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Fission widths of hot nuclei from Langevin dynamics

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The fission dynamics of excited nuclei is studied in the framework of the Langevin equation. One-body wall-and-window friction is used as the dissipative force in the Langevin equation. In addition to the usual wall formula friction, the chaos-weighted wall formula developed earlier is also considered here. The fission rate calculated with the chaos-weighted wall formula is found to be larger by about a factor of 2 compared to that obtained with the usual wall friction. The systematic dependence of the calculated fission width on temperature and spin of the fissioning nucleus is investigated and a simple parametric form of the fission width is obtained.

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I. INTRODUCTION

The fission of highly excited compound nuclei formed in heavy-ion-induced fusion reactions has emerged as a topic of considerable theoretical and experimental interest in recent years. Multiplicity measurements of light particles and photons emitted in the prescission stage strongly suggest [1] that fission is a much slower process for hot nuclei than that determined from the statistical model of Bohr and Wheeler [2] based on phase space arguments. This led to a revival of theoretical studies based on the original work of Kramers [3] who considered the fission of excited nuclei as a consequence of thermal fluctuations. Dynamical models for fission based on the Fokker-Planck equation [4,5] and Langevin equation [6,7] were subsequently developed.

The most extensive application of the Langevin equation to study fission dynamics was made by Fröbrich and Gontchar [7]. A combined dynamical and statistical model for fission was employed in their calculations in which a switching over to a statistical model description was made when the fission process reached the stationary regime. The required fission widths for the statistical branch of the calculation were obtained from the stationary limit of the fission rates as determined by the Langevin equation [8]. The dissipative property of nuclei is an important input to such Langevin dynamical calculations. Though a complete theoretical understanding of the dissipative force in fission dynamics is yet to be developed, a detailed comparison [9] of the calculated fission probability and prescission neutron multiplicity excitation functions for a number of nuclei with the experimental data led to a phenomenological shape-dependent nuclear friction. The phenomenological friction turned out to be considerably smaller ($\sim$10%) than the standard wall formula value for nuclear friction for compact shapes of the fissioning nucleus whereas a strong increase of this friction was found to be necessary at large deformations. A clear physical picture for such a friction is yet to be developed and the present work is an effort in this direction.

The wall formula for nuclear friction was developed by Blocki *et al.* [10] in a simple classical picture of one-body dissipation. It was also derived from a formal theory based on classical linear response theory [11]. One crucial assumption of the wall formula concerns the randomization of the particle (nucleon) motion due to the successive collisions it suffers at the nuclear surface. In other words, a complete mixing in the classical phase space of particle motion is required. It was early realized [10,11] that any deviation from this randomization assumption would give rise to a reduced strength of the wall formula. Further, Nix and Sierk suggested [12,13] in their analysis of mean fragment kinetic energy data that the dissipation is about 4 times weaker than that predicted by the wall-plus-window formula of one-body dissipation. However, it is only recently that a modification of the wall formula has been proposed [14] in which the full randomization assumption is relaxed in order to make it applicable to systems in which mixing in the phase space is partial. Considering only those chaotic particle trajectories which arise due to irregularity in the shape of the one-body potential and which are responsible for irreversible energy transfer, a modified friction coefficient was obtained in Ref. [14]. In what follows, we shall use the term ‘‘chaos-weighted wall formula’’ (CWWF) for this modified friction in order to distinguish it from the original wall formula (WF) friction. As was shown in Ref. [14], the CWWF friction coefficient η_{cwwf} will be given as

$$\eta_{cwwf}=\mu\,\eta_{wf}, \tag{1.1}$$

where η_{wf} is the friction coefficient as was given by the original wall formula [10] and μ is a measure of chaos (chaoticity) in the single-particle motion and depends on the instantaneous shape of the nucleus. The value of chaoticity μ changes from 0 to 1 as the nucleus evolves from a spherical shape to a highly deformed one. The CWWF friction is thus much smaller than the WF friction for compact nuclear shapes while they become closer at large deformations. The CWWF friction was subsequently found [15,16] to describe satisfactorily the collective energy damping of cavities containing classical particles and undergoing time-dependent shape evolutions. Thus suppression of the strength of the wall formula friction achieved in the chaos-weighted wall formula suggests that chaos in single-particle motion (rather

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a lack of it) can provide a physical explanation for the reduction in strength of friction for compact nuclear shapes as required in the phenomenological friction of Ref. [9] and this has motivated us to apply CWWF friction to fission dynamics in the present work.

We shall present in this paper a systematic study of fission rates by using both CWWF and WF frictions in the Langevin equation. The aim of our study is twofold. First, we would like to find the effect of introducing the chaoticity factor in friction on fission rates at different excitation energies and spins of the compound nucleus. The second one concerns a parametric representation of the fission width, the need for which arises as follows. The fission width is an essential input along with the particle and γ widths for a statistical theory in the stationary branch of compound nucleus decay. Kramers [3] obtained a simple expression for the stationary fission width assuming a large separation between the saddle and scission points and a constant friction. Gontchar *et al.* [8] later derived a more general expression, taking the scission point explicitly into account, but still assuming a constant friction coefficient. The CWWF friction, however, is not constant and is strongly shape dependent and hence the corresponding stationary fission width cannot be analytically obtained. Thus it becomes necessary to find a suitable parametric form of the numerically obtained stationary fission widths using CWWF friction in order to use them in the statistical regime of compound nucleus decay. We shall concentrate upon the parametric representation of fission widths in the present work while application of CWWF friction in a full dynamical plus statistical model will be reported in a future publication.

We shall describe the Langevin equation along with the necessary input as used in the present calculation in the next section. The calculated fission rates and the systematic behavior of the stationary fission widths will be given in Sec. III. A summary of the results will be presented in the last section.

II. LANGEVIN EQUATION FOR FISSION

A. Nuclear shape, potential, and inertia

In order to specify the collective coordinates for a dynamical description of nuclear fission, we will use the shape parameters c,h, and α as suggested by Brack *et al.* [17]. We will consider only symmetric fission ($\alpha=0$) and will further neglect the neck degree of freedom ($h=0$) in order to simplify the calculation. The surface of a nucleus of mass number A will then be defined as

$$\rho^2(z)=\left(1-\frac{z^2}{c_0^2}\right)(a_0c_0^2+b_0z^2), \tag{2.1}$$

where

$$c_0=cR,$$

$$R=1.16A^{1/3}$$

and

$$a_0=\frac{1}{c^3}-\frac{b_0}{5},$$

$$b_0=\frac{c-1}{2},$$

in cylindrical coordinates for the elongation parameter c. Considering c and its conjugate momentum p as the dynamical variables, the coupled Langevin equations in one dimension will be given [18] as

$$\frac{dp}{dt}=-\frac{p^2}{2}\frac{\partial}{\partial c}\left(\frac{1}{m}\right)-\frac{\partial F}{\partial c}-\eta\dot{c}+R(t),$$

$$\frac{dc}{dt}=\frac{p}{m}. \tag{2.2}$$

In the above equations, m and η are the shape-dependent collective inertia and friction coefficients, respectively. The free energy of the system is denoted by F while $R(t)$ represents the random part of the interaction between the fission degree of freedom and the rest of the nuclear degrees of freedom considered collectively as a thermal bath in the present picture.

We will make the Werner-Wheeler approximation for incompressible irrotational flow to calculate the collective inertia [19]. The driving force in a thermodynamic system should be derived from its free energy for which we will use the following expression [9] considering the nucleus as a noninteracting Fermi gas:

$$F(c,T)=V(c)-a(c)T^2, \tag{2.3}$$

where T is the temperature of the system and $a(c)$ is the coordinate-dependent level density parameter which is given as [20]

$$a(c)=a_vA+a_sA^{2/3}B_s(c). \tag{2.4}$$

The values for the parameters a_v, a_s and the dimensionless surface area B_s are chosen following Ref. [9].

The potential energy $V(c)$ enters into our calculation through its dependence on the deformation coordinate. This deformation-dependent potential energy is obtained from the finite-range liquid drop model [21] where we calculate the generalized nuclear energy by double folding the uniform density within the surface [Eq. (2.1)] with a Yukawa-plus-exponential potential. The Coulomb energy is obtained by double folding another Yukawa function with the density distribution. The various input parameters are taken from Ref. [21] where they were determined from fitting fission barriers of a wide range of nuclei. The centrifugal part of the potential is calculated using the rigid-body moment of inertia.

The instantaneous random force $R(t)$ plays a very crucial role in the Langevin description of nuclear fission. As a result of receiving incessant random kicks, the fission degree of freedom can finally pick up enough kinetic energy to overcome the fission barrier. This random force is modeled after that of a typical Brownian motion and is assumed to have a stochastic nature with a Gaussian distribution whose average is zero [6]. It is further assumed that $R(t)$ has extremely short correlation time, implying that the intrinsic nuclear dynamics is Markovian. Consequently the strength of the random force can be obtained from the fluctuation-dissipation theorem and the properties of $R(t)$ can be written as

$$\langle R(t)\rangle = 0,$$

$$\langle R(t)R(t')\rangle = 2\eta T\delta(t-t'). \tag{2.5}$$

B. One-body dissipation

One-body dissipation was used more successfully in fission dynamics than two-body viscosity in the past [6,18]. Accordingly, we shall consider the one-body wall-and-window dissipation [10] to account for the friction coefficient η in the Langevin equation. For the one-body wall dissipation, we shall use the chaos-weighted wall formula [Eq. (1.1)] introduced in the preceding section. At this point we will briefly recall the physical arguments made in order to arrive at this expression. The particle trajectories moving in a one-body nuclear potential were first identified as either regular or chaotic depending on their nature of time evolution [14,15]. Originating from a given point near the nuclear surface and moving in a given direction, a regular trajectory closes smoothly in phase space. On the other hand, another trajectory leaving the same point but in a different direction could be a chaotic one that does not close in phase space. Considering the contributions of these two types of trajectories separately, it was argued in Refs. [14,15] that only chaotic trajectories give rise to irreversible energy transfer and the resulting friction coefficient acting on the wall motion will be as given in Eq. (1.1). The chaoticity μ is a measure of chaos in the single-particle motion of the nucleons within the nuclear volume, and in the present classical picture, this will be given as the average fraction of the trajectories that are chaotic when the sampling is done uniformly over the nuclear surface.

The chaoticity is a specific property of the nonintegrability of the nuclear shape. Thus it is required to be calculated for all possible shapes up to the scission configuration. A typical calculation of chaoticity for a given shape proceeds as follows. The initial coordinates of a classical trajectory starting from the nuclear surface are chosen by sampling a suitably defined set of random numbers such that all initial coordinates follow a uniform distribution over the nuclear surface. The initial direction of the trajectory is also chosen randomly and its Lyapunov exponent is then obtained by following the trajectory for a considerable length of time. Each trajectory is identified either as a regular or as a chaotic one by considering the magnitude of its Lyapunov exponent and the nature of its variation with time. The details of this procedure are given in Ref. [22].

We have calculated the chaoticity over a range of shapes from oblate to the scission configuration (at $c=2.08$ where the neck radius becomes zero) at small steps of c, the elongation coordinate. Figure 1 shows the calculated values of chaoticity which will be subsequently employed to obtain the chaos-weighted wall formula friction. It is important to note here that chaoticity is very small for near-spherical shapes ($c\sim 1$). This immediately implies, through Eq. (1.1), a strong suppression of the original wall formula friction for compact shapes of the compound nucleus. Chaoticity, however, increases as the shape becomes more oblate or changes towards the scission configuration. We find here that the full chaotic regime ($\mu=1$) in the single-particle dynamics is not reached even at the scission configuration. This observation is specific to the parametric form of the shape [Eq. (2.1)] used in the present calculation. It was observed earlier that the value of chaoticity reaches 1 near the scission point when the Legendre-polynomial P_2-deformed quadrupole shapes were considered [16]. It is difficult to speculate to what extent the final calculated observables will be sensitive to the choice of the shape parametrization.

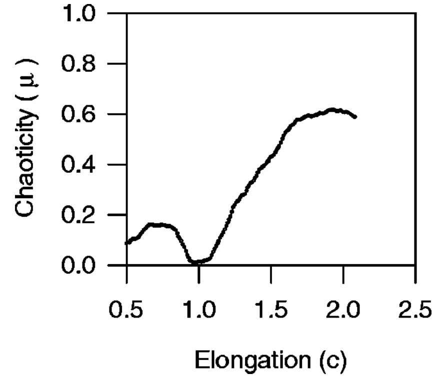

FIG. 1. Variation of chaoticity with elongation.

We shall use the following expression to calculate the wall formula friction coefficient [23]:

$$\eta_{wf}(c) = \frac{1}{2}\pi\rho_m\bar{v}\left\{\int_{z_{\min}}^{z_N}\left(\frac{\partial\rho^2}{\partial c}+\frac{\partial\rho^2}{\partial z}\frac{\partial D_1}{\partial c}\right)^2 \times\left[\rho^2+\left(\frac{1}{2}\frac{\partial\rho^2}{\partial z}\right)^2\right]^{-1/2} dz + \int_{z_N}^{z_{\max}}\left(\frac{\partial\rho^2}{\partial c}+\frac{\partial\rho^2}{\partial z}\frac{\partial D_2}{\partial c}\right)^2 \times\left[\rho^2+\left(\frac{1}{2}\frac{\partial\rho^2}{\partial z}\right)^2\right]^{-1/2} dz\right\}, \tag{2.6}$$

where ρ_m is the mass density of the nucleus, $\bar{v}$ is the average nucleon speed inside the nucleus, and D_1, D_2 are positions of the centers of mass of the two parts of the fissioning system relative to the center of mass of the whole system. $z_{\min}$ and $z_{\max}$ are the two extreme ends of the nuclear shape along the z axis and z_N is the position of the neck plane that divides the nucleus into two parts. The chaos-weighted wall formula friction is subsequently obtained from Eq. (1.1) as $\eta_{cwwf}(c)=\mu(c)\eta_{wf}(c)$. Defining a quantity $\beta(c)=\eta(c)/m(c)$ as the reduced friction coefficient, its dependence on the elongation coordinate is shown in Fig. 2 for

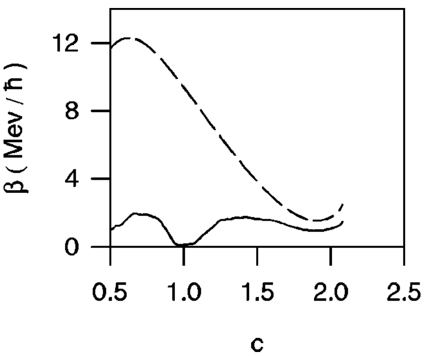

FIG. 2. Variation of the reduced friction coefficient β with elongation c for chaos-weighted wall formula (solid line) and wall formula (dashed line) frictions.

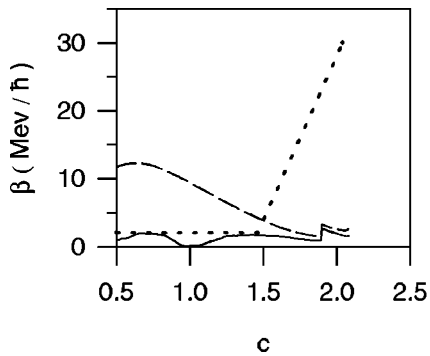

FIG. 3. Reduced one-body friction coefficient β with chaos weighted wall formula (solid line) and wall formula (dashed line) frictions. The phenomenological reduced coefficient (dotted line) from Ref. [9] is also shown.

both WF and CWWF frictions for the ^{200}Pb nucleus. The reduction in the strength of the wall friction due to chaos considerations is evident from this figure.

We shall now consider the role of window friction in one-body dissipation. The window friction is expected to be effective after a neck is formed in the nuclear system [23]. Further, the radius of the neck connecting the two future fragments should be sufficiently narrow in order to enable a particle that has crossed the window from one side to the other to remain within the other fragment for a sufficiently long time. This is necessary to allow the particle to undergo a sufficient number of collisions within the other side and make the energy transfer irreversible. It therefore appears that the window friction should be very nominal when neck formation just begins. Its strength should increase as the neck becomes narrower, reaching its classical value when the neck radius becomes much smaller than the typical radii of the fragments. We however know very little regarding the detailed nature of such a transition. We shall therefore refrain from making any further assumption regarding the onset of window friction. Instead, we shall define a critical elongation coordinate c_{win} beyond which the window friction will be switched on. The window friction coefficient will then be given as

$$\eta_{win}(c)=\theta(c-c_{win})\frac{1}{2}\rho_m\bar{v}\left(\frac{\partial R}{\partial c}\right)^2\Delta\sigma, \tag{2.7}$$

where

$$\theta(c-c_{win})=0 \quad (\text{for} \quad c<c_{win})$$
$$=1 \quad (\text{for} \quad c\geqslant c_{win}),$$

and R is the distance between centers of mass of future fragments and $\Delta\sigma$ is the area of the window between the two parts of the system. The full one-body friction will now be written as

$$\eta(c)=\eta_{wall}(c)+\eta_{win}(c), \tag{2.8}$$

and in what follows, we will use either η_{wf} or η_{cwwf} for η_{wall} in the above expression. For the window friction, the value of c_{win} is taken as 1.9 where the neck radius is half of the fragment radius. Figure 3 shows the reduced one-body friction coefficients. The phenomenological reduced friction obtained in Ref. [9] is also shown in Fig. 3. Though the one-body friction with CWWF agrees qualitatively with the phenomenological friction for $c<1.5$, it is beyond its scope to explain the steep increase of phenomenological friction for $c>1.5$. It may be noted, however, that the compulsion of having a very strong friction at large deformations was to allow a sufficient number of neutrons to evaporate during the saddle-to-scission transition (i.e., after fission has taken place) in order to fit the experimental prescission neutron multiplicities for very heavy nuclei [9]. Therefore, the role of a very strong friction beyond the saddle point will not be significant for fission rates which is of our present concern.

III. RESULTS

With all the necessary input defined as above, the Langevin equation (2.2) is numerically integrated following the procedure outlined in Ref. [6]. A very small time step of $0.005\hbar/\text{MeV}$ for numerical integration is used in the present work. The numerical stability of the results is checked by repeating a few calculations with still smaller time steps. The initial distribution of the coordinates and momenta is assumed to be close to equilibrium and hence the initial values of (c,p) are chosen from sampling random numbers following the Maxwell-Boltzmann distribution. Starting with a given total excitation energy (E^*) and angular momentum (l) of the compound nucleus, energy conservation in the form

$$E^*=E_{int}+V(c)+p^2/2m \tag{3.1}$$

gives the intrinsic excitation energy E_{int} and the corresponding nuclear temperature $T=(E_{int}/a)^{1/2}$ at each integration step. The centrifugal potential is included in $V(c)$ in the above equation. A Langevin trajectory will be considered as having undergone fission if it reaches the scission point (c_{sci}) in the course of its time evolution. The calculations are repeated for a large number (typically 100 000 or more) of trajectories and the number of fission events is recorded as a function of time. Subsequently the fission rates can be easily evaluated [24].

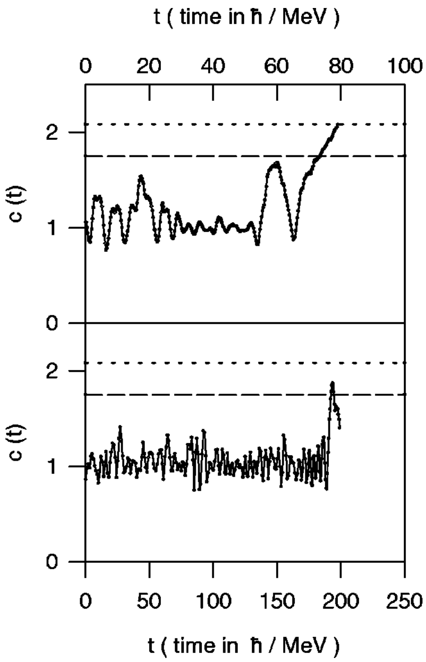

FIG. 4. The upper panel shows a typical Langevin trajectory reaching the scission point (dotted line). The lower panel shows a trajectory which returns to the potential pocket after crossing the saddle point (dashed line).

A typical Langevin trajectory that has reached the scission point and has ended up as a fission event is shown in Fig. 4 (upper panel). Another trajectory, the kind which is less frequent, is shown in the lower panel of the same figure. The Langevin trajectory in the latter case crosses the saddle point and after spending some time beyond the saddle point drifts back into the potential pocket. Such a trajectory may or may not finally reach the scission point within the observation time and corresponds to a to-and-fro motion across the saddle. This point is further illustrated in Fig. 5 where fission rates are plotted as a function of time. Two different criteria are used to define a fission event here. The solid circles correspond to fission events defined by those trajectories that reach the scission point whereas the open circles correspond to those which cross the saddle point. The fission rate is very small for both cases at the beginning when the compound nucleus is just formed and the Langevin dynamics has just been turned on. Subsequently the fission rate grows with time and after a certain equilibration time it reaches a stationary value that corresponds to a steady flow across the fission barrier. The fission rate defined at the saddle point reaches a stationary value earlier than that defined at the scission point. The time difference between them gives the average time of descent from the saddle to the scission. This observation was also made in earlier works [25]. The main purpose of the present discussion is to investigate the role of backstreaming in the fission process. It is observed in Fig. 5 that the stationary fission rate at the saddle point is higher than that at the scission point. The difference between these two stationary rates can be regarded as due to backstreaming. The backstreaming is thus small compared to the steady outward flow though it is not negligible. This also shows that

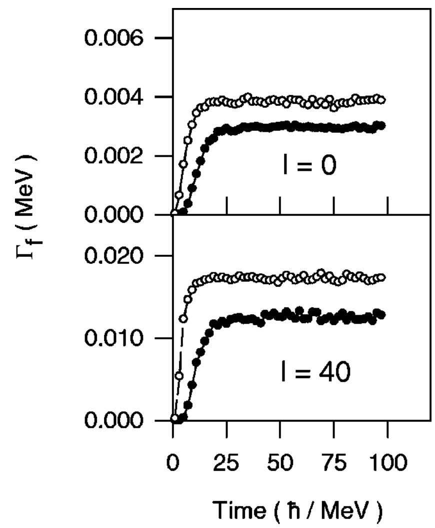

FIG. 5. Time development of fission widths for compound nuclear spins of 0 and 40 (in units of $\hbar$). Open circles correspond to trajectories for which the saddle point crossing is considered as fission. Solid circles represent trajectories which reach the scission point.

crossing the saddle point is not an adequate criterion for fission in stochastic calculations and can lead to an overestimation of the fission rate.

We shall now compare fission rates calculated with chaos-weighted wall-and-window friction with those obtained with wall-and-window friction [Eq. (2.8)]. Figure 6 shows the fission widths at three spins of the compound nucleus ^{200}Pb.

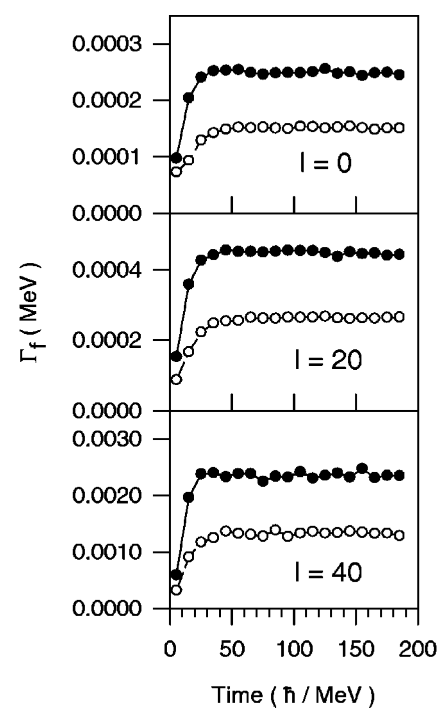

FIG. 6. Time development of fission widths calculated with chaos-weighted wall formula (solid circles) and wall formula (open circles) frictions for different compound nuclear spins l (in units of $\hbar$).

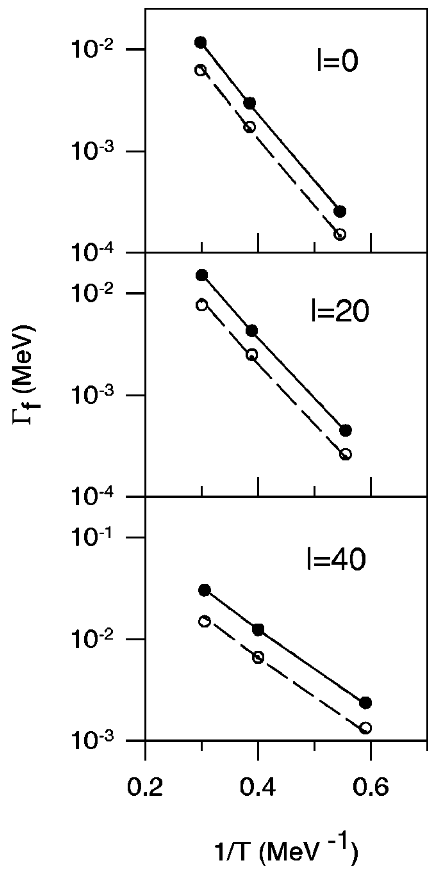

FIG. 7. Temperature dependence of stationary fission widths calculated with chaos-weighted wall formula (solid circles) and wall formula (open circles) frictions for different compound nuclear spins l (in units of $\hbar$). The lines are fitted as explained in the text.

The effect of suppression in the chaos-weighted wall formula shows up as an enhancement by about a factor of 2 of the stationary fission rates. A similar enhancement of the stationary fission rate calculated with chaos-weighted wall-and-window friction in comparison with that obtained with wall-and-window friction is also observed over a wide range of compound nuclear spin and temperature. The enhancement factor (of about 2) remains almost the same when different choices of c_{win} are used in the window friction [Eq. (2.7)].

We shall next extract the stationary fission widths systematically at different temperatures for a given spin of the compound nucleus. This is done by taking the average of the fission rates in the plateau region. These fission rates are essentially the Kramers' limit of the Langevin equation under consideration and we expect the stationary fission widths Γ_f to depend upon the temperature T as $\Gamma_f(l,T) = A_l \exp(-b_f/T)$ for a given spin (l) of the compound nucleus where b_f is the height of the fission barrier in the free energy profile and A_l is a parameter. Such a dependence of stationary fission widths on temperature is indeed found and is shown in Fig. 7. The parameter A_l can now be extracted by fitting the calculated fission widths with the above expression. Subsequently we shall look into the dependence of the parameter A_l on l, a few typical plots of which are shown in Fig. 8. Using these values of A_l, one can now obtain the value of this parameter for any arbitrary spin by interpolation. Even with a limited number of calculated values, the interpolated values will be quite accurate because A_l depends on l rather weakly as can be seen in Fig. 8. Consequently it will now be possible to extract the fission width of a compound nucleus at any given temperature and spin from a set of a limited number of calculated widths. This fact will be very useful in statistical model calculations where fission widths are required at numerous values of temperature and spin which are encountered during the evolution of a compound nucleus. Therefore in such cases where analytical expressions for fission widths are not available, the above systematic behavior can generate fission widths from a limited set of calculations.

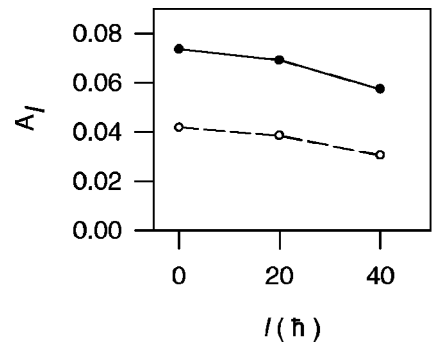

FIG. 8. Variation of the parameter A_l with compound nuclear spin l.

Two time scales are of physical significance in the Langevin description of the dynamics of fission. One is the equilibration time τ_{eq}, the time required to attain a steady flow across the barrier. The other is the fission lifetime $\tau_f = \hbar/\Gamma_f$. Figure 9 shows these time intervals for different values of spin of the compound nucleus ^{200}Pb. At very small values of spin, the fission life time is many times longer than the equilibration time. This means that a statistical theory for compound nuclear decay is applicable in such cases. On the other hand, τ_{eq} and τ_f become comparable at higher values of the compound nuclear spin and this corresponds to a dynamics dominated decay of the compound nucleus. Statistical models are not meaningful in these cases and dynamical descriptions such as the Langevin equation become essential for the fission of such compound nuclei.

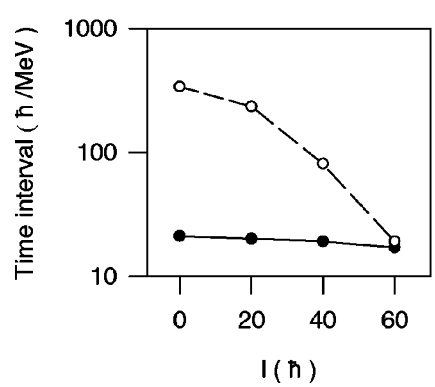

FIG. 9. Dependence of the equilibration time τ_{eq} (solid circles) and the fission lifetime τ_f (open circles) on compound nuclear spin l.

In the above, we restricted our investigations to the ^{200}Pb nucleus as a representative example. This nucleus has been experimentally formed in ^{19}F+ ^{181}Ta reactions at a number of excitation energies [26–28]. The fission probability of the compound nucleus and the prescission neutron and γ multiplicities were measured in these experiments. These quantities can be calculated in a statistical model which requires the fission width as well as the neutron and γ emission widths as input to the calculation. In particular, the input fission width plays a critical role in order to reproduce the experimentally determined prescission neutron and γ multiplicities at high excitation energies (typically a few tens of MeV or higher) in statistical model calculations [27,28]. While the neutron and γ widths can be obtained from the Weisskopf formula [6], a dynamical theory is required to calculate the fission width of the hot compound nucleus. The present work is a step in this direction and the fission widths obtained here can thus serve as input to statistical model calculations.

IV. SUMMARY AND OUTLOOK

In the preceding sections, we have presented a systematic study of fission dynamics using the Langevin equation. Among the various physical input required for solving the Langevin equation, we paid particular attention to the dissipative force for which we chose the wall-and-window one-body friction. We used a modified form of wall friction, the chaos-weighted wall formula, in our calculation. The chaos-weighted wall formula took into account the nonintegrabilty of single-particle motion in the nucleus and it resulted in a strong suppression of friction strength for near-spherical shapes of the nucleus. The fission widths calculated with the chaos-weighted wall formula turned out to be about twice the widths calculated with the normal wall formula friction. Chaos-weighted wall friction thus enhances the fission rate substantially compared to that obtained with normal wall friction.

We further made a parametric representation of the calculated fission widths in terms of the temperature and spin of the compound nucleus. It was found that this parametric form can be well determined from the fission widths calculated over a grid of spin and temperature values of limited size. This fact would make it possible to perform statistical model calculations of the decay of a highly excited compound nucleus where the fission widths are to be determined from a dynamical model such as the Langevin equation. When the friction form factor has a strong shape dependence as in the chaos-weighted wall formula, the corresponding fission widths cannot be obtained in an analytic form. In such cases, the frequently required values of the fission width in a statistical model calculation can be made economically accessible through a parametric representation of the fission width which has to be obtained in a separate calculation similar to the present one. We shall report on such applications of the parametrized fission widths in compound nuclear decay in our future works.

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