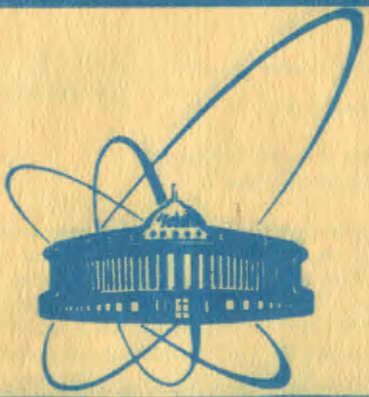

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KINETIC ENERGY DISSIPATION IN HEAVY ION COLLISIONS AND NUCLEAR MATTER DENSITY VIBRATIONS

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INTRODUCTION

Deep inelastic heavy ion collisions are accompanied by a great kinetic energy loss. The value of loss (100 ÷ 300 MeV[1]) tells us that the excited nuclear states can be of complicated structure. However, since the time of the energy loss is extremely short (~10^{-22} s [1,2]) some coherent mechanism of the energy damping may exist.

It means that the doorway states can be the collective ones. From the large value of the energy loss it follows that first of all it is necessary to take into account high frequency modes of excitation. Among them the multipole nuclear matter density vibrations (giant resonances) are the most important[3-5].

Based on these assumptions, we have got equations[5], describing the relative motion of the interacting nuclei in a classic approximation. These equations have a number of peculiarities, distinguishing them from analogous equations, usually used in the phenomenological models[2]. Since the time of the kinetic energy loss to the inner degrees of freedom and the relaxation time are of the same order, the explicit time dependence of frictional coefficient and of the varying (renormalized due to excitation of inner degrees of freedom) nucleus-nucleus potential and mass parameters plays an essential role in the equations of motion. Second, the asymptotic value of the radial frictional coefficient is 50-100 times as small as phenomenological one[2]. These two facts made it possible to suppose that the source of the irreversible kinetic energy losses is not only the radial friction but also the explicit time-dependent interaction potential. Preliminary estimations (except trajectory calculations) showed that the significant part of the experimental kinetic energy loss might be explained in this way.

The aim of the present paper is to study possible ways of explanation of the experimental data on the energy loss and energy angle (E-θ) correlations [1,2] on the basis of our model, having done detailed trajectory calculations.

The physical results strongly depend on the spectrum of the nuclear inner excitations. We shall point, if possible, to the changes of the results while using another microscopic model.

EQUATION OF MOTION

In paper[5] we get the following equations for the description of the relative motion in the collision plane:

$$\mu\ddot{R} - L^2/\mu R^3 + \frac{\partial U}{\partial R} = -\frac{\partial \langle \delta V \rangle}{\partial R},$$

$$L = \mu R^2 \dot{\theta}, \qquad (1a)$$

$$\dot{L} = -\frac{\partial \langle \delta V \rangle}{\partial \theta},$$

where μ is the reduced mass; U, the nucleus-nucleus potential, containing the nuclear[2] and Coulomb [6] terms. The connection between the relative motion and the inner degrees of freedom is described by the interaction

$$\langle \delta V \rangle = \sum_{\alpha} \gamma^{*}_{\alpha}(\bar{R})\, \xi_{\alpha}(t), \qquad (1b)$$

$\gamma^{*}_{\alpha} = \gamma^{*}_{n\ell m}(\bar{R}) = G_{n\ell}(R)\, Y^{*}_{\ell m}(\hat{R})$ is the form factor (see also Appendix A), $\xi_{\alpha} = \xi_{n\ell m}$ are amplitudes, and (n, ℓ, m) are quantum numbers of the density oscillations. In the hydrodynamic approximation linear damped oscillations are described in the presence of an external force, which appears due to the connection of oscillations with relative motion:

$$\xi_{\alpha}(t) = -\frac{2\omega_{\alpha}}{h\Omega_{\alpha}} \int_{0}^{t} dt' \cdot \exp(-\Gamma_{\alpha}(t-t')/2) \sin(\Omega_{\alpha}(t-t'))\, \gamma_{\alpha}(t'), \qquad (2)$$

$\omega_{\alpha} = \omega_{n\ell}$ are natural frequencies $\Omega_{\alpha} = \sqrt{\omega_{\alpha}^2 - \Gamma_{\alpha}^2/4}$, $\Gamma_{\alpha} = \Gamma_{n\ell}$ is the damped term, describing phenomenologicaly the connection of collective vibrations with the other inner degrees of freedom.

Substituting (2) into (1) we get the system of integro-differential equations for the trajectory $R(t)$, $\theta(t)$ in the collision plane. We succeeded in getting these equations of a closed type owing to the interaction linearity in density-oscillation amplitudes. It gave us a possibility of writing (2). For a more complicated interaction we must solve a large number of coupled differential equations for relative and inner motions. In principle the system (1.2) can be solved numerically, for example, by using the finite-difference method. However, for comparison with the phenomenological

models it is useful to reduce the equations to an approximate differential form, separating the induced potential, friction, mass parameters

$$
\begin{aligned}
&(\mu + \tilde{M})\ddot{R} + \frac{1}{2}\dot{R}^2 \frac{\partial \tilde{M}}{\partial R} - \frac{L^2}{\mu R^3}\left(1 + \frac{1}{2\mu R}\,\frac{\partial \Psi}{\partial R}\right) + \\
&+ \frac{\partial (U + \tilde{U})}{\partial R} + \dot{R}\tilde{\chi}_R = 0 , \\
&\frac{d}{dt}\left(L\left(1 + \frac{\tilde{\Psi}}{\mu R^2}\right)\right) = \frac{-L}{\mu R^2}\tilde{\chi}_\theta , \\
&\dot{\theta} = \frac{L}{\mu R^2} .
\end{aligned}
\tag{3}
$$

In paper [5] the induced potential ($\tilde{U}$) and the radial friction ($\tilde{\chi}_R$) were briefly considered. Let us write the expressions for all functions, entering into equations (3) (see also Appendix A):

$$
\begin{aligned}
&\tilde{U}(R,t) = -\gamma^2(R)\, J_R^{(0)}(t) , \\
&\tilde{\chi}_R(R,t) = 2\left|\frac{\partial \gamma(R)}{\partial R}\right|^2 J_R^{(1)}(t) , \\
&\tilde{M}(R,t) = -\left|\frac{\partial \gamma(R)}{\partial R}\right|^2 J_R^{(2)}(t) , \\
&\tilde{\chi}_\theta(R,t) = \gamma^2(R)\, J_\theta^{(1)}(t) , \\
&\tilde{\Psi}(R,t) = -\frac{1}{2}\gamma^2(R)\, J_\theta^{(2)}(t) , \\
&J_R^{(n)}(t) = i\int_0^t s^n \chi_{RR}(s)\, ds , \\
&J_\theta^{(n)}(t) = i\int_0^t s^n \chi_{\theta\theta}(s)\, ds ,
\end{aligned}
\tag{4}
$$

$$i\chi_{RR}(s) = \sum_{n\ell} M_{n\ell} \cdot \sin(\Omega_{n\ell}\, s) \cdot \exp(-\Gamma_{n\ell} s/2),$$

$$i\chi_{\theta\theta}(s) = \sum_{n\ell} M_{n\ell} \cdot \ell(\ell+1) \cdot \sin(\Omega_{n\ell}\, s) \cdot \exp(-\Gamma_{n\ell}\, s/2), \qquad (4)$$

$\tilde{\chi}_\theta$ is the tangential frictional coefficient, $\tilde{M}$ and $\tilde{\Psi}$ inertial parameters. Equations of motion are completely determined by the radial form factor and structure time functions $i\chi_{RR}, i\chi_{\theta\theta}$, which are analogous to the response functions [7,8] for radial and tangential motions. It is to be noted that they are deduced without averaging over the grand canonical ensemble (as is usually done in Linear Response Theory [7,8,9]). Besides, we did not suppose the localization of interaction in angle [8] while separating the angular part, but expended the interaction over spherical functions. The renormalized potential, friction, etc., depend not only on distance but also on time. In Linear Response Theory this dependence is neglected and asymptotic values are used. But for such fast processes as the dissipation in deep inelastic heavy ion collisions to do the same is wrong, to our mind [5]. Let us analyse this statement in detail. In order to estimate the characteristic relaxation times we calculate the response function $i\chi_{RR}$ (Fig. 1) and its initial moments $J_R^{(0)}$ (Fig. 2) and $J_R^{(1)}$ (Fig. 3) defining the explicit dependence of the induced potential $\tilde{U}$ and radial friction $\tilde{\chi}_R$ on time.

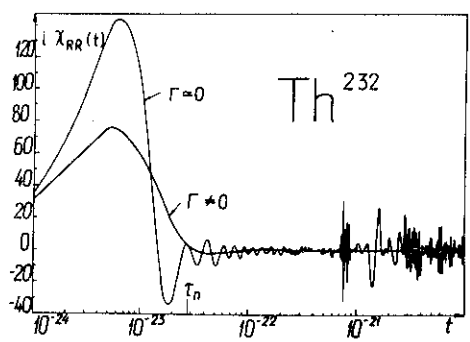
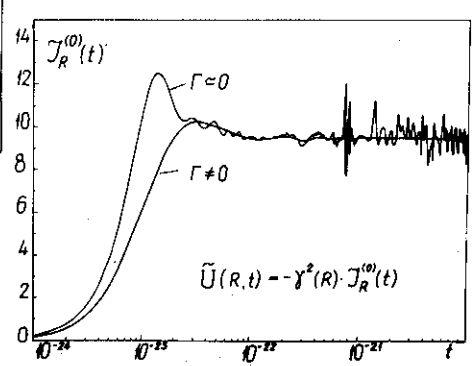


Fig.1. The radial response function of the nucleus Th.

Fig.2. The zeroth moment of the response function $J_R^{(0)}(t)$ defining the explicit time dependence of the induced potential $\tilde{u}$.



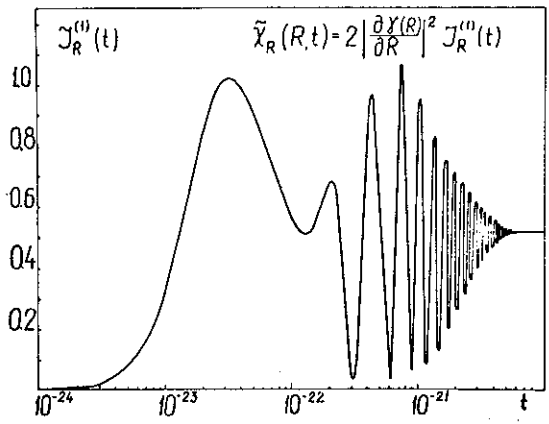


Fig.3. The first moment of the response function $J_R^{(1)}(t)$ defining the explicit time dependence of the coefficient of radial friction $\tilde{\chi}_R$.

The following conclusions can be drawn:

i) The time of taking asymptotic value by the potential $\tau_{as,p} \simeq (0.5\text{-}2.0)\cdot 10^{-22}$ (Fig. 2) is close to the kinetic-energy-loss time, following from the calculation according to the classic friction models /2/.

ii) The analogous value for the radial frictional coefficient $\tau_{as,fr} \simeq 50\cdot 10^{-22}$ (Fig. 3) is comparable with the whole reaction time /1/. Strong oscillations $J_R^{(1)}$ for times lesser than $\tau_{as,fr}$ are connected with the response-function oscillations. Although they are not practically seen in Fig. 1, they strengthen owing to the integration with the weight function proportional to t .

iii) The memory time (the time of two largest response function oscillations (Fig. 1) $\tau_m \sim 3\cdot 10^{-23}$s is several times as small as the analogous one ($\tau_m \sim 2\cdot 10^{-22}$s) got in ref./8/ by assuming the one-body dissipation mechanism. It is explained by the fact that we consider a higher-frequency branch of nuclear excitations than in ref./8/ (See Appendix A).

iv) Damped terms, introduced phenomenologically to take into account the decay of collective modes into noncollective ones, are important for smoothing characteristic pulsations in the response function at times larger than ($\sim 5\text{-}7\cdot 10^{-21}$s). Besides, this damping allows us to neglect the artifical procedure of cutting while integrating over time used in ref./8/.

The asymptotic values ($t \to \infty$) of radial and tangential frictional coefficients are proportional to damping values.

So, the damping is significant for the description of irreversible energy losses caused by friction.

v) The potential, friction, etc., are factorized into radial form factor $\gamma(R)$ (or its derivatives) and time functions $J^{(n)}_{R,\theta}(t)$ (4), defined by response functions $i\chi_{RR,\theta\theta}$ and having recurrent correlations:

$$J^{(n)}(t) = t \cdot J^{(n-1)}(t) - \int_0^t J^{(n-1)}(s)\, ds.$$

The form of response functions depends on the nature of states excited in the nuclear collision. The main features are the memory time and amplitudes. Let us substitute the response function by the effective step function:

$$i\chi(t) = \frac{\chi^{00}}{\tau_m} \begin{cases} 1 & t \le \tau_m \\ 0 & t > \tau_m \end{cases}. \qquad (5a)$$

We arrive at the following result:

$$J^{(n)}(t) = \frac{\chi^{00}}{\tau_m (n+1)} \begin{cases} t^{n+1} & t \le \tau_m \\ \tau_m^{n+1} & t > \tau_m \end{cases}. \qquad (5b)$$

Varying χ^{00} and τ_m we can roughly estimate the influence of the concrete choice of a microscopic model on the renormalized terms. The amplitude χ^{00} and the memory time τ_m are completaly defined by the excitation spectrum and its contribution into the connection term with relative motion $\langle \delta V \rangle$ (1b). But the amplitude is an integral characteristic defined by the whole spectrum, then the memory time strongly depends on the distribution of the connection term $\langle \delta V \rangle$ over spectrum. So, later on we shall keep χ^{00} value the same as in ref.[5] and shall vary only τ_m. As a result, we get a suitable parametrization for all time-dependent functions in the Equations of Motion.

RADIAL FRICTION

In phenomenological trajectory calculations the radial friction is usually chosen as $\chi_R = C_R \cdot f(R)$, where C_R is the free parameter and $f(R)$ is the form factor that is most frequently

taken as the square of the first or second derivative of a nucleus-nucleus potential with respect to the distance [2] : $f(R) \sim (\partial U/\partial R)^2$ or $(\partial^2 U/\partial R^2)^2$

$$\tilde{\chi}_R(R,t) = 2\left(\frac{\partial \gamma}{\partial R}\right)^2 J_R^{(1)}(t) = \left(\frac{\partial^2 U}{\partial R\ \partial R}\right)^2 \cdot \tilde{C}_R(t) .$$

An analogous form factor can also be separated in our case. However, the effective frictional strength $\tilde{C}_R$ explicitly depends on time and as was mentioned above, the time of taking the asymptotic value appears to be rather large ($\sim 5 \cdot 10^{-21}$ s) and comparable with the whole reaction time. Preliminary estimations of the frictional force showed [5] that it is too small (50-100 times as small as phenomenological one [2]). We checked this statement in trajectory calculations in this work also choosing different nucleus-nucleus potentials. As a result, we can draw the conclusion that the obtained radial frictional forces in our model are too small (on the average, 100 times as less as phenomenological ones) *. It means that significant kinetic energy losses cannot be explained only by the radial friction. Earlier we mentioned the rough proportionality of asymptotic frictional forces to damping and duration of the response function. We selected damping as to describe the widths of the lowest excitations, therefore its considerable variations are not allowed. It is possible to enlarge the effective memory time of the response function by applying a more realistic model for the description of density oscillations (e.g., taking into account deformation of the form of a nucleus, diffusion of the nuclear surface and considering the density oscillations in the double nuclear system). We intend to investigate the influence of the form and nuclear surface diffusion parameters on the response function.

POTENTIAL

In our previous paper we supposed that irreversible kinetic energy losses can be caused not only by friction but also by the explicitly time-dependent potential:

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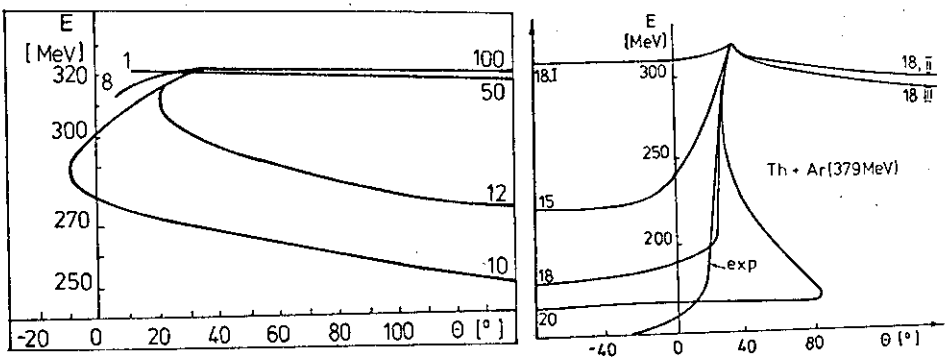
* We should mention the analogous result obtained in Linear Response Theory under the assumption about the one-body dissipation mechanism [8].



Figs.4,5. (E–θ) correlation for the reaction Th+Ar. There are effective memory times (in units 10^{-22} s) near all the curves (see the text).

$$\tilde{U}(R,t) = -\gamma^2(R) \cdot J_R^{(0)}(t) \equiv -\tilde{U}_0(R) \cdot J_R^{(0)}(t) . \qquad (6)$$

In this case the irreversibility of energy losses is ensured by the threshold form of $J_R^{(0)}$ (Fig. 2), which follows from the coherent superposition of nuclear density oscillations. Energy losses are defined by the potential $\tilde{U}$ averaged over the time of interaction (τ_{int}):

$$\Delta E(\tau_{int}) = E_0 - E(\tau_{int}) = \int_0^{\tau_{int}} \tilde{U}_0 \cdot \frac{dJ_R^{(0)}}{dt} \cdot dt . \qquad (7)$$

This correlation can be easily obtained from the equation for radial motion with the explicitly time-dependent force. Parametrizing $J_R^{(0)}$ in the simplest way (5b) we get:

$$\Delta E(\tau_{int}) = \frac{\chi_{RR}^{00}}{\tau} \int_0^{\tau} \tilde{U}_0(R(t))\, dt , \qquad (7a)$$

where $\tau \equiv \min(\tau_{int}, \tau_m)$, τ_m is the effective memory time (5). From this it follows that τ_m must be comparable with the interaction time (energy losses will be small both for small τ_m (the integral (7a) itself is small) and for large τ_m). In Fig. 4 this is illustrated by the calculation of (E–θ) correlation for the reaction Th + Ar (380 MeV) at several values

of time τ_m. We neglected all the renormalized terms (except the potential $\tilde{U}$). It is seen that we can expect the considerable energy losses for $\tau_m \sim 10^{-21}$ s. The losses are insignificant both for $\tau_m \sim 10^{-22}$ s and for $\tau_m \sim 10^{-20}$ s.

TANGENTIAL FRICTION, INERTIAL PARAMETERS

In our calculations the value of tangential friction is near to the phenomenological one [2].

The ratio of the value of the renormalized mass to the reduced one does not exceed several per cent:

$$|\tilde{M}/\mu| < 0.09 .$$

However, this ratio strongly oscillates in time. The analogous result can be obtained for the renormalized inertial moment

$$|\tilde{\psi}/R^2\mu| < 0.001 .$$

Asymptotic value $\tilde{M}$ and $\tilde{\psi}$ are very small but times of taking these values are about 10^{-19} s.

TRAJECTORY CALCULATIONS RESULTS

One can see the results of our calculations of the energy angle $(E-\theta)$ [1,2] correlation of the reaction Th+Ar (380 MeV) products in Figs. 4-6. The calculations have been performed on the basis of equations (3,4). The response functions are chosen in form (5). The amplitudes $\chi^{00}_{RR}, \chi^{00}_{\theta\theta}$ are defined within our model. The only parameter of the spectrum is the memory time $\tau^R_m = \tau^\theta_m \equiv \tau_m$ *.

First, let us analyse the result of calculations of correlation when only the induced potential $\tilde{U}$ (among the renormalized terms) is taken into account (Fig. 4). It is seen that in such a way one can describe considerable kinetic energy losses (about 70 MeV). However, the whole experimental [1] correlation cannot be thus described.

* The model used defines also $\tau_m \simeq 3 \cdot 10^{-23}$ s.

The results of calculation of $(E-\theta)$ correlation in the case, when the time-dependent potential $\tilde{U}$ and friction $\tilde{\chi}_R$ are taken into account are shown in Fig. 5. For $\tau_m = 18 \cdot 10^{-22}$s (the variant of the best description of experimental data) one can observe the results of calculation in the cases: i) the only radial friction (I); ii) the only potential $\tilde{U}$ (II); iii) the tangential friction and potential (III). One can see that it is very important to take into account the time-dependent potential and friction simultaneously. At the beginning of collision the main contribution is due to the renormalized time-dependent potential, at a later stage the role of friction enlarges. This statement is confirmed both by the results of trajectory calculations and by the peculiarities of time-dependent forms of the potential and friction (Figs.2, 3).

The results of our calculations are compared with the results of phenomenological friction calculation [2] and experimental data [1] in Fig. 6. It is shown that the results of both theoretical calculations are similar. The difference between experimental and theoretical results can be eliminated by including deformation in the exit channel [11].

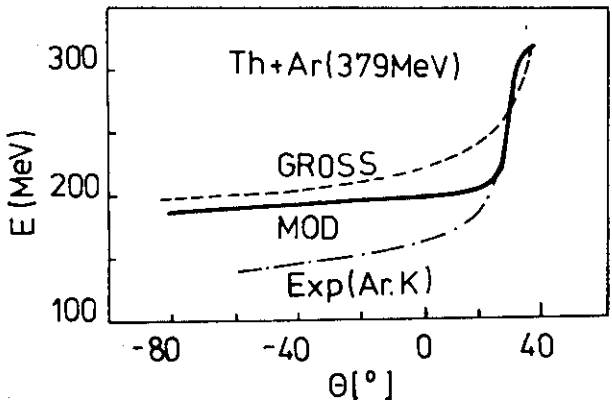


Fig.6. Comparison of $(E-\theta)$ correlations, calculated in the framework of our model (MOD) and the phenomenological friction model (GROSS) with the experimental data.

APPENDIX A

The connection of the relative motion of two colliding nuclei with the density vibrations is described by the potential

$$\delta V(\vec{R}, \xi) \equiv \frac{\partial U(\vec{R}, R_p, R_t)}{\partial R_t} \cdot \delta R_t(\vec{R}, \xi),$$

where U is the nuclear part of the statical potential [2]; $R_t(R_p)$, the average radius of the target (the projectile)

$$\delta R_t = R_t \sum_{n\ell m} F_{n\ell} \frac{j'(z_{n\ell})}{z_{n\ell}} Y^*_{\ell m}(\hat{R}) \sqrt{\frac{h\omega_{n\ell}}{2m u_s^2 \rho_0}}\, \xi_{n\ell m}(t),$$

the variables $\xi_{n\ell m}$ describe the density vibrations, j_ℓ, j'_ℓ are the spherical Bessel function and its derivative

$$F_{n\ell}^{-2} = \int_0^{R_t} j_\ell^2\left(\frac{z_{n\ell}}{R_t} r\right) r^2 dr,$$

$\rho_0 = 0.17\ \mathrm{fm}^{-3}$, $u_s = 0.2 \cdot c$ is the "sound" velocity in nuclear matter, m is the nucleon mass, $\omega_{n\ell} = u_s z_{n\ell}/R_t$, $z_{n\ell}$ can be found from equations [12]:

$$\beta_{n\ell} = \frac{j_\ell(z_{n\ell})}{j'_\ell(z_{n\ell})} = \frac{17(\ell-1)(\ell+2)}{3m u_s^2 A_t^{1/3} z_{n\ell}}.$$

In the explicit form:

$$\delta V \equiv \sum_{n\ell m} Y^*_{n\ell m} \cdot \xi_{n\ell m} = \sum_{n\ell m} G_{n\ell}(R)\, Y^*_{\ell m}(\hat{R})\, \xi_{n\ell m},$$

$$G_{n\ell}(R) = \gamma(R) \sqrt{\frac{4\pi h \Omega_{n\ell} M_{n\ell}}{(2\ell+1)\omega_{n\ell}}} =$$

$$= \gamma(R) \sqrt{\frac{4\pi h}{z_{n\ell}} \left(1 + \beta_{n\ell}/z_{n\ell} + (1 - \ell/\ell+1)/z_{n\ell}^2\right) \beta_{n\ell}^2)^{-1}},$$

$$\gamma(R) = \frac{\partial U(R)}{\partial R_t} \sqrt{\frac{R_t}{3A_t m u_s}}.$$

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