

TWO-BODY RELAXATION TIMES IN HEATED NUCLEI**V.A. Plujko¹, O.M. Gorbachenko, M.O. Kavatsyuk***Nuclear Physics Department, Taras Shevchenko National University, Pr. Acad. Glushkova, 2,
Bldg. 11, 03022 Kiev, Ukraine*

Received 21 June 2001, in final form 2 July 2001, accepted 3 July 2001

The retardation and temperature effects in two-body collisions are studied. The collision integral with retardation effects is obtained on the base of the Kadanoff-Baym equations for the Green functions in a form with allowance for reaching the local equilibrium system. The collisional relaxation times of collective vibrations are calculated using both the transport approach and the doorway state mechanism with hierarchy of particle-hole configurations in heated nuclei. The relaxation times of the kinetic method are rather weakly dependent on multipolarity of the Fermi surface distortion and on the mode of the collective motion. The dependence of the relaxation times on temperature as well as on frequency of collective vibrations is considered and compared. It is shown that variations of the in-medium two-body cross-sections with energy lead to non-quadratic dependence of the collisional relaxation time both on temperature and on the frequency of the collective motion.

PACS: 05.20.+w, 21.60.Ev, 21.65.+f, 24.10.Pa, 24.30.Cz

1 Introduction

The damping of the collective excitations as well as transport coefficients for viscosity and heat conductivity are strongly governed by the particle collisions. The relaxation time method is widely used as the simplest and rather accurate approach for simulation of the collisional relaxation rate $\nu_c \propto 1/\tau$, where τ is the so-called relaxation time [1–3]. Relaxation time method can be applied to the description of the decay rates of arbitrary modes of motion, but an explicit form of the relaxation time depends on specific features of the mode. In this contribution, the collisional relaxation times responsible for the width of the collective vibrations are studied.

2 Semiclassical kinetic equation approach

The collisional relaxation times can be calculated using the collision integral of the transport equation. In studies of the damping widths of collective excitation in the Fermi liquid, they are

¹E-mail address: plujko@mail.univ.kiev.ua

determined by the coefficients $\tau_\ell^{(\pm)}$ of the multipole expansion of the total number $\mathcal{N}^{(\pm)}(\hat{\mathbf{p}})$ of the collisions in the direction $\hat{\mathbf{p}} = \mathbf{p}/p$ of the momentum space [2, 4]

$$\mathcal{N}^{(\pm)}(\hat{\mathbf{p}}) \equiv \int_0^\infty d\epsilon J^{(\pm)}(\hat{\mathbf{p}}, \epsilon) = \sum_{\ell \geq \ell_0^{(\pm)}} \sum_{m=-\ell}^{\ell} \frac{\phi_{\ell m}^{(\pm)}}{\tau_\ell^{(\pm)}} Y_{\ell m}(\hat{\mathbf{p}}). \quad (1)$$

Here, $J^{(\pm)}(\hat{\mathbf{p}}, \epsilon)$ are the linearized collision integrals

$$J^{(\pm)}(\hat{\mathbf{p}}, \epsilon) \equiv J^{(\pm)}(\mathbf{p}, \mathbf{r}, t) = (J_p(\mathbf{p}, \mathbf{r}, t) \pm J_n(\mathbf{p}, \mathbf{r}, t))/2, \quad (2)$$

where the signs (+) and (−) stand for isoscalar and isovector modes of vibrations, respectively, and the subscripts p and n stand for protons and neutrons, respectively; $J_\alpha(\mathbf{p}, \mathbf{r}, t)$ is a collision integral in phase space (\mathbf{p}, \mathbf{r}) , when a nucleon of the sort $\alpha = (p \text{ or } n)$ with momentum \mathbf{p} is scattered; ϵ is nucleon energy. The collision integrals are linearized with respect to the dynamical component of the phase space distribution function $\delta f_\alpha(\mathbf{p}, \mathbf{r}, t)$. The values $\ell_0^{(\pm)}$ determine the initial components of the multipole expansion of the total number of the collisions. The functions $\nu_{\ell m}^{(\pm)}$ are the partial components of the energy-integrated distribution function $\delta f^{(\pm)}(\mathbf{p}, \mathbf{r}, t) = (\delta f_p \pm \delta f_n)/2 \equiv \delta f^{(\pm)}(\hat{\mathbf{p}}, \epsilon, \mathbf{r}, t)$,

$$\int_0^\infty d\epsilon \delta f^{(\pm)}(\hat{\mathbf{p}}, \epsilon, \mathbf{r}, t) = \sum_{\ell \geq 0} \sum_{m=-\ell}^{\ell} \phi_{\ell m}^{(\pm)}(\mathbf{r}, t) Y_{\ell m}(\hat{\mathbf{p}}), \quad (3)$$

where $Y_{\ell m}(\hat{\mathbf{p}})$ is the spherical harmonic function. In the approximation of a small difference between the chemical potentials of protons and neutrons and assuming $\bar{f}_p = \bar{f}_n = \bar{f}$, where $\bar{f} \equiv \bar{f}(\mathbf{p}, \mathbf{r})$ is the equilibrium distribution function, the dynamical distortions $\delta f^{(\pm)}(\mathbf{p}, \mathbf{r}, t)$ of the phase space distribution functions are solutions of the linearized Landau-Vlasov equation

$$\frac{\partial \delta f^{(\pm)}}{\partial t} + \frac{\mathbf{p}}{m} \frac{\partial \delta f^{(\pm)}}{\partial \mathbf{r}} - \frac{\partial \bar{\epsilon}^{(\pm)}}{\partial \mathbf{r}} \frac{\partial \delta f^{(\pm)}}{\partial \mathbf{p}} - \frac{\partial \delta U^{(\pm)}}{\partial \mathbf{r}} \frac{\partial \bar{f}}{\partial \mathbf{p}} = J^{(\pm)}(\mathbf{p}, \mathbf{r}, t), \quad (4)$$

where $\delta U^{(\pm)} \equiv \delta U^{(\pm)}(\mathbf{p}, \mathbf{r}, t)$ is the Wigner transform of the variation of the self-consistent potential with respect to the equilibrium value $\bar{\epsilon}^{(\pm)}$. In the nuclear interior the mean field variation $\delta U^{(\pm)}$ can be expressed in terms of the Landau interaction amplitude $F^{(\pm)}(\mathbf{p}, \mathbf{p}')$ as

$$\delta U^{(\pm)} = \frac{g}{N_F} \int \frac{d\mathbf{p}'}{(2\pi\hbar)^3} F^{(\pm)}(\mathbf{p}, \mathbf{p}') \delta f^{(\pm)}(\mathbf{p}', \mathbf{r}, t), \quad (5)$$

where $N_F = 2p_F m^*/(g\pi^2\hbar^3)$, p_F is the Fermi momentum, m^* is the effective mass of nucleon and g is the spin degeneracy factor. The quantity $F^{(\pm)}(\mathbf{p}, \mathbf{p}')$ is usually parameterized in terms of the Landau constants $F_0^{(\pm)}$ and $F_1^{(\pm)}$ as

$$F^{(\pm)}(\mathbf{p}, \mathbf{p}') = F_0^{(\pm)} + F_1^{(\pm)}(\hat{\mathbf{p}} \cdot \hat{\mathbf{p}}'). \quad (6)$$

In the isoscalar case, the Landau constants are related to the incompressibility modulus K [5] of matter and the effective mass m^* [6] by

$$K = 6\mu(1 + F_0^{(+)}), \quad m^* = m \left(1 + F_1^{(+)} / 3\right). \quad (7)$$

Here m is the mass of free nucleon and μ is the chemical potential. We have $\mu \approx \epsilon_F = p_F^2/2m^*$ for $T \ll \epsilon_F$, where ϵ_F is the Fermi energy and T is the temperature. In the isovector case, the Landau parameter $F_0^{(-)}$ is related to the nuclear symmetry energy b_{symm} . Namely [7, 8],

$$b_{\text{symm}} = \frac{1}{3} \mu (1 + F_0^{(-)}). \quad (8)$$

The quantities $\tau_\ell^{(\pm)}$ in Eq. (1) can be considered as the partial collective relaxation times because they determine a collisional contribution to the damping widths resulting from the two-body collisions in the layer of the momentum space with multipolarity ℓ ,

$$\frac{1}{\tau_\ell^{(\pm)}} \equiv \int_0^\infty d\epsilon \int d\Omega_p J^{(\pm)}(\hat{\mathbf{p}}, \epsilon) Y_{\ell 0}(\hat{\mathbf{p}}) / \int_0^\infty d\epsilon \int d\Omega_p \delta f^{(\pm)} Y_{\ell 0}(\hat{\mathbf{p}}). \quad (9)$$

These times are proportional to the relaxation times $\tau_c^{(\pm)}$ defining the damping widths $\Gamma_c^{(\pm)}(L)$ of the isoscalar and the isovector vibrations with frequency ω [2, 4, 9, 10] in the regime of rare collisions with $\omega \tau_c^{(\pm)} \gg 1$ in the Fermi liquid. In particular, the collisional damping widths of giant resonances with dipole ($L = 1$) and quadrupole ($L = 2$) multiplicities resemble the widths in the relaxation rate approach

$$\Gamma_c^{(\pm)}(L) = \hbar / \tau_c^{(\pm)}(L), \quad \tau_c^{(-)}(L = 1) = \tau_{\ell=1}^{(-)}, \quad \tau_c^{(+)}(L = 2) = \tau_{\ell=2}^{(+)}, \quad (10)$$

in the case when nuclear fluid dynamical model with relaxation is used [4, 9–12]. The collisional damping width [2] of zero sound in the Fermi liquid with its relative velocity $S_r \simeq 1$ is also given by Eq.(10) but with the use of the $\tau_{\ell \rightarrow \infty}^{(+)} \propto \tau_2^{(+)}$ for $\tau_c^{(\pm)}(L)$. The time $\tau_{\ell=2}^{(+)}$ at $\omega = 0$ is the thermal relaxation time determining the viscosity coefficient of the Fermi liquid [13].

The variations of the mean field and of the dynamical component of the phase-space distribution function change rapidly in the systems with high frequency collective vibrations. This leads to the memory (retardation, i.e. non-Markovian) effects in the collision term. There are different expressions for memory-dependent collision integral in the Fermi liquid ([14]- [20]).

The non-Markovian collision term of the semiclassical Landau-Vlasov equation was studied in Born approximation with the use of the Kadanoff-Baym equations for the Green functions in Refs. [17, 19]. The one-component Fermi liquid was considered with the periodic time variation of the nonequilibrium distribution function $\delta f = \delta f_n = \delta f_p$, $\delta f \propto \exp(-i\omega t)$. As a result, the linearized collision integral consists of two components (see Eqs.(42),(43) and (45),(46) of Ref. [19]), i.e.,

$$J(\mathbf{p}, \mathbf{r}, t) = J^{(1)}(\mathbf{p}, \mathbf{r}, t) + J^{(2)}(\mathbf{p}, \mathbf{r}, t), \quad (11)$$

where the components $J^{(1)}(\mathbf{p}, \mathbf{r}, t)$ and $J^{(2)}(\mathbf{p}, \mathbf{r}, t)$ are determined by the variations of the distribution function and the mean field, respectively, and

$$J^{(j)}(\mathbf{p}, \mathbf{r}, t) = 2 \int \frac{d\mathbf{p}_2 d\mathbf{p}_3 d\mathbf{p}_4}{(2\pi\hbar)^6} W(\{\mathbf{p}_i\}) \delta(\Delta\mathbf{p}) B^{(j)}(\mathbf{p}, \mathbf{r}, t). \quad (12)$$

Here, $W(\{\mathbf{p}_i\}) = (d\sigma/d\Omega)4(2\pi\hbar)^3/m^2$ is the probability of two-body collisions with the initial momenta $\mathbf{p}_1 = \mathbf{p}, \mathbf{p}_2$ and final ones $\mathbf{p}_3, \mathbf{p}_4$, ($i = 1 \div 4$); $d\sigma/d\Omega$ is the in-medium differential

cross-section (in the Born approximation);

$$B^{(1)}(\mathbf{p}, \mathbf{r}, t) = \sum_{k=1}^4 \delta f_k(t) \frac{\partial Q(\{\bar{f}_j\})}{\partial \bar{f}_k} [\delta_+(\Delta\bar{\epsilon} + \hbar\omega) + \delta_-(\Delta\bar{\epsilon} - \hbar\omega)], \quad (13)$$

$$B^{(2)}(\mathbf{p}, \mathbf{r}, t) = Q(\{\bar{f}_j\}) \frac{\Delta(\delta U(t))}{\hbar\omega} \{[\delta_+(\Delta\bar{\epsilon} + \hbar\omega) - \delta_+(\Delta\bar{\epsilon})] - [\delta_-(\Delta\bar{\epsilon} - \hbar\omega) - \delta_-(\Delta\bar{\epsilon})]\},$$

where $\bar{f}_k \equiv \bar{f}(\mathbf{p}_k, \mathbf{r})$; $\partial Q(\{\bar{f}_j\})/\partial \bar{f}_k$ are the derivatives of the is the Pauli blocking factor Q with respect to \bar{f}_k ,

$$Q(\{\bar{f}_j\}) = (1 - \bar{f}_1)(1 - \bar{f}_2)\bar{f}_3\bar{f}_4 - \bar{f}_1\bar{f}_2(1 - \bar{f}_3)(1 - \bar{f}_4). \quad (14)$$

The $\bar{\epsilon}_i = \bar{\epsilon}(\mathbf{p}_i, \mathbf{r})$ is the equilibrium single-particle energy and δU_j the variation of the mean field, and $\Delta\bar{\epsilon} = \bar{\epsilon}_1 + \bar{\epsilon}_2 - \bar{\epsilon}_3 - \bar{\epsilon}_4$, $\Delta(\delta U) \equiv \delta U_1 + \delta U_2 - \delta U_3 - \delta U_4$, $\Delta\mathbf{p} = \mathbf{p}_1 + \mathbf{p}_2 - \mathbf{p}_3 - \mathbf{p}_4$. The equilibrium distribution function $\bar{f}_k \equiv \bar{f}(\mathbf{p}_k, \mathbf{r})$ depends on the equilibrium single-particle energy $\bar{\epsilon}_k \equiv \bar{\epsilon}(\mathbf{p}_k, \mathbf{r})$: $\bar{f}_k = \bar{f}(\bar{\epsilon}_k)$. It equals the Fermi function evaluated at the temperature T , $\bar{f}(\bar{\epsilon}_k) = 1/[1 + \exp((\bar{\epsilon}_k - \mu)/T)]$. The nonequilibrium component δf of the distribution function can be presented in the the form

$$\delta f(\mathbf{p}, \mathbf{r}, t) = -\nu(\mathbf{p}, \mathbf{r}, t) \frac{\partial \bar{f}(\bar{\epsilon})}{\partial \bar{\epsilon}}. \quad (15)$$

With the use of this relation the expression for quantity $B^{(1)}$ can be transformed to the form

$$\begin{aligned} B^{(1)}(\mathbf{p}, \mathbf{r}, t) &= - \sum_{k=1}^4 \nu_k \frac{\partial Q(\{\bar{f}_j\})}{\partial \bar{\epsilon}_k} [\delta_+(\Delta\bar{\epsilon} + \hbar\omega) + \delta_-(\Delta\bar{\epsilon} - \hbar\omega)] = \\ &= \Delta\nu Q(\{\bar{f}_j\}) \frac{\partial}{\partial \hbar\omega} [\delta_+(\Delta\bar{\epsilon} + \hbar\omega) + \delta_-(\Delta\bar{\epsilon} - \hbar\omega)] - \delta B^{(1)}, \end{aligned} \quad (16)$$

where $\Delta\nu \equiv \nu_1 + \nu_2 - \nu_3 - \nu_4$, $\nu_k = \nu(\mathbf{p}_k, \mathbf{r}, t)$ and

$$\begin{aligned} \delta B^{(1)} &= \sum_{k=1}^4 \frac{\partial}{\partial \bar{\epsilon}_k} \{ \nu_k Q(\{\bar{f}_j\}) [\delta_+(\Delta\bar{\epsilon} + \hbar\omega) + \delta_-(\Delta\bar{\epsilon} - \hbar\omega)] \} \\ &+ \sum_{k=1}^4 Q(\{\bar{f}_j\}) [\delta_+(\Delta\bar{\epsilon} + \hbar\omega) + \delta_-(\Delta\bar{\epsilon} - \hbar\omega)] \frac{\partial \nu_k}{\partial \bar{\epsilon}_k}. \end{aligned} \quad (17)$$

The first component in the Eq.(17) determines a probability flux of colliding particles which is connected with possibility of variation of the energy $\bar{\epsilon}_k$ when the values of other energies ($\bar{\epsilon}_{j \neq k}$ and $\hbar\omega$) are fixed. This term should be equal zero because of fixing the total energy in initial or final states and therefore it does not contribute to the total number of the collisions $\mathcal{N}^{(\pm)}$, Eq.(1). The last statement can be easily verified by direct calculation of this contribution to the $\mathcal{N}^{(\pm)}$ with the use of the procedure proposed by Abrikosov and Khalatnikov (see Eqs.(25) to (30)) for evaluation of the manifold energy integrals. A relative dynamical component ν_k of the distribution function is slowly dependent on energy and it can be considered (at least for low

temperatures $T \ll \epsilon_F$) as a function of the momentum direction rather than of the momentum: $\nu_k \equiv \nu(\mathbf{p}_k, \mathbf{r}, t) = \nu(\hat{\mathbf{p}}_k, \epsilon_F, \mathbf{r}, t)$. Therefore the second component in the Eq.(17) is also negligible and the term $\delta B^{(1)}$ in the Eq.(16) should be rejected,

$$\delta B^{(1)} = 0. \quad (18)$$

Note that the generalized functions δ_+ , δ_- appearing in Eqs. (13), (16) and (17) include also integral contribution,

$$\delta_+(x) = \frac{1}{2\pi} \int_{-\infty}^0 d\tau e^{-ix\tau} = \frac{i}{2\pi} \frac{1}{x+i0} = \frac{1}{2} \delta(x) - \frac{1}{2\pi i} \mathcal{P}\left(\frac{1}{x}\right), \quad \delta_-(x) = \delta_+^*(x), \quad (19)$$

where $\delta(x)$ is the delta function and the symbol \mathcal{P} denotes the principal value of integral contribution. The integral terms of the δ_{\pm} , corresponding to virtual transitions, are usually omitted in the J because they assumed to be included by renormalizing the interactions between particles [21]. This corresponds to substituting $\delta(x)/2$ for δ_{\pm} in Eqs. (13), i.e., to taking into account real transitions with conservation of energy. The shift in energy $\Delta\bar{\epsilon}$ by $\hbar\omega$ in the arguments of the δ -functions of the expressions for the collision integral agrees with the interpretation of the collisions in the presence of the collective excitations proposed by Landau [22]. According to this interpretation, high-frequency oscillations in Fermi liquid can be considered as phonons, that are absorbed and created at the two-particle collisions.

In the one-component Fermi liquid the nonequilibrium distribution function $\delta f(\mathbf{p}, \mathbf{r}, t) \equiv f(\mathbf{p}, \mathbf{r}, t) - \bar{f}(\bar{\epsilon}(\mathbf{p}, \mathbf{r}))$ is a solution of the linearized Landau-Vlasov equation in the form

$$\frac{\partial \delta f}{\partial t} + \frac{p}{m} (\hat{\mathbf{p}} \cdot \frac{\partial}{\partial \mathbf{r}}) \delta \bar{f} = J. \quad (20)$$

Here, $\delta \bar{f}$ is a linear deviation of the distribution function from its local equilibrium value $f_{l.e.}$, where a function $f_{l.e.}$ is equal to the the Fermi function $\bar{f}(\epsilon)$ evaluated with actual one-particle energy $\epsilon = \bar{\epsilon} + \delta U$, $f_{l.e.} = \bar{f}(\epsilon(\mathbf{p}, \mathbf{r}, t))$,

$$\delta \bar{f} = \delta f - \frac{d\bar{f}}{d\bar{\epsilon}} \delta U = f(\mathbf{p}, \mathbf{r}, t) - f_{l.e.} = -\chi \frac{d\bar{f}}{d\bar{\epsilon}}, \quad \chi = \nu + \delta U, \quad (21)$$

where relationship for δU has the form of the Eq.(5) but with the interactions and distribution function for one-component Fermi liquid.

According to Eq.(20), a Fermi system tends to the local equilibrium (when $\partial \delta f / \partial t = 0$) if the collision integral is a functional of the $\delta \bar{f}$, $J = \Phi(\delta \bar{f})$. The collision integral given by the expressions (11)-(13) and (16), (18) has the following general form: $\delta J = \delta J^{(1)}(\delta f) + \delta J^{(2)}(\delta U)$. Therefore it does not lead to local equilibrium of a system. The condition of the existence of the local equilibrium of a system is general property of the Landau-Vlasov equation in the Fermi liquid [23, 24] at $\partial \delta f / \partial t = 0$. Therefore the Born approximation (11)-(13) for collision integral is poor approach without additional modification, and a revision of the derivation method of the collision integral expression is needed.

It should be initially noted that the foregoing relationships for collision integral are obtained with the use of the perturbation theory in nearly nonhomogeneous systems with week interaction between particles. A week interaction can not change rapidly the trajectory of the particle and due to this it can not lead to rapid variations of the distribution function. It means that retardation

effects are overestimated in the expression for collision integral in Born approximation where it was assumed that distribution function was varied very quickly during all possible interval of the time changing ($-\infty \leq t' \leq t$). Consequently, the collision integral given by Eqs. (11)-(13), (16), (18) can be in fact correct in the case of small retardation, i.e., for small values of the $\hbar\omega$.

With this in mind, we replace the derivatives of the form $\partial \delta_+(\Delta\bar{\epsilon} + \hbar\omega)/\partial \hbar\omega$ and $\partial \delta_-(\Delta\bar{\epsilon} - \hbar\omega)/\partial \hbar\omega$ in the Eq.(16) by the finite differences $(\delta_+(\Delta\bar{\epsilon} + \hbar\omega) - \delta_+(\Delta\bar{\epsilon}))/\hbar\omega$ and $(\delta_-(\Delta\bar{\epsilon}) - \delta_-(\Delta\bar{\epsilon} - \hbar\omega))/\hbar\omega$, respectively. Then we combine the resulting expression together with contribution $B^{(2)}$ arising from mean-field variation and obtain the linearized collision integral for one-component Fermi liquid in the following form

$$J(\mathbf{p}, \mathbf{r}, t) = \int \frac{d\mathbf{p}_2 d\mathbf{p}_3 d\mathbf{p}_4}{(2\pi\hbar)^6} W(\{\mathbf{p}_i\}) \delta(\Delta\mathbf{p}) \Delta\chi Q \frac{\delta(\Delta\epsilon + \hbar\omega) - \delta(\Delta\epsilon - \hbar\omega)}{\hbar\omega}. \quad (22)$$

With the use of the algebraic relation [23]

$$[(1 - \bar{f}_1)(1 - \bar{f}_2)\bar{f}_3\bar{f}_4 - \bar{f}_1\bar{f}_2(1 - \bar{f}_3)(1 - \bar{f}_4) \exp\left(\frac{\mp\hbar\omega}{T}\right)] \delta(\Delta\epsilon \pm \hbar\omega) = 0, \quad (23)$$

Eq. (22) can be presented as

$$J(\mathbf{p}, \mathbf{r}, t) = \int \frac{d\mathbf{p}_2 d\mathbf{p}_3 d\mathbf{p}_4}{(2\pi\hbar)^6} W(\{\mathbf{p}_i\}) \delta(\Delta\mathbf{p}) \Delta\chi \bar{f}_1 \bar{f}_2 (1 - \bar{f}_3)(1 - \bar{f}_4) [\Phi(\hbar\omega, T) - \Phi(-\hbar\omega, T)], \quad (24)$$

where $\Phi(\hbar\omega, T) = \delta(\Delta\epsilon + \hbar\omega)[\exp(-\hbar\omega/T) - 1]/\hbar\omega$.

The collision integral of the form (22) or (24) provides driving distribution function towards its local equilibrium value because it depends on the variation δf , $J \equiv J(\delta f)$. This behaviour is in line with general properties of the Landau- Vlasov equation in the Fermi-liquid [23, 24] at $\partial \delta f / \partial t = 0$. The expressions (22), (24) depend only on the occupation probability $\mathcal{P}_{2p2h} \equiv \bar{f}_1 \bar{f}_2 (1 - \bar{f}_3)(1 - \bar{f}_4)$ of the 2p-2h states in the phase space. This fact leads to interpretation of the collisional damping with linearized collision term as the relaxation process due to the coupling of one-particle and one-hole states to more complicated $2p - 2h$ configurations.

The form of the collision integral (24) in the Markovian limit ($\omega \rightarrow 0$) coincides with the standard expression for the collision integral in Fermi-liquid without retardation effects [23, 24] because in this case the term in square brackets of Eq.(24) tends to the value $-\delta(\Delta\epsilon)/T$.

The equation (22) for some special explicit form of the quantity χ_j was used at first in Refs. [16, 25, 26]. The derivation of the collision integral (22) is performed in Ref. [14] within framework of the extended time-dependent Hartree-Fock model. The expressions for the distortion functions χ_j corresponding to a perturbation approach on collision term and including the amplitudes of the random phase approximation were used in this method.

The expression for the collision integral in two-component Fermi-system is obtained from Eq.(24) in the same manner as done in Ref. [4] under the assumption that chemical potentials and the equilibrium distribution functions are the same for protons and neutrons.

The analytical expressions for partial collective relaxation times τ_ℓ (Eq.(9)) can be obtained in low-temperature and low-frequency limits ($T, \hbar\omega \ll \epsilon_F$). In this case the momentum integrals are calculated using the Abrikosov- Khalatnikov procedure [23, 27-29] which is based

on the assumption that particles are scattered near Fermi surface with the momentum values p_i approximately equal to the Fermi momentum p_F . In this case the probability $W(\{\mathbf{p}_i\})$ of two-body collisions can be taken as a function of two scattering angles ϕ and θ , where ϕ is the angle between the momenta \mathbf{p}_1 and \mathbf{p}_2 , and θ is the angle between the $(\mathbf{p}_1\mathbf{p}_2)$ and $(\mathbf{p}_3\mathbf{p}_4)$ planes; that is

$$\begin{aligned}\cos \phi &= (\hat{\mathbf{p}}_2 \cdot \hat{\mathbf{p}}_1), \\ \cos \theta &= [\hat{\mathbf{p}}_1 \times \hat{\mathbf{p}}_2] \cdot [\hat{\mathbf{p}}_3 \times \hat{\mathbf{p}}_4] / |[\hat{\mathbf{p}}_1 \times \hat{\mathbf{p}}_2]| |[\hat{\mathbf{p}}_3 \times \hat{\mathbf{p}}_4]|.\end{aligned}\quad (25)$$

It allows to separate the angular and the energy integrations in the collision integral at arbitrary scattering angle [29]. The integrals with respect to momenta in expression for the collision integral are calculated employing the transformation [23, 27, 29]

$$\begin{aligned}& \int d\mathbf{p}_2 d\mathbf{p}_3 d\mathbf{p}_4 \delta(\Delta\mathbf{p})(\dots) = \\ &= \frac{m^{*3}}{2} \int_0^\pi d\phi \int_0^\pi d\theta \frac{\sin \phi}{\cos \frac{\phi}{2}} \int_0^{2\pi} d\varphi \int_0^\infty \int_0^\infty \int_0^\infty d\epsilon_2 d\epsilon_3 d\epsilon_4 (\dots).\end{aligned}\quad (26)$$

Here, φ is the azimuthal angle of the momentum \mathbf{p}_2 in the coordinate system with the Z axes along \mathbf{p}_1 .

The integration with respect to the azimuthal angle φ is performed by the relation [28]

$$\int_0^{2\pi} \frac{d\varphi}{2\pi} Y_{lm}(\hat{\mathbf{p}}_j) = Y_{lm}(\hat{\mathbf{p}}_1) P_l(\hat{\mathbf{p}}_j \hat{\mathbf{p}}_1), \quad (27)$$

where P_l is a Legendre polynomial, and

$$\begin{aligned}(\hat{\mathbf{p}}_3 \hat{\mathbf{p}}_1) &= \cos^2(\phi/2) + \sin^2(\phi/2) \cos \theta, \\ (\hat{\mathbf{p}}_4 \hat{\mathbf{p}}_1) &= \cos^2(\phi/2) - \sin^2(\phi/2) \cos \theta.\end{aligned}\quad (28)$$

To perform over energies in the collision integral the following expressions are used [23]

$$\begin{aligned}I_\nu(y) &= \int_{-\infty}^{+\infty} dx_1 dx_2 \dots dx_\nu n(x_1) n(x_2) \dots n(x_\nu) \delta(x_1 + x_2 + \dots + x_\nu) \equiv \\ &\equiv \int_{-\infty}^{+\infty} dx_\nu n(x_\nu) I_{\nu-1}(x_\nu + y) \equiv \int_{-\infty}^{+\infty} dt n(t - y) I_{\nu-1}(t),\end{aligned}\quad (29)$$

where $n(x) = 1/(1 + \exp(x))$, $n(x) + n(-x) = 1$ and

$$I_3(y) = \frac{1}{2} \frac{y^2 + \pi^2}{1 + \exp(-y)}, \quad I_4(y) = \frac{1}{6} \frac{y(y^2 + 4\pi^2)}{(1 - \exp(-y))}. \quad (30)$$

Finally we get the following relation for relaxation times using the collisional integral of the form given by Eq.(24):

$$\frac{\hbar}{\tau_\ell^{(\pm)}} = \mathcal{R}(\omega, T) \left[\langle \sigma'_{av} \Phi_\ell^{(+)} \rangle + \langle \sigma'_{pn} \Phi_\ell^{(\pm)} \rangle \right], \quad (31)$$

where $\sigma'_{av} = (\sigma'_{nn} + \sigma'_{pp})/2$; $\sigma'_{jj'} \equiv d\sigma_{jj'}/d\Omega$ is in-medium differential cross-section for scattering of the nucleons j and j' (here, $j = n$ or p , and similarly $j' = p$ or n). The quantity $\mathcal{R}(\omega, T)$ result from the energy integrations and has the following form in low-temperature and low-frequency limits ($T, \hbar\omega \ll \epsilon_F$) in the approximation $m^* \simeq m$

$$\mathcal{R}(\omega, T) = \frac{2}{3\pi} \frac{m}{\hbar^2} \{(\hbar\omega)^2 + (2\pi T)^2\}. \quad (32)$$

The symbol $\langle \dots \rangle$ in Eq.(31) denotes averaging over angles of the relative momenta of the colliding particles,

$$\langle (\dots) \rangle = \frac{1}{\pi} \int_0^\pi d\phi \sin(\phi/2) \int_0^\pi d\theta (\dots). \quad (33)$$

The functions $\Phi_\ell^{(\pm)}$ in (31) define the angular constraints on nucleon scattering within the distorted layers of the Fermi surface with multipolarity ℓ :

$$\Phi_\ell^{(\pm)} = 1 \pm P_\ell(\cos \phi) - P_\ell((\hat{\mathbf{p}}_3 \hat{\mathbf{p}}_1)) \mp P_\ell((\hat{\mathbf{p}}_4 \hat{\mathbf{p}}_1)), \quad (34)$$

where the scalar products $(\hat{\mathbf{p}}_3 \hat{\mathbf{p}}_1)$ and $(\hat{\mathbf{p}}_4 \hat{\mathbf{p}}_1)$ are given by Eq.(28). It follows $\Phi_{\ell=0}^{(+)}(\phi, \theta) = \Phi_{\ell=1}^{(+)}(\phi, \theta) = \Phi_{\ell=0}^{(-)}(\phi, \theta) = 0$. These relations lead to possibility of the two-body collisions in layers of the Fermi surface distortion with multipolarity beginning with the value $\ell_0^{(+)} = 2$ in the isoscalar case and $\ell_0^{(-)} = 1$ for the isovector vibrations. As a result, the isovector dipole relaxation time $\tau_{\ell=1}^{(-)}$ has a finite value, that means a nonconservation of the isovector current in the presence of $n-p$ collisions [30].

Due to the momentum conservation and conditions $p_i = p_F$, the angle θ agrees with the scattering angle in the center-of-mass reference frame of two nucleons. The angle ϕ defines the magnitudes of the relative momenta $\mathbf{k}_i = (\mathbf{p}_2 - \mathbf{p}_1)/2$ and $\mathbf{k}_f = (\mathbf{p}_4 - \mathbf{p}_3)/2$ before and after collision, respectively. The value of total momentum, $\mathbf{P} = \mathbf{p}_1 + \mathbf{p}_2$, also depends on a magnitude of the ϕ . We have

$$\mathbf{k}_i \mathbf{k}_f = k^2 \cos \theta, \quad k^2 = k_i^2 = k_f^2 = p_F^2 \sin^2(\phi/2), \quad \mathbf{P}^2 = 4p_F^2 \cos^2(\phi/2). \quad (35)$$

Therefore the relative kinetic energy E_{rel} of two nucleons as well as the energy of centrum mass motion E_{cm} are dependent on angle ϕ

$$E_{rel} = k^2/m = 2\epsilon_F \sin^2(\phi/2), \quad E_{cm} = P^2/2m = 2\epsilon_F \cos^2(\phi/2) \quad (36)$$

and the total energy $E_{tot} = E_{rel} + E_{cm}$ holds only fixed, $E_{tot} = 2\epsilon_F$. Therefore the in-medium differential cross-sections $\sigma'_{j,m}$ of two nucleon scattering depend on the relative momenta \mathbf{k}_i and \mathbf{k}_f at fixed total energy rather than at fixed relative kinetic energy E_{rel} , because the magnitude of E_{rel} changes with angle ϕ between colliding particles. The transfer momenta $\mathbf{q} = \mathbf{k}_i - \mathbf{k}_f = \mathbf{p}_3 - \mathbf{p}_1$ and $\mathbf{q}' = -(\mathbf{k}_i + \mathbf{k}_f) = \mathbf{p}_1 - \mathbf{p}_4$ for scattering due to direct and exchange interactions respectively are also functions of the ϕ and θ : $q = 2k(\phi) \sin(\theta/2)$ and $q' = 2k(\phi) \cos(\theta/2)$.

Now we estimate the collisional relaxation times in the case of the isotropic scattering with independent of energy the angle-integrated cross sections $\sigma_{jj'}$. Performing angular integration in (31) with the use of Eqs.(33) and (34) we find that $1/\tau_{\ell < \ell_0^{(\pm)}}^{(\pm)} = 0$ and

$$\frac{\hbar}{\tau_\ell^{(\pm)}} = \frac{1}{\alpha_\ell^{(\pm)}} [(\hbar\omega/2\pi)^2 + T^2], \quad \frac{1}{\alpha_\ell^{(\pm)}} = \frac{8m}{3\hbar^2} [c_\ell \sigma_{av} + d_\ell^{(\pm)} \sigma_{np}], \quad (37)$$

$$c_\ell = 1 - \frac{2 - (-1)^\ell}{2\ell + 1}, \quad d_\ell^{(-)} = \frac{1 - (-1)^\ell}{2\ell + 1}, \quad d_\ell^{(+)} = d_{\ell=0}^{(-)} = c_{\ell=0} = c_{\ell=1} = 0,$$

where $\sigma_{av} = [\sigma_{pp} + \sigma_{nn} + 2\sigma_{np}]/4$ is the in-medium spin-isospin averaged nucleon-nucleon cross section. The magnitude of the in-medium cross section $\sigma_{jj'}$ is taken usually proportional to the value of the free space cross section $\sigma_{jj'}^{(free)}$ with a factor $F = \sigma_{jj'}/\sigma_{jj'}^{(free)}$, so that the parameter $\alpha_\ell^{(\pm)}$ can be rewritten in the form

$$\alpha_\ell^{(\pm)} = \tilde{\alpha}_\ell^{(\pm)}/F, \quad \tilde{\alpha}_\ell^{(\pm)} = 4.18 / \left[c_\ell + 1.3d_\ell^{(\pm)} \right], \quad MeV. \quad (38)$$

Here, the values $\sigma_{av}^{(free)} = 3.75 \text{ fm}^2$ and $\sigma_{np}^{(free)} = 5 \text{ fm}^2$ are adopted [10,16]; they correspond to the free space cross sections near Fermi energy.

The relative relaxation times $\tau_\ell^{(\pm)}/\tau_{\ell=2}^{(+)}$ with the free space cross sections are shown on Fig.1 in relation to the multipolarity ℓ of the distorted layers of the Fermi surface which are accessible to particle collisions. Solid and dashed lines connect the values which correspond to isoscalar and isovector modes of vibrations respectively. The magnitudes of the relaxation times are different for isoscalar and isovector modes of vibrations and they are dependent on the multipolarity ℓ . As seen from the Fig.1, the collisional relaxation times rather slowly vary with multipolarity ℓ and with collective motion mode at isotropic scattering with energy independent free cross sections. In particular, parameters $\tilde{\alpha}_\ell^{(\pm)}$, which define relaxation times by the Eq.(37), take the same value at $\ell \rightarrow \infty$, $\tilde{\alpha}_{\ell \rightarrow \infty}^{(\pm)} \equiv \tilde{\alpha} = 4.18 \text{ MeV}$, and $\tilde{\alpha}/\tilde{\alpha}_{\ell=1}^{(-)} \simeq 0.9$, $\tilde{\alpha}_{\ell=2}^{(-)}/\tilde{\alpha}_{\ell=1}^{(-)} \simeq 1.1$; $\tilde{\alpha}/\tilde{\alpha}_{\ell=2}^{(+)} \simeq 0.8$, $\tilde{\alpha}_{\ell=3}^{(+)}/\tilde{\alpha}_{\ell=2}^{(+)} \simeq 1.4$.

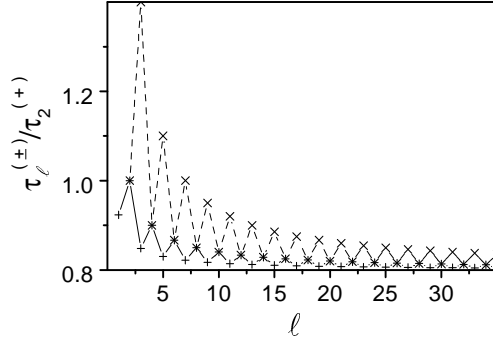


Fig.1: The relative relaxation times $\tau_\ell^{(\pm)}/\tau_2^{(+)}$ versus multipolarity ℓ . Solid and dashed lines connect the values which correspond to isoscalar and isovector modes of vibrations respectively.

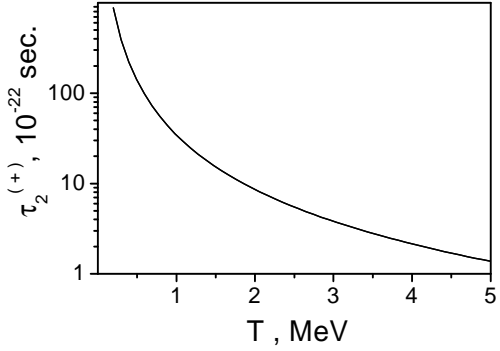


Fig.2: The relaxation time $\tau_2^{(+)}$ in dependence of T at $\hbar\omega = 0$ in units of 10^{-22} sec.

The dependence of the viscosity relaxation time $\tau_{\ell=2}^{(+)}(\omega = 0, T)$ on the temperature is shown in Fig.2. The value of $\alpha_{\ell=2}^{(+)}$ is used from Eq.(38). The temperature dependence arises from smearing out the equilibrium distribution function near the Fermi momentum in heated nuclei. The collisional rates $1/\tau_\ell^{(\pm)}$ given by Eqs.(31) and (37) are quadratic both in temperature and in frequency with the same relationship between the components much as in the zero sound

attenuation factor of heated Fermi liquid within the Landau prescription [9, 14, 16, 22]. The relaxation times $\tau_\ell^{(\pm)}$ depend on frequency ω due to the memory effects in the collision integral.

3 Doorway state mechanism in heated nuclei

The relationship (10) gives the possibility to evaluate the relaxation time in system with weak damping in an independent way from decay rates $\lambda_c^{(\pm)} = \Gamma_c^{(\pm)}/\hbar$. We adopt the following physical notion: $\lambda_c^{(\pm)} = \Gamma_c^{(\pm)}/\hbar$ is the spreading decay rate of the initial state $|\nu_i^{(\pm)}\rangle$ to the final state $|\nu_f^{(\pm)}\rangle$ within first-order approximation of the time-dependent perturbation theory, as given by the golden rule

$$\lambda_c^{(\pm)} \equiv 1/\tau_c^{(\pm)} = \frac{2\pi}{\hbar} \overline{|M^{(\pm)}|^2} \rho_f^{(\pm)}, \quad (39)$$

where $\rho_f^{(\pm)}$ is the density of the available final states.

The quantity $\overline{|M^{(\pm)}|^2}$ is the mean square matrix element for transitions due to residual interaction V_{res}

$$\overline{|M^{(\pm)}|^2} = \overline{|\langle \nu_f^{(\pm)} | V_{res} | \nu_i^{(\pm)} \rangle|^2}, \quad (40)$$

where the line over symbols denotes an average over final states² [31].

The initial state should describe giant collective vibration in heated system at given temperature T . It is taken as a mixture of a collective state (GR) and a thermal state which is approximated by uncorrelated superposition states of $mp - mh$ configuration with m excited particles and holes corresponding to the most probable number of excitons $\bar{n} = 2m$ at given temperature T . The excitation energy of the system is the sum of collective energy $\hbar\omega$ and thermal excitation energy $U = \bar{n}\varepsilon$ with $\varepsilon = \pi^2 T / (12 \ln 2)$ for the average excitation energy per thermal exciton [34]:

$$|\nu_i\rangle = |\{GR\}, \{mpmh\}\rangle, \quad E = \hbar\omega + U, \quad U = aT^2, \quad \bar{n} = 2m = 2gT \ln 2, \quad (41)$$

where the expression for \bar{n} is taken from [34]; $a = \pi^2 g / 6$. The quantity g is the single nucleon state density at the Fermi surface and the same values of g are taken for neutrons and protons.

Next we accept common feature that giant resonance state (GR) is formed by coherent superposition of many (predominantly correlated) one-particle one-hole configurations and due to this fact wave function of initial state can be presented as the sum of wave functions $|\{(m+1)p(m+1)h\}\rangle \equiv |\{n_i\}; k_i\rangle$ of incoherent $(m+1)p - (m+1)h$ configurations with $n_i = 2 + \bar{n}$ excitons, k_i stands for other quantum numbers

$$|\nu_i\rangle = \sum_{k_i} C_{k_i}^{\nu_i} |\{n_i\}; k_i\rangle, \quad (42)$$

where the quantity $C_{k_i}^{\nu_i}$ defines the magnitude of the admixture of different components of quasi-particle eigenstates.

²To simplify the presentation, we will omit in the following the superscript (\pm) and include them only when it is necessary to avoid confusion.

Because of two-body character of the residual interaction V_{res} , the final state can consist of configurations with $n_f = n_i, n_i \pm 2$ excitons. The averaged squared matrix elements $\overline{|M|^2}$ of the transitions to states with fixed number of excitons can be rewritten as

$$\begin{aligned} \overline{|M|^2} &= \sum_{k_i, k'_i} \overline{C_{k_i}^{\nu_i} C_{k'_i}^{\nu_i, *} \langle \{n_f\}; k_f | V_{res} | \{n_i\}; k_i \rangle \langle \{n_i\}; k'_i | V_{res} | \{n_f\}; k_f \rangle} \\ &\sum_{k_i} \overline{|C_{k_i}^{\nu_i}|^2 |\langle \{n_f\}; k_f | V_{res} | \{n_i\}; k_i \rangle|^2} \simeq \mathcal{M}^2(n_i \rightarrow n_f, E), \end{aligned} \quad (43)$$

where $\mathcal{M}^2(n_i \rightarrow n_f, E) = \overline{|\langle \{n_f\} | V_{res} | \{n_i\} \rangle|^2}$ is effective mean square matrix element for transition between incoherent particle-hole states.

This transformation is performed by the use of the following assumption and properties:

- i) The compensation of the binary products of the matrix elements coupling together incoherent exciton states with different values k_i is assumed to take place due to very complicated character of the final state;
- ii) Approximate normalization of the factors $C_{k_i}^{\nu_i}$ is used, $\sum_{k_i} |C_{k_i}^{\nu_i}|^2 \simeq 1$;
- iii) The mean square matrix elements for transitions between incoherent exciton states with different values of the numbers k_i, k_f are taken as equal to the same magnitude $\mathcal{M}^2(n_i \rightarrow n_f, E)$ which is dependent only on numbers of excitons and the total excitation energy. We also assume that effective mean square matrix elements $\mathcal{M}^2(n_i \rightarrow n_f, E)$ for interactions between different kinds of nucleons are equal in magnitude [32–35].

With the use of (43), the collisional relaxation rate λ_c (Eq.(39)) coincides with the particle interactions rate of the exciton model starting from the n_i configuration [34]. The relation (39) for the collisional relaxation time $\tau_c = \tau_c(\omega, T)$ is

$$\begin{aligned} \frac{\hbar}{\tau_c} &= 2\pi \mathcal{M}^2(n_{<} = \bar{n} + 2, E) \rho_c(E) + 2\pi \mathcal{M}^2(n_{<} = \bar{n}, E) \rho_a(E), \\ E &= \hbar\omega + U, \quad U = aT^2, \quad \bar{n} = bT, \quad a = \pi^2 g/6, \quad b = 0.843a, \end{aligned} \quad (44)$$

when processes of creation and of annihilation of the particle-hole pairs are included. The matrix elements for both processes are taken to be determined by the number excitons $n_{<}$ in the simplest state [36] and ρ_c (ρ_a) is the density of the final accessible states corresponding to the pair creation (annihilation).

The transitions to final configuration with $n_f = n_i + 2 \equiv \bar{n} + 4$ dominate at low excitation energies. Using the simplest expression within the exciton model [34] for density of final accessible states, $\rho_c(E) = (g^3/2)(E^2/(n_i + 1))$, the Eq.(44) is given by

$$\frac{\hbar}{\tau_c(\omega, T)} = \pi g^3 \mathcal{M}^2(n_{<} = \bar{n} + 2, E = \hbar\omega + U) \frac{(\hbar\omega + aT^2)^2}{3 + bT}. \quad (45)$$

According to the exciton model studies ([33]- [40]) the effective mean square matrix elements $\mathcal{M}^2(n_i, E)$ is energy-independent at low excitation energies and it is inversely proportional to energy at higher excitations. The energy-independent estimation \mathcal{M}^2 was obtained with the use of the Fermi gas model as [37, 38] $\mathcal{M}^2 = \mathcal{K}_M/A^3$, $\mathcal{K}_M \simeq 15.3 \text{ MeV}^2$, where A is the mass number. The behaviour of collisional relaxation, as given by Eq.(45), with such magnitude of the mean square matrix element agrees with estimation (37) based on kinetic equation approach at low temperatures $T \ll \hbar\omega$.

There are different estimates for the mean square matrix element with dependence on energy and number of excitons [33, 39, 40]. The fulfillment of the condition of equiprobability of all particle-hole configurations is assumed in most of them and therefore they can not be used in the considered case of collective (predominantly 1p-1h) state overlapped with temperature-fixed background particle-hole states. The expression for $\mathcal{M}^2(n, E)$ without assumption on a uniform sharing of the excitation energy E into n excitons was proposed in Ref. [40]:

$$\mathcal{M}^2(n_{<}, E) = \frac{n_{<} + 1}{4} \frac{\mathcal{K}_B}{A^3 E}, \quad (46)$$

where quantity \mathcal{K}_B is not changed with E and n but can be dependent on numbers of protons and neutrons, $\mathcal{K}_B = 190 \text{ MeV}^3$. If this value of \mathcal{M}^2 is employed as the squared intronuclear matrix element, the collisional relaxation time τ_c is a linear function of the collective energy $\hbar\omega$ and thermal energy U . The corresponding expressions (Eqs.(45) and (46)) for the relaxation time have the same form as that one obtained within test particle approach, when collisions were simulated by modeling s-wave scattering between pseudoparticles [41, 42]:

$$\frac{\hbar}{\tau_c(\omega, T)} = \frac{1}{\alpha_e} (\hbar\omega + U), \quad \frac{1}{\alpha_e} = \frac{\mathcal{K}_B \pi}{4} (g/A)^3. \quad (47)$$

The relaxation times given by Eqs.(37) and (47) can agree together at zero temperature if the magnitude of \mathcal{K}_B is equal to the value $\mathcal{K}_0 = (\hbar\omega/\alpha)(A/(g\pi))^3$. Here, $\mathcal{K}_0 \equiv 70.9\hbar\omega/\alpha \simeq 220 \text{ MeV}^3$ for giant isovector dipole resonances in heavy nuclei, when $\hbar\omega \simeq 13 \text{ MeV}$, $g = A/13$ and $\alpha = \tilde{\alpha} = 4.18 \text{ MeV}$. This value of \mathcal{K}_B is rather close to the $\mathcal{K}_B = 190 \text{ MeV}^3$. It means that in cold nuclei the relaxation times for the GDR within doorway state mechanism are not too different from those obtained within the transport approach.

The dependence of the collisional relaxation times $\tau_c(\omega, T)$ on temperature and energy $\hbar\omega$ is demonstrated on Figs.3-6. Solid and dashed lines correspond to the relaxation times $\tau_c(\omega, T)$ within doorway state mechanism with the mean square matrix elements $\mathcal{M}^2(n_{<}, E) \propto 1/E$ and $\mathcal{M}^2(n_{<}, E) = \text{const}$ respectively. Dot-dashed lines correspond to the collisional relaxation times $\tau_{\ell=1}^{(-)}$ within the framework of the transport approach with the value $\alpha_1^{(-)}$ determined by free $n - p$ cross-section (see, Eq.(38)). The factors \mathcal{K}_M and \mathcal{K}_B of the mean square matrix elements are fixed from the condition of the coincidence of the relaxation times $\tau_c(\omega = E_{GDR}/\hbar, T = 0)$ and $\tau_{\ell=1}^{(-)}(\omega = E_{GDR}/\hbar, T = 0)$ in cold nuclei at a frequency corresponding to the GDR energy. The magnitude of this energy is taken as equals to the GDR energy in ^{208}Pb : $E_{GDR} = 13.43 \text{ MeV}$. The values of the relaxation times are given in units of 10^{-22} sec .

Figures 3, 4 show relaxation times at $\omega = 0$ (Fig.3) and $\omega = E_{GDR}/\hbar$ (Fig.4) in relation to the temperature.

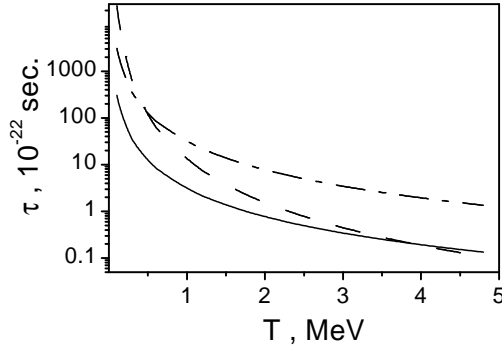


Fig.3: The relaxation time τ in dependence of T at $\hbar\omega = 0$ in units of 10^{-22} sec. Solid and dashed lines denote τ within doorway state mechanism with $M^2 \propto 1/E$ and $M^2 \propto const.$ Dot-dashed line is τ within the framework of the transport approach.

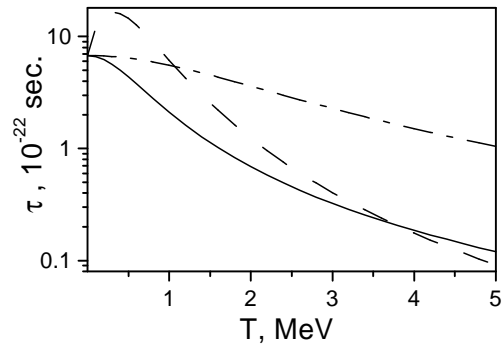


Fig.4: The relaxation time in dependence of T at $\hbar\omega = E_{GDR}$ in units of 10^{-22} sec. Notations are the same as in Fig. 3.

Figures 5, 6 demonstrate dependence of the relaxation times on energy $\epsilon = \hbar\omega$ in cold (Fig.3, $T = 0$) and heated (Fig.4, $T = 2$ MeV) nuclei.

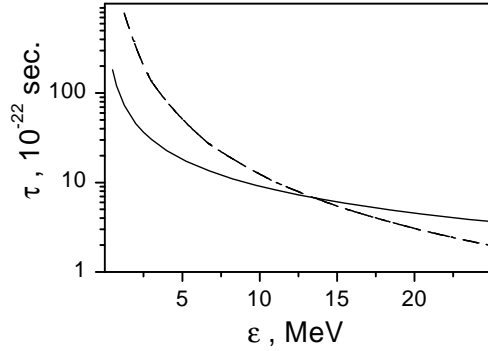


Fig.5: The relaxation time in dependence of $\epsilon \equiv \hbar\omega$ at $T = 0$ in units of 10^{-22} sec. Notations are the same as in Fig. 3.

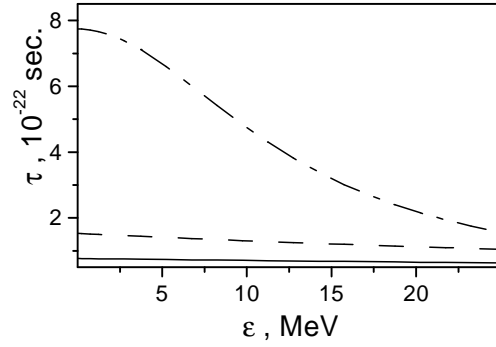


Fig.6: The relaxation time in dependence of $\epsilon \equiv \hbar\omega$ at $T = 2$ in units of 10^{-22} sec. Notations are the same as in Fig. 3.

The collective relaxation times in heated nuclei can be approximately presented by the expression

$$1/\tau_c(\omega, T) = q_1[\omega^\beta + q_2 T^2]^\gamma / [q_3 + q_4 T]^\delta, \quad (48)$$

where q_j are some constants and the exponents are the functions of frequency and temperature and they are $\beta = 2, \gamma = 1, \delta = 0$ in the transport method; the β, γ are changed from 2 to 1 and δ varies from 1 to 0 with growing of the excitation energy in the doorway state approach with allowance for pair creation.

4 Results and conclusions

The retardation and temperature effects in two-body collisions in heated Fermi systems were studied. An expression for non-Markovian collision integral of the Landau-Vlasov transport equation was obtained in a form which allows for reaching the local equilibrium in system. It is based on the Kadanoff- Baym equations for Green functions within the limit of small retardations.

The expressions for collisional relaxation times of the collective vibration in heated nuclei are derived with the use of the non-Markovian collision integral as well as of decay rates of exciton model. The relaxation times depend on frequency of the collective vibrations and the temperature. The temperature dependence arises from smearing out the equilibrium distribution function near the Fermi momentum in heated nuclei. The frequency dependence results from the retardation (memory) effects in the collisions. Analytical expressions for relaxation times of the isoscalar and isovector modes of the collective motion are derived in the case of the energy independent isotropic cross-sections in the two-body collisions. The relaxation times vary rather slowly with multipolarity of the Fermi surface distortions governed by collective motion and two-body collisions. It enables to use the approximation of the relaxation time ansatz for the collision integral. The relaxation times depend on type of collective motion mode, such as the lifetimes of the particle-hole configurations in two-component exciton model of the Ref. [32].

New approach for calculations of the collision relaxation time in heated nuclei are proposed using the formulae for the transition rates of the intranuclear transition between thermal state with collective vibrations and incoherent particle-hole configurations. This method leads to the same results as the transport approach in the case of low temperatures and energy-independent mean square matrix element of particle-particle collisions. It makes possible to take into account the energy dependence of the in-medium cross-sections in a simple phenomenological way by the use of the parametrization of the mean square matrix element \mathcal{M}^2 of particle-particle collisions from exciton model of nuclear reactions. The dependence of the matrix element \mathcal{M}^2 (i.e., the in-medium cross-section) on energy leads to non-quadratic dependence of the relaxation times on temperature and collective vibration frequency.

Acknowledgments This work is supported in part by the IAEA(Vienna) under contract 302-F4-UKR-11567.

References

- [1] G.A. Brooker, J. Sykes: *Ann. of Phys* **61** (1970) 387
- [2] V.M. Kolomietz, A.G. Magner, V.A. Plujko: *Z. Phys. A* **345** (1993) 131
- [3] M.M. Abu-Samreh, H.S. Koehler: *Nucl. Phys. A* **552** (1993) 101
- [4] V.A. Plujko: *Acta Phys. Pol. B* **30** (1999) 1383
- [5] J.P. Blaizot: *Phys. Rep.* **64** (1980) 171
- [6] P. Ring, P. Schuck: *The nuclear many-body problem*, Springer, Berlin 1980
- [7] H. Krivine, J. Treiner, O. Bohigas: *Nucl. Phys. A* **336** (1980) 155
- [8] A.B. Migdal, A.A. Lushnikov, D. F. Zaretsky: *Nucl. Phys.* **66** (1965) 193
- [9] V.M. Kolomietz, V.A. Plujko, S. Shlomo: *Phys. Rev. C* **52** (1995) 2480
- [10] V.M. Kolomietz, V.A. Plujko, S. Shlomo: *Phys. Rev. C* **54** (1996) 3014

- [11] V.A.Plujko: *Acta Phys. Pol. B* **31** (2000) 435
- [12] V.A. Plujko: in *Proceedings of the 9th Inter. Conf. Nucl. Reaction Mechanisms*, Varenna, June 5-9, 2000 (Ed. E. Gadioli). Univ. Studi Milano, Suppl. N. 115, 2000, p.113,
- [13] G. Bertsch: *Z. Phys. A* **289** (1978) 103
- [14] S. Ayik, O. Yilmaz, A. Gokalp, P. Schuck: *Phys. Rev. C* **58** (1998) 1594
- [15] H. Reinhardt, P.G. Reinhard, K. Goeke: *Phys. Lett. B* **151** (1985) 1985
- [16] S. Ayik, D. Boiley: *Phys. Lett. B* **276** (1992) 263 ; *Phys. Lett. B* **284** (1992) 482E
- [17] V.M. Kolomietz, A.G. Magner, V.A. Plujko: *Yad. Fiz.* **55** (1992) 2061 [*Sov. J. Nucl. Phys.* **55** (1992) 1143]
- [18] V.M. Kolomietz, A.G. Magner, V.A. Plujko: *Yad. Fiz.* **56** (1993) 110 [*Phys. At. Nucl.* **56** (1993) 209]; *Nucl. Phys. A* **545** (1992) 99c
- [19] V.M. Kolomietz, V.A. Plujko: *Yad. Fiz.* **57** (1994) 992 [*Phys. At. Nucl.* **57** (1994) 931]
- [20] K. Morawetz, U. Furmann: *RIEKEN Review No.* **23** (1999) 90 ;
LANL e-print: (<http://xxx.lanl.gov/abs/nucl-th/9907012>)
- [21] P.G. Reinhard, H.L. Yadov, C.Toepffer: *Nucl. Phys. A* **458** (1986) 301
- [22] L.D. Landau: *Zh. Eksp. Teor. Fiz.* **32** (1957) 59 [*Sov. Phys. JETP* **5** (1957) 101]
- [23] G. Baym, Ch. Pethick: *Landau Fermi-liquid theory*, John Wiley and Sons, New York 1991
- [24] E.M. Lifshitz, L. Pitajevski: *Physical Kinetics*, Nauka, Moscow 1979
- [25] S. Ayik, M.Belkacem, A.Bonasera: *Phys. Rev. C* **51** (1995) 611
- [26] M. Belkacem, S. Ayik, A. Bonasera: *Phys. Rev. C* **52** (1995) 2499
- [27] A.A. Abrikosov, I.M. Khalatnikov: *Usp. Fiz. Nauk.* **66** (1958) 177 [*Rep. Prog. Phys.* **22** (1959) 329]
- [28] J. Sykes, G.A. Brooker: *Ann. of Phys.* **56** (1970) 1
- [29] J. Vogel, E. Vogel, C. Toepffer: *Ann. of Phys.* **164** (1985) 463
- [30] K. Ando, A. Ikeda, G. Holzwarth: *Z. Phys. A* **310** (1983) 223
- [31] M. Danos, W. Greiner: *Phys. Rev.* **138** (1965) 876
- [32] J. Dobes, E.Betak: *Z. Phys. A* **310** (1983) 329
- [33] C. Kalbach: *Phys. Phys. C* **33** (1986) 818
- [34] E. Gadioli, P.E. Hodgson: *Pre-Equilibrium Nuclear Reactions*, Clarendon Press, Oxford 1992
- [35] E. Betak, P.E. Hodgson: *Rep. Prog. Phys.* **61** (1998) 483
- [36] C. Kalbach: *Z. Phys. A* **287** (1978) 319
- [37] G.M. Braga-Marcazzan, E. Gadioli-Erba, L. Milazzo-Colli, P.G. Sona: *Phys. Rev. C* **6** (1972) 1398
- [38] E. Gadioli, L. Milazzo-Colli: *Lect. Not. Phys.* **22** (1973) 84
- [39] C. Kalbach: *Z. Phys. A* **287** (1978) 319
- [40] P. Oblozinsky: *Phys. Phys. C* **35** (1987) 407
- [41] A. Bonasera, G.F. Burgio, M. Di Toro: *Phys. Lett. B* **221** (1989) 233
- [42] A. Smerzi, A. Bonasera, M. Di Toro: *Phys. Rev. C* **44** (1991) R1713