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STATISTICAL DESCRIPTION OF THE RADIATIVE STRENGTH FUNCTIONS

V.A. Plujko, S.N.Ezhov[1], A.S.Mikulyak[1]

[1] *Taras Shevchenko University, Kiev, Ukraine*

A closed-form thermodynamic pole approach is developed for average description of the E1 radiative strength functions using the microcanonical ensemble for initial states. A semiclassical description of the collective excitation damping in the method is based on modern physical notion on the relaxation processes in Fermi systems. It is shown that the model is able to cover a relatively wide energy interval, ranging from zeroth gamma-ray energy to values above GDR peak energy. It gives rather accurate means of simultaneous description of the γ-decay and photoabsorption strength functions in the medium and heavy nuclei. For gamma-ray energies near neutron binding energies the calculations within the proposed model describe experimental data somewhat better for heavy nuclei with A >150 as compared to other closed-form approaches.

1. Introduction

Gamma-emission is one of the most universal channel of the nuclear decay, because it generally may attend any nuclear reaction. This process as well as absorption of the gamma-rays and electron-positron decay are described in the many-body systems by the radiative strength functions [1 - 3]. These functions are important for the study of the nuclear structure models, γ-decay mechanisms, deformation and fluctuation of the nuclear shape, energies and widths of the collective excitations [4 - 6, 8, 9]. Besides this fundamental importance from a theoretical point of view, the strength functions are necessary to generate the data for the energy and non-energy applications. It is critically important to have a simple closed-form expression for the γ-ray strength function because in most cases this function is an auxiliary quantity used in calculations of different nuclear characteristics and processes. The theory-based approaches for γ-strength are preferred over the empirical ones to improve the reliability and accuracy of such calculationsand to understand the physical sense of used parameters.

According to Brink hypothesis [10, 11], the Lorentzian line shape with the energy-independent width (SLO model) is widely used for calculations of the dipole E1 radiative strength. This approach is most appropriate simple method for the description of the photoabsorption data on medium and heavy nuclei [12 – 14]. The situation is more complicated in the case of the gamma-emission. The SLO model strongly underestimates the gamma-ray spectra at low energies $\varepsilon_\gamma \leq 1\,\mathrm{MeV}$ [15, 16]. A global description of the gamma-spectra by the Lorentzian can be obtained rather satisfactorilyin the range $1 \leq \varepsilon_\gamma \leq 8$ MeV but with use of the giant dipole resonance (GDR) parameters which are different from those based on photoabsorption data. On the whole, SLO approach overestimates the integral experimental data (the capture cross sections, the average radiative widths) in heavy nuclei [14 - 21]. The models for description of the E1 strengths at low energies ε_γ were proposed in refs. [22, 23]. An enhanced generalized Lorentzian model (EGLO) was used and analyzed in refs.[24, 25] for a unified description of the low energy and integral data. The EGLO radiative strength function consists of two components (for spherical nuclei): a Lorentzian with the energy and temperature dependent width $\Gamma_k(\varepsilon_\gamma, T)$ and finite value term from [22] corresponding to zero value of γ-ray energy. An empirical expression for width $\Gamma_k(\varepsilon_\gamma, T)$ was used with two additional parameters. The dependence of the parameters on mass number was obtained to fit EGLO calculations to the experimental data. Nowadays the EGLO method is recommended by the IAEA [25] as the best practical model for calculation of the dipole gamma-ray strength function when the experimental data are unavailable.

It should be noted that the SLO and EGLO expressions for the gamma-decay strength function of heated nuclei are in fact the parametrizations of the experimental data, namely:

1. these expressions are not consistent with general relation between strength function and the imaginary part of the response function of the heated nuclei (see [26 - 28] and Sect.2);
2. the EGLO damping width Γ_k has the empirical dependences on γ-energy and temperature T. It is similar to the one of the zero sound damping in the infinite fermi- liquid when the collisional (two-body) dissipation is taken into account only. It is well known that an important contribution to the total width in heavy nuclei is given by the fragmentation (one-body) width which determines a redistribution of the particle-hole excitations near of the collective state [29 - 31]. The latter component of the width is almost independent of the nuclear temperature. The width of the SLO model can be identified with this fragmentation component.

The statistical description of the average γ-decay strength of excited states is presented below using the microcanonical ensemble for initial states. The contributions to the relaxation width resulting from the both interparticle collisions with retardation effects and fragmentation are taken into account in a semiclassical way. The dependences of the γ-decay and photoabsorption strength functions on the initial excitation energy, the gamma-ray energy and mass number are investigated within the thermodynamic pole approximation (TPA method). The TPA calculations are compared with those ones within EGLO, SLO models and with the experimental data. It is shown that the TPA model is able to cover a relatively wide energy interval, ranging from zeroth gamma-ray energy to values above GDR peak energy. It gives rather accurate method of simultaneous description of the γ-decay and photoabsorption strength functions in the medium and heavy nuclei.

2. Gamma-ray strength functions in heated nuclei

We shall consider the radiative strength function averaged over spins of initial states for γ-emission of the electric type with the energy ε_γ and multipolarity λ. The general expression for this function can be obtained from the relation for the average radiative width $d\Gamma_\lambda / d\varepsilon_\gamma$ per unit of the γ-ray energy interval [32]. Within statistical mechanics the width $d\Gamma_\lambda / d\varepsilon_\gamma$ is defined in standard way as the quantity averaged over states with slightly different values of the total initial energy E and numbers of protons Z and neutrons N

$$\frac{d\Gamma_\lambda}{d\varepsilon_\gamma} = \sum_{J,M,\Delta E,\Delta Z,\Delta N,J_f} \frac{P_{if}}{N} = \sum_{J,M,J_f} \frac{\langle P_{if}\rangle}{\Omega(U,Z,N)} \tag{1}$$

where $N = \Omega(U,Z,N)\Delta E \Delta Z \Delta N$ is the total number of initial states; $\Omega(U,Z,N)$ is the density of states; $U=E-E_0$ is the initial excitation energy; E_0 is the ground state energy. The quantities J and M are the spin of initial states and its projection on the Z axis, respectively; ΔE, ΔZ, ΔN are the small-scale intervals of the dispersion in values of the energy, numbers of protons and neutrons near the average values E, Z, N.

The quantity $P_{if}(\varepsilon_\gamma) = d_\lambda(\varepsilon_\gamma) \cdot B_{if}^{(\lambda)} \cdot \delta(E - E_f - \varepsilon_\gamma)$ is the γ-transition probability with the energy ε_γ from an initial state i to the final state f; $d_\lambda(\varepsilon_\gamma) = (\varepsilon_\gamma / \hbar c)^{2\lambda+1} 8\pi(\lambda+1) / \lambda(2\lambda+1)^2$, and

$B_{if}^{\lambda} = \sum_{M_f,\mu} \left|\langle J_f M_f E_f | Q_{\lambda\mu} | J_i M_i E_i\rangle\right|^2$ is the reduced transition probability with the multipole operator $Q_{\lambda\mu}$ for $E\lambda$ radiation, $Q_{\lambda\mu} = \sum_k e_k(\lambda) r_k^{\lambda} Y_{\lambda\mu}(r_k)$, The quantities $e_n(\lambda) = e(-1)^k Z/A^{\lambda}$ and $e_p(\lambda) = e\left((A-1)^{\lambda} + (-1)^{\lambda}(Z-1)\right)/A^{\lambda}$ are the effective kinematic charge of the neutrons and protons in nucleus, respectively.

Eq.1 can be represented in the following form

$$d\Gamma_{\lambda}/d\varepsilon_{\gamma} = \langle D(\varepsilon_{\gamma})\rangle / \Omega(U,Z,N) \tag{2}$$

Here, $\Omega(U,Z,N) = \sum_J \Omega(U,Z,N,J)$ is the total density of the initial states; the symbol <...>denotes an average over the energies and numbers of the protons and neutrons with the unit weight functions in the intervals ΔE, ΔZ and ΔN, respectively

$$D(\varepsilon_{\gamma}) = d_{\lambda}(\varepsilon_{\gamma}) \cdot \sum_{N^1, Z^1, J,, M, J_f, M_f, \nu} \delta(E - E_{\nu}) \cdot \delta(N - N^1) \cdot \delta(Z - Z^1) \cdot \times \left|\langle J_f M_f E_f | Q_{\lambda} | JME\rangle\right|^2 \delta\left(E - E_{\nu} - \varepsilon_{\gamma} - \gamma_1 (N - N^1) - \gamma_1 (Z - Z^1)\right) \tag{3}$$

where $Q_{\lambda} = \sum_{\mu} Q_{\lambda\mu}$. The identical changing the arguments is made in the δ-function depending on energy. The additional constants γ_i, defined below fix the numbers of the protons and neutrons.

In the region of high excitation energies being discussed the density of states in the intervals ΔE, ΔZ, ΔN is almost constant. Therefore, we can assume that the D varies a little and $\langle D\rangle = D$. In this case the quantity $\Gamma_k(\varepsilon_{\gamma}, T) = D/\Omega$ coincides with the width of the γ- decay of states of the microcanonical ensemble with the given constants of motion E, Z and N. Using the integral representation of the δ-functions

$$\delta(x) = \frac{1}{2\pi} \int_{-\infty}^{+\infty} e^{itx} dt \tag{4}$$

and completeness relation for wave functions one obtains

$$D(\varepsilon_{\gamma}) = \frac{1}{(2\pi)^3} \int_{-\infty}^{+\infty} d\alpha_1 \cdot d\alpha_2 \cdot d\alpha_3 \cdot \exp(\alpha_3 U - \alpha_1 N - \alpha_2 Z) \cdot Z(\{\alpha_j\}) \cdot \bar{\Gamma}(\{\alpha_j\}, \varepsilon_{\gamma}). \tag{5}$$

Here, $Z(\{\alpha_j\}) = Sp(\exp(-\beta \mathit{H}))$ is the partition function of the grand canonical ensemble characterized by three constants α_1, α_2, $\beta = \alpha_3$; $\mathit{H} = H - \gamma_1 \hat{N} - \gamma_2 \hat{Z}$, $\gamma_j = \alpha_j/\beta$, H is the nuclear

Hamiltonian and $\hat{N}$, $\hat{Z}$ are the operators of the neutron and proton numbers. The quantity $\bar{\Gamma}\left(\{\alpha_j\},\varepsilon_\gamma\right)$ is the mean width per unit energy of the γ-decay of the states of the canonical ensemble with the parameters $\{\alpha_j\}$

$$\bar{\Gamma}\left(\{\alpha_j\},\varepsilon_\gamma\right)=\frac{d_\lambda\left(\varepsilon_\gamma\right)}{\pi}I_{Q_\lambda^*,Q_\lambda}\left(\{\alpha_j\},\omega\right),\ \omega=\varepsilon_\gamma/\hbar \tag{6}$$

where

$$I_{Q_\lambda^*,Q_\lambda}\left(\{\alpha_j\},\omega\right)=\frac{1}{2\pi}\int_{-\infty}^{+\infty}dt\cdot Sp\left(\rho(H)\,Q_\lambda(0)\,Q_\lambda^*(t)\right)\exp(i\omega t) \tag{7}$$

is the spectral intensity for the expectation value of the product $Q_\lambda(0)Q_\lambda^*(t)$ the multipole operator $Q_\lambda(t)$ in the Heisenberg representation $Q_\lambda(t)=\exp(itH/\hbar)Q_\lambda\exp(-itH/\hbar)$, $Q_\lambda=Q_\lambda(0)$ i.e. the time-depended correlation function for the operator $Q_\lambda(t)$. The canonical average $Sp(\rho(H)\ldots)$ is taken over the Gibbs ensemble with the density matrix $\rho(H)=\exp(-\beta H)/Z(\{\alpha_j\})$.

Taking into account the fluctuation- dissipative relation between the spectral density of nonsymmetrized correlation function and response function [33, 34], Eq.(6) can be written as

$$\bar{\Gamma}\left(\{\alpha_j\},\varepsilon_\gamma\right)=s_\lambda\left(\varepsilon_\gamma,\{\alpha_j\}\right)\exp\left(-\beta\,\varepsilon_\gamma\right) \tag{8}$$

where

$$s_\lambda\left(\varepsilon_\gamma,\{\alpha_j\}\right)=\frac{d_\lambda\left(\varepsilon_\gamma\right)}{\pi}I_{Q_\lambda^*,Q_\lambda}\left(\{\alpha_j\},\omega\right)\exp(\beta\hbar\omega)=-\frac{d_\lambda\left(\varepsilon_\gamma\right)}{\pi}\left(1-\exp\left(-\beta\varepsilon_\gamma\right)\right)^{-1}\operatorname{Im}\chi_\lambda\left(\omega,\{\alpha_j\}\right) \tag{9}$$

The quantity χ_λ is the linear response function given by

$$\chi_\lambda\left(\omega,\{\alpha_j\}\right)=Z\,e_p(\lambda)\,\chi_\lambda^{(p)}\left(\omega,\{\alpha_j\}\right)+N\,e_n(\lambda)\,\chi_\lambda^{(n)}\left(\omega,\{\alpha_j\}\right),$$
$$\chi_\lambda^k\left(\omega,\{\alpha_j\}\right)=Sp\left(r^\lambda\sum_\mu Y_{\lambda\mu}^*(r)\,\delta n_k\right)\Big/q_\omega(t) \tag{10}$$

where $\delta n_k(t)$ is the change of the single-particle density matrix n_k induced by the external field $V_{ext}^k=q_\omega(t)\,e_k(\lambda)\,r^\lambda\sum_\mu Y_{\lambda\mu}(r)$, $q_\varpi(t)=q_0\exp(-i(\omega+i\eta)t)$, $\eta\to+0$, $q_0\ll 1$ for protons (k=p) and neutrons (k=n).

The integral in Eq.(5) can be evaluated by means of the saddle- point method. The parameters α_j of the saddle point are found from the condition of an extremum of the logarithm of the integrand. They are the solutions of the equations

$$\partial S_\gamma\left(\{\alpha_j\}\right)/\partial\alpha_k = 0,\ k = 1 \div 3 \tag{11}$$

with

$$S_\gamma\left(\{\alpha_j\}\right) = \ln Z\left(\{\alpha_j\}\right) - \alpha_1 N - \alpha_2 Z + \alpha_3 U + \ln\bar{\Gamma}\left(\{\alpha_j\}, \varepsilon_\gamma\right). \tag{12}$$

Next it is assumed that the dependence of the function s_λ on the saddle point parameters $\{\alpha_j\}$ is more smooth as that of the partition function $Z\left(\{\alpha_j\}\right)$, namely,

$$\left|\partial^n \ln s_\lambda / \partial^n \alpha_j\right| << \left|\partial^n \ln Z / \partial^n \alpha_j\right|,\ n = 1,2. \tag{13}$$

This assumption is in agreement with many investigations of the linear response function properties in medium and heavy heated nuclei (see, for example кefs. [27, 28, 35]. One finally obtains[1] for the average radiative width $\Gamma_\lambda\left(\varepsilon_\gamma\right)$:

$$\frac{d\Gamma_\lambda}{d\varepsilon_\gamma} = \left(\frac{\varepsilon_\gamma}{\pi\hbar c}\right)^2 o_\lambda\left(\varepsilon_\gamma, T_f\right)\frac{\Omega\left(U - \varepsilon_\gamma, Z, N\right)}{\Omega(U, Z, N)} \tag{14}$$

where

$$\sigma_\lambda\left(\varepsilon_\gamma, T_f\right) = -\left(\frac{\pi\hbar c}{\varepsilon_\gamma}\right)^2 \frac{d_\lambda\left(\varepsilon_\gamma\right)}{\pi} \frac{\operatorname{Im}\chi_\lambda\left(\omega, T_f\right)}{1 - \exp\left(-\varepsilon_\gamma/T_f\right)} \tag{15}$$

Here, for simplicity, the designations a dependence of the functions s and χ on the parameters α_1 and α_2 is not indicated in Eq.(14) and below. The quantity $\Omega\left(U_f, Z, N\right)$ is the total density of the final states with the energy $U_f = U - \varepsilon_\gamma$ and the temperature $T_f = \beta$,

$$\Omega\left(U_f, Z, N\right) = \exp\left(S_f\left(\{\alpha_j\}\right)\right) / 2\pi^2\left|\det\left(a_{kl}\right)\right|^{1/2} \tag{16}$$

where $\det\left(a_{kl}\right)$ is the determinant of the matrix with the elements

$$a_{kl} = \partial^2 S_f / \partial\alpha_k \partial\alpha_l\ \ k, l = 1 \div 3 \tag{17}$$

The parameters $\alpha_1, \alpha_2, \alpha_3 = 1/\beta = 1/T_f$ are the solutions of equations of the thermodynamical state in final nucleus, i.e. they are the solutions of the Eq. [11] with the entropy of final states

$$S_f\left(\{\alpha_j\}\right) = S_\gamma\left(\{\alpha_j\}\right) - \ln\Gamma\left(\{\alpha_j\}, \varepsilon_\gamma\right) = \ln\Gamma\left(\{\alpha_j\}\right) - \alpha_1 N - \alpha_2 Z + \alpha_3 U$$

instead of the function S_γ.

[1] Note that the expression for the average gamma-width, with use of the microcanonical distribution and an assumption on a weak dependence of the width on ε_γ, was first considered in [36]. In this case the average radiative width is equal to $d\Gamma_\lambda / d\varepsilon_\gamma = -d_\lambda\left(\varepsilon_\gamma\right) / \left(\pi \operatorname{Im}\chi_\lambda(\omega, T)\right)$

The relation (14) is the same one as those given by detailed balance principle [4] with the photoabsorption cross-section $\sigma_\lambda(\varepsilon_\gamma, T_f)$ given by Eq.(15). It follows from Eq.(14) and (7), (9) that the rate of the γ-transitions between excited states depends mainly on thermal fluctuations of the multipole moments in the final states.

General form of the $\sigma_\lambda(\varepsilon_\gamma, T)$ coincides with that one from refs. [26 - 28],[37, 38] obtained by making use of the averaging over the canonical ensemble with a constant temperature T.

The emission and absorption processes for γ-rays are generally connected with different radiative strengths [1, 14]. The gamma- decay (downward) strength function $\overleftarrow{f}_{E\lambda}$ determines the γ-emission of heated nuclei. It is associated with the average radiative widths $\Gamma_\lambda(\varepsilon_\gamma)$. The photoexcitation (upward) strength function $\overrightarrow{f}_{E\lambda}$ is connected with photoabsorption cross-section at fixed temperature T. For the dipole transitions these functions have the form

$$\overleftarrow{f}_{E1}(\varepsilon_\gamma, T) = \frac{d\Gamma_{\lambda=1}/d\varepsilon_\gamma}{3\varepsilon_\gamma^3} \frac{\Omega(U,Z,N)}{\Omega(U-\varepsilon_\gamma,Z,N)} = F(\varepsilon_\gamma, T_f) \tag{18}$$

and

$$\overrightarrow{f}_{E1}(\varepsilon_\gamma, T) = \frac{\sigma_{\lambda=1}(\varepsilon_\gamma, T)}{3\varepsilon_\gamma(\pi\hbar c)^2} = F(\varepsilon_\gamma, T). \tag{19}$$

Here, spectral function $F(\varepsilon_\gamma, T)$ is given by

$$F(\varepsilon_\gamma, T) = \frac{s_{\lambda=1}(T,\varepsilon_\gamma)}{3\varepsilon_\gamma^3} = -2e^2\frac{NZ}{A}\left(\frac{2}{3\hbar c}\right)^3 L(\varepsilon_\gamma, T)\,\mathrm{Im}\,\chi^{(-)}(\omega, T), \tag{20}$$

$$L(\varepsilon_\gamma, T) = \frac{1}{1-\exp(-\varepsilon_\gamma/T)} \tag{21}$$

where $\chi^{(-)}(\omega,T) = \mathrm{Sp}\left(r\sum_\mu Y_{1\mu}^*(r)\delta n^{(-)}\right)\Big/q_\omega(t)$ is the response function of the heated nuclei to the field $q_\omega(t) r\sum_\mu Y_{1\mu}(r)$ and $\delta n^{(-)} = \delta n_p(t) - \delta n_n(t)$ the variation of the isovector single-particle density matrix.In the case of the spherical nuclei, we have $\chi^{(-)}(\omega,T) = 3\mathrm{Sp}\left(r\,Y_{10}(r)\delta n^{(-)}\right)/q_\omega(t)$ with the isovector density perturbation $\delta n^{(-)}$ induced by the dipole field $q_\omega(t) r\,Y_{10}(r)$.

Note that the γ-decay strength function depends on temperature T_f of the final states. This temperature is a function of the γ-ray energy in contrast to the initial states temperature T.

In the case of cold nuclei the radiative strength functions is also connected with the response function by Eqs.(18) - (20) but with factor $L = 1$. The scaling factor $L(\varepsilon_\gamma, T)$ (21), defines the enhancement of magnitude of the radiative strength functions in heated nuclei with temperature T as compared to the cold nuclei. This factor can be interpreted as average number of the 1p-1h excited states in heated system placed in an external field with energy $\hbar\omega$,

$$L(\varepsilon_\gamma, T) = N_{1p-1h} = \frac{1}{\hbar\omega}\int_{-\infty}^{+\infty}\int_{-\infty}^{+\infty} d\varepsilon_1 d\varepsilon_2 n(\varepsilon_1)(1 - n(\varepsilon_2))\delta(\varepsilon_1 - \varepsilon_2 + \hbar\omega) \quad (22)$$

where $n(\varepsilon) = 1/(1 - \exp((\varepsilon - \mu)/T))$ is the Fermi distribution function for occupation of the single-particle states.

In order to get a simple closed-form expression for the response function, one assums as usual that gamma-decay is determined by the collective motion mode which excited in the associated photoabsorption process. Due to this the E1 transitions are considered as corresponding to the giant dipole excitations. Next the hydrodynamic model with damping ([39]) is applied for description of the collective motion of the neutrons against the protons which corresponds to the GDR in the classical picture. This approach is an extension of the Steinwedel- Jensen (SJ) model and provides a simple description of the GDR excitation simultaneously with its damping.In common with the other classical hydrodynamics models SJ model with dampingcorresponds to the semiclassical description of the Fermi systems by means of the Landau-Vlasov kinetic equation with truncation of the Fermi sphere distortion by the layers of monopole and dipole multipolarities only [40]. Note that the SJ model describes volume oscillations of the transition density and these oscillations are almost unaffected by the dynamical distortion of the Fermi surface with multipolarities $l > 1$ [41]. The SJ mode plays most important role in heavy nuclei [42].

We make use of the expression for the induced dipole moment within the extended SJ model from $14.4 of Ref.[39] and combine it with the relation for classical absorption cross-section and Eqs.(15), {20). Then we get the spectral function

$$F(\varepsilon_\gamma, T) = 8.674 \cdot 10^{-7}\frac{NZ}{A}\alpha_0\beta_0^2\frac{Y(z)}{1-\exp(-\varepsilon_\gamma/T)}, \text{ MeV}^{-3}, \quad (23)$$

where

$$Y(z) = \mathrm{Im}\left(\frac{j_2(z)}{z\, j_1(z) - z^2 j_2(z)}\right) = \mathrm{Im}\left(\frac{1}{z^2}\left(\frac{\tan(z) - z}{\varphi(z)}\right)\right) = \frac{A}{NZ\beta_0^2}\sum_{n\geq 1} f_n \frac{\varepsilon_\gamma \Gamma}{(\varepsilon_\gamma^2 - \varepsilon_n^2)^2 + (\Gamma\varepsilon_\gamma)^2}. \quad (24)$$

Here, $\varepsilon_n = z_n/\beta$ and $f_n = (NZ/A)^2 \cdot 2/(z^2 - 2)$ are the energy and classical oscillator strength of the n-resonance, respectively; z_n are solutions of the equation $\varphi(z) = (z^2 - 2)^2\tan(z) + 2z = 0$;

$$z = z(\varepsilon_\gamma, T) = \pm\beta_0\sqrt{\frac{\varepsilon_\gamma}{2\tilde{\varepsilon}_\gamma}} \cdot (\tilde{\varepsilon}_\gamma + i\Gamma), \quad \tilde{\varepsilon}_\gamma = \varepsilon_\gamma + \sqrt{\varepsilon_\gamma^2 + \Gamma^2} \quad (25)$$

and $\alpha_0 = 4\pi e^2\hbar/mc = 0.305$, $\beta_0 = R_0/\hbar u$, $u = (4 b_{vol}/m)(ZN/A^2)$ is the isovector sound velocity with the volume symmetry energy coefficient b_{vol} entering the semi-empirical mass formula; $j_n(z)$ are the spherical Bessel functions.

The quantity Γ in Eqs.(24) and (25) is damping width of the isovector velocity $\mathbf{v} = \mathbf{v}_p - \mathbf{v}_n$, where $\mathbf{v}_p$, $\mathbf{v}_n$ are velocities of the proton and neutron fluids, respectively. It determines the reduced friction force $\mathbf{v}\Gamma$ in the equation for isovector velocity, Eq. (14. 60) from ref. [39]. It can be seen that this term corresponds to the expression $v\Gamma = 2\hbar\int d\mathbf{p}(\mathbf{p}/m)J(\mathbf{p},\mathbf{r},t)/(2\pi\hbar)^3$ in semiclassical picture given by Landau-Vlasov equation with a source term $J(\mathbf{p},\mathbf{r},t)$ for relaxation processes [43,44]. The source term is taken as the sum of two components [43, 44], $J = J_c + J_s$: the first one, J_c, is the isovector collision integral with retardation (memory) effects and the second one, J_c, is connected with fragmantation contribution to the damping width (isovector one-body relaxation). The latter component is described within the framework of the relaxation time approximation with the relaxation time τ_s different from infinity only in the distorted layer of the Fermi surface with dipole multipolarity. The expression for damping width is obtained by the use of relation (26) and the expression for isovector collision integral from [44]. The width is the energy-dependent one and has the following form[2]

$$\Gamma = \Gamma(\varepsilon_\gamma, T) = \Gamma_c(\varepsilon_\gamma, T) + \Gamma_s, \quad \Gamma_c(\varepsilon_\gamma, T) = (m^*/m)k_s(\varepsilon_\gamma)\Gamma_\omega \tag{27}$$

$$\Gamma_s = (m^*/m)\hbar/\tau_s = (m^*/m)k_s(\varepsilon_\gamma)\Gamma_\omega, \quad \Gamma_\omega = \hbar\,\bar{\upsilon}/R \tag{28}$$

Here, Γ_s and Γ_c are the collisional and one-body contributions to the total width, respectively. The quantity m^* is an effective mass of nucleon; we will use m^*=m. The quantities R and $\bar{\upsilon} = 3\upsilon_F/4$ are the nuclear radius and average velocity of the nucleon, respectively; v_F is the Fermi velocity.

The quantity $\tau(\varepsilon_\gamma,T)$ is the collisional relaxation time for the isovector dipole distortion of the Fermi surface. It is associated with two-body collisions in the heated nucleus which is placed in the electric field with the frequency $\omega = \varepsilon_\gamma/\hbar$. For the isotropic collision probabilities it is given by [43, 44]

$$\frac{\hbar}{\tau_c(\varepsilon_\gamma, T)} = \frac{T^2}{\alpha}\left(1 + (\varepsilon_\gamma/2\pi T)^2\right) \tag{29}$$

The dependence of the relaxation time τ_c on the energy ε_γ results from memory effects in the collision integral and follows Landau's prescription [45,47-49]. The temperature dependence arises from the smeared out behavior of the equilibrium distribution function near the Fermi momentum in the heated nuclei.

The quantity α in (29) is defined by the magnitude of the in-medium neutron-proton cross section σ(np) near the Fermi surface

$$\alpha = \frac{\text{const}}{\sigma(np)}, \qquad \text{const} = \frac{9\hbar^2}{16m^*} = 23.514, \ \alpha \text{ in MeV}, \ \sigma\text{ (np) in fm}^2. \tag{30}$$

The magnitude of the in-media cross section σ(np) is usually taken proportional to a value of the free space cross section σ_{free}(np) with a factor F,

$$\sigma(np) = F \cdot \sigma_{free}(np) \tag{31}$$

then the parameter α can be rewritten in the form

[2]Here we do not use a normalization of the damping width to the value which corresponds to the zero-sound magnitude for the collisional width in infinite matter with arbitrary multipolarities of the Fermi sphere distortion [43, 45, 46].

$$\alpha = a_{free}/F = 4.7/F,\ \alpha_{free} = 23.514/\sigma_{free}(np) \tag{32}$$

when the value $\sigma_{free}(np) = 5\ fm^2$ is adopted in line with refs.[43, 48, 50, 51].

The quantities α and F determine two-body contribution B_c to the damping width of the GDR at zero temperature and the thermal relaxation time $\tau(T)$ for nuclear viscosity [52] in heated nuclei:

$$B_c = \Gamma_c(\varepsilon_\gamma = E_r)/\Gamma_r = q/\alpha = F\cdot q/\alpha_{free}, \qquad q = E_r^2/4\pi\Gamma_r$$

$$\tau(T)/\hbar = ab\,T^2, \qquad b = (5/6)\sigma(np)/\bar{\sigma} \approx 0.833\,\sigma_{free}(np)/\sigma_{free} \approx 1.11, \tag{33}$$

where the E_r and Γ_r=$\Gamma(\varepsilon_\gamma = E_r, T=0)$ are respectively the GDR energy and width at zero temperature; $\bar{\sigma} = (\sigma(pp)+\sigma(n)+2\sigma(np))/4$ is the in-media spin-isospin averaged nucleon-nucleon cross section near the Fermi surface; $\bar{\sigma}_{free} = 3.75\,fm^2$ free averaged cross section.

The values of the B_c should be located in the interval from 0 to 1. This condition determines the limiting values of α and F,

$$\alpha \geq \alpha_{free} = q,\quad F \leq F_{max} = \alpha_{freee}/q = \alpha_{free}/\alpha_{min} \tag{34}$$

The isovector one-body relaxation width Γ_s in (28) is taken to be similar to the wall formula expression Γ_w [30, 31, 41] but scaled with an energy-dependent coefficient $k_s(\varepsilon_\gamma)$ [43, 44, 53, 54]. The quantal calculations within framework of a simplified RPA [53] show significant reduction of the one-body width in comparison with the wall value,in particular $k_s \approx 0.1$ in the range where collective phonon energy exceeds the nuclear binding energy and $k_s \approx 0.7$ if collective phonon energy is negligibly small. The value k_s=0.62 was adopted in ref.[54].
The spectral function *F* given by Eqs. (23), (24) can be written in more convinient form

$$F(\varepsilon_\gamma, T) = 8.674\cdot 10^{-8}\frac{\sigma_r\Gamma_r}{1-\exp(-\varepsilon_\gamma/T)}\sum_{n=1}^{K}\omega_n\frac{\varepsilon_\gamma\Gamma(\varepsilon_\gamma, T)}{(\varepsilon_\gamma^2-\varepsilon_n^2)^2+(\varepsilon_\gamma\Gamma(\varepsilon_\gamma, T))^2} \tag{35}$$

Here, $w_n=f_n/f_1=(z_1^2-2)/(z_n^2-2)$; $K\to\infty$ and σ_r is the peak value of the photoabsorption cross-section

$$\sigma_r = 10.0\,\alpha_0 f_{n=1}/\Gamma_r = 8.4(NZ/A)\alpha_0/\Gamma_r = 0.5\pi\sigma_{TRK} \tag{36}$$

where $\sigma_{TRK} \approx 60NZ/A$ is the classical sum rule in Mev mb.

The first term in expansion (35) corresponds to the excitation of the GDR with the energy $E_r = \varepsilon_{n=1}$ and therefore the quantities β_0, $\varepsilon_{n>1}$ can be defined in the term of the GDR energy as

$$\beta_0 = z_{n=1}/E_r,\ \varepsilon_n = x_n\cdot E_r, \qquad x_n = z_n/z_1, \qquad z_1 = 2.08 \tag{37}$$

The values of the parameters w_n and x_n at $n \leq 4$ are following: $w_1 = 1$, $w_2 = 0.070$, $w_3 = 0.028$, $w_4 = 0.015$ and $x_1 = 1$, $x_2 = 2.86$, $x_3 = 4.42$, $x_4 = 5.97$.

The imaginary part of the dipole response function $\chi_{\lambda=1}$ associated with the Eq.(35) exhibits the resonance behaviour, in which the individual resonances have a Lorentzian shape with energy-dependent width. In the cold nuclei the first term of the expression (35) for the $\mathrm{Im}\chi_{\lambda=1}$ was obtained within the random-phase approximation [55]. This first term of (35) is also in close agreement with the imaginary part of the response function of the heated Fermi-liquid drop on an external pressure, when approximation of the dissipative nuclear fluid dynamics is used for description of the system [46].

The γ-decay and photoexcitation dipole strength functions, (18),(19), have the same temperature-dependent limiting value for vanishing gamma-ray energy and it is equal to

$$\overleftarrow{f}_{E1}(\varepsilon_\gamma = 0, T_f = T) = \overrightarrow{f}_{E1}(\varepsilon_\gamma = 0, T = 0) = 8.674\cdot 10^{-8} q_{K\to\infty}\sigma_r\Gamma_r T\Gamma(\varepsilon_\gamma = 0, T)/E_r^4, \quad (38)$$

where $q_K = \sum_{n=1}^{K} q_n/x_n^4$ and $q_{K\to\infty} = (4/175)z_1^4/(z_1^2 - 2)$. All value of the sum $q_{K\to\infty}$=.008656 z_1=2.0815 is practically contained in the first term $q_{K=1} = 1$.

3.Testing of the closed-form models for E1 strength

Here, the calculations of the E1 radiative strength functions are compared within the framework of the SLO, EGLO models and the approach described in Sect.2. For not very high excitation energies the main contribution to the spectral function *F* results from the first term of the Eq.(35). Therefore, in what follows we use this approximation, and the expression (35) with K = 1 is referred to as the thermodynamic pole approximation [56, 57] (TPA model), $F = F_{TPA}$,

$$F_{TPA}(\varepsilon_\gamma, T) = 8.674\cdot 10^{-8}\frac{\sigma_r\Gamma_r}{1-\exp(-\varepsilon_\gamma/T)}\sum_{n=1}^{K}\omega_n\frac{\varepsilon_\gamma\Gamma(\varepsilon_\gamma, T)}{(\varepsilon_\gamma^2-\varepsilon_n^2)^2 + (\varepsilon_\gamma\Gamma(\varepsilon_\gamma, T))^2} \quad (39)$$

The SLO spectral function, $F=F_{SLO}$, has the Lorentzian form but with the energy independent width Γ_r rather than $\Gamma(\varepsilon_\gamma,T)$,

$$F_{SLO}(\varepsilon_\gamma, T) = 8.674\cdot 10^{-8}\sigma_r\Gamma_r\frac{\varepsilon_\gamma\Gamma_r}{(\varepsilon_\gamma^2-E_r^2)^2 + (\varepsilon_\gamma\Gamma_r)^2} \quad (40)$$

The EGLO dipole spectral function is given by [24,25], $F=F_{EGLO}$,

$$F_{EGLO} = 8.674\cdot 10^{-8}\sigma_r\Gamma_r\left[\frac{\varepsilon_\gamma\Gamma_k(\varepsilon_\gamma, T)}{(\varepsilon_\gamma^2-E_r^2)^2 + (\varepsilon_\gamma\Gamma_k(\varepsilon_\gamma, T))^2} + 0.7\frac{\Gamma_k(\varepsilon_\gamma = 0, T)}{E_r^3}\right] \quad (41)$$

where the energy-dependent width $\Gamma_k(\varepsilon_\gamma,T)$

$$\Gamma_k(\varepsilon_\gamma, T) = \left[k_0(A) + (1 - k_0(A))\frac{\varepsilon_\gamma - \varepsilon_0}{E_r - \varepsilon_0}\right]\cdot(\varepsilon_\gamma^2 + 4\pi^2T^2)\frac{\Gamma_r}{E_r^2} \quad (42)$$

where $k_0(A)$ is the empirical factor; ε_0 = 4.5 MeV. For the case $k_0(A)$ = 1, the quantity Γ_k corresponds to the expression for the collisional damping width in the infinite matter.

The values $k_0(A)$ are mainly obtained by fitting of the low energy experimental data and they depend on the models used for calculations of the temperature and level density. The two expressions for quantity $k_0(A)$ are recommended in [25]

$$k_0(A)=\begin{cases}1, A<148\\ 1+0.09(A-148)^2\exp(-0.18\cdot(A-148)), A\geq 148\end{cases} \tag{43}$$

when the backshifted Fermi gas model [58] (BSFG) is taken for level densities, and

$$k_0(A)=\begin{cases}1.5, A<145\\ 1+0.131(A-145)^2\exp(-0.154\cdot(A-145)), A\geq 145\end{cases} \tag{44}$$

for level density from [59].

Below the backshifted Fermi gas model [58] is adopted and the Eq.(43) for $k_0(A)$ is used. The equations for the temperatures T, T_f have the following form

$$T=\frac{1+\sqrt{1+4a(U-\Delta)}}{2a}, T_f=\frac{1+\sqrt{1+4a(U-\varepsilon_\gamma-\Delta)}}{2a}=\frac{1+\sqrt{1+4a(aT^2-T)-4a\,\varepsilon_\gamma}}{2a} \tag{45}$$

where Δ the energy shift parameter and a the level density parameter. The values of the a$ and Δ are taken from the data file *beijing_bs1.dat* at the rigid-body value for moment inertia (see, the RIPL Handbook [25]}, Ch.5), and from the global fitting in [60], namely,

$$a=0.21A^{0.87}, MeV^{-1}, \Delta=-6.6A^{-0.32}+12\chi A^{-0.5}, MeV, \tag{46}$$

when data in ref.[25] are absent. Here, χ= 0, 1 and 2 for odd-odd, odd-even(even-odd) and even-even nuclides, respectively.

For very small temperature T and with negative values of Δ the BSFG model can lead to negative values of the initial excitation energy $U=aT^2-T+\Delta$ In this case usual Fermi- gas model is used for calculation of the energy, $U=a\,T^2$.

The values of the GDR energy, E_r, and width, Γ_r, and the peak of the E1 absorption cross-section, σ_r, are considered as the temperature- independent and taken from photonuclear data file *beijing_gdr.dat* [25] (when the data exist) or from the global systematics at zero temperature. In the last case they equal

$$E_r=31.2\,A^{-1/3}+20.6\,A^{-1/6},\ \Gamma_r=0.026\,E_r^{1.91},\ o_r=1.2\cdot 120\,NZ/(A\pi\Gamma_r) \tag{47}$$

for spherical nuclei and

$$\begin{aligned}&E_{r,1}=E_r/(1+2\beta/3), && E_{r,2}=E_r/(1-\beta/3),\\ &\Gamma_{r,1}=0.026\,E_{r,1}^{1.91}, && \Gamma_{r,2}=0.026\,E_{r,2}^{1.91}\\ &\sigma_{r,1}=\sigma_r/3, && \sigma_{r,2}=2\sigma_r/3\end{aligned} \tag{48}$$

for deformed nuclei, where β is the quadrupole deformation parameter. All deformed nuclei are considered as the axially symmetric spheroids with the effective quadrupole deformation parameters β. For every nucleus the quantity β is founded as an effective quadrupole deformation parameter which gives the same value of quadrupole moment (Q) as it is in the case when the general expression for the Q ([61], Eq.6.53) is used allowing for the deformations of multipolarities L=2, 4, 6 with parameters β_2, β_4 and β_6, respectively. The values of the last parameters were taken from *moller.dat* file for nuclear ground-state masses and deformations from RIPL\cite [25], Ch.1, (see also [62]).

For axially deformed nuclei the E1 strengths (18) and (19) are the sum of two spectral function *F* with the parameters $E_{r,1}$, $\Gamma_{r,1}$, $\sigma_{r,1}$ and $E_{r,1}$, $\Gamma_{r,1}$, $\sigma_{r,1}$, respectively.

The two approaches are used for the scaling coefficient $k_s(\varepsilon_\gamma)$ in Eq.(28) for one-body isovector width Γ_s:

1.the energy-independent value

$$k_s(\varepsilon_\gamma)=k_s(\varepsilon_\gamma=E_r)=(\Gamma_r-\Gamma_c(\varepsilon_\gamma=E_r,T=0))/\Gamma_\omega \qquad (49)$$

i.e., the quantity $k_s(E_r)=k_s(\varepsilon_\gamma=E_r)$ is obtained from fitting of the the GDR width at zero temperature and it defines the one-body contribution near the GDR resonance energy;

2.the energy-dependent value in a power approximation of the form

$$k_s(\varepsilon_\gamma)=\begin{cases} k_s(E_r)+(k_s(0)-k_s(E_r))\left|(\varepsilon_\gamma-E_r)/E_r\right|^n, \varepsilon_\gamma\le 2E_r \\ k_s(0), \varepsilon_\gamma\ge 2E_r \end{cases} \qquad (50)$$

where the quantity $k_s(0)= k_s(\varepsilon_\gamma=0)$ defines the fragmentation width at zero energy.

The quantity B_c, (33) appears as more convenient than F in calculations of the α,

$$\alpha=q/B_c,\ q=0.02533\,E_r^2/\Gamma_r \qquad (51)$$

determining two-body component of the width by Eqs.(27), (29). The reason is that maximal value of the F is changed from nucleus to nucleus as opposite to B_c, see Eqs.(33) and (34).

It should be noted that the damping component Γ_c given by the Eqs.(27), (29) with α from (32), (33), is two-body relaxation width only in some specific medium where multipolarity of the Fermi sphere distortion does not exceed dipole multipolarity and two-body collisions are isotropic ones with energy-independent cross-sections $\sigma = \sigma_{free}$. Therefore the parameters B_c, (33), and F, (31), are respectively two-body contribution to width and proportionality factor of the in-medium cross-section to free-space magnitude in the specific medium and their values account for the difference between real system and specific medium too.

The dipole γ-decay strength functions $\overleftarrow{f}_{E1}$ considered as a function of mass number are shown in fig.1. The experimental data taken from *kopecky.dat* file of the RIPL-handbook ([25]). Calculations were performed for nuclei from this data file (50 nuclei corresponding to (n,γ) reaction) and at excitation energies and gamma-ray energies are equal to mean energy of E1 transitions. These energies are rather close to the corresponding neutron binding energies. The different lines connect the values calculated within framework of given model. Hereafter the values n = 0.5, $k_s(0)$ = 0.3, B_c = 0.35 and n = 3, $k_s(0)$ = 0.7, B_c = 0.7 are used in TPA calculations. The values of index n are only indicated in the figure for short.

As it can be seen from this figure, for gamma-ray energies near neutron binding energies the calculations within the TPA model describe experimental data in somewhat better way for heavy nuclei with A > 150 as compared with other approaches. The parameters n = 3, $k_s(0)$ = 0.7, B_c = 0.7 can be recommended as more appropriate set in TPA calculations.

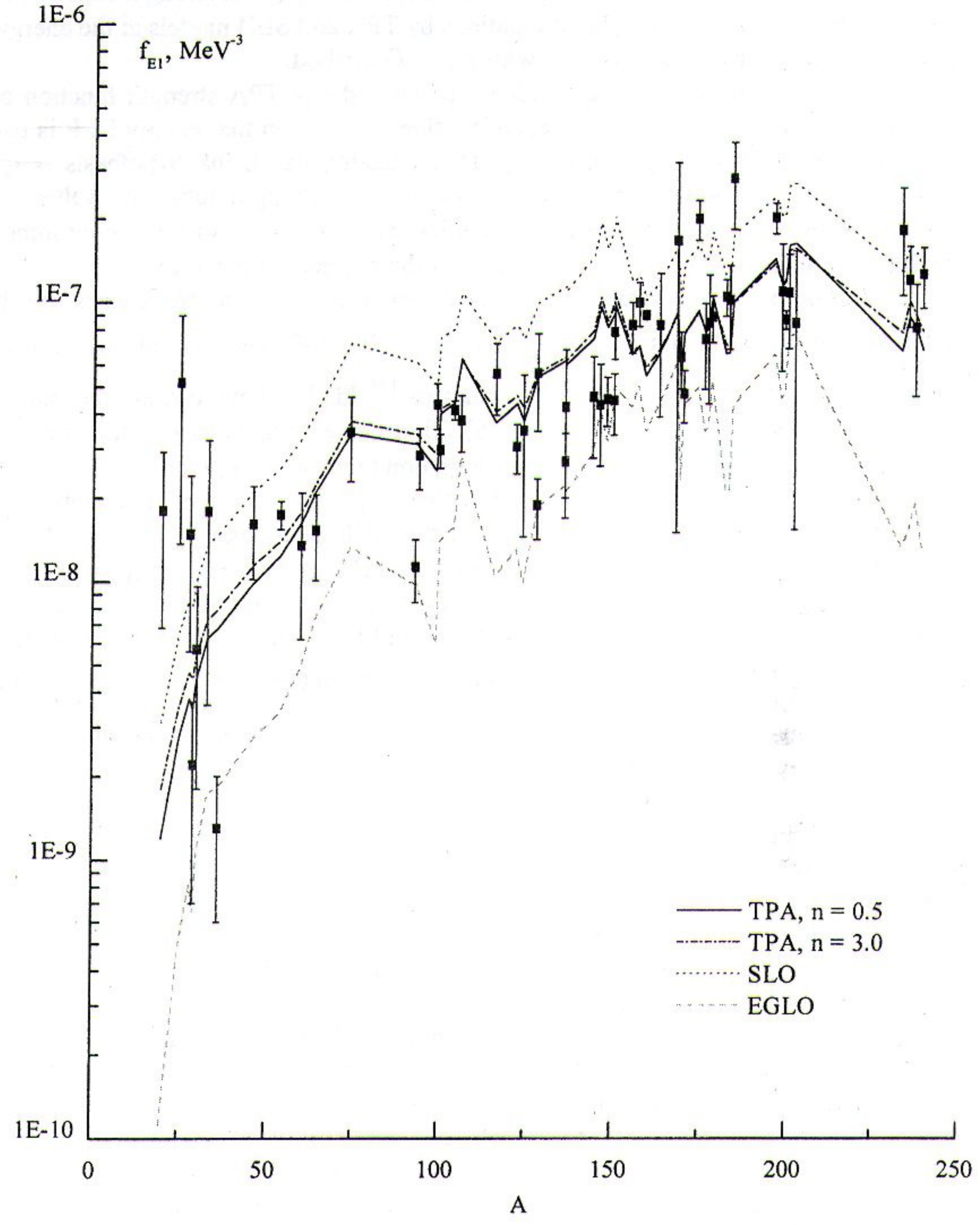


Fig.1. The E1 gamma-decay strength functions versus mass number A.

In fig.2 the results of the calculations of the strength functions $\overleftarrow{f}_{E1}$ in ^{144}Nd with the initial excitations energy E which is equal to the neutron binding energy $B_n \approx 7.8$ MeV are shown. The experimental data are taken from ref.[15].

The results obtained by EGLO and TPA approaches are almost the same at low energies $\varepsilon_\gamma \leq 3$MeV. In this range the EGLO and TPA models describe experimental data much better than the SLO model and give a non-zero temperature-dependent limit of the strength function for vanishing gamma-ray energy, see Eq.(38). The calculations by TPA and SLO models at the energies $\varepsilon_\gamma \geq 5$MeV lie more close to experimental data than within EGLO method.

Figure 3 demonstrates the dependence of the γ-decay TPA strength function on the initial excitation energy U. The E1 strength depends rather strongly on the energy U. It is usually named as a breakdown of Brink hypothesis [14]. This violation of Brink hypothesis is growing with increasing excitation energy. The difference of the E1 strength function values calculated at different U is increased with decreasing γ-energies and these deviations are more important for the γ-transitions with energies under or of the order of the nuclear temperature T.

In figs.4, 5 the comparison is shown between different approaches in the case the photoabsorption strength function $\overrightarrow{f}_{E1}$, Eqs.(19), (39), (40) and (41) at different values of the temperature T = 0.01, 2 MeV of absorbing nucleus ^{144}Nd. The notations are the same as in figs.2. The experimental data are taken from ref.[15]. They correspond to (n,γ) reaction at ε_γ = 6-8 MeV and were obtained from photoabsorption cross-section in the range $\varepsilon_\gamma >$ 8 MeV.

The behaviour of the E1 strength functions calculated by the TPA method is almost in coincidence with SLO model in the vicinity of the GDR peak energy. It is mainly resulted from account of the one-body relaxation width Γ_s, (28), which is practically independent of the gamma-ray energy. Note that the SLO approach is probably the most appropriate simple method for the estimation of the E1 photoabsorption strength for cold nuclei in the range of giant resonance peak energy. The strength function $\overrightarrow{f}_{E1}$ depends only weakly on the temperature if a magnitude of the T is much smaller than the gamma-ray energy. The form of the strength is rather sensitive to the excitation energy of absorbing nucleus at low energies of the γ-rays.

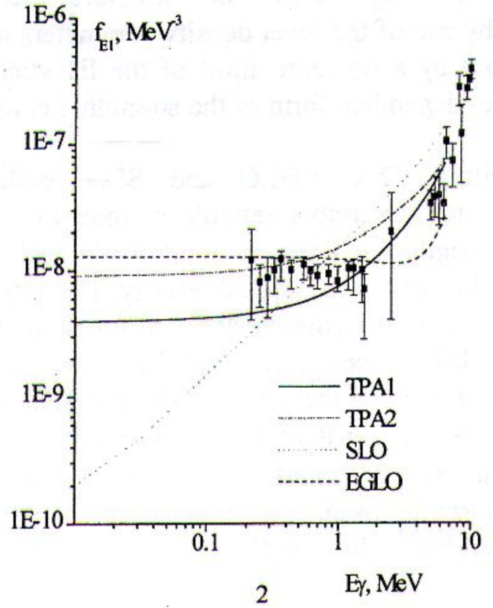
4.Conclusions

A closed-form TPA approach is developed for average description of the E1 radiative strength functions. This method is not time consuming and is applicable for calculations of the statistical contribution to the dipole strengths for processes of the gamma-decay as well as photoabsorption with compound system formation. It has the following main features:

1.The general expression between radiative strength function and imaginary part of the temperature response function is used. This relationship is based on microcanonical ensemble for initial excited states and it is in line with a detailed balance principle.

2.The form of the temperature response function is taken within framework of the Steinwedel-Jensen hydrodynamic model with damping.The response function has the Lorentzian line shape (two for axially deformed nuclei) with width depending on γ-ray energy. The Landau-Vlasov kinetic approach with the monopole and dipole distortions of the Fermi sphere is employed to calculate the damping width which is proportional to friction coefficient of the isovector velocity of the relative motion of the protons over neutrons.

3.Description of damping in the TPA method is based on modern physical understanding of the relaxation processes in Fermi systems. The contributions to the Lorentzian width resulting from the interparticle collisions as well as fragmentation component caused by interaction of particles with time dependent self-consistent mean field are included. A method of independent source of relaxation is employed to account for all contributions to width. The energy dependence of the collisional contribution is arisen from memory effects in the collision integral.



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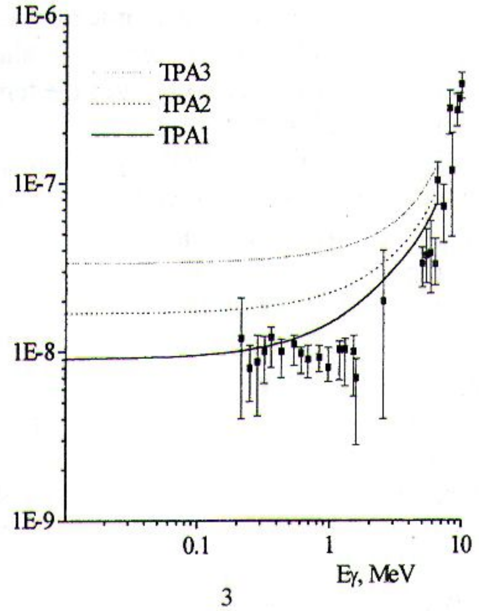

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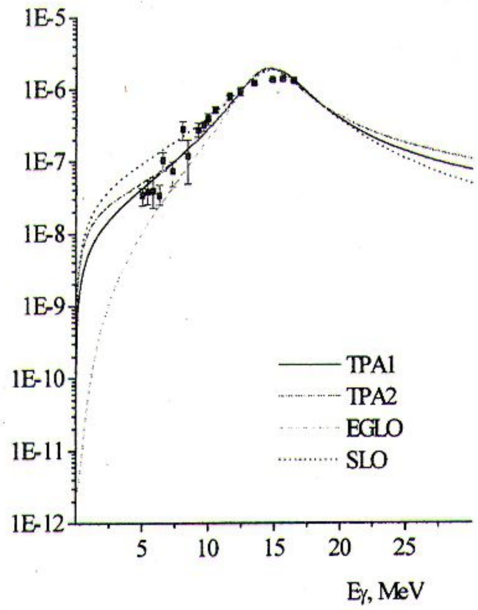

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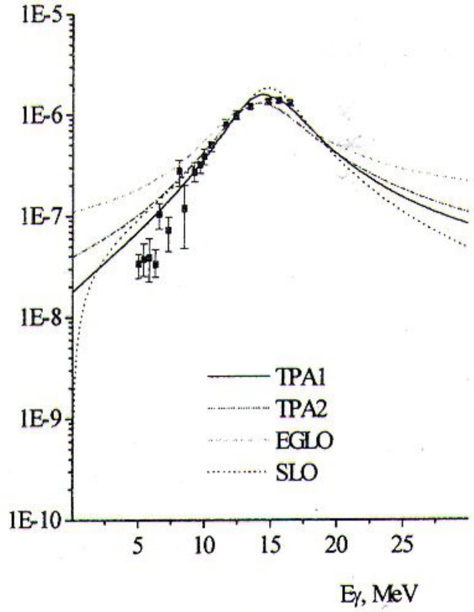

5

Figures

2. The E1 gamma-decay strength functions for Nd ^{144} at U=Bn;
3. Dependence of the gamma-decay strength functions at different excitation energies, n=3.0;
4, 5. The photoabsorption strength functions for Nd-144 at T=0.01 MeV(4) and T=2 MeV(5).

4.The form of the E1 radiative strength function within framework of the TPA model is determined by both the Lorentzian shape of the response function with energy dependent width and an average number of the excited 1p-1h states at given γ-ray energy. Shell structure and pairing correlations are included in phenomenological way by use of the level density parameters allowing for these effects. The TPA approach is characterized by a non-zero limit of the E1 strength for vanishing gamma-ray energy. It gives the temperature- dependent form of the strength, i.e. leads to a breakdown of Brink hypothesis.

The comparison between calculations within TPA, EGLO and SLO models and experimental data showed that the TPA approach provides rather reliable method of a unified description of the γ-decay and photoabsorption strength functions in a relatively wide energy interval, ranging from zeroth gamma-ray energy to values above GDR peak energy. The TPA model will be useful for the prediction of the downward and upward radiative strength functions for cold and heated nuclei.. Values $n = 3$, $k_s(0) = 0.7$ and $B_c = 0.7$ can be recommended as best suited set to calculations of the E1 strengths in medium and heavy nuclei by the TPA model. It should be noted that a behaviour of the TPA strength functions is rather sensitive to the type of γ-ray energy dependence of scaling coefficient $k_s(\varepsilon_\gamma)$ in Eq.(28) for one-body isovector width. The phenomenological approximations (49) and(50) are currently used. The further investigations of the fragmentation width are necessary to refine the form of the scaling coefficient.

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СТАТИСТИЧНИЙ ОПИС РАДІАЦІЙНИХ СИЛОВИХ ФУНКЦІЙ

В.А.Плюйко, С.М.Єжов, А.С.Микуляк

Розроблено термодинамічний підхід для опису усереднених E1 радіаційних силових функцій з використанням мікроканонічного ансамбля для початкових станів. Затухання колективних збуджень розглянуто напівкласично з використанням сучасних уявлень про процеси релаксації у Фермі-системах. Показано, що модель здатна описати дані по силовим функціям γ-розпаду та фотопоглинання у середніх та важких ядрах у відносно широкому інтервалі енергій, а саме від нульової енергії гамма-квантів до значень біля ГДР. Для ядер з A>150 і при енергіях гамма-квантів поблизу енергії зв'язку нейтронів запропонована модель описує експериментальні дані краще, ніж інші аналітичні підходи.

СТАТИСТИЧЕСКОЕ ОПИСАНИЕ РАДИАЦИОННЫХ СИЛОВЫХ ФУНКЦИЙ

В.А.Плюйко, С.М.Ежов, А.С.Микуляк

Развит термодинамический подход описания средних E1 радиационных силовых функций с использованием микроканонического ансамбля для начальных состояний. Затухание коллективных возбуждений рассматривается полуклассически на основе современных представлений о процессах релаксации в Ферми-системах. Показано, что модель описывает данные по силовым функциям γ-распада и фотопоглощения в средних и тяжелых ядрах в относительно широком интервале энергий, простирающемся от нулевых значений до энергий вблизи ГДР. Для ядер с A>150 и при енергиях гамма-квантов вблизи енергий связи нейтронов предложенная модель описывает експериментальные данные лучше, чем другие аналитические подходы