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SHELLS, ORBITS AND TRANSPORT COEFFICIENTS OF THE NUCLEAR COLLECTIVE DYNAMICS

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To the 75-th anniversary of birth of Vilen Mitrofanovich Strutinsky

The Strutinsky shell correction method is applied to the nuclear collective dynamics within the periodic orbit theory extended to dissipative phenomena through some effective residue interactions. The macroscopic limit of the semiclassical transport coefficients with the known wall formula for the friction is obtained. The shell corrections to the stiffness, the inertia and the friction parameters are calculated analytically as functions of the particle number and temperature for the low-energy excitations of heavy nuclei. It is shown that they are approximately proportional to the free-energy shell corrections, with the same sign for the inertia and the opposite one for the stiffness and the friction. The shell oscillating components of the transport coefficients disappear exponentially at temperatures larger than the ones of the free energy. The reduced friction and the energy excitation estimations are in agreement with some experimental data and theoretical calculations.

1. Introduction

Low energetic nuclear collective dynamics can be treated within the Landau-Migdal theory by using the notion of elementary excitations as defined through the concept of quasi-particles [1]. For practical calculations based on phenomenological access it would be highly desirable to separate the quasi-particle contributions from the background ones and parametrize the latter phenomenologically through some macroscopic models.

As is well-known, such a goal has been successfully met by Strutinsky in connection to *static* nuclear properties like the total energy (or the free energy for heated nuclei), for instance for fission barrier calculations [2, 3]. There, the shell energy (or free energy) plays the role of the quasi-particle component in sense of the Landau-Migdal theory. The phenomenological macroscopic part is given by the nuclear liquid drop energy (liquid-drop free energy, respectively).

A first attempt to generalize these ideas to the problems of the collective dynamics has been made in [4] in the so called gas-liquid (or "liquid-particle") model for the time evolution of the one-body density matrix starting from von Neumann equation. Practically, it was reduced to the lowest order of the semiclassical approximation ($\hbar \to 0$) for the quasi-particle part of the density matrix [5]. The latter was defined as the solution of the Landau-Vlasov equation, and applied for description of the giant resonances [6, 7] and temperature dependence of the macroscopic transport coefficients, such as the stiffness, the inertia and the friction parameters within the response function theory, see [8 – 10]. By employing this response theory, more general and realistic proposals which include the quantum shell effects, as considered for instance within the Periodic Orbit Theory (POT) [11–14], for study of the quasi-particle motion at low energy excitations were suggested in [15]. In this way, the collective variables were introduced explicitly through a deformed single-particle field and parameterization of the nuclear excitations by making use of the adequate collective response functions.

In the present paper we wish to develop these ideas directly for observable quantities like the transport coefficients related to the collective response function. They are more simple functions of the particle number and temperature in close analogy to the free energy. Therefore, we hope to more easy and stright extension of the Strutinsky shell correction method to the collective dynamics in the case of the low energy excitations of enough heavy nuclei. We will

start with a general discussion of the response function formalism following basically [9, 16] in Section 2. In Section 3 we will explain the semiclassical calculation of the quasi-particle (shell) response functions in the POT by presenting more complete solution than in [15]. The main focus will be concentrated to discussion of the transport coefficients including the corresponding shell ingradients along with more specific macroscopic parts. These results for the stiffness and the inertia in the case of the low-energy collective quadrupole excitations depending on particle number and temperature are discussed in comparison with experimental data on fission and collective modes as well as with the traditional microscopic methods [3, 9, 16–21] in Section 4. Summary will end the paper, see section 5.

2. Response Function Theory

Many-body excitations are conveniently described in terms of the response to an external perturbation $V_{ext} = q_{ext}(t)\,\hat{F}$, $(q_{ext}(t) = q^{\omega}_{ext}e^{-i\omega t})$ with $\hat{F}$ being some one-body operator. Its quantal average $\langle\hat{F}\rangle_t$, at time t can be calculated through the Fourier transform $\langle\hat{F}\rangle_\omega$ obtained within the linear response theory as [8, 9]

$$\langle\hat{F}\rangle_\omega = -\chi^{coll}_{FF}(\omega)\,q^{\omega}_{ext}, \qquad \hat{F} = \left(\frac{\partial H(Q)}{\partial Q}\right)_{Q=0} = \left(\frac{\partial V}{\partial Q}\right)_{Q=0}. \tag{1}$$

Here, $\chi^{coll}_{FF}(\omega)$ is the collective response function. The Hamiltonian $H(Q)$ at $q_{ext} = 0$ depends on a collective variable Q through the potential $V(Q)$, $H(Q) = K + V(Q)$, where K is the particle kinetic energy. The variable Q can be defined as a time-dependent deformation parameter of the nuclear surface, say, $R(Q) = R\,(1 + Q(t)\,Y_{\Lambda 0}(\hat{\mathbf{r}}))$, where $Q(t) = Q_\omega e^{-i\omega t}$, $\hat{\mathbf{r}} = \mathbf{r}/r$ is the radial unit vector at the nuclear surface for multipole vibrations of potential $V(Q)$ near the spherical shape $r = R$ (we omit everywhere the unperturbed quantities). With help of the consistency condition,

$$\langle\hat{F}\rangle_\omega = \kappa\,Q_\omega, \qquad \kappa = -\chi_{FF}(0) - \left(\frac{\partial^2 F}{\partial Q^2}\right)_{Q=0}, \tag{2}$$

where κ is the coupling constant, F the free energy of the system, the collective response $\chi^{coll}_{FF}(\omega)$ can be expressed in terms of the "intrinsic" response function (8) as

$$\chi^{coll}_{FF}(\omega) = \kappa\frac{\chi_{FF}(\omega)}{\chi_{FF}(\omega) + \kappa}, \qquad \chi_{FF}(\omega) = -\langle\hat{F}\rangle_\omega/\,(Q_\omega + q^{\omega}_{ext})\,. \tag{3}$$

For low excitation energies we may use the Taylor expansion of the $\chi_{FF}(\omega)$ in ω at $\omega = 0$,

$$\chi_{FF}(\omega) \equiv \chi(\omega) = \chi(0) + i\gamma(0)\omega + B(0)\omega^2 + ..., \qquad \chi(0) = -\kappa - C(0), \tag{4}$$

where $C(0) = (\partial^2 F/\partial Q^2)_{Q=0}$, $\gamma(0)$ and $B(0)$ are the so called the stiffness, the friction and the inertia (transport) coefficients in the "zero frequency limit". As explained in [9, 17], it is convinient to consider the response function $\chi^{coll}_{QQ}(\omega)$ in the Q-mode, which can be now written in an inversed form as

$$\frac{1}{\chi^{coll}_{QQ}(\omega)} = \frac{1}{\chi_{QQ}(\omega)} + \kappa, \tag{5}$$

with $\chi_{QQ}(\omega) = \chi_{FF}(\omega)/\kappa^2$. The inversed collective response functions for low frequencies can be expanded then in ω at $\omega = 0$ and approximated by the response function of a damped harmonic oscillator with parameters C, B and γ,

$$\frac{1}{\chi_{QQ}^{coll}(\omega)} \approx -B\omega^2 - i\gamma\omega + C. \tag{6}$$

They can be called as the collective transport coefficients which are related to the corresponding "intrinsic" parameters $C(0)$, $B(0)$ and $\gamma(0)$ of Eq. (4) in the "zero frequency limit" mentioned above as [9]

$$C = \left[1 + \frac{C(0)}{\chi(0)}\right] C(0), \quad \gamma = \left[1 + \frac{C(0)}{\chi(0)}\right]^2 \gamma(0), \quad B = \left[1 + \frac{C(0)}{\chi(0)}\right]^2 \left(B(0) + \frac{\gamma(0)^2}{\chi(0)}\right). \tag{7}$$

The poles of the collective response function $\chi_{QQ}^{coll}(\omega)$ (6) are the solutions of the known Langevin equation $-B\omega^2 - i\gamma\omega + C = 0$.

The "intrinsic" response $\chi_{FF}(\omega)$ in Eqs. (3) can be expressed in terms of the "dressed" one-body Green function $G_\Sigma(e)$ as (see Refs.[9, 15, 16, 22, 23])

$$\begin{aligned} \chi_{FF}(\omega) &= \frac{1}{\pi} \int_0^\infty \mathrm{d}e\, n(e) \int \mathrm{d}\mathbf{r}_1 \int \mathrm{d}\mathbf{r}_2\, \hat{F}(\mathbf{r}_1)\, \hat{F}(\mathbf{r}_2) \times \\ &\times \left[\bar{G}_\Sigma(\mathbf{r}_1, \mathbf{r}_2, e - \hbar\omega) + G_\Sigma(\mathbf{r}_1, \mathbf{r}_2, e + \hbar\omega)\right] \mathrm{Im}\, G_\Sigma(\mathbf{r}_1, \mathbf{r}_2, e), \end{aligned} \tag{8}$$

where $n(e)$ is the occupation number at the energy e. For $G_\Sigma(\mathbf{r}_1, \mathbf{r}_2, e)$ we may use the energy spectral representation

$$G_\Sigma(e) = \sum_\nu \frac{a_\nu |\nu\rangle\langle\nu|}{e - \mathrm{Re}\,\Sigma_\nu + i\mathrm{Im}\,\Sigma_\nu}. \tag{9}$$

The poles of this Green function are laying near the real axis for enough small $\mathrm{Im}\,\Sigma_\nu$ ($\mathrm{Im}\,\Sigma_\nu << e_F$, e_F is the Fermi energy), and roughly $\mathrm{Re}\,\Sigma_\nu \approx e_\nu$, $a_\nu \approx 1$. Then, e_ν and $|\nu\rangle$ are the eigenvalues and eigenfunctions of the corresponding quantum problem for an unperturbed Hamiltonian $H(Q)$ at $Q = 0$ (bar above G_Σ in Eq. (8) means the complex conjugation).

Usually, the POT is considered as a mean field approximation with $\mathrm{Im}\,\Sigma_\nu \to +0$. We shall try to extend this theory to a small (in sense mentioned above) but finite damping related to the residue interaction by using a simple model for $\mathrm{Im}\,\Sigma_\nu$ based mainly on the Landau Fermi-liquid theory, see for instance [9],

$$\mathrm{Im}\,\Sigma_\nu \approx \mathrm{Im}\,\Sigma(\omega, T) = \frac{1}{2}\Gamma(\omega, T) = \frac{1}{2\,\Gamma_0} \frac{(\hbar\omega)^2 + \pi^2 T^2}{1 + \left[(\hbar\omega)^2 + \pi^2 T^2\right]/c^2}, \tag{10}$$

Such residue interaction can be associated with the two-body collisional damping. In our calculations we use the traditional consistent parameters of the response function theory, $\Gamma_0 = 33.3$ MeV presenting matrix elements of the residue interaction, $c = 20$ MeV for a so-called "cut-off" parameter, see [9, 17–19]. As seen from Eq. (9) with the approximation (10), one formally can consider in further derivations the complex energies $e \to e + i\mathrm{Im}\,\Sigma(\omega, T)$ like for averaging of the level density with the Lorentzian weight function in a folding integration [12].

3. Semiclassical Response

Expression (8) is ready for a semiclassical calculation in terms of the Feyman path integral representation of one-body Green function by the stationary phase method [11, 13, 15]. According to Ref. [11, 13] the semiclassical version of the Green function is given by:

$$G(\mathbf{r}_1, \mathbf{r}_2, e) = \sum_\alpha \mathcal{A}_\alpha(\mathbf{r}_1, \mathbf{r}_2, e) \exp\left[\frac{i}{\hbar} S_\alpha(\mathbf{r}_1, \mathbf{r}_2, e) - i\frac{\pi}{2}\mu_\alpha\right]. \tag{11}$$

The index α covers all classical paths inside the potential well $V(Q)$ at $Q = 0$, which connect the two spatial points $\mathbf{r}_1$ and $\mathbf{r}_2$ there for a given energy e. The $\mathcal{A}_\alpha$ is a Green function amplitude depending on stability of the trajectory α. The S_α is the classical action along such trajectory α, $S_\alpha = \int_{\mathbf{r}_1}^{\mathbf{r}_2} \mathrm{d}l\, p(\mathbf{r}, e)$, where $p = |\mathbf{p}| = \sqrt{2M(e - V)}$, $\mathbf{p}$ is the momentum, M the mass of nucleon. The μ_α in Eq. (11) is the phase related to the Maslov index. The latter can be determined by the number of the caustic and turning points [24].

Among all possible trajectories we single out α_0 which directly connects $\mathbf{r}_1$ and $\mathbf{r}_2$ without reflections from the potential well edge in intermidiate points. For G this leads to the separation $G = G_0 + G_{osc}$, see [11–14]. The first term G_0 for $\mathbf{r}_1$ close to $\mathbf{r}_2$ writes:

$$G_0(\mathbf{r}_1, \mathbf{r}_2, e) = -\frac{M}{2\pi\hbar^2 s} e^{iks}, \qquad s = |\mathbf{r}_1 - \mathbf{r}_2|, \qquad k = \frac{p}{\hbar}. \tag{12}$$

Its trace corresponds to the Thomas-Fermi model for the level density, see [13]. The "oscillating" part G_{osc} in Green function G is determined by other trajectories $\alpha \neq \alpha_0$.

The Thomas-Fermi friction $\gamma_{ThF}(0)$ and inertia $B_{ThF}(0)$ in the "zero frequency limit" can be found from the "intrinsic" response function $\chi_{00}(\omega)$, see Eq. (8) with $G = G_0$ Eqs. (12), by using expansion (4). Due to the δ-function of $\hat{F} = -V_0 R\delta(r - R)Y_{\Lambda 0}$ the spatial points $\mathbf{r}_1$ and $\mathbf{r}_2$ in Eq. (8) are taken at the potential surface. The boundary condition there can be applied for the infinite square well potential ($V_0 \to \infty$), for instance, like in [15, 17, 22]. In order to obtain now the leading term of the friction and the inertia in the semiclassical parameter $k_F R >> 1$ ($k_F = p_F/\hbar$, p_F is the Fermi momentum equal approximately to the particle momentum p at the energy $e = \lambda \approx e_F$, λ the chemical potential) we note that small distances s between $\mathbf{r}_1$ and $\mathbf{r}_2$ give the main contribution into the integrals over the potential surface in $\chi_{00}(\omega)$ of Eq. (8), according to (12). By making use of such local approximation $s \to 0$ for the double normal derivatives of Green function $\mathrm{Im}\, G_0$ of Eqs. (12) at the potential surface which come from the boundary conditions one has $\partial^2 \mathrm{Im}\, G_0/\partial r_1 \partial r_2 \approx -(Mk^3/2\pi\hbar^2)\, j_1(ks)/(ks)$ at $r_i = R$ ($i = 1, 2$), where $j_1(x)$ is the spherical Bessel function. This expression becomes exact for the semi-infinite fermionic system with a flat potential surface, see [22, 25]. Such approximation is similar to the δ-function of the distance s between the two surface points $\mathbf{r}_1$ and $\mathbf{r}_2$. In order to remove quantum fluctuations we consider now averaging of $\gamma_{ThF}(0)$ and $B_{ThF}(0)$ over the particle number or the energies, i.e. over the single common parameter $k_F R$. Exchanging this averaging and the surface integrations over $\mathbf{r}_1$ and $\mathbf{r}_2$ we note that the double surface integrand essentially depends on $k_F s$ such that large $k_F R$ corresponds to the dominating contribution of small s as mentioned above. The averaging of the integrand in $k_F R$ can be considered as folding integral with a Gaussian weight function, for instance. This integral over $k_F R$ can be naturally transformed to the surface variable $k_F s$. The averaging over $k_F s$ is the same as the one over the variable $\mathbf{r}_1$ at the potential surface. That leads to $\mathbf{r}_1$-independence of the internal integral over $\mathbf{r}_2$ with respect to the external integration variable $\mathbf{r}_1$ on the surface in Eq. (8). Formally, for the surface integration over $\mathbf{r}_2$ we may equivalently use the spherical coordinate system with the polar axis crossing $\mathbf{r}_1$ at the potential surface. Such statistical arguments are

similar to the ones used in [17, 22, 25, 26] in the derivation of the so-called "wall formula" for the friction. In this way one obtains

$$\gamma_{ThF}(0) = \left(\frac{\partial \chi''_{00}}{\partial \omega}\right)_{\omega=0} = \frac{\hbar\,(k_F R)^4}{4\pi^2}\left[1+\frac{\pi^2}{3}\left(\frac{T}{\lambda}\right)^2\right]\left(1-\frac{\pi\Gamma(0,T)}{3\,\lambda}\right), \tag{13}$$

$$B_{ThF}(0) = \frac{1}{2}\left(\frac{\partial^2 \chi'_{00}}{\partial \omega^2}\right)_{\omega=0} = \frac{M\,R^2\,(k_F\,R)^2}{2\,\pi}\,B_\Gamma\left(\frac{\Gamma(0,T)}{4\lambda}\right), \tag{14}$$

where $\chi = \chi' + i\chi''$,

$$B_\Gamma(x) = \frac{2}{\pi}\left[\mathrm{arctg}(1/x) - \frac{x}{2}\ln\left(1+\frac{1}{x^2}\right)\right], \qquad B_\Gamma(x) \to 1 + \frac{2}{\pi}\,x\ln\left(\frac{x}{e}\right) \tag{15}$$

at $x \to 0$. The last brackets in (13) and (14) account for small corrections linear in the damping "width" $\Gamma(\omega = 0, T)$ of Eq. (10) for simplicity, $\Gamma/4\lambda \approx \Gamma/4e_F << 1$ (these damping factors can be presented exactly in terms of elementary functions but were ommited because of rather lengthy expressions). We took into account explicitly the spin degeneracy factor 2. Note that the friction $\gamma_{ThF}(0)$ (13) without small Γ corrections is exactly the same as obtained in [17, 22, 25, 26]. The next terms in expansion of $\gamma_{ThF}(0)$ in $1/k_F R$ are relatively proportional $1/(k_F R)^2$. Generally speaking, the expansions of smooth (macroscopic) $\gamma(0)$ and $B(0)$ are really in small parameter $1/(k_F R)^2$ related to $\hbar^2$ corrections of the extented Thomas-Fermi model. Note also that we found that the Thomas-Fermi friction $\gamma_{ThF}(0)$ is exactly identical to the wall formula for the case of the semi-infinite potential well with the flat surface without the additional statistical asumption concerning averaging in $k_F R$ or relatively, in the surface integration variable $\mathbf{r}_1$ in agreement with the results obtained in [22, 25, 26] as mentioned above. The contributions of other trajectories α with reflections from the potential surface in the intermidiate points between $\mathbf{r}_1$ and $\mathbf{r}_2$ coming from G_{osc} are discussed below. Note that within the same statistical approximation one obtains also the mass parameter $B_{ThF}(0)$. Its main semiclassical term (14) is of the order of $1/k_F R \sim A^{-1/3}$ with respect to the irrotational-flow inertia $B_{irrot} \propto MR^2(k_F R)^3 \sim MR^2 A$ in the hydrodynamic model (A is the particle number in nucleus).

For the macroscopic part of the stiffness $\tilde{C}(0)$ related to the smooth (liquid-drop) free energy $\tilde{F}$, $\tilde{C}(0) = \partial^2\tilde{F}/\partial Q^2$ at $Q = 0$, it is obvious to use the liquid drop (hydrodynamic) value C_{LD}, $\tilde{C}(0) = C_{LD} = b_s A^{2/3}(\Lambda-1)(\Lambda+2)/4\pi$ where b_s the energy surface constant, and for smooth coupling constant too. It is not so obvious for a smooth macroscopic inertia $\tilde{B}(0)$ and a friction $\tilde{\gamma}(0)$ though many authors use the irrotational-flow inertia B_{irrot} and the hydrodynamic friction of the frequent collision regime γ_{HD} at least for comparison of different approaches, see for instance, Ref. [10, 26]. We are going to compare these two approaches for the macroscopic components of the friction and the inertia within the POT, the local "wall" statistical Thomas-Fermi (13), (14) and the hydrodynamical approximations "in the zero frequency limit" mentioned above.

We have written down all these well-known steps because they serve as starting point for our separation into smooth parts and deviations according to the quasi-particle or the shell effects. For this purpose we will first isolate quantities which are determined by excitations near the Fermi level from the ones which can be replaced by statistical (macroscopic) averages. The latter shall be defined in close analogy to Strutinsky's original ideas [2, 3].

Applying now the averaging procedure in particle number through the occupation numbers in spurious of the Strutinsky shell correction method [2] one has

$$n = \tilde{n} + \delta n \quad \to \quad F = \tilde{F} + \delta F, \qquad \chi(\omega) = \tilde{\chi}(\omega) + \delta\chi(\omega), \qquad \kappa = \tilde{k} + \delta\kappa \tag{16}$$

etc. Here $\tilde{\kappa}$ is the macroscopic part of the coupling constant which can be evaluated in the nuclear effective surface approximation [5, 7], $\tilde{\kappa} = -Kb_s r_0^5 A^{4/3}/54\beta$, K is the incompressibility, $K = 6\lambda(1+F_0) \approx 200\,MeV$, $\lambda \approx e_F = 40\,MeV$ for the chemical potential, F_0 the Landau interaction constant ($F_0 \approx -0.2$). The quantity b_s in (16) stands for the surface constant, part of the second equation in Eqs. (16), $b_s \approx 17\,MeV$, $r_0 = (3/4\pi\tilde{\rho})^{1/3} \approx 1.2fm$, $\tilde{\rho}$ the average (nuclear-matter) particle density, β is proportional to the surface coefficient b_s, $\beta \approx 90$ MeV·fm^5, see Ref. [10]. All numbers for the constants mentioned above and used below in the specific calculations are typical for the physics of heavy nuclei.

By substituting the "zero frequency" expansion (4) into Eq. (5) and making use of the last two equations in Eqs. (16) one easy can express $\chi_{QQ}(\omega)$ and $\chi^{coll}_{QQ}(\omega)$ through the "intrinsic" transport coefficients $C(0)$, $B(0)$ and $\gamma(0)$, which can be presented in the usual splitting form of Eqs. (16) as

$$C(0) = \tilde{C}(0) + \delta C(0), \qquad B(0) = \tilde{B}(0) + \delta B(0), \qquad \gamma(0) = \tilde{\gamma}(0) + \delta\gamma(0). \tag{17}$$

Note that the correction $C(0)/\chi(0) \approx -C(0)/\tilde{\kappa}$ in Eqs. (7) is always small for the standard nuclear parameters mentioned above and similar splitting for the collective stiffness C and friction γ is valid. For the case of the stiffness $C(0)$ or C such shell correction splitting is obvious because of the definition of C through the free energy F and the splitting of the shell-correction method for F mentioned in Eqs. (16), though it is not so obvious for the friction. However, the collective inertia B and the corresponding response function $\chi^{coll}_{QQ}(\omega)$ (6) are non-linear with respect to the shell components of the transport coefficients and coupling constant due to its friction dependent component in B of Eqs. (7). Note also that the "intrinsic" inertia $B(0)$, which is helpful but having no direct physical meaning, can be negative at finite temperature in contrast to the collective mass parameter B of Eqs. (7) which becomes positive because of the positive dominating friction dependent correction and $\chi(0) > 0$ in line of the microscopic quantum calculations discussed in [9]. Note also that for the macroscopic collective response functions one has the relation similar to Eq. (5) by putting there zero all shell corrections. A similar relation of the macroscopic collective transport coefficients and the response function to the macroscopic parts $\tilde{C}(0)$, $\tilde{B}(0)$ and $\tilde{\gamma}(0)$ of the "intrinsic" transport coefficients of Eqs. (17) can easily be found. For smooth transport coefficients $\tilde{\gamma}$ and $\tilde{B}$ one may test one of the macroscopic models such as the "wall" Thomas-Fermi (13), (14) or hydrodynamic approaches. In both these approximations the static smooth stiffness coefficient $\tilde{C}(0)$ and coupling constant $\tilde{\kappa}$ (see below) can be derived within the effective surface approximation of the liquid-drop model [5, 7]. They are related to a smooth free energy $\tilde{F}$ considered usually as the liquid-drop one, according to the shell correction method.

The shell component of the total free energy δF can be written as

$$\delta F = \sum_\beta \delta F_\beta, \qquad \delta F_\beta = \left(\frac{\hbar}{m\tau_\beta}\right)^2 \Phi(z_\beta)\, g^{(\beta)}_{osc}(\lambda), \tag{18}$$

where

$$\Phi(z) = \frac{z}{\sinh(z)}, \qquad z_\beta = \frac{\pi\tau_\beta T}{\hbar}. \tag{19}$$

The sum runs the isolated periodic orbits β or families, $g^{(\beta)}_{osc}(\lambda)$ is the β-component of the shell density $g_{osc}(e)$, $g_{osc}(e) = \sum_\beta g^{(\beta)}_{osc}(e)$ at the chemical potential $e = \lambda$, τ_β the time of particle motion along the periodic orbit β at this energy e, T the temperature. The temperature cut-off factor (see Eqs. (19)) decreasing exponentially with T appears here and below through the derivatives of the Fermi distribution $n(x) = 1/(1+e^x)$ with $x = (e-\lambda)/T$ after integration over the energy e by parts.

Substituting Eqs. (16) and $G = G_0 + G_{osc}$ for the Green function G into Eq. (8) for the intrinsic response $\chi(\omega)$ one gets a sum of the two leading terms for its shell component $\delta\chi(\omega)$, $\delta\chi(\omega) \approx \delta\chi_{10}(\omega) + \delta\chi_{11}(\omega)$, where

$$\delta\chi_{1,\{{}^{0}_{1}\}}(\omega) = \frac{1}{\pi}\int_0^\infty \mathrm{d}e\,\delta n(e)\int d\mathbf{r}_1\int d\mathbf{r}_2\,\hat{F}(\mathbf{r}_1)\,\hat{F}(\mathbf{r}_2)\times$$
$$\times\left[\bar{G}_{osc}(\mathbf{r}_1,\mathbf{r}_2,e-\hbar\omega)+G_{osc}(\mathbf{r}_1,\mathbf{r}_2,e+\hbar\omega)\right]\mathrm{Im}\,G_{\{{}^{0}_{osc}\}}(\mathbf{r}_1,\mathbf{r}_2,e)\,. \qquad (20)$$

The component $\delta\chi_{11}(\omega)$ was derived in Ref. [15].

More explicit analytical expressions for $\delta\chi_{10}(\omega)$ can be obtained for the quadrupole vibration near the spherical shape of the infinitely deep square-well potential. By making use of the radial $\delta(r-R(Q))$-function form for the $\hat{F}(r)$ operator of Eqs. (1) and the boundary conditions at this potential surface, see [15, 17, 22, 23], we find the leading semiclassical approximation for the normal derivatives $\partial^2\mathrm{Im}G_0/\partial r_1\partial r_2$ at the potential surface, which are just the radial ones with the same sharp bell-like shape as function of $k_F s$, see just before Eqs. (13),(14).

We transform now the integral over the rest tangent-to-surface (angle) variables of $\mathbf{r}_2$ to the ones of $\mathbf{r}_1 - \mathbf{r}_2$. The main contribution into the integral over these angle variables can be evaluated by the stationary phase method like in [15]. However, the integration is simplified because the closing of orbits $\mathbf{r}_2 \to r_1$ $(s \to 0)$ means the periodicity condition, and we get identically such integral as a sum over closed periodic orbits connecting these two points without change of the order in $\hbar$. Therefore, the external integration over the angle variables of $\mathbf{r}_1$ is simply reduced to the spatial region occupied by the periodic orbit families. Expanding now the exponent phases as function of the energy e to the first order in $e-\lambda$, and then, in $\hbar\omega/\lambda$ we may take analytically the integral over e in (20) like in [15]. Finally, one obtains

$$\delta\chi_{10}(\omega) = -\frac{5}{32}\,f_\Lambda\,(k_F R)^3\sum_\beta a\,m\sin\phi\;\delta F_\beta\;e^{i\omega t_\beta}, \qquad (21)$$

where

$$f_\Lambda = \int_0^1 \mathrm{d}x\,\sqrt{1-x^2}\,P_\Lambda^2(x) = \left\{\begin{array}{ll}\pi/4, & \Lambda=0,\\ 5\pi/128, & \Lambda=2\end{array}\right\}, \qquad (22)$$

for the monopole and quadrupole vibrations, respectively, $P_\Lambda(x)$ is the Legendre polynomial. The sum is taken over the periodic orbit families $\beta=(a,b,m)$, m is the period number, a and b are the corner and winding numbers for one period $m=1$, $\phi=\pi b/a$. We keep here only linear terms in ω/λ in exponent and zero order ones in the pre-exponent for simplicity. Next order terms in ω/λ can be easily found like in [15].

The contribution of planar orbits with the classical degeneracy parameter $\mathcal{K}=3$ into the free-energy correction δF_β into (21) is given by $\delta F_\beta = \mathrm{Im}\,\delta\mathcal{F}_\beta$ with

$$\delta\mathcal{F}_\beta = 4E_0\cot\phi\sqrt{\frac{(k_F R)^5\sin\phi}{\pi\,(a\,m)^5}}\,\Phi\,(z_\beta)\,\exp\left[i\,k_F\,L_\beta - i\,\frac{\pi}{2}\,\mu_\beta^{(3)}\right]\exp\left(-\frac{k_F\,L_\beta\Gamma(\omega,T)}{4\,\lambda}\right),\ \mathcal{K}=3. \qquad (23)$$

Here, $E_0=\hbar^2/2MR^2$, μ_β is the Maslov index for planar orbits with $\mathcal{K}=3$, $\mu_\beta = 3\,a\,m+2\,(m\,b-1)$, and we took explicitly into account the spin-degeneracy factor 2. The action S_β is given by $S_\beta(\lambda)=p_F L_\beta$, where $p_F=\sqrt{2\,M\,\lambda}$, L_β the length of the periodic orbits β, $L_\beta = 2\,R\,a\,m\sin\phi$. The temperature cut-off factor $\Phi\,(z_\beta)$ is function of $z_\beta=\pi\tau_\beta T/\hbar$ at $e=\lambda$ in accordance with Eqs. (19). Last exponential cut-off (damping) factor in Eq. (23) is originated from $\mathrm{Im}\,\Sigma(\omega,T)$ (10) of the self-energy Σ in the Green function G Eq. (9) through

$$\delta\mathcal{F}_\beta \to \delta\mathcal{F}_\beta \exp\left(-\frac{\tau_\beta\,\Gamma(\omega,T)}{2\hbar}\right) \approx \delta\mathcal{F}_\beta \exp\left(-\frac{k_F\,L_\beta\Gamma(\omega,T)}{4\,\lambda}\right) \qquad \text{for} \qquad \Sigma_\nu \to \mathrm{e}_\nu \tag{24}$$

as related to $e \to e + i\mathrm{Im}\,\Sigma = e + i\Gamma/2$, where $\Gamma(\omega,T)$ is given by (10). For the diameter orbits $\mathcal{K}=2$ one has respectively

$$\delta\mathcal{F}_\beta = E_0\,\frac{(k_F R)^2}{8\,\pi\,m^3}\,\Phi\,(z_\beta)\,\exp\left[i\,k_F\,L_d - i\,\frac{\pi}{2}\,\mu_d\right]\,\exp\left(-\frac{k_F\,L_\beta\Gamma(\omega,T)}{4\,\lambda}\right),\ \ \mathcal{K}=2, \tag{25}$$

where L_d is the length of the diameters ($a=2, b=1, \phi=\pi/2$), $L_d = 4mR$, μ_d is related to the Maslov index for the diameters, $\mu_d = 2\pi m + 2$. According to these expressions, the contribution of the planar orbits into the $\delta\chi_{10}(\omega)$ is larger in factor $\sqrt{k_F R}$ than for the diameters because of larger degeneracy parameter $\mathcal{K}$ in one unit. However, because shorter orbits might be important for larger temperature T we keep the diameter terms too along with more degenerated planar-orbit components in $\delta\chi_{10}$ and $\delta\chi_{11}$.

By making use of (21) (with next quadratic terms in exponent and linear ones in pre-exponent in $\hbar\omega/\lambda$) one obtains the stiffness δC, the inertia δB and the friction $\delta\gamma$ shell corrections to the macroscopic quantities $\tilde{C}(0)$, $\tilde{B}(0)$ and $\tilde{\gamma}(0)$ in the zero frequency limit [9], respectively,

$$\delta C = -\delta\kappa - \delta\chi(0) = \left(\frac{\partial^2 F}{\partial Q^2}\right)_{Q=0} = -\frac{5}{16\pi}\,(k_F R)^2 \sum_\beta\,(a\,m\sin\phi)^2\ \delta F_\beta, \tag{26}$$

$$\delta B_{10} = \frac{1}{2}\left(\frac{\partial^2\delta\chi'_{10}}{\partial\omega^2}\right)_{\omega=0} = \frac{5\,f_\Lambda\,k_F R\,(M\,R^2)}{32\,E_0}\sum_\beta\,(a\,m\sin\phi\,)^3\ \delta F_\beta, \tag{27}$$

$$\delta\gamma_{10} = \left(\frac{\partial\delta\chi''_{10}}{\partial\omega}\right)_{\omega=0} = -\frac{5\,f_\Lambda\,\hbar(k_F R)^2}{32\,E_0}\sum_\beta\,(a\,m\sin\phi)^2\ \delta F_\beta, \tag{28}$$

where $\delta\chi_{10}$ is the shell response function part of the χ_{10} from Eq. (21), $\delta\mathcal{F}_\beta$ the β-component of the free-energy shell correction δF, see Eqs. (19). We included here and below the damping factor related to the $\mathrm{Im}\,\Sigma(\omega=0,T)$ (10) into δF, according to Eq. (24).

The components of the inertia δB_{11} and the friction $\delta\gamma_{11}$ related to the response function term $\delta\chi_{11}$ are given by the response function shell correction obtained in [15], $\delta B_{11} = \delta B_{11}^{(-)} + \delta B_{11}^{(+)}$, $\delta\gamma_{11} = \delta\gamma_{11}^{(-)} + \delta\gamma_{11}^{(+)}$, where

$$\delta B_{11}^{(-)} = \frac{(2\Lambda+1)f_\Lambda\,k_F R\,M\,R^2}{24\,E_0}\,\mathrm{Re}\sum_\beta\,a^3\,m^2\,(2\,m+1)\,(m+1)\sin^4\phi\cos\phi\ \delta\mathcal{F}_\beta, \tag{29}$$

$$\delta B_{11}^{(+)} = -\frac{(2\Lambda+1)f_\Lambda\,k_F R\,M\,R^2}{4\sqrt{2}\,E_0}\,\mathrm{Re}\sum_\beta\sum_{n=1}^{m-1}\frac{a^3\,m^{3/2}\,(m^2+2\,n^2-2\,m\,n)}{\sqrt{4\,a\,n\,(m-n)-m}}\,\sin^4\phi\cos\phi\ \delta\mathcal{F}_\beta, \tag{30}$$

$$\delta\gamma_{11}^{(-)} = -\frac{(2\Lambda+1)\hbar\,f_\Lambda\,(k_F R)^2}{8\,E_0}\,\mathrm{Re}\sum_\beta\,a^2\,m^2\,(m+1)\sin^3\phi\cos\phi\ \delta\mathcal{F}_\beta, \tag{31}$$

$$\delta\gamma_{11}^{(+)} = \frac{(2\Lambda+1)\hbar\,f_\Lambda\,(k_F R)^2}{4\sqrt{2}\,E_0}\,\mathrm{Re}\sum_\beta\sum_{n=1}^{m-1}\frac{a^2\,m^{3/2}\,(2\,n-m)}{\sqrt{4\,a\,n\,(m-n)-m}}\sin^3\phi\cos\phi\ \delta\mathcal{F}_\beta, \tag{32}$$

$\delta\mathcal{F}_\beta$ is given by (23) for planar $\mathcal{K}=3$ orbits. Note that the diameter contributions into the POT sum for the $\delta\chi_{11}$, see [15], are small of the order of $1/k_F R$ with respect the leading terms (29) - (32) and therefore, can be neglected practically in the specific calculations.

The convergence of the POT sum for the shell corrections in the response function $\chi(\omega)$ (21) and the transport coefficients (26) - (28), (29) - (32) is garanteed by the temperature cut-off factor $\Phi(z_\beta)$ in (19) as mentioned above. Due to this factor all these shell corrections exponentially decrease with increasing temperature T and disappear practically at somewhat larger T about $3-4$ MeV than the critical value $T=T_{sh}$, $T_{sh}\approx\hbar\Omega/\pi=2-3$ MeV of disappearance of the free-energy shell correction δF, see the next section. The distance between gross shells $\hbar\Omega$ is determined by the average period $\bar{T}_\beta$ of the most important, shortest and more degenerated orbits occupied larger phase-space volume, $\hbar\Omega=\hbar v_F/R\approx 2\pi\hbar/\bar{T}_\beta\approx\lambda/A^{1/3}$, see Ref. [13]. The typical values of $\hbar\Omega$ at $\lambda\approx 40$ Mev are $6-10$ MeV for heavy nuclei with $A\sim 60-240$. The additional convergence factor depending also on temperature (and on frequency in $\chi(\omega)$) is due to the factors originated from $\mathrm{Im}\,\Sigma$ (10) through the last exponent in Eq. (24).

4. Discussion

The shell corrections to the transport coefficients $\delta C(0)$, $\delta B(0)$, and $\delta\gamma(0)$ Eqs. (26) - (28), (29) - (32) as function of the particle number, $A^{1/3}$, are shown in Fig. 1 for the case of the

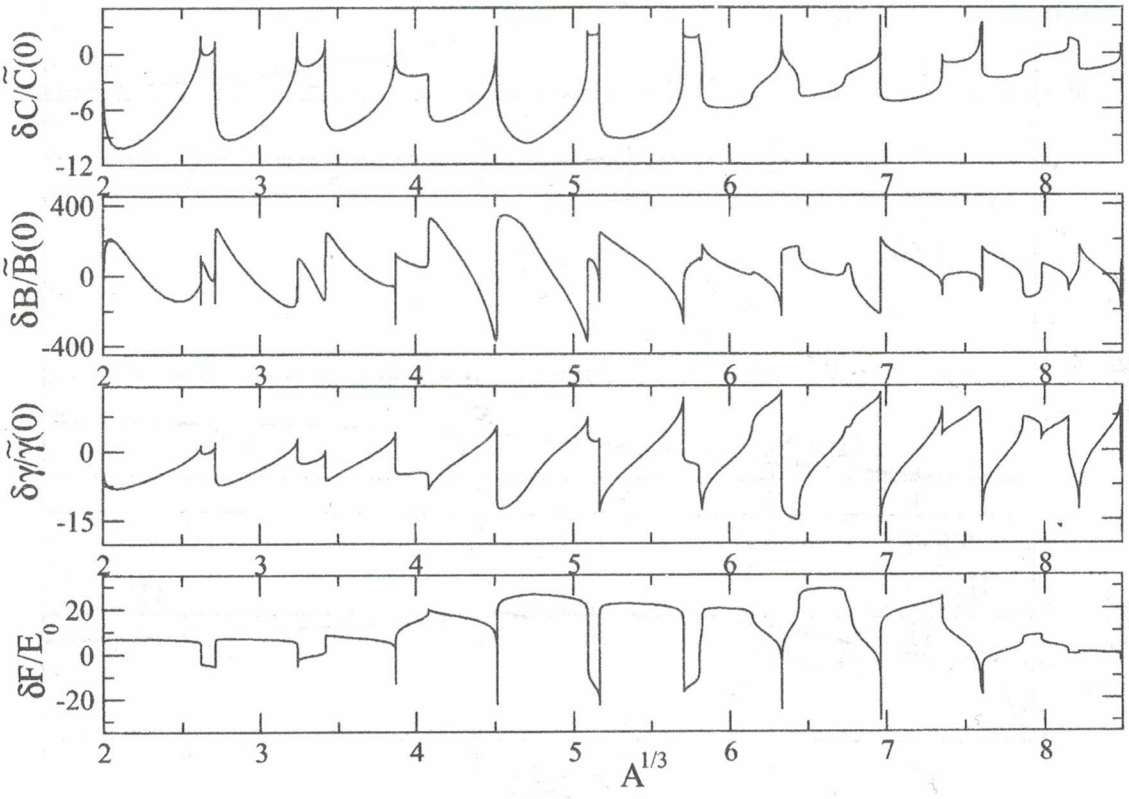

Fig. 1. The stiffness, the inertia, the friction, and the free energy shell corrections as function of particle number ($A^{1/3}$) in units of the macroscopic transport coefficients in the "zero frequency limit"; T=0.5 MeV; λ=40 MeV, r_0=1.2 fm, Γ_0=33.3 MeV, c=20 MeV.

quadrupole ($\Lambda=2$) vibrations ($\int_0^\infty n(e)g(e)de=A$). As seen from this Figure, they are largely proportional to the shell free-energy component δF (18), with the same sign for the case of the inertia and with the opposite sign for the stiffness and the friction. We found that in a qualitative

agreement with the realistic calculations of the inertia and the stiffness, see Fig. IX-I on p. 387 in Ref. [3]. As the so called plateaux condition for the shell corrections [2, 3] is fulfilled for the shell free-energy δF (18), all corrections δC, δB and $\delta\gamma$ satisfy this condition. That confirms their correct extraction from the total quantities in the "zero frequency limit".

Fig. 2 shows the quadrupole friction γ, the accociated reduced friction γ/B, and the effective damping $\gamma/2\sqrt{BC}$ as functions of temperature T for the hydrodynamic (HD) and "wall" Thomas-Fermi (TF) approaches to a smooth (macroscopic) transport coefficients with the semiclassical shell corrections (SC) and without ones, respectively. The calculations are performed at the minimum $k_F R = 9.84$ of the free energy δF Eqs. (18), where the particle number is close to the ^{208}Pb nucleus, in a small temperature region. The HD+SC friction (the top pannel) is monotonically decreasing function of temperature in contrast to the TF+SC approach for which it has a wide maximum at small temperatures due to the shell effects. The quantum shell effects disappear exponentially at temperatures about $3-4$ MeV, a little larger than that for the free energy δF, namely $2-3$ MeV. The TF+SC friction approaches the wall formula asymptotics at large temperatures. The latter is similar to the results for $\gamma(0)/\hbar$ presented in Fig. 5.3 in Ref. [17] with including the diagonal ($e_\nu = e_\mu$) matrix elements in the double sum over quantum

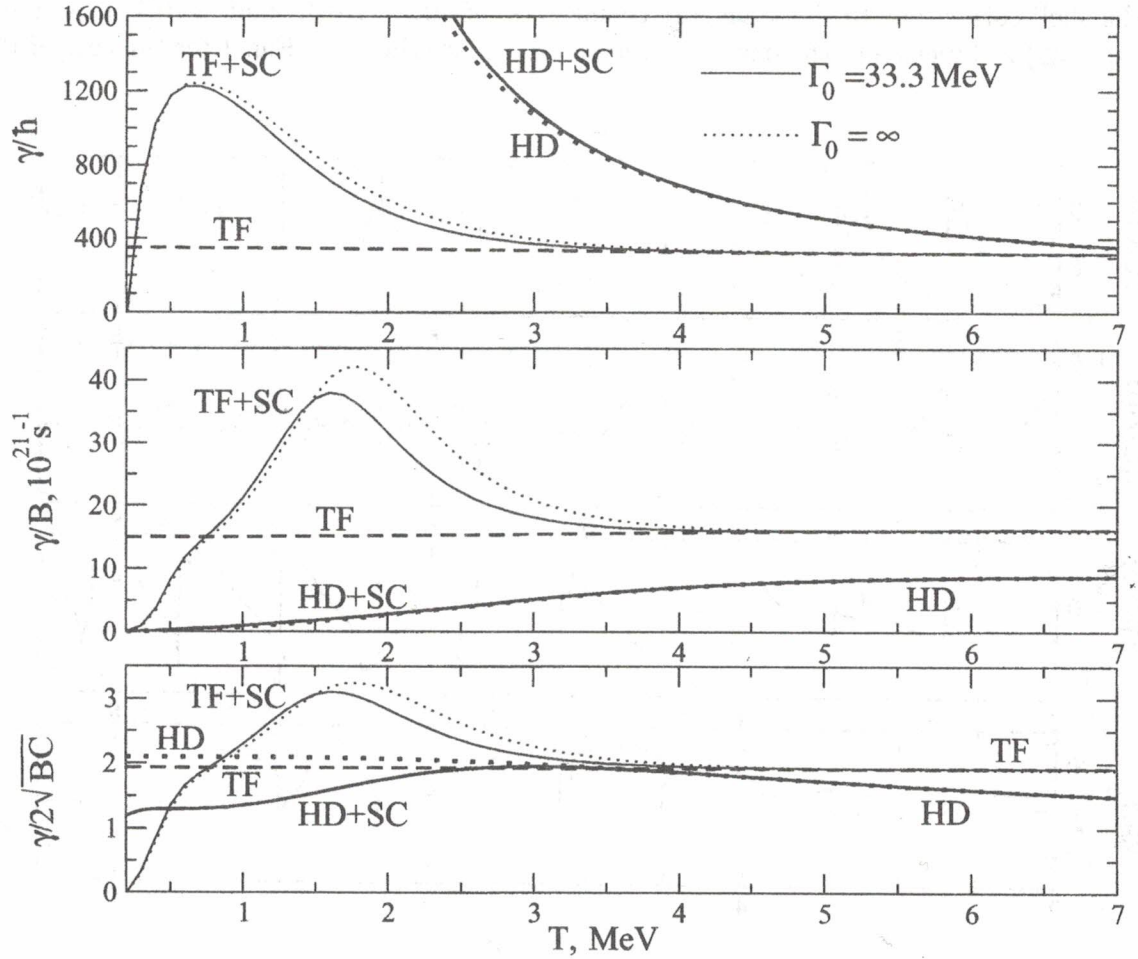

Fig. 2. The friction, the reduced friction, and the effective damping (7) versus temperature; HD+SC (full fat) is the hydrodynamical, and TF+SC (full thin) the Thomas-Fermi components with shell corrections; heavy dots and dashed are the corresponding macroscopic parts; frequent dots show the TF+SC approach for $Im\ \Sigma$=0; $k_F R = 9.84$; parameters are the same as in Fig. 1.

states in the intrinsic response function $\chi(\omega)$ like Eq. (8) ($\gamma \approx \gamma(0)$ because $C(0)/\chi(0)$ is a small

correction as mentioned above). The reduced friction γ/B (middle in Fig. 2) for the HD+SC case increases monotonically from almost zero to about 10 (in units $10^{21}\ s^{-1}$) with growing temperature. In contrast, the TF+SC reduced friction has a maximum of about 40 at temperature about 1.5 MeV. At smaller temperatures the TF+SC γ/B is sharply increasing on left of this maximum and turns into almost constant about 20 on its right for larger temperatures. These values are in agreement with the fission experimental data collected in Ref. [21] in order of the magnitude. For both macroscopic models one has mainly the overdamped motion, $\gamma/2\sqrt{BC} > 1$, for temperatures larger about 0.4 MeV, see the bottom of Fig. 2. In contrast to the HD+SC model, for the TF+SC case at very small temperatures the collective motion becomes underdamped, in accordance with the experimental data mentioned above. The difference between frequent dotted and thin solid lines shows a rather big effect of the residue (damping) interaction related to the $Im\Sigma$, especially for the reduced friction at the intermeadiate temperatures.

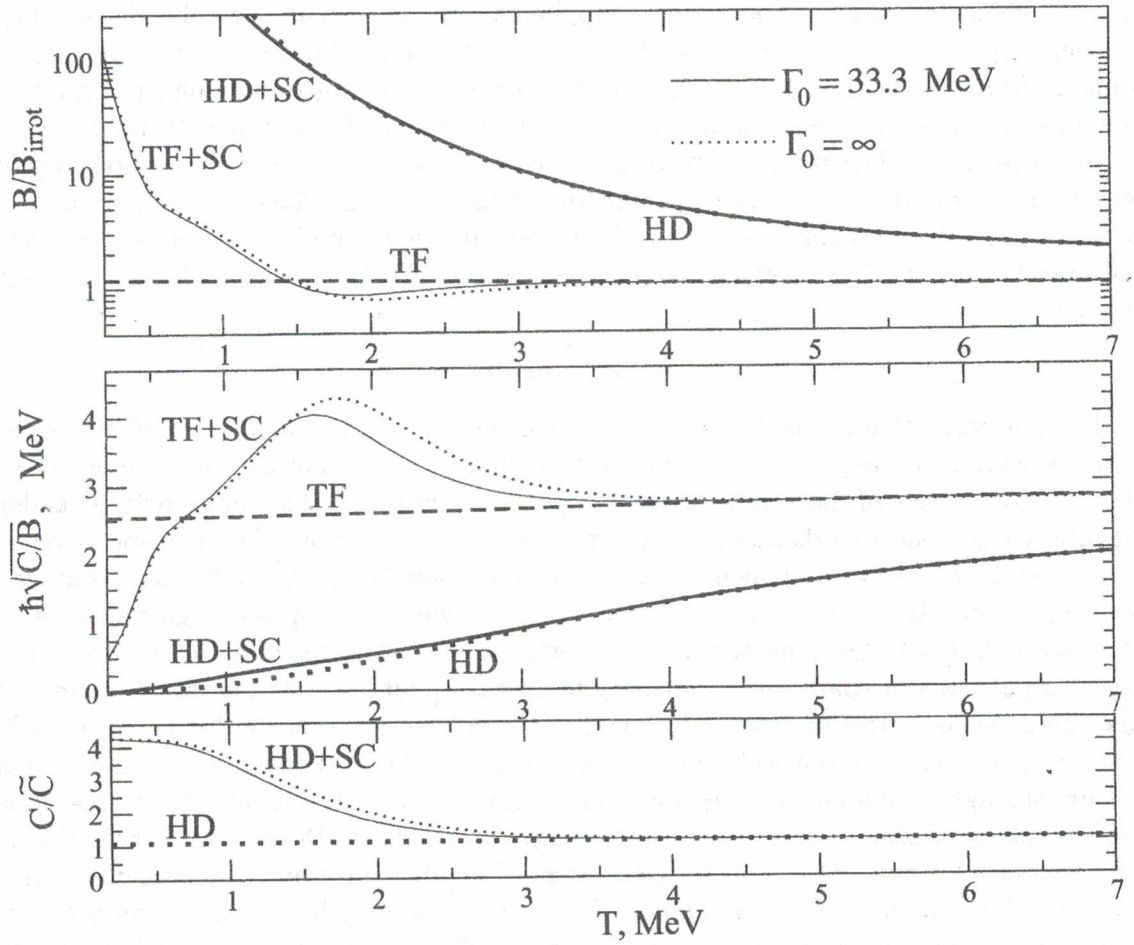

Fig. 3. The inertia, the excitation energy parameter, and the stiffness as function of temperature with the same denotations and parameters as in Fig. 2.

The inertia B, the excitation energy $\hbar\sqrt{C/B}$, and the stiffness C are shown in Fig. 3. The HD+SC mass parameter B (top) is monotonically decreasing function of temperature, while the TF+SC inertia decreases at small temperatures, and approach asymptotically the Thomas-Fermi (TF) approximation with slight shell oscillations around a constant for large ones. Again, we obtain a large shell effect in the TF+SC mass parameter at smaller temperatures. Note that the collective TF inertia shown in this figure is of the order of the well-known value B_{irrot} of the irrotational flow. This is much larger than found from Eq. (14), because of the transformation

(7) to the collective inertia at zero shell components. In particular, it depends on the TF friction (13). Therefore, the high temperature limits for the inertia in the both compared models are rather close to the same B_{irrot} like in some semi-microscopic calculations, see Fig. 5.1.2 of Ref. [9]. The excitation energy parameter $\hbar\sqrt{C/B}$, see the middle, for TF+SC model is much larger than for the HD+SC model having a maximum of about 4.5 MeV at temperature about 1.5 MeV. This parameter of the TF+SC model becomes close to the hydrodynamical value HD at large temperatures. Notice, a large enhancement of the stiffness at small temperatures due the quantum shell effects is clearly seen in the bottom. The quantum (shell) effects in the friction γ and the inertia B (the reduced friction γ/B, the effective damping $\gamma/2\sqrt{B\,|C|}$, and the excitation energy parameter $\hbar\sqrt{C/B}$) at smaller temperatures are relatively much greater for the TF+SC approach than in the HD+SC case. Both compared approximations are largely in agreement with the corresponding theoretical calculations [9, 17–19, 21].

We emphasize here that we should expect that the "zero frequency" approximation for the transport coefficients discussed above can be applied for enough low-laying collective excitations. More general approach which includes high nuclear collective excitations like nuclear giant resonances is based on a local fitting of the individual peaks of the strength function $\mathrm{Im}\chi^{coll}_{QQ}(\omega)$ by the harmonic oscillator one of a well-known Lorenzian-like form, see for instance [9, 16, 17, 20]. The collective transport coefficients are defined in this case as parameters of such harmonic-oscillator response function locally near the strength function peaks. The semiclassical calculations of these transport coefficients and comparison with their "zero frequency limit" as well as more detailed discussion of the relation to the semi-microscopic approach [9, 17–19] will be discussed in further publications.

5. Conclusions

The shell correction method is applied for calculations of the transport coefficients for low collective excitations in the "zero frequency limit" within the POT. For the quadrupole vibrations near the spherical shape of the infinite squre-well potential we found that these coefficients depend significantly on the effective damping related to a residue interaction. The transport coefficient shell components in the "zero frequency limit", which are significant at smaller temperatures, are largely proportional to the free-energy shell corrections, with the opposite sign for the stiffness and the friction but with the same sign for the inertia. The shell corrections to the transport coefficients disappear exponentially with increasing temperature for larger temperatures than that for the free-energy shell correction, especially for the inertia TF+SC with the "wall" Thomas-Fermi macroscopic part. The friction with the hydrodynamical (HD+SC) model for a smooth macroscopic part strongly monotonically decreases in contrast to the TF+SC approach. In the latter case we have the significant oscillations at small temperatures due to the quantum shell effects, and almost constant asymptotics for large temperatures. The reduced friction is in agreement with the experimental data on the fission probabilities in order of the magnitude. We found a significant quantum shell enhancement effect in the inertia and stiffness parameters at small temperatures where we found the underdamped motion such that we have much larger excitation energies for the TF+SC model in contrast to the HD+SC model. All versions of the macroscopic models are mainly overdamped for enough large temperatures. They are in reasonable agreement with some experimental fission data and the theoretical results obtained within similar approximations.

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ЯДЕРНІ ОБОЛОНКИ, ПЕРІОДИЧНІ ОРБІТИ ТА ТРАНСПОРТНІ КОЕФІЦІЄНТИ КОЛЕКТИВНОЇ ДИНАМІКИ

О. Г. Магнер, А. М. Гжебінський, С. М. Федоткін

Для розрахунку транспортних коефіцієнтів колективного ядерного руху в рамках теорії періодичних орбіт, узагальненої на випадок дисипативних процесів, використовується метод оболонкових поправок. Одержано макроскопічну границю квазікласичних транспортних коефіцієнтів з відомою формулою "стінкової в'язкості" для тертя. Обчислено оболонкові поправки до коефіцієнтів жорсткості, тертя та масового параметра як функцій числа частинок і температури при низьких енергіях збудження у досить важких ядрах. Вони пропорційні оболонковим поправкам до вільної енергії з таким же знаком для масового параметра, а для жорсткості та тертя – з протилежним знаком. Показано, що оболонкові компоненти транспортних коефіцієнтів експоненційно зникають при температурах більших, ніж для вільної енергії. Результати знаходяться в узгодженні з деякими експериментальними даними й мікроскопічними квантовомеханічними розрахунками.

ЯДЕРНЫЕ ОБОЛОЧКИ, ПЕРИОДИЧЕСКИЕ ОРБИТЫ И ТРАНСПОРТНЫЕ КОЭФФИЦИЕНТЫ КОЛЛЕКТИВНОЙ ДИНАМИКИ

А. Г. Магнер, А. Н. Гжебинский, С. Н. Федоткин

Для расчетов транспортных коэффициентов коллективной ядерной динамики в рамках теории периодических орбит, обобщенной на случай диссипативных процессов, используется метод оболочечных поправок. Получен макроскопический предел квазиклассических транспортных коэффициентов с известной формулой "стеночной вязкости". Вычислены оболочечные поправки для коэффициентов жесткости, трения и массового параметра как функции числа частиц и температуры для низколежащих коллективных состояний в тяжелых ядрах. Они пропорциональны оболочечным поправкам к свободной энергии с таким же знаком для осцилляций массового параметра, а для осцилляций жесткости и коэффициента трения – с противоположным знаком. Показано, что оболочечные компоненты транспортных коэффициентов экспоненциально исчезают при температурах больших, чем для свободной энергии. Результаты находятся в согласии с некоторыми экспериментальными значениями и микроскопическими квантовомеханическими расчетами.

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