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Instytut Fizyki UMCS
GSI Darmstadt

P. ROZMEJ*, W. NÖRENBERG

Dynamical Diabatic Hindrance of Heavy-Ion Fusion

The fusion process for heavy nuclei has been extensively studied within a generalized formulation of the diabatic approach to collective nuclear motion [1]. It was shown there that coupled equations of motion decouple during the approach if appropriate diabatic channels are used. Thus the mean value and the variance of the fusion barrier are completely determined from the ensamble of diabatic barriers being defined by initial correlations. The mean shift ΔB with respect to the adiabatic barrier and the variance σ_B^2 have been calculated from the splitting of the diabatic single particle levels at the barrier and the initial occupation probabilities given by the pairing wavefunctions of the separated nuclei.

The microscopic model applied in [1] reproduces both the general (smooth) dependence of the mean shift ΔB on the mean fusibility parameter x_m and the strong opposite isotopic dependence due to shell structure present in experimental data [2,3]. The accuracy of the theoretical predictions decreases, however, with the increasing fusibility. This is a signature that the existing theoretical model does not contain a physical mechanism which becomes important for heavier systems and sets on for $x_m \gtrsim 0.75$.

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The principal results of the paper [1] are presented in figure 1.

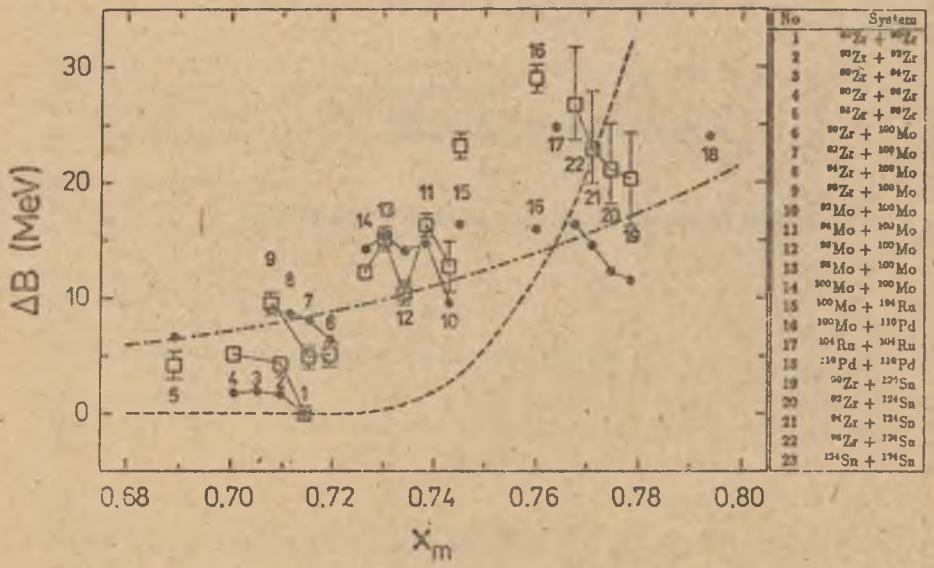

Fig. 1. Experimental (squares) and theoretical barrier shifts for 23 fusion systems as functions of the mean fusibility. The full dots give the values for the pure diabatic shift. The extra push obtained from the one–body dissipation model [4] and the surface–friction model [5] are shown by the dashed and dash–dotted lines, respectively.

The model used in [1] can be considered as the diabatic limit. Barrier shifts and fluctuations are calculated under the assumption that the single–particle motion during the approach phase of the collision is purely diabatic. For heavier systems, as an overlap of the colliding nuclei at the barrier distance increases with fusibility, one can expect the increasing influence of two–body collisions. These collisions produce dissipation that generally broadens the barrier distribution. The mean shift is also increased because the additional energy ('diabatic extra–push') is required by the system to reach the fusion configuration. Two–body collisions are incorporated in DDD theory (Dissipative Diabatic Dynamics) [6]. The present paper aims to formulate DDD in a way convenient to describe the fusion process of heavy nuclei. As there is no single fusion barrier and contrary there is a definite distribution of fusion barriers (with the mean value and the variance calculated in [1]) we introduce reaction channels occuring with a definite probability. For simplicity we limit our considerations to only one collective degree of freedom q (in our case it is the distance between centers of colliding nuclei). Then the equation of the collective motion

for a particular channel ν reads as

$$B(q)\ddot{q} + \frac{1}{2}\left(\frac{dB}{dq}\right)\dot{q}^2 = F_\nu(q,t)\,, \tag{1}$$

where B is mass parameter and $F_\nu(q,t)$ is the force for the given channel. At the time $t = 0$ the force is given by the derivative of the potential that is the sum of adiabatic part and diabatic shift

$$F_\nu(q, t = 0) = -\frac{d}{dq}\left[V^{\rm ad}(q) + \Delta V_\nu^{\rm diab}(q)\right]\,. \tag{2}$$

The adiabatic potential $V^{\rm ad}$ determining the equilibrium force is taken from the frozen density Hartree–Fock (FDHF) calculations [1]. This paper supply also the distribution of the diabatic shift of the fusion barrier $\Delta V^{\rm diab}$. Choosing the different values of the diabatic shift one can define a particular reaction channel that is populated with the determined probability. Figure 2 shows the example of initial distribution of the fusion barriers obtained in static approach [1]. The mean value represent the sum of the adiabatic (FDHF) barrier and the mean value of the diabatic shift.

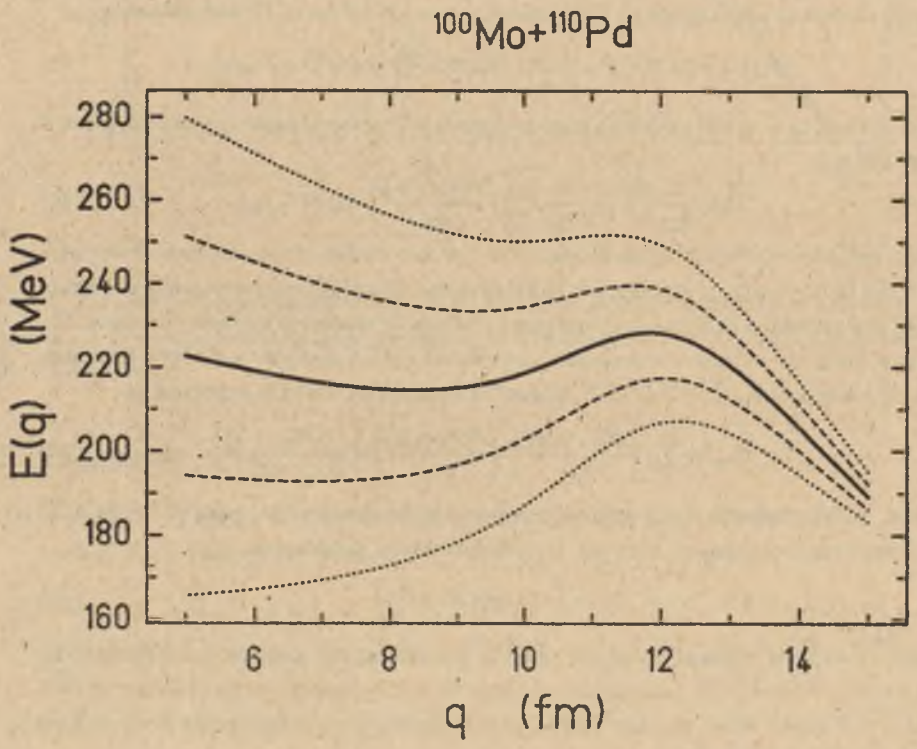

Fig. 2. Initial distribution of fusion barriers for the $^{100}Mo+^{110}Pd$ system. The thick solid line displays the mean barrier and the dashed and dotted lines correspond to barriers different from the mean value by one and 2 standard deviations, respectively.

Due to two–body collisions the force $F_\nu(q,t)$ fluctuates and is driven to its thermal equilibrium value. We may describe the stochastic behaviour of the force by a distribution function which satisfies a Fokker–Planck equation. Alternatively we introduce the stochastic variable $F_\nu(q,t)$ which satisfies the Langevin–type equation

$$\frac{\partial F_\nu(q,t)}{\partial t} = -\frac{1}{\tau_{loc}} \left[F_\nu(q,t) - \overline{F}(q)\right] + F_\nu^{fluct}(q,t)\,, \tag{3}$$

where τ_{loc} denotes local equilibration time.

This time has been estimated by Bertsch [7] within Fermi–gas model as

$$\tau_{loc} = \frac{2 \cdot 10^{-22}\mathrm{s} \cdot \mathrm{MeV}}{\epsilon^*} \tag{4}$$

and depends on the excitation energy ϵ^* per particle (which is related to the temperature). $\overline{F}(q)$ stands for the adiabatic force. The fluctuating force $F_\nu^{fluct}(q,t)$ is an irregular function in time with zero mean and a white spectrum (δ – functions in time) [8]. It can be related to the diffusion coefficient D of the corresponding Fokker–Planck equation [9]. For the actual numerical calculations we discretize equation (3) following [9] and obtaining

$$F_\nu(t+\tau) - F_\nu(t) = -\frac{1}{\tau_{loc}} \left[F_\nu(t) - \overline{F}\right] \tau + \sqrt{D}\,\sqrt{\tau}\,\omega(t)\,. \tag{5}$$

Here τ is a time step and $\omega(t)$ is a gaussian distributed random number, such that $\langle\omega\rangle = 0$ and $\langle\omega^2\rangle = 2$.

$$D = \frac{2}{\tau_{loc}}\,\overline{\sigma_F^2} \cong \frac{2}{\tau_{loc}} \sum_a \left(\frac{\partial(\varepsilon_a - \mu)}{\partial q}\right) \overline{n}_a\,(1 - \overline{n}_a) \tag{6}$$

is the diffusion coefficient which is related to the mean value of the variance of the force $\overline{\sigma_F^2}$. The last quantity is evaluated in the right hand side of (6) microscopically according to [8] and presented in figure 3. It appears to be closely related to stiffness parameter C_{qq} (tensor for more collective coordinates) introduced within dissipative diabatic dynamics [10,11] via equation $\overline{\sigma_F^2} = T \cdot C(T)$. Alternatively stiffness tensor is expressed as

$$C_{nm} = \sum_a \left\{\frac{\partial(\varepsilon_a - \mu)}{\partial q_n}\right\} \left\{\frac{\partial(\varepsilon_a - \mu)}{\partial q_m}\right\} \left\{\frac{\partial \overline{n}_a}{\partial(\varepsilon_a - \mu)}\right\} \tag{7}$$

The ε_a are the diabatic single–particle energies, μ is the chemical potential and $\overline{n}_a(q,\mu,T)$ are equilibrium occupation numbers according to Fermi distribution :

$$\overline{n}_a = \left\{1 + \exp\left(\frac{\varepsilon_a - \mu}{T}\right)\right\}. \tag{8}$$

Figure 3 displays stiffness parameter (sum of the proton and neutron contributions) for the system ^{100}Mo+^{110}Pd calculated from diabatic single–particle states obtained in two–centre shel model which parameters had been adjusted to reproduce initial Hartree–Fock single–particle energies [1].

Numerical calculations of the fusion probability are not finished yet for the dynamical model. Preliminary results (which estimate the mean dynamical shift of the fusion barrier) idicate that the dynamics of the process can indeed be responsible for the previous discrepancy between the theoretical and experimental results.

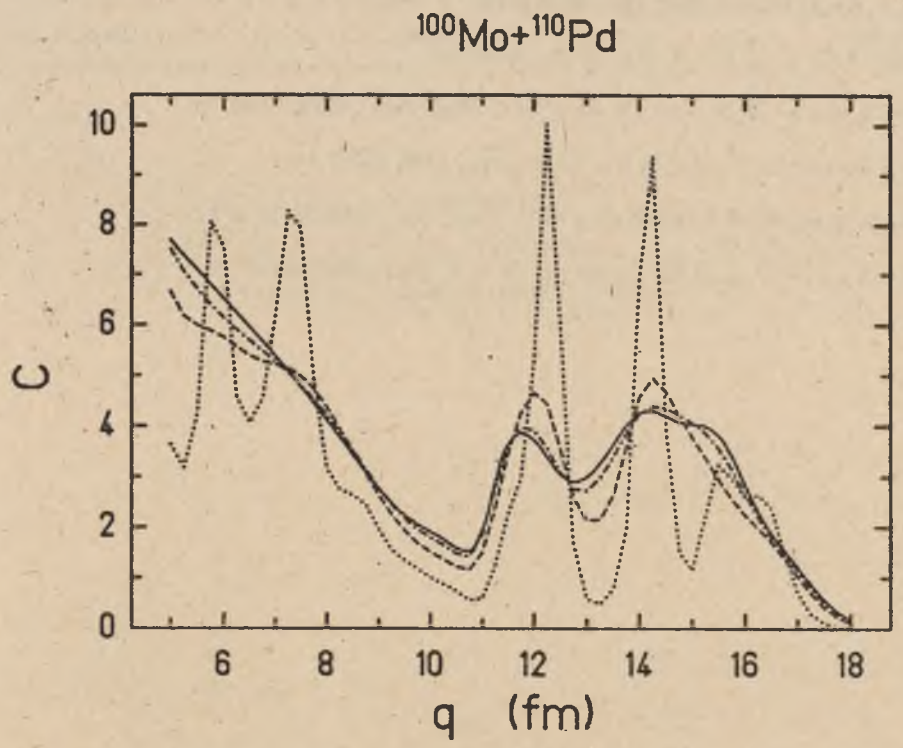

Fig. 3. The stiffness parameter for $^{100}Mo+^{110}Pd$ system at different temperatures. Dotted, dashed, dash–dotted and solid lines correspond to temperatures $T = 0.2$, 0.6, 1.0 and 1.4 MeV respectively.

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