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-A-UR -84-1566

CONF - 8405179--1

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LA-UR--84-1566

DE84 012441

TITLE: DYNAMICS OF FISSION AND HEAVY ION REACTIONS

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SUBMITTED TO: to be presented at the International Conference on Theoretical Approaches to Heavy Ion Reaction Mechanisms, Paris, France, May 14-18, 1984

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DYNAMICS OF FISSION AND HEAVY ION REACTIONS

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We discuss recent advances in a unified macroscopic-microscopic description of large-amplitude collective nuclear motion such as occurs in fission and heavy ion reactions. With the goal of finding observable quantities that depend upon the magnitude and mechanism of nuclear dissipation, we consider one-body dissipation and two-body viscosity within the framework of a generalized Fokker-Planck equation for the time dependence of the distribution function in phase space of collective coordinates and momenta. Proceeding in two separate directions, we first solve the generalized Hamilton equations of motion for the first moments of the distribution function with a new shape parametrization and other technical innovations. This yields the mean translational fission-fragment kinetic energy and mass of a third fragment that sometimes forms between the two end fragments, as well as the energy required for fusion in symmetric heavy-ion reactions and the mass transfer and capture cross section in asymmetric heavy-ion reactions. In a second direction, we specialize to an inverted-oscillator fission barrier and use Kramers' stationary solution to calculate the mean time from the saddle point to scission for a heavy-ion-induced fission reaction for which experimental information is becoming available.

1. INTRODUCTION

We have already heard from the other speakers at this International Conference on Theoretical Approaches to Heavy Ion Reaction Mechanisms about the many complementary aspects displayed by the atomic nucleus. With its relatively small number of degrees of freedom, the nucleus is both microscopic and macroscopic on the one hand and both quantal and classical on the other, which gives it a rich dynamical behaviour ranging from elastic vibrations of solids to long-mean-free-path dissipative fluid flow with statistical fluctuations. Experimental clues to this challenging many-body problem continue to be provided by fission and heavy ion reactions. Yet our major goal of determining whether in large-amplitude collective nuclear motion the nucleons interact primarily through the mean field generated by the remaining nucleons, or whether two-particle collisions play a substantial role, has proved elusive.

Our challenge is not to explain the experimental data in terms of some model with adjustable parameters--since often several models with widely different physical bases are capable of doing this equally well--but instead to find and calculate physical observables that depend sensitively upon the magnitude and

*This work was supported by the U.S. Department of Energy.

mechanism of nuclear dissipation. The difficulty arises because many of the gross experimental features of fission and heavy ion reactions are determined primarily by a competition between the attractive nuclear force and the repulsive Coulomb and centrifugal forces, and any theoretical approach that includes correctly these relatively trivial forces reproduces the data with fair accuracy. Also, the final effects on observable quantities caused by dissipation are often very similar to the final effects caused by collective degrees of freedom.

A possible starting point for a theory of nuclear dynamics is the time-dependent mean-field (Hartree-Fock) approximation, in which nucleons interact only through the mean field generated by the other nucleons, with two-particle collisions neglected entirely¹. The approximate validity of this approach stems from the Pauli exclusion principle and the details of the nucleon-nucleon interaction, which at low excitation energies lead to a nucleon mean free path that is long compared to the nuclear radius. In this approximation, the many-body wave function for a system of A nucleons is represented at all times by a single Slater determinant consisting of A single-particle wave functions.

However, critical comparisons of two predictions of this approach with experimental data suggest that in real nuclei the type of nuclear dynamics predicted by the time-dependent mean-field approximation is modified significantly by residual interactions arising from two-particle collisions. These comparisons involve the experimental demonstration that nuclei do not penetrate through each other in nearly central collisions as predicted by this approximation, and the substantially smaller predicted energy loss at large angles in heavy ion reactions than is observed experimentally^{1,2}.

Some important steps have been taken to incorporate two-particle collisions into the time-dependent mean-field approximation^{3,4}. Although certain conceptual problems remain and computational difficulties have precluded comparisons of such extended mean-field approximations with experimental results, these studies have nevertheless shown that two-particle collisions cannot be neglected. This motivates us to take the opposite tack and study nuclear dynamics by use of a macroscopic-microscopic method. Our purpose is to calculate for two radically different dissipation mechanisms observable quantities in fission and heavy ion reactions and confront these predictions with experimental data in an attempt to determine the magnitude and mechanism of nuclear dissipation.

2. MACROSCOPIC-MICROSCOPIC METHOD

We focus from the outset on those few collective coordinates that are most relevant to the phenomena under consideration. In particular, for a system of A nucleons, we separate the $3A$ degrees of freedom representing their center-of-mass motion into N collective degrees of freedom that are treated explicitly and

3A - N internal degrees of freedom that are treated implicitly.

2.1. Collective coordinates

In our earlier dynamical studies we have usually described the nuclear shape in terms of smoothly joined portions of three quadratic surfaces of revolution, with three symmetric and two independent asymmetric shape coordinates⁵⁻⁹. Although suitable for many purposes, this three-quadratic-surface parametrization breaks down in the later stages of many heavy-ion fusion calculations, is unable to describe division into more than two fragments and leads to very complicated expressions for the forces involved.

Because of these disadvantages, we have switched to a more suitable parametrization in which an axially symmetric nuclear shape is described in cylindrical coordinates by means of the Legendre-polynomial expansion¹⁰

$$\rho_s^2(z) = R_0^2 \sum_{n=0}^N q_n P_n[(z-\bar{z})/z_0] .$$

In this expression, z is the coordinate along the symmetry axis, ρ_s is the value on the surface of the coordinate perpendicular to the symmetry axis, z_0 is one-half the distance between the two ends of the shape, \bar{z} is the value of z at the midpoint between the two ends, R_0 is the radius of the spherical nucleus, P_n is a Legendre polynomial of degree n , and q_n for $n \neq 0$ and 1 are $N - 1$ shape coordinates. Since the nucleus is assumed to be incompressible, the quantity q_0 is not independent but is instead determined by volume conservation. Also, q_1 is determined by fixing the center of mass. Throughout this paper we use $N = 11$, corresponding to five independent symmetric and five independent asymmetric shape coordinates. In addition, we include an angular coordinate $\theta \equiv q_{N+1}$ to describe the rotation of the nuclear symmetry axis in the reaction plane, which leads to a total of N degrees of freedom that are considered.

2.2. Potential energy

In terms of the N collective coordinates $q = q_2, \dots, q_{N+1}$ we calculate the potential energy of deformation $V(q)$ as the sum of repulsive Coulomb and centrifugal energies and an attractive Yukawa-plus-exponential potential¹¹, with constants determined in a recent nuclear mass formula¹². This generalized surface energy takes into account the reduction in energy arising from the nonzero range of the nuclear force in such a way that saturation is ensured when two semi-infinite slabs are brought into contact.

2.3. Kinetic energy

The collective kinetic energy is given by

$$T = \frac{1}{2} M_{ij}(q) \dot{q}_i \dot{q}_j = \frac{1}{2} [M(q)^{-1}]_{ij} p_i p_j ,$$

where the collective momenta p are related to the collective velocities \dot{q} by

$$p_i = M_{ij}(q) \dot{q}_j .$$

In these equations and the remainder of this paper we use the convention that repeated indices are to be summed over from 2 to $N + 1$. We calculate the inertia tensor $M(q)$, which is a function of the shape of the system, for a superposition of rigid-body rotation and incompressible, nearly irrotational flow by use of the Werner-Wheeler method, which determines the flow in terms of circular layers of fluid⁵⁻⁹.

2.4. Dissipation mechanisms

The coupling between the collective and internal degrees of freedom gives rise to a dissipative force whose mean component in the i -th direction may be written as

$$F_i = - \eta_{ij}(q) \dot{q}_j = - \eta_{ij}(q) [M(q)^{-1}]_{jk} p_k .$$

For the calculation of the shape-dependent dissipation tensor $\eta(q)$ that describes the conversion of collective energy into single-particle excitation energy, we consider both one-body dissipation^{8,9,13-15} and ordinary two-body viscosity^{6,8,9}, whose dissipation mechanisms represent opposite extremes. In the former case dissipation arises from collisions of nucleons with the moving nuclear surface and when the neck is smaller than a critical size also from the transfer of nucleons through it, with a magnitude that is completely specified by the model. In the latter case dissipation arises from collisions of nucleons with each other, but the coefficient of two-body viscosity must be determined from an adjustment to experimental results.

Compared to our previous calculations with one-body dissipation^{8,9,13}, our present calculations incorporate two improvements. First, to describe the transition from the wall formula that applies to mononuclear shapes to the wall-and-window formula that applies to dinuclear shapes we now use the smooth interpolation

$$\eta = \sin^2\left(\frac{\pi}{2}\alpha\right) \eta_{\text{wall}} + \cos^2\left(\frac{\pi}{2}\alpha\right) \eta_{\text{wall-and-window}} ,$$

where

$$\alpha = (r_{\text{neck}}/R_{\text{min}})^2$$

is the square of the ratio of the neck radius r_{neck} to the transverse semi-axis R_{min} of the end fragment with the smaller value. Second, in determining the drift velocities of the end fragments relative to which velocities in the wall-and-window formula are measured, we now require the conservation of linear and angular momentum rather than using the velocities of the centers of mass. However, the results calculated with both prescriptions for the drift velocity are nearly identical.

2.5. Generalized Fokker-Planck equation

In addition to the mean dissipative force, the coupling between the collective and internal degrees of freedom gives rise to a residual fluctuating force. When this stochastic force is treated with classical statistical mechanics under the Markovian assumption that it does not depend upon the system's previous history, we are led to the generalized Fokker-Planck equation

$$\begin{aligned} \frac{\partial f}{\partial t} + (M^{-1})_{ij} p_j \frac{\partial f}{\partial q_i} - \left[\frac{\partial V}{\partial q_i} + \frac{1}{2} \frac{\partial (M^{-1})_{jk}}{\partial q_i} p_j p_k \right] \frac{\partial f}{\partial p_i} \\ = \eta_{ij} (M^{-1})_{jk} \frac{\partial}{\partial p_i} (p_k f) + \tau \eta_{ij} \frac{\partial^2 f}{\partial p_i \partial p_j} \end{aligned}$$

for the dependence upon time t of the distribution function $f(q,p,t)$ in phase space of collective coordinates and momenta. The last term on the right-hand side of this equation describes the spreading of the distribution function in phase space, with a rate that is proportional to the dissipation strength and the nuclear temperature τ , which is measured here in energy units.

2.6. Generalized Hamilton equations

Because of the practical difficulty of solving the generalized Fokker-Planck equation exactly except for special cases, in some of our studies we use equations for the time rate of change of the first moments of the distribution function, with the neglect of higher moments. These are the generalized Hamilton equations

$$\dot{q}_i = (M^{-1})_{ij} p_j$$

and

$$\dot{p}_i = - \frac{\partial V}{\partial q_i} - \frac{1}{2} \frac{\partial (M^{-1})_{jk}}{\partial q_i} p_j p_k - \eta_{ij} (M^{-1})_{jk} p_k,$$

which we solve numerically for each of the N generalized coordinates and momenta.

3. FISSION

As our first application, we consider the fission process, with particular emphasis on the mean translational fission-fragment kinetic energies of nuclei throughout the periodic table. Although similar to earlier studies^{6,8}, our present calculations are performed, as discussed above, with a more flexible shape parametrization, with a more realistic set of constants, and with two improvements in our treatment of one-body dissipation. Also, our initial conditions at the fission saddle point now incorporate the effect of dissipation on the fission direction¹⁶ and are calculated for excited nuclei with nuclear

temperature $\tau = 2$ MeV by determining the mean velocity of all nuclei that pass per unit time through the saddle point with positive velocity¹⁷.

3.1. Dynamical trajectories

In our fission calculations we specialize to reflection-symmetric shapes and zero angular momentum, so that only five coordinates are considered explicitly. We then project out of this five-dimensional space the two most important symmetric degrees of freedom, which are conveniently defined in terms of the central moments⁶⁻⁹

$$r = 2\langle z \rangle$$

and

$$\sigma = 2\langle (z - \langle z \rangle)^2 \rangle^{1/2},$$

where the angular brackets $\langle \rangle$ denote an average over the half volume to the right of the midplane of the reflection-symmetric shape. The moment r gives the distance between the centers of mass of the two halves of the dividing nucleus and σ measures the elongation of each half about its center of mass. Our calculations are performed for nuclei with atomic number Z related to mass number A according to Green's approximation to the valley of beta stability¹⁸.

As shown in fig. 1, the mean dynamical trajectories for light nuclei correspond to short descents from dumbbell-like saddle-point shapes to compact scission shapes, whereas those for heavy nuclei correspond to long descents from cylinder-like saddle-point shapes to elongated scission shapes. The trajectories for one-body dissipation lie inside those for no dissipation and two-

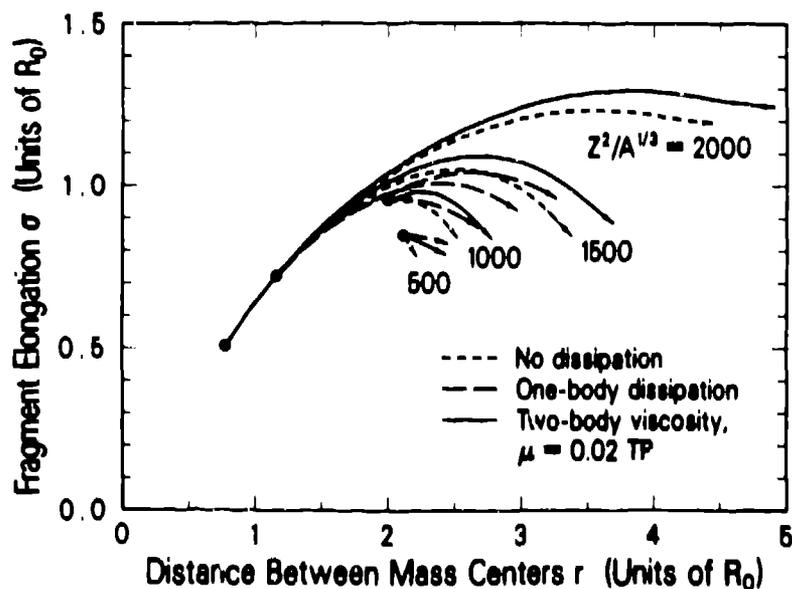


FIGURE 1

Effect of dissipation on mean dynamical trajectories for four fissioning nuclei. Saddle points are indicated by solid circles and scission points by arrowheads.

body viscosity for heavy nuclei, but for very light nuclei lie outside. In contrast, the trajectories for two-body viscosity always lie somewhat above those for no dissipation, leading to more elongated scission shapes. Our calculations for two-body viscosity are performed with viscosity coefficient

$$\mu = 0.02 \text{ TP} = 1.25 \times 10^{-23} \text{ MeV s/fm}^3 ,$$

which as we see later is the value required to optimally reproduce experimental mean fission-fragment kinetic energies.

3.2. Ternary division

An exciting new aspect of these dynamical calculations is the formation of a third fragment between the two end fragments for sufficiently heavy nuclei with either no dissipation or two-body viscosity. As shown in fig. 2, the mass of this third fragment increases with increasing $Z^2/A^{1/3}$ above a critical value that is slightly lower for two-body viscosity than for no dissipation. Since no third fragment is formed with one-body dissipation, accurate experimental information concerning such true ternary-fission processes should help decide the nuclear-dissipation issue. Further theoretical aspects of this problem are currently being studied at Los Alamos by Cârjan¹⁹.

3.3. Fission-fragment kinetic energies

In calculating the mean fission-fragment translational kinetic energy at infinity, we treat the post-scission dynamical motion in terms of two spheroids, with initial conditions determined by keeping continuous the values of r , σ , \dot{r} ,

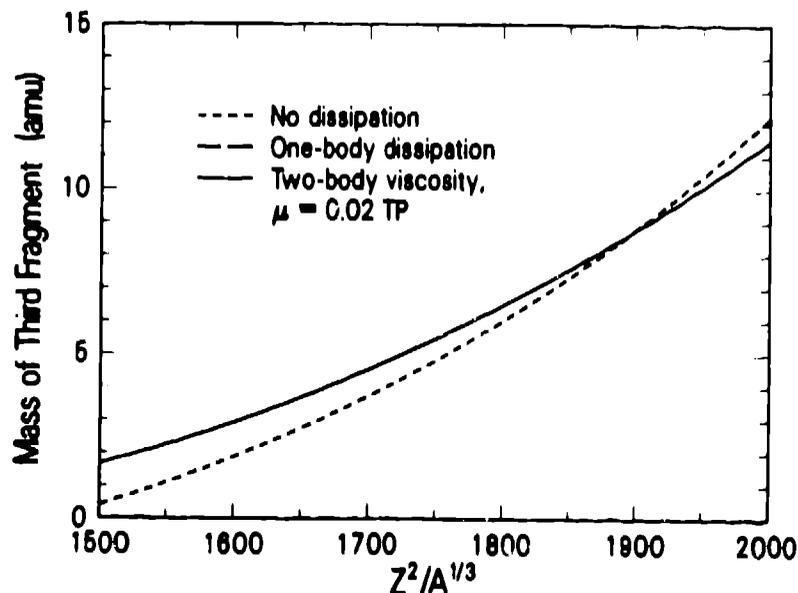


FIGURE 2

Effect of dissipation on the formation of a third fragment between the two end fragments.

and σ at scission. When a small third fragment is formed in a realistic situation off the symmetry axis and/or with some transverse velocity, it moves away and contributes less to the kinetic energy of the two larger end fragments than it would in our idealized calculation, where it remains stationary at its origin. In the presence of a third fragment, we obtain a lower limit to the fission-fragment kinetic energy by calculating the post-scission separation of the end fragments in the absence of the middle fragment. Also, we estimate an upper limit in terms of the kinetic energy at scission of the two end fragments plus the Coulomb interaction energy of three spherical nuclei positioned at their respective centers of charge.

We compare in figs. 3 and 4 our mean kinetic energies calculated in this way with experimental values for the fission of nuclei at high excitation energy⁶, where single-particle effects have decreased in importance. As shown by the short-dashed curves in both figures, the results calculated with no dissipation are for heavy nuclei substantially higher than the experimental values. Dissipation of either type lowers the calculated kinetic energy. However, as shown by the long-dashed curve in fig. 3, one-body dissipation with a magnitude that is specified by the theory predicts for heavy nuclei values that lie below the experimental data. This underprediction arises because the highly dissipative descent from the saddle point damps out much of the pre-scission kinetic energy

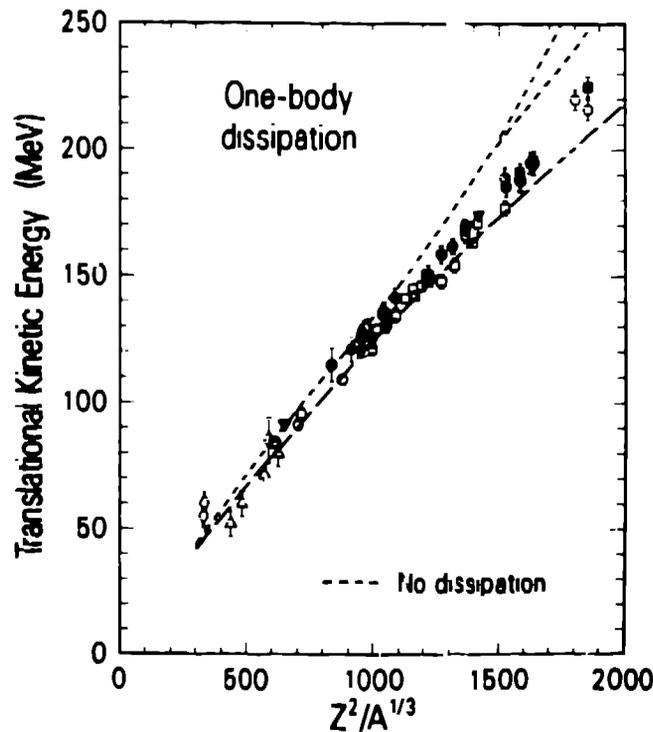


FIGURE 3
Reduction of mean fission-fragment kinetic energies by one-body dissipation, compared to experimental values.

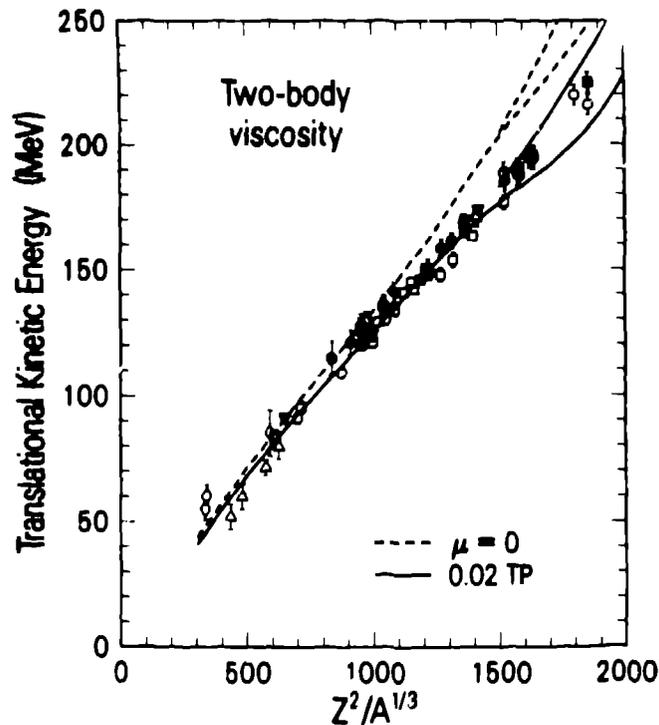


FIGURE 4

Reduction of mean fission-fragment kinetic energies by two-body viscosity, compared to experimental values.

and our improved parametrization leads to moderately elongated scission shapes with lower Coulomb repulsion. We regard this discrepancy as experimentally demonstrating that one-body dissipation as presently formulated is not the complete dissipation mechanism in large-amplitude collective nuclear motion.

In contrast, as shown by the solid curves in fig. 4, when the two-body viscosity coefficient is adjusted to the value $\mu = 0.02$ TP, the experimental data for heavy nuclei lie between the calculated lower and upper limits and are adequately reproduced throughout the rest of the periodic table. For two-body viscosity, the dynamical trajectories lead to elongated scission shapes with less Coulomb repulsion, but this is supplemented by some pre-scission kinetic energy. These results calculated with several improvements demonstrate that mean fission-fragment kinetic energies are capable after all of distinguishing between dissipation mechanisms.

4. HEAVY ION REACTIONS

Even better prospects for determining the dissipation mechanism reside with heavy ion reactions, where we are able to choose the total mass of the combined system, the mass asymmetry of the entrance channel and the bombarding energy with foresight. This permits us to select for study those dynamically interesting cases that involve large distances in deformation space.

4.1. Energy for fusion

A necessary condition for compound-nucleus formation is that the dynamical trajectory of the fusing system pass inside the fission saddle point in a multi-dimensional deformation space. For heavy nuclear systems and/or large impact parameters, the fission saddle point lies inside the contact point and the center-of-mass bombarding energy must exceed the maximum in the one-dimensional zero-angular-momentum interaction barrier by an amount ΔE in order to form a compound nucleus.

This additional energy ΔE has been calculated both by solving the generalized Hamilton equations numerically with the three-quadratic-surface shape parametrization and realistic forces^{7,9} and approximately with the two-sphere-plus-conical-neck shape parametrization and schematic forces^{20,21}. Since we have not yet performed such calculations with our present shape parametrization and other improvements, we use the results from ref. 9 as an illustration.

Figure 5 compares calculated and experimental values of the additional center-of-mass bombarding energy ΔE as a function of a scaling variable $(Z^2/A)_{\text{mean}}$ defined in terms of the atomic numbers and mass numbers of the target and projectile⁹. Solid symbols denote values extracted from measurements of evaporation residues, which require the formation of true compound nuclei. Open symbols

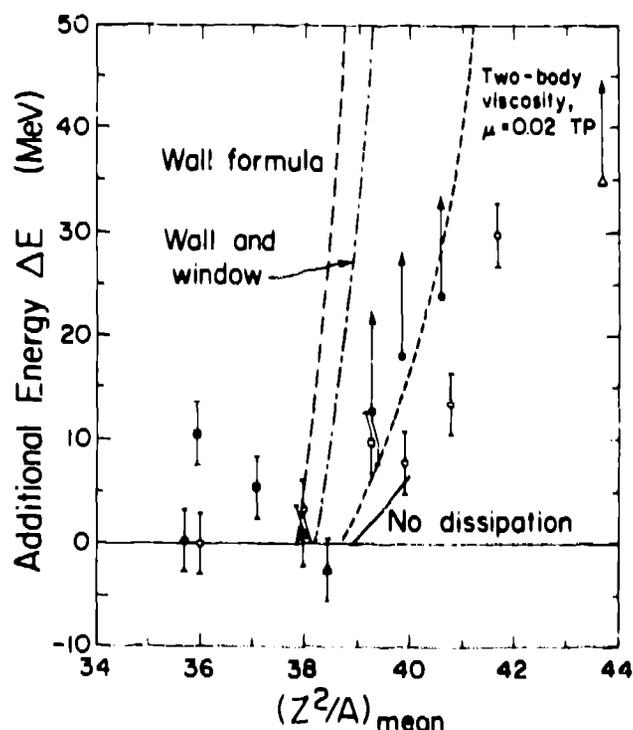


FIGURE 5
Comparison of calculated and experimental values of the additional energy ΔE required for compound-nucleus formation.

single values extracted from measurements of nearly symmetric fission-like fragments, where fast-fission processes involving significant mass transfer but not true compound nucleus formation also contribute. Because the open symbols do not represent cases where a true compound nucleus has necessarily been formed, and because the error bars for the three solid symbols with the largest values of $(I^2/A)_{\text{mean}}$ extend to infinity, we are not able to reach definitive conclusions from comparisons of this type. It is necessary instead to calculate such observable quantities as the mass transfer and capture cross section for asymmetric systems, to which we now turn our attention.

4. Mass transfer

We consider the reaction $^{208}\text{Pb} + ^{58}\text{Fe}$ at five bombarding energies corresponding to cases that have been studied experimentally by Bock et al.²², examining in particular the dependence upon angular momentum of the mass transfer A_{tr} to the lighter ^{58}Fe nucleus from the heavier ^{208}Pb nucleus. Our calculations with one body dissipation are not yet completed, but fig. 6 shows our results for two body viscosity. For a given bombarding energy the mass transfer is zero for angular momenta larger than a critical value at which the nuclei first come into contact, at which point it jumps discontinuously in our calculation to a nonzero value. The mass transfer then increases with decreasing

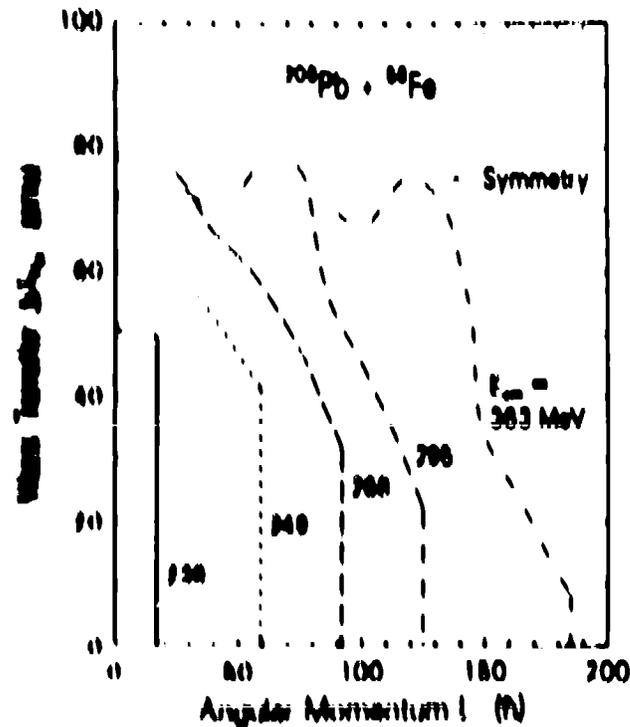


FIGURE 6

Dependence of mass transfer A_{tr} on angular momentum I for two body viscosity, with viscosity coefficient $\eta = 10^{-24}$ g cm² s⁻¹.

angular momentum unit values corresponding to a symmetric nuclear system are reached, below which some oscillations occur.

4.3. Capture cross section

In the experimental study of the $^{208}\text{Pb} + ^{58}\text{Fe}$ reaction, a cross section σ_c corresponding to capture, or symmetric fragmentation, was determined as the average of two separate procedures²². With the first procedure all mass transfers greater than 40 amu were included, whereas with the second procedure Gaussian distributions centered at symmetry were adjusted to the data. The resulting experimental points, with error bars reflecting the differences between the two procedures²², are shown in fig. 7. In calculating a corresponding theoretical capture cross section, we use the first procedure involving mass transfers greater than 40 amu. Other than at the lowest bombarding energy, our calculated curve for two-body viscosity lies substantially above the experimental points, with the deviation increasing to almost a factor of 2 at the highest bombarding energy. We regard this important discrepancy as experimentally demonstrating that two-body viscosity is also not the complete dissipation mechanism in large-amplitude collective nuclear motion.

Our next step is to perform analogous calculations for one-body dissipation, taking into account the dissipation associated with a time rate of change of the mass asymmetry degree of freedom in the completed wall-and-window formula¹⁵. Like the rest of you at this Conference, we are eager to find out if one-body dissipation can quantitatively account for experimental capture cross sections.

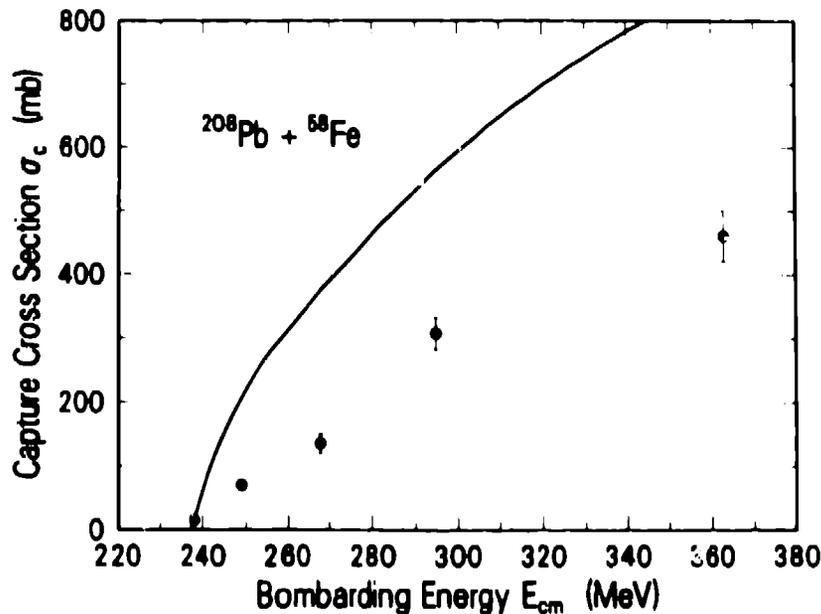


FIGURE 7
Comparison of experimental capture cross sections with results calculated for two-body viscosity, with viscosity coefficient $\mu = 0.02$ TP.

5. SADDLE-TO-SCISSION TIME

We now proceed in a second direction and apply the generalized Fokker-Planck equation to a one-dimensional inverted-oscillator fission barrier with frequency ω . The inertia with respect to the deformation coordinate q is assumed constant with value m and the dissipation coefficient is assumed constant with value η . It is natural to measure the dissipation strength in terms of the dimensionless ratio $\gamma = \eta/(2m\omega)$, defined so that unity corresponds to critical damping in the inverted oscillator turned upright.

Except for extremely small values of γ , Kramers' stationary solution of the Fokker-Planck equation for the inverted oscillator can be used to derive an analytical expression for the mean time \bar{t} required for the system to move from the saddle point at $q = 0$ to the scission point at $q = q_c$. The result is¹⁷

$$\bar{t} = \frac{2}{\omega} [(1 + \gamma^2)^{1/2} + \gamma] R\left[\left(\frac{1}{2}m\omega^2 q_c^2/\tau\right)^{1/2}\right],$$

where

$$R(z) = \int_0^z \exp(y^2) dy \int_y^\infty \exp(-x^2) dx$$

is a readily computed function studied and tabulated by Rosser. Because of its strong dependence upon the dissipation strength, the mean saddle-to-scission time provides a direct method for determining the magnitude of nuclear dissipation. As an example of this possibility, we consider the reaction $^{16}_0\text{O} + ^{142}_{84}\text{Nd} \rightarrow ^{158}_{64}\text{Er}$ at a laboratory bombarding energy $E_{\text{lab}} = 208$ MeV, which is being studied experimentally by the Los Alamos-Oak Ridge collaboration²³. For three values of angular momentum spanning the window that contributes to fission²³, we show our calculated mean saddle-to-scission times in fig. 8. The constants of the inverted oscillator representing the fission barrier for each angular momentum are determined by equating two quantities calculated for a parabolic barrier to the corresponding quantities calculated with our dynamical model described earlier. The nuclear temperature τ is determined from the excitation energy E^* at the saddle point by use of a Fermi-gas relationship.

Experimental data on the spectra of neutrons emitted in this reaction and on their angular distributions will be analyzed to yield the number of neutrons emitted prior to scission and the number of neutrons emitted from the fission fragments²³. These quantities may in turn be related by means of a statistical model to the sum of the time required to build up the quasi-stationary probability flow over the fission barrier²⁴ and the mean saddle-to-scission time. The former is currently being calculated for this reaction by Grangé and Weidenmüller²⁵. When fully completed and analyzed, this experiment should set stringent limits on the magnitude of nuclear dissipation.

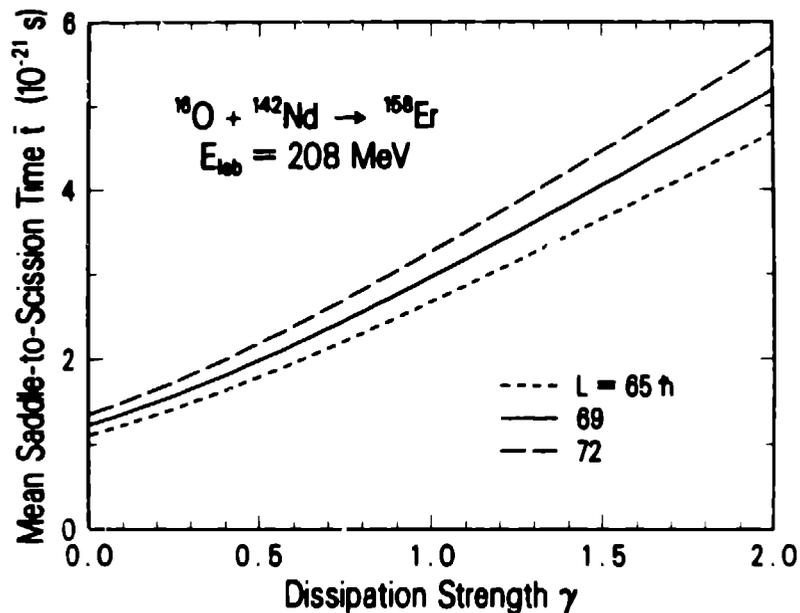


FIGURE 8

Increase of the mean saddle-to-scission time with dissipation strength.

6. OUTLOOK

We are entering a new era in fission and heavy ion reactions. Up to now theoretical approaches with vastly different pictures of the underlying nuclear dynamics have reproduced many of the gross experimental features of fission and heavy ion reactions because they include correctly the dominant nuclear, Coulomb and centrifugal forces. However, calculations are now being designed specifically to test the dissipation mechanism. When compared with mean fission-fragment kinetic energies, these calculations demonstrate that one-body dissipation is not the complete dissipation mechanism. Also, when compared with capture cross sections, they demonstrate that two-body dissipation is not the complete dissipation mechanism.

We are led experimentally to the suggestion that dissipation in large-amplitude collective nuclear motion is intermediate between these two extremes, arising from both mean-field effects and two-particle collisions. Further comparisons of the type made here, together with forthcoming experimental information on true ternary fission and the mean saddle-to-scission time, offer the exciting prospect of finally determining the magnitude and mechanism of nuclear dissipation.

ACKNOWLEDGEMENT

We are grateful to N. Cârjan, A. Gavron, H. Hofmann, F. Plasil, J. Randrup, W.J. Swiatecki and H.A. Weidenmüller for stimulating discussions.

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