,8], where it was applied to isovector GR. A phenomenological mean field and the Landau-Migdal particle-hole interaction are used in the calculations together with some partial self-consistency conditions. The isoscalar part of the interaction is somewhat modified in order to reproduce, in the calculations, the above-mentioned spurious state at zero energy. Thus, our approach is not fully self-consistent and we check its abilities in the description of both the GR excitation energy and the GR relative strength [as compared with the relevant energy-weighted sum rule (EWSR)]. Then the approach is applied to the description of the direct nucleon decay of the considered GRs. Earlier attempts along this line were undertaken in Refs. [9,10] in calculations of the direct neutron decay of the ISGMR in ^{208}Pb. Basic relationships, used in the approach adopted in this work, are given in Sec. II. The values of model parameters and the results of the calculation are presented in Sec. III, together with the corresponding experimental data. Discussion of the results and a summary are given in Sec. IV.

II. BASIC RELATIONSHIPS

The CRPA equations are given below in the form accepted within the finite Fermi-system theory [11]. Let $\hat{V}e^{-(i/\hbar)\omega t}+\text{H.c.}$, where $\hat{V}=\Sigma_a V(x_a)$, be an external periodic-in-time single-particle field, acting upon the nucleus in the process of GR excitation. For consideration of the isoscalar GR of multipolarity L this field can be chosen in the form $V_L(x)=V_L(r)Y_{L0}(\vec{r}/r)$. Within the CRPA the strength function corresponding to field $\hat{V}_L$ is determined by the following expression:

$$S_L(\omega)=-\frac{1}{\pi}\text{Im}\, P_L(\omega),$$

$$P_L(\omega)=\sum_\alpha \int V_L^*(r)A_\alpha^L(r,r',\omega)\tilde{V}_{L,\alpha}(r',\omega)drdr'. \quad (1)$$

Here, $P_L(\omega)$ is the nuclear polarizability corresponding to the above-chosen external field, $(r^2r'^2)^{-1}A_\alpha^L(r,r',\omega)$ is the radial part of the free particle-hole propagator of multipolarity L for neutrons $(\alpha=n)$ and protons $(\alpha=p)$, ω is the excitation energy, and $\tilde{V}_{L,\alpha}$ are the so-called effective fields. They satisfy the system of integral equations

$$\tilde{V}_{L,\alpha}(r,\omega)=V_L(r)+\sum_\beta \frac{F_{\alpha\beta}(r)}{r^2}\int A_\beta^L(r,r',\omega)\tilde{V}_{L,\beta}(r')dr'. \quad (2)$$

Here, $F_{\alpha\beta}(r)=\{[F(r)+F'(r)]\delta_{\alpha\beta}+[F(r)-F'(r)]\tau_{\alpha\beta}^{(1)}\}$ is the matrix of the spin-independent part of the Landau-Migdal particle-hole interaction $(F+F'\vec{\tau}\cdot\vec{\tau}')\delta(\vec{r}-\vec{r}')$, where F and F' are phenomenological quantities and $\vec{\tau}=(\tau^{(1)},\tau^{(2)},\tau^{(3)})$ are the isospin Pauli matrices.

The expression for the propagator $A_\alpha^L(r,r',\omega)$ can be presented in the form

$$A_{\alpha}^{L}=\sum_{\mu,(\lambda)}(t_{(\lambda)(\mu)}^{L})^{2}n_{\mu}^{\alpha}\chi_{\mu}^{\alpha}(r)\chi_{\mu}^{\alpha}(r')[g_{(\lambda)}^{\alpha}(r,r',\varepsilon_{\mu}+\omega)+g_{(\lambda)}^{\alpha}(r,r',\varepsilon_{\mu}-\omega)], \quad (3)$$

which allows one to take into account exactly the single-particle continuum [12]. Here, $n_{\mu}=N_{\mu}(2j_{\mu}+1)^{-1}$ is the occupation number, N_{μ} is the number of nucleons occupying the single-particle level $\mu(=\{\varepsilon_{\mu},j_{\mu},l_{\mu}\})$, $r^{-1}\chi_{\mu}(r)$ is the radial (real) bound-state wave function, $g_{(\lambda)}(r,r',\varepsilon)$ is the Green's function of the radial single-particle Schrödinger equation with $(\lambda)=j_{\lambda},l_{\lambda}$, and $t_{(\lambda)(\mu)}^{L}=(2L+1)^{-1/2}\langle(\lambda)||Y_{L}||(\mu)\rangle$ is the kinematic factor.

An alternative representation of $S_{L}(\omega)$, which is more convenient for the description of the GR nucleon decay, follows from Eqs. (1) and (2):

$$S_{L}(\omega)=-\frac{1}{\pi}\mathrm{Im}\sum_{\alpha}\int\tilde{V}_{L,\alpha}^{*}(r,\omega)A_{\alpha}^{L}(r,r',\omega)\times\tilde{V}_{L,\alpha}(r',\omega)drdr'=\sum_{c}|M_{c}^{L}|^{2}, \quad (4)$$

where the expression for the amplitude M_{c}^{L} has the form

$$M_{c}^{L}(\omega)=(n_{\mu}^{\alpha})^{1/2}t_{(\lambda)(\mu)}^{L}\int\chi_{\varepsilon(\lambda)}^{\alpha}\tilde{V}_{L,\alpha}(r,\omega)\chi_{\mu}^{\alpha}dr. \quad (5)$$

Here, $r^{-1}\chi_{\varepsilon(\lambda)}^{\alpha}$ is the radial (real, normalized to the δ function of the energy) continuum-state wave function of the escaping nucleon with energy $\varepsilon=\varepsilon_{\mu}+\omega$; $c=\mu,(\lambda),\alpha$ is the set of the decay-channel quantum numbers. Two comments can be made concerning the above-given equations. (i) In the description of the GR gross properties, the nucleon pairing can be neglected within an accuracy of $(2\Delta_{pair}/\bar{\omega})^{2}$, where Δ_{pair} is the energy gap and $\bar{\omega}$ is the GR mean excitation energy; (ii) in the description of the GR direct nucleon decay, the modification of the single-hole strength in the product nucleus due to the nucleon pairing and (or) to the coupling to low-lying collective states can be phenomenologically taken into account by using in Eq. (5) the experimental reduced spectroscopic factor s_{μ} instead of occupation factor n_{μ}.

The GR relative strength is characterized by the ratio

$$x_{L}=\int_{(\delta)}\omega S_{L}(\omega)d\omega/(\mathrm{EWSR})_{L}, \quad (6)$$

where $\delta=\omega_{2}-\omega_{1}$ is a rather large excitation energy interval in the vicinity of the GR, and

$$(\mathrm{EWSR})_{L}=(\hbar^{2}/2m)\int[(dV_{L}/dr)^{2}+L(L+1)\times(V_{L}/r)^{2}]\rho^{(+)}(r)r^{2}dr \quad (7)$$

is the only slightly model-dependent energy-weighted sum rule (m is the nucleon mass, $\rho^{(+)}=\rho^{n}+\rho^{p}$ is the nuclear density) [13]. The mean excitation energy $\bar{\omega}_{L}$ and the root-mean-square (rms) energy dispersion Δ_{L} are determined by the strength function $S_{L}(\omega)$ and the considered excitation energy interval δ.

We choose the radial dependence of the external field $V_{L}(r)$ in such a way that the considered GR would exhaust the major part of the corresponding $(\mathrm{EWSR})_{L}$. Therefore, for consideration of the ISGMR, the above-mentioned spurious state (s.s.) and the ISGDR we choose the fields $V_{L=0}(r)=(r/R)^{2}$, $V_{L=1}^{s.s.}(r)=(r/R)$, and $V_{L=1}(r,\eta)=(r/R)^{3}-\eta(r/R)$ (R is the radius of the nucleus), respectively. The parameter η is determined by the condition that the spurious state have zero strength relative to the field $V_{L=1}(r,\eta)$. The calculations can be considered as sufficiently self-consistent provided that the spurious state has zero energy and exhausts the major part of the $(\mathrm{EWSR})_{L=1}^{s.s.}$, and the above parameter η is close to $5\langle r^{2}\rangle/3R^{2}$, which appears in the EWSR of the ISGDR [14].

In the vicinity of the GR with not-too-large excitation energy the RPA amplitudes $P_{L}(\omega)$ and $M_{c}^{L}(\omega)$ exhibit one or several rather narrow resonances. They correspond to the particle-hole-type doorway states forming the GR. The Breit-Wigner parametrization of the mentioned amplitudes,

$$P_{L}(\omega)=\sum_{g}\frac{R_{g}}{\omega-\omega_{g}+\frac{i}{2}\Gamma_{g}^{\uparrow}},$$

$$|M_{c}^{L}(\omega)|=\frac{1}{\sqrt{2\pi}}\left|\sum_{g}\frac{R_{g}^{1/2}(\Gamma_{gc}^{\uparrow})^{1/2}}{\omega-\omega_{g}+\frac{i}{2}\Gamma_{g}^{\uparrow}}\right|, \quad (8)$$

allows one to deduce the doorway-state parameters: excitation energy ω_{g}, partial strength R_{g}, and partial and total escape widths $\Gamma_{gc}^{\uparrow}$ and $\Gamma_{g}^{\uparrow}$, respectively. The possibility of using the above parametrization can be checked by satisfying the equality $\Gamma_{g}^{\uparrow}=\Sigma_{c}\Gamma_{gc}^{\uparrow}$, which follows from Eqs. (1), (4), and (8). Note that the signs of $R_{g}^{1/2}$ and $(\Gamma_{gc}^{\uparrow})^{1/2}$ are not determined; only their product is used below. In accordance with Eq. (8) the relative and double-relative strengths of the doorway-state resonance equal $x_{L,g}=\omega_{g}R_{g}/(\mathrm{EWSR})_{L}$ and $y_{L,g}=x_{L,g}/x_{L}$, respectively. The sums $x_{L}=\Sigma_{g}x_{L,g}$ and $\bar{\omega}_{L}=\Sigma_{g}\omega_{g}R_{g}/(\Sigma_{g}R_{g})$, taken over δ, define, respectively, the GR relative strength and the GR mean excitation energy.

The expansion of the effective fields (2) in terms of the doorway-state resonances (see, e.g., Ref. [11]),

$$\tilde{V}_{L,\alpha}(r,\omega)=V_{L}(r)+\sum_{g}\frac{R_{g}^{1/2}v_{g\alpha}^{L}(r)}{\omega-\omega_{g}+\frac{i}{2}\Gamma_{g}^{\uparrow}}, \quad (9)$$

allows one to calculate the transition potentials $v_{g\alpha}^{L}(r)=\Sigma_{\beta}F_{\alpha\beta}(r)\rho_{g\beta}^{L}(r)$ and, as a result, the transition density $\rho_{g}^{L}(r)=\Sigma_{\alpha}\rho_{g\alpha}^{L}(r)$ for each doorway state with strength $R_{g}^{1/2}=\int\rho_{g}^{L}(r)V_{L}(r)r^{2}dr$. Note that for $L=0$, $\int\rho_{g}^{L}(r)r^{2}dr=0$, as follows from Eqs. (2) and (3).

The doorway-state coupling to many-quasiparticle configurations leads to the formation of the GR as a single resonance in the energy dependence of energy-averaged reaction cross sections. Following Refs. [7,8] we take this coupling into consideration phenomenologically by adopting an independent spreading width for each doorway-state resonance. It means that the transition to the energy-averaged amplitudes $\overline{P}_L(\omega)$ and $\overline{M}^L_c(\omega)$ can be realized by the following substitution in Eq. (8): $\Gamma^{\uparrow}_g \to \Gamma_g = \Gamma^{\uparrow}_g + \Gamma^{\downarrow}$. The doorway-state spreading width $\Gamma^{\downarrow}$ is considered as the only adjustable parameter in the approach. It can be found by equating the total width Γ_{tot} (dependent on $\Gamma^{\downarrow}$) of the energy-averaged strength function of the GR,

$$\overline{S}_L(\omega) = -\frac{1}{\pi}\,\mathrm{Im}\sum_g \frac{R_g}{\omega-\omega_g+\frac{i}{2}\Gamma_g}, \tag{10}$$

to the total width Γ^{expt}_{tot} of the GR in the experimental inclusive reaction cross section.

Because each doorway-state resonance in the energy dependence of the amplitudes $\overline{M}^L_c(\omega)$ becomes rather broad, as a rule it is also necessary to take into account the variation of the penetrability of the potential barrier for escaping nucleons over the resonance. It can be done as follows [8]:

$$|\overline{M}^L_c(\omega)| = \frac{1}{\sqrt{2\pi}}\left|\sum_g \frac{R_g^{1/2}(\overline{\Gamma}^{\uparrow}_{gc})^{1/2}}{\omega-\omega_g+\frac{i}{2}\Gamma_g}\right|,$$

$$\overline{\Gamma}^{\uparrow}_{gc} = \Gamma^{\uparrow}_{gc}\overline{P}_{g,\alpha,(\lambda)}/P_{\alpha,(\lambda)}(\varepsilon_{g\mu}). \tag{11}$$

Here, $\varepsilon_{g\mu} = \omega_g + \varepsilon_\mu$, $\overline{P}_{g,\alpha,(\lambda)}$ is the penetrability averaged over the resonance:

$$\overline{P}_{g,\alpha,(\lambda)} = \frac{1}{\sqrt{2\pi}\sigma_g}\int P_{\alpha,(\lambda)}(\varepsilon)\exp\left(-\frac{(\varepsilon-\varepsilon_{g\mu})^2}{2\sigma_g^2}\right)d\varepsilon,$$

$$\sigma_g = \Gamma_g/2.35. \tag{12}$$

Thus, the energy-averaged partial ''cross section'' $\overline{\sigma}^L_\mu(\omega) = \Sigma_{(\lambda)}|\overline{M}^L_c(\omega)|^2$ can be calculated without the use of free parameters. The summation in the last equation is performed over quantum numbers of the escaping nucleon, which are compatible with the selection rules for L transitions. ''Cross section'' $\overline{\sigma}^L_\mu(\omega)$ corresponds to the population of the single-hole state μ^{-1} in the product nucleus after the GR nucleon decay.

The GR branching ratios and partial widths for direct nucleon decay are defined as follows:

$$b^L_\mu = \int \overline{\sigma}^L_\mu(\omega)d\omega \Big/ \int \overline{S}_L(\omega)d\omega \tag{13}$$

and $\overline{\Gamma}^{\uparrow}_\mu = b^L_\mu \Gamma_{tot}$. As follows from Eq. (13) when only one doorway state corresponds to the considered GR (for instance, in the case of the ISGMR) one can calculate the average partial escape widths of this GR directly with the help of Eqs. (10)–(12).

III. VALUES OF THE MODEL PARAMETERS AND RESULTS

The nuclear mean field and particle-hole interaction are the input data for any CRPA approach. In the calculations presented in this work the phenomenological nuclear mean field is taken as the sum of isoscalar, spin-orbit, isovector, and Coulomb parts:

$$U(x) = U_0(r) + U_{so}(r)\vec{\sigma}\cdot\vec{l} + \frac{1}{2}\tau^{(3)}v(r) + \frac{1}{2}(1-\tau^{(3)})U_c(r). \tag{14}$$

Here, $U_0(r) = -U_0 f_{ws}(r,R,a)$,

$$U_{so}(r) = -U_{so}\frac{1}{r}\frac{df_{ws}}{dr},$$

where f_{ws} is the Woods-Saxon function with $R = r_0A^{1/3}, r_0 = 1.24$ fm, and $a = 0.63$ fm. The values of the parameters used in the calculations are the following:

$$U_0 = 54\ \text{MeV},\quad U_{so} = 13.9\left(1+2\frac{N-Z}{A}\right)\ \text{MeV fm}^2. \tag{15}$$

The symmetry potential $v(r)$ in Eq. (14) was determined in a self-consistent way (see, e.g., Ref. [15]): $v(r) = 2F'\rho^{(-)}(r)$, where $\rho^{(-)} = \rho^n - \rho^p$ is the neutron excess density. The Coulomb part in Eq. (14) was calculated in a Hartree approximation via the proton density ρ^p. The intensities of the phenomenological Landau-Migdal particle-hole interaction are parametrized in the standard form (see, e.g., Refs. [11,10]),

$$F(r) = C[f^{ex} + (f^{in} - f^{ex})f_{ws}(r,R,a)],\quad F'(r) = Cf', \tag{16}$$

with the following parameters:

$$C = 300\ \text{MeV fm}^3,\quad f' = 1.0,$$

$$f^{in} = -0.0875,\quad f^{ex} \simeq -(2.3--2.5). \tag{17}$$

The above-given mean field allows one to reproduce satisfactorily the nucleon separation energies B_α for closed-shell nuclei using the Landau-Migdal parameter f'. The parameter f^{ex} in Eqs. (17) was slightly varied to make the spurious-state energy close to zero for each considered nucleus (see Table I). In accordance with the spectral expansion of the polarizability (see, e.g., Ref. [11]), in the limit $\omega \to \omega_{s.s.} \to 0$, it can be presented in the form

$$P_{L=1}(\omega) \to \frac{2\omega_{s.s.}R_{s.s.}}{\omega^2-\omega^2_{s.s.}},\quad x_{s.s.} = \frac{\omega_{s.s.}R_{s.s.}}{(\text{EWSR})^{s.s.}_{L=1}}. \tag{18}$$

The calculated relative strengths $x_{s.s.}$ are also listed in Table I.

The parameter η in the expression for the field $V_{L=1}(r,\eta)$ was taken initially to be $5\langle r^2\rangle/3R^2$, where $\langle r^2\rangle = 4\pi A^{-1}\int r^4\rho^{(+)}(r)dr$ with $\rho^{(+)} = \rho^n + \rho^p$. For this choice of η the spurious state exhausts less than 0.1% of the $(\text{EWSR})_{L=1}$. This fact, which allows us to use the above

TABLE I. Calculated and experimental nucleon separation energies for some closed-shell nuclei. The isoscalar strength of the particle-hole interaction and the relative strength of the spurious state are also given.

	^{90}Zr	^{116}Sn	^{124}Sn	^{144}Sm	^{208}Pb
B_n^{calc} [MeV]	11.64			11.59	7.24
B_n^{expt} [MeV]	11.98			10.55	7.37
B_p^{calc} [MeV]		9.25	12.58		7.99
B_p^{expt} [MeV]		9.28	12.11		8.01
$-f^{ex}$	2.36	2.42	2.40	2.41	2.51
$x_{s.s.}$ [%]	96	93	93	96	92

expression for η, is an additional argument in favor of the rather good self-consistency of the model used in this work. Thus, all the model parameters used below in the calculations of some characteristics of isoscalar monopole and dipole resonances in several medium-heavy nuclei are fixed from independent data.

We start by describing the results obtained for the ISGMR in ^{90}Zr, ^{124}Sn, and ^{208}Pb. In the resonance region the energy dependence of the monopole strength function of these nuclei, calculated by Eqs. (1)–(3), do not exhibit any gross structure. The results of the approach for the ISGMR energy (obtained within 10% of the experimental value) and the relative strength are given in Table II, where a comparison with corresponding experimental data is given. Partial nucleon escape widths of the ISGMR were calculated using Eqs. (10)–(12). To take the averaged potential barrier penetrability more accurately into consideration in accordance with Eq. (12) we used the experimental decay-channel energies $\varepsilon_{g\mu}^{expt}=\omega_g^{expt}+\varepsilon_\mu^{expt}$. The calculated partial widths are also multiplied by the experimental reduced spectroscopic factors s_μ of the corresponding single-hole state. The calculation results together with available experimental data for the above-mentioned nuclei are given in Tables III–V.

The transition density calculated by Eq. (9) for the ISGMR in ^{208}Pb is shown in Fig. 1 in comparison with the quantity $\rho^{tr}(r)\sim[3\rho+rd\rho/dr]$ obtained within the scaling model [18] ($\rho=\rho^{(+)}$).

The mean energy, rms energy dispersion, and relative strength calculated with the use of the corresponding CRPA

TABLE II. Calculated and experimental parameters of the ISGMR.

	^{90}Zr	^{124}Sn	^{208}Pb
$\omega_1-\omega_2$ [MeV]	13.5–16.5	13–16	11–14
$x_{L=0}$ [%]	90	92	92
ω_g[MeV]	14.9	14.6	12.7
ω_g^{expt} [MeV]	16.1±0.4 [a]	15.40±0.35 [b]	13.9±0.3 [c]
	17.9±0.4 [d]		14.2±0.3 [d]
Γ_{tot}^{expt} [MeV]	3.1±0.4 [a]	3.8±0.5 [b]	2.4±0.3 [c]

[a]Values from Ref. [1].
[b]Values from Ref. [2].
[c]Values from Ref. [3].
[d]Values from Ref. [4].

TABLE III. Calculated and experimental partial nucleon escape widths of the ISGMR in ^{90}Zr.

μ^{-1}	E_x[MeV]	$\overline{\Gamma_\mu^\uparrow}$ [keV][a]	$\Gamma_\mu^\uparrow$ [keV][b]
Neutron			
9/2$^+$	0.00	59	155±31[c]
1/2$^-$	0.59	358	∼220
Proton			
1/2$^-$	0.00	221	∼175
9/2$^+$	0.91	9	
3/2$^-$	1.51	187	∼80
5/2$^-$	1.74	10	

[a]$s_\mu=1$ is taken for all states μ^{-1}.
[b]The experimental values deduced without extraction of the statistical decay contribution [5].
[c]Value from Ref. [1].

strength function for the ISGDR in several nuclei are given in Table VI together with some experimental data.

The CRPA strength function for the ISGDR exhibits an essential gross structure. We attempted to expand the strength function of the ISGDR in ^{208}Pb in terms of doorway-state resonances of the Breit-Wigner type (see Fig. 2). A specific feature of this expansion is the appearance of a rather broad resonance-like structure with total escape width $\Gamma_g^\uparrow\simeq1.9$ MeV. As has been checked by direct calculations, performed in the spirit of Ref. [20], valence particle-hole transitions are mainly responsible for the appearance of this resonance. In view of the rather strong coupling of the ISGDR to the continuum, the above-mentioned expansion of the CRPA strength function for ^{208}Pb is considered to be a crude approximation. In particular, the deviation from equality $\Gamma_g^\uparrow=\Sigma_c\Gamma_{gc}^\uparrow$ for the main doorway-state resonances is about 30–40 %. For this reason the accuracy of calculations of partial nucleon branching ratios by Eqs. (10)–(13) is expected to be not too high. These branching ratios were evaluated using the experimental one-hole energies ε_μ^{expt}, the reduced spectroscopic factors s_μ for the corresponding single-hole states and mean doorway-state spreading width $\Gamma^\downarrow=3.2$ MeV. This value was found by equating the total width the energy-averaged strength function to $\Gamma_{tot}^{expt}=5$ MeV (see, e.g., Ref. [19]). Calculated partial branching ratios for the main decay channels are given in Table VII.

TABLE IV. Calculated and experimental partial neutron escape widths of the ISGMR in ^{124}Sn.

μ^{-1}	E_x[MeV]	$(2j_\mu+1)s_\mu$ [a]	$\overline{\Gamma_\mu^\uparrow}$ [keV]	$\Gamma_\mu^\uparrow$ [keV] [a]
11/2$^-$	0.00	9.3±1.5	29	} 100±40 (11/2$^-$, 3/2$^+$, 1/2$^+$)
3/2$^+$	0.00	2.8±0.5	69	
1/2$^+$	0.15	1.6±0.2	27	
7/2$^+$	1.06	1.4±0.2	44	
7/2$^+$	1.16	9.0±3.1	275	
5/2$^+$	1.19	2.9±0.5	69	} <250±60 (5/2$^+$, 5/2$^+$)
5/2$^+$	1.50	2.8±0.2	65	
$\Gamma^\uparrow$ [keV]			578	570±190

[a]Values from Ref. [2].

TABLE V. Calculated and experimental partial neutron escape widths of the ISGMR in ^{208}Pb.

μ^{-1}	Ex [MeV]	s_μ [a]	$\overline{\Gamma}^{\uparrow}_{\mu}$ [keV]	$\Gamma^{\uparrow}_{\mu}$ [b] [keV]	$\Gamma^{\uparrow}_{\mu}$ [c] [keV]
$(1/2)^-$	0.00	1.0	31	incl. $(13/2)^+(19\pm27)$	140 ± 35
$(13/2)^+$	1.63	0.91	4	$75\pm35(73\pm33)$	incl. $(1/2)^-$
$(5/2)^-$	0.57	0.98	180	$<35(65\pm43)$	70 ± 15
$(3/2)^-$	0.89	1.0	57	$75\pm40(133\pm45)$	50 ± 10
$(7/2)^-$	2.34	0.7	135	$<140\pm30(155\pm33)$ incl. $(5/2)^+(7/2)^+$	165 ± 40
$(9/2)^-$	3.41	0.61	3		
$\Gamma^{\uparrow}$ [keV]			410	$325\pm105(445\pm181)$	425 ± 100

[a]Values from Ref. [16].
[b]Values from Ref. [3].
[c]Values from Ref. [17].

TABLE VI. Calculated and experimental parameters of the ISGDR.

	^{90}Zr	^{116}Sn	^{144}Sm	^{208}Pb
$\omega_1-\omega_2$ [MeV]	18–31	16–30	15–27	15–30
$x_{L=1}$ [%]	79	77	75	76
$x^{expt}_{L=1}$ [%] [a]	91 ± 11	89 ± 10	105 ± 12	95 ± 13
$\overline{\omega}_{L=1}$ [MeV]	25.05	24.10	23.08	21.84
ω^{expt} [MeV] [a]	26.3 ± 0.4	24.3 ± 0.3	23.0 ± 0.3	20.3 ± 0.3 22.4 ± 0.5 [b]
Δ^{CRPA} [MeV]	2.7	2.6	2.2	2.3
Δ^{expt} [MeV] [a]	3.2 ± 0.2	3.5 ± 0.2	3.2 ± 0.2	2.5 ± 0.2

[a]Values from Ref. [4].
[b]Values from Ref. [19].

IV. DISCUSSION AND SUMMARY

The calculated mean excitation energies of the ISGMR in ^{90}Zr, ^{124}Sn, and ^{208}Pb were found to be lower than the corresponding experimental energies by about 10% (Table II). This is not surprising since our calculation is not fully self-consistent and no attempt was made to reproduce the experimental results. Nevertheless, such a description seems satisfactory because a rather simple phenomenological mean field and particle-hole interaction are used in the calculations and the above quantities are only partially self-consistent (concerning self-consistent calculations of the ISGMR strength function see Refs. [9,14,22,23,25]). Figure 1 shows that the CRPA and scaling transition densities of the ISGMR, calculated for ^{208}Pb, have different behaviors only in the interior region. Within the CRPA approach a qualitative description of partial nucleon escape widths of the ISGMR in the mentioned nuclei was obtained without the use of any adjustable parameters (see Tables III–V). It is interesting to note that the use of $f^{ex}=-2.275$ from Ref. [10] for the ISGMR in ^{208}Pb leads to $\omega^{s.s.}_{L=1}\simeq2.5$ MeV and to increasing the calculated total neutron escape width by factor $\simeq1.5$.

The calculated mean excitation energy of the ISGDR in

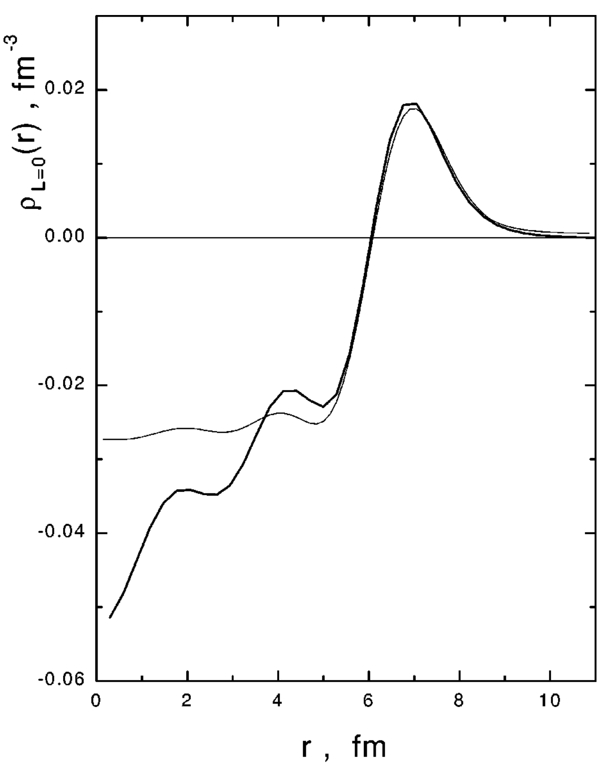

FIG. 1. Calculated transition density of the ISGMR in ^{208}Pb (thick line). The thin line corresponds to the scaling model calculation [18]. Both densities have the same strength R_g.

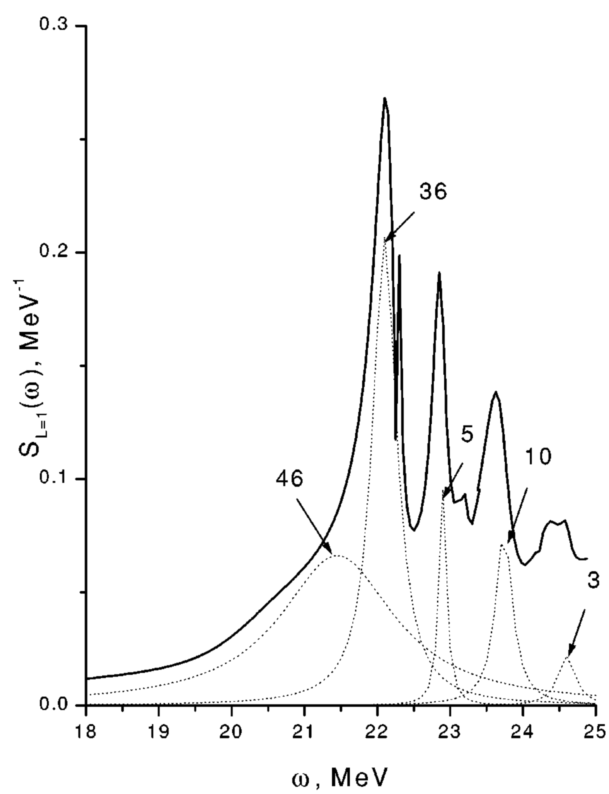

FIG. 2. Calculated CRPA strength function of the ISGDR in ^{208}Pb (solid line) and its expansion in terms of doorway-state resonances (dotted lines). The double-relative strengths of the resonances $y_{L=1,g}$ (in %) are also given.

TABLE VII. Calculated partial branching ratios for the direct nucleon decay of the ISGDR in ^{208}Pb. The branching ratios for population of neutron deep-hole states $(3/2)^+$, $(11/2)^-$, and $(5/2)^+$ are, respectively, 0.31, 0.78, and 0.80 (we used $s_\mu = 0.6$).

		Neutron				Proton	
μ^{-1}	$(5/2)^-$	$(13/2)^+$	$(7/2)^-$	$(9/2)^-$	$(1/2)^+$	$(3/2)^+$	$(5/2)^+$
E_x[MeV]	0.57	1.63	2.34	3.41	0.00	0.35	1.68
s_μ [a]	0.98	0.91	0.70	0.61	0.545	0.57	0.535
b_μ [%]	0.17	4.47	0.39	1.21	0.24	0.60	1.43
$\Sigma_\mu b_\mu$ [%]		8.13				2.27	

[a]Values from Refs. [16,21].

^{90}Zr, ^{124}Sn, and ^{208}Pb was found to be in good agreement with the experimental energy (Table VI). The ISGDR exhausts 75–80 % of the $(\mathrm{EWSR})_{L=1}$ (Table VI). The major part of the rest is exhausted by the ''pygmy'' resonance within the excitation energy interval 8–15 MeV [24]. Note that the calculated ISGDR strengths are somewhat lower than the corresponding experimental results [4] also given in Table VI. The rms energy dispersion Δ^{CRPA} is a characteristic of the Landau damping of the considered GR. The calculated values for Δ^{CRPA} are only slightly smaller than the experimental values (Table VI). The calculated partial branching ratios for the direct nucleon decay of the ISGDR in ^{208}Pb (see Table VII) should be considered as rough estimates in view of the rather poor expansion of the CRPA strength function and reaction amplitudes in terms of isolated Breit-Wigner resonances. Nevertheless, the calculated branching ratios can be used as a guide in experimental searches of the direct nucleon decay of the ISGDR in ^{208}Pb.

In conclusion, we have reported results of the CRPA calculations of some parameters of the isoscalar monopole and dipole giant resonances in several medium-heavy nuclei. In spite of the absence of full self-consistency between the mean field and particle-hole interaction, the calculation results are found to be in at least qualitative agreement with the corresponding values deduced from experimental data. An attempt to determine qualitatively the values of the partial branching ratios for the direct nucleon decay of the ISGDR in ^{208}Pb has also been undertaken.

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