LA-UR --1748

CONF-890491--7

Received by OSTI

JUN 07 1989

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LA-UR--89-1748

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TITLE SUMMARY: OUR 50-YEAR ODYSSEY WITH FISSION

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SUBMITTED TO Presented at the International Conference on Fifty Years Research in Nuclear Fission, Berlin, West Germany, April 3-7, 1989

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SUMMARY: OUR 50-YEAR ODYSSEY WITH FISSION

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On the occasion of this International Conference on Fifty Years Research in Nuclear Fission, we summarize our present understanding of the fission process and the challenges that lie ahead. The basic properties of fission arise from a delicate competition between disruptive Coulomb forces, cohesive nuclear forces, and fluctuating shell and pairing forces. These static forces are primarily responsible for such experimental phenomena as deformed ground-state nuclear shapes, fission into fragments of unequal size, sawtooth neutron yields, spontaneously fissioning isomers, broad resonances and narrow intermediate structure in fission cross sections, and cluster radioactivity. However, inertial and dissipative forces also play decisive roles in the dynamical evolution of a fissioning nucleus. The energy dissipated between the saddle and scission points is small for low initial excitation energy at the saddle point and increases with increasing excitation energy. At moderate excitation energies, the dissipation of collective energy into internal single-particle excitation energy proceeds largely through the interaction of nucleons with the mean field and with each other in the vicinity of the nuclear surface, as well as through the transfer of nucleons between the two portions of the evolving dumbbell-like system. These unique dissipation mechanisms arise from the Pauli exclusion principle for fermions and the details of the nucleon-nucleon interaction, which make the mean free path of a nucleon near the Fermi surface at low excitation energy longer than the nuclear radius. With its inverse process of heavy-ion fusion reactions, fission continues to yield surprises in the study of large-amplitude collective nuclear motion. Future challenges include devising experiments to unambiguously distinguish dissipative effects from analogous effects caused by collective degrees of freedom and computing fission directly from the underlying hadronic interaction.

1 INTRODUCTION

As this enjoyable conference comes to an end, it is time for us to summarize what we have learned from our 50-year odyssey with fission. Most of you have participated first-hand in this odyssey, which began here in Berlin with an unexpected experimental discovery by Hahn and Straßmann[1]. Like other epic journeys, fission research has been full of surprises from the outset. As vividly recounted by Herrmann[2] in his talk, the very idea that a heavy nucleus could split into two nuclei of intermediate mass was so astonishing that in 1934 Fermi[3] misinterpreted as transuranium elements the fission fragments that he had produced in Rome. Commenting on Fermi's experiments, Noddack[4] immediately suggested that "One could think that in the bombardment of heavy nuclei with neutrons these nuclei disintegrate into several large fragments." Yet her revolutionary idea went unheeded, even though she was also working in Berlin, and four years later Hahn and Straßmann[1] [illegible] accepted their own chemical evidence that barium had been produced.

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

Fission has continued to surprise us with division into fragments of unequal size, sawtooth neutron yields, spontaneously fissioning isomers, broad resonances and narrow intermediate structure in fission cross sections, cluster radioactivity, and unique dissipation mechanisms quite unlike those responsible for viscosity in fluids. Such phenomena provide an invaluable testing ground for many-body theories of large-amplitude collective nuclear motion. The atomic nucleus is a unique quantal many-body system consisting of nucleons and mesons at one level of approximation, or quarks and gluons at another. With its relatively small number of degrees of freedom, the nucleus displays aspects that are both microscopic and macroscopic on the one hand, and both quantal and classical on the other. The interplay between these complementary aspects leads to a static nuclear potential-energy landscape containing multiple minima, saddle points, and valleys, as well as to a rich dynamic behavior ranging from elastic vibrations of solids to long-mean-free-path dissipative fluid flow with statistical fluctuations.

Despite this underlying complexity, Meitner and Frisch[5] quickly recognized that fission is basically a large-amplitude collective rearrangement of nucleons, much like the division of a charged drop of liquid. On the basis of the liquid-drop model, Bohr and Wheeler[6] provided a satisfactory account of many of the early properties of fission, as discussed by Vandenbosch[7] in his introduction to the conference. This model, where fission represents a delicate competition between disruptive Coulomb forces and cohesive nuclear forces idealized as surface tension, was studied extensively[8–15] for the next quarter century. Yet many experimental phenomena lay outside its predictive capabilities, and it became increasingly clear that single-particle effects also play decisive roles.

A quantitative method for treating single-particle shell and pairing corrections was developed in 1966 by Strutinsky[16], whom we had the pleasure of hearing from at this conference[17]. The idea of a macroscopic-microscopic method, in which smooth trends are taken from a macroscopic model and local fluctuations from a microscopic model, had been introduced by Swiatecki[18] in 1963 with a simplified procedure for calculating the microscopic fluctuations that worked only for small deformations. These developments revolutionized the calculation of fission barriers. This renaissance in our understanding of static fission properties was followed in 1974 by another renaissance associated with the dynamics of fission. Through the work of Gross[19], Swiatecki[20], and others, it was realized that at low and moderate excitation energies the mechanism of nuclear dissipation is intimately connected with the long mean-free path of nucleons inside a nucleus.

Simultaneously, fully selfconsistent microscopic methods were also developed and applied to both statics and dynamics. For a given effective nucleon-nucleon interaction, the constrained static Hartree-Fock approximation[21–23] was used to calculate fission barriers, the time-dependent Hartree-Fock approximation[24–27] was used to calculate the dynamical evolution of a fissioning nucleus, as well as the inverse process of heavy-ion fusion reactions, and an imaginary-time mean-field approximation[28] was used to calculate spontaneous

Many new results were presented at the conference that contribute to our understanding of this vast field. In attempting to summarize these contributions, I have decided to concentrate on what fission has taught us concerning the three fundamental quantities that characterize large-amplitude nuclear shape changes: the nuclear potential energy of deformation, the nuclear inertia, and nuclear dissipation. Since these three quantities enter the equations of motion for the collective coordinates of a fissioning nucleus, the properties of fission can be computed at some level of approximation once they are known. Experimental results will be interspersed with the theoretical developments at appropriate points, even though they are affected to some extent by all three quantities. I will be focusing on the underlying physical principles involved rather than on the mathematical or experimental details. Also, I will not have time to mention the work of everyone who spoke, but instead will use those results that best illustrate the particular points I am trying to make. For continuity, I will occasionally mention recent work that was not explicitly presented at the conference.

2 NUCLEAR POTENTIAL ENERGY OF DEFORMATION

2.1. Selfconsistent Hartree-Fock calculations

As Berger[21] told us, the most fundamental way to calculate the nuclear potential energy of deformation is to start with a realistic interaction between the underlying nuclear constituents and solve the appropriate many-body equations in some approximation. To put what is actually done in proper perspective, let us remind ourselves that since nucleons themselves are composed of quarks and gluons, one should ideally start with the Lagrangian of quantum chromodynamics[29] and solve that somehow. To get to something tractable, one must make several approximations. First, one must approximate the confined quarks and gluons by nucleons and mesons. Second, one must replace the relativistic equation satisfied by spin-$\frac{1}{2}$ nucleons with a nonrelativistic equation. Third, one must replace the realistic interaction between nucleons with an effective interaction.

For simple effective interactions, such as those of Gogny and Skyrme, the constrained Hartree-Fock approximation has been used to compute the fission barriers of ^{240}Pu and other heavy nuclei. Results computed[22,23] with the finite-range interaction of Gogny are shown in Fig. 1. Axial asymmetry is taken into account at the first barrier, and reflection asymmetry is taken into account beyond the second minimum. The upper curve, which is computed with the original values of the constants in the Gogny D1 interaction, lies significantly higher at the first peak, second minimum and second peak than experimental values[4], which are indicated by the vertical arrows. However, when the surface energy constant a_s is reduced from its original value of 20.2 MeV to 19.0 MeV, the calculated and experimental values agree to within about 2 MeV accuracy, as indicated by the lower curve.

The constrained Hartree-Fock fission barrier for ^{240}Pu has also been computed[4] with the zero-range Skyrme M* interaction, for which[23] the surface energy constant a_s = 17.2

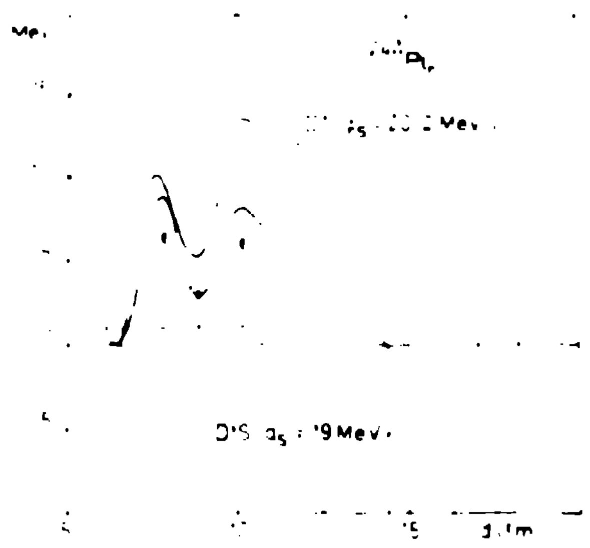

Fig. 1. Fission barriers[22-25] computed with the Gogny finite-range interaction.

MeV. These calculations impose axial symmetry and reflection symmetry, whose removal would lower the second peak to below the experimental value. The smaller value of a_s in this interaction compared to that in the Gogny interaction arises partly from the zero range of the Skyrme interaction.

2.2. Macroscopic-microscopic method

Because of computational difficulties and other reasons, selfconsistent Hartree-Fock calculations of the fission barrier have been carried out for only a few nuclei. Instead, nearly all fission barriers have been calculated by means of an alternative approach the macroscopic-microscopic method. This method synthesizes the best features of two complementary approaches. The smooth trends of the potential energy (with respect to particle numbers and deformations) are taken from a macroscopic model, and the local fluctuations are taken from a microscopic model. The method in its present form was developed in 1966 by Strutinsky[17] and has since revolutionized the calculation of fission barriers. The idea of a macroscopic-microscopic method had been introduced earlier by Swiatecki[18] and others.

In this method, which is suitable for treating nuclear systems that contain a large number of particles, the total nuclear potential energy of deformation is written as the sum of two terms,[17] namely

$$V = V_{\mathrm{macro}} + \Delta V_{\mathrm{micro}} \qquad (1)$$

The first term is a smoothly varying macroscopic energy that reproduces the broad trends of the potential energy. The second term contains oscillating microscopic corrections that arise because of the discreteness of the individual particles. The most important of these purely microscopic contributions are the shell and pairing corrections.

The calculation of the nuclear potential energy of deformation by means of this approach is usually performed in five steps: (1) Specify the nuclear shape. (2) Calculate the [illegible]

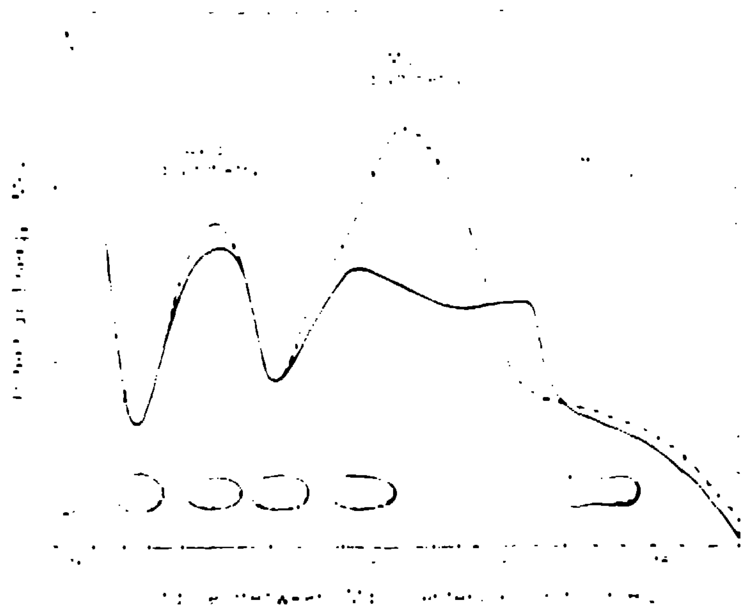

Fig. 2. Fission barrier computed with the macroscopic-microscopic method, illustrating the effects of axial asymmetry and mass asymmetry[32].

scopic energy. (3) Generate the single-particle potential felt by the nucleons. (4) Solve the Schrödinger equation for the single-particle energies. (5) Calculate the microscopic (shell and pairing) corrections. The total potential energy is then given by the sum of the macroscopic energy calculated in step 2 and the microscopic corrections calculated in step 5.

As illustrated in Fig. 2, the fission barrier of a heavy actinide nucleus calculated by this method contains a secondary minimum surrounded by two peaks[32]. The secondary minimum arises from shell effects associated with a nuclear shape whose axes are in the ratio 2:1. The dashed curve gives the potential energy for symmetric deformations as a function of the distance r between the centers of mass of the two nascent fragments, in units of the radius R of the spherical shape. The inclusion of axial asymmetry at the first peak lowers the energy by approximately 1 MeV, whereas the inclusion of mass asymmetry at the second peak lowers the energy by approximately 4 MeV. Such mass-asymmetric second saddle points are intimately connected with asymmetric fission-fragment mass distributions in heavy actinide nuclei at low excitation energies.

2.3 Experimental consequences

Single-particle effects of the type we have been discussing lead to several interesting experimentally observed phenomena. As discussed by Habs[33], the presence of a deep secondary minimum in the fission barrier leads to some states of the system in which the wave function is concentrated primarily at this larger deformation. Such spontaneously fissioning isomers were discovered experimentally in 1962 by Polikanov and coworkers[34] at Dubna, and have now been observed in 33 nuclei ranging from [illegible]U to [illegible]Bk. In [illegible]U, [illegible]U and [illegible]Np gamma decay back to the first well has been observed in addition to spontaneous fission. R[illegible]

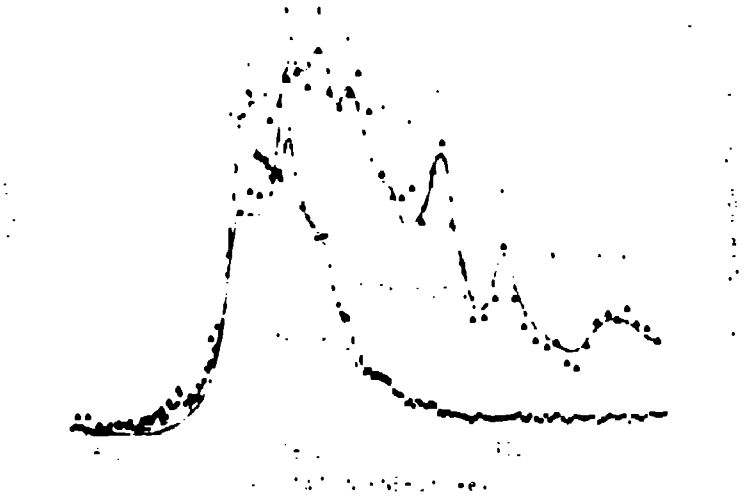

Fig. 3. Analysis of fission cross sections in terms of rotational bands of opposite parity[36]

gamma decay back to the first well has been tentatively identified also in ^{233}Th.

From detailed spectroscopic studies in the second minimum[33], the extracted moments of inertia and quadrupole moments correspond to nuclear shapes whose axes are in the ratio of approximately 2:1, in accordance with theoretical expectations. Seven high-spin superdeformed states corresponding to shapes with axes in the ratio of approximately 3:1 have been observed[33] in rare-earth nuclei ranging from ^{146}Gd to ^{152}Dy.

When a nucleus undergoes fission through a barrier containing two peaks separated by a secondary minimum, two distinct types of resonances in the fission cross section become possible. A classic example of the first type appears in the neutron-induced fission cross section of ^{23}Th, which was measured at Harwell by James, Lynn, and Earwaker[35]. The relatively broad resonance of width 40 keV at 720 keV incident neutron energy has traditionally been associated with a quasistationary vibrational level in the secondary minimum of the fission barrier.

However, Blons[36] described recent evidence indicating that this resonance is actually associated with a *third* minimum in the barrier corresponding to a mass-asymmetric deformation. This possibility had been pointed out in 1973 by Möller[32] as a resolution of the thorium anomaly that existed at that time. The recent evidence, shown in Fig. 3, consists of a simultaneous analysis of the fission cross sections for the ^{23}Th(n,f) and [illegible]Th(d,pf) reactions in terms of transition states corresponding to two rotational bands of opposite parity. Freiesleben[17] presented additional tentative evidence for a third asymmetric minimum in terms of a forward-backward anisotropy in fission-fragment angular distributions.

As discussed by Weigmann[38], the second type of resonance structure arising from a two-peaked fission barrier occurs when the excitation energy is close to a compound [illegible] level for the nucleus situated at the secondary minimum. This is illustrated in Fig. [illegible] [illegible] fission cross section of [illegible] Pu which was measured at Geel by [illegible]

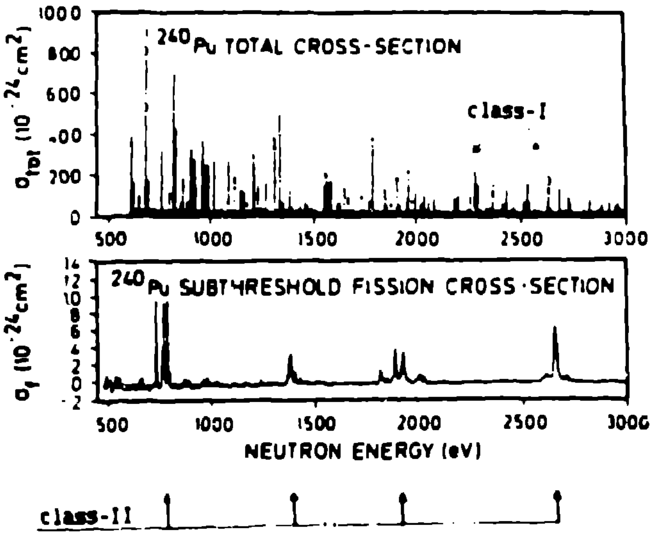

Fig. 4. Resonances corresponding to class-I states at the ground-state minimum and to class-II states at the secondary minimum[38,39].

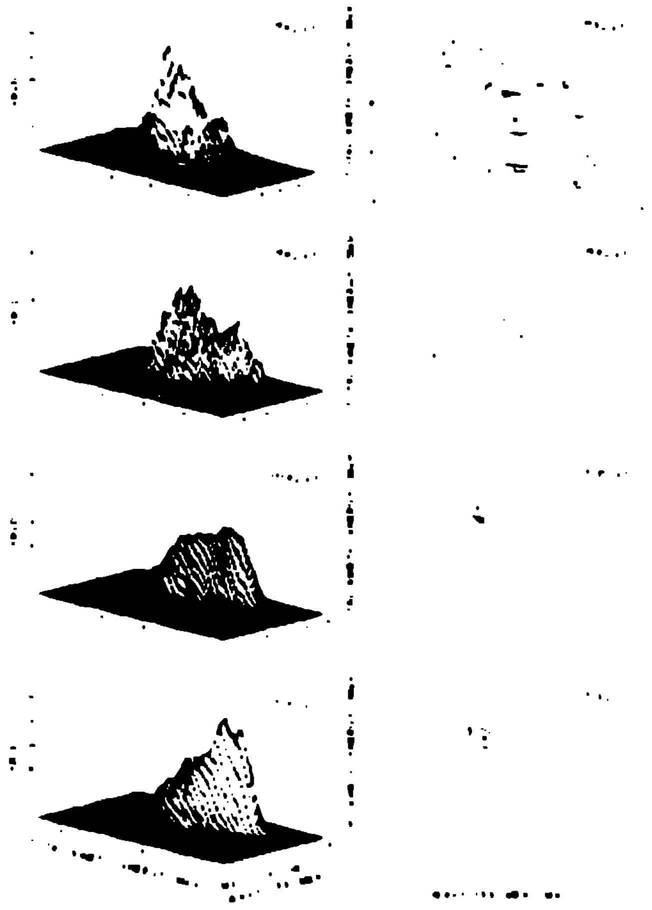

Fig. 5. Dramatic effect of neutron number on fission-fragment mass and kinetic-energy distributions[44]

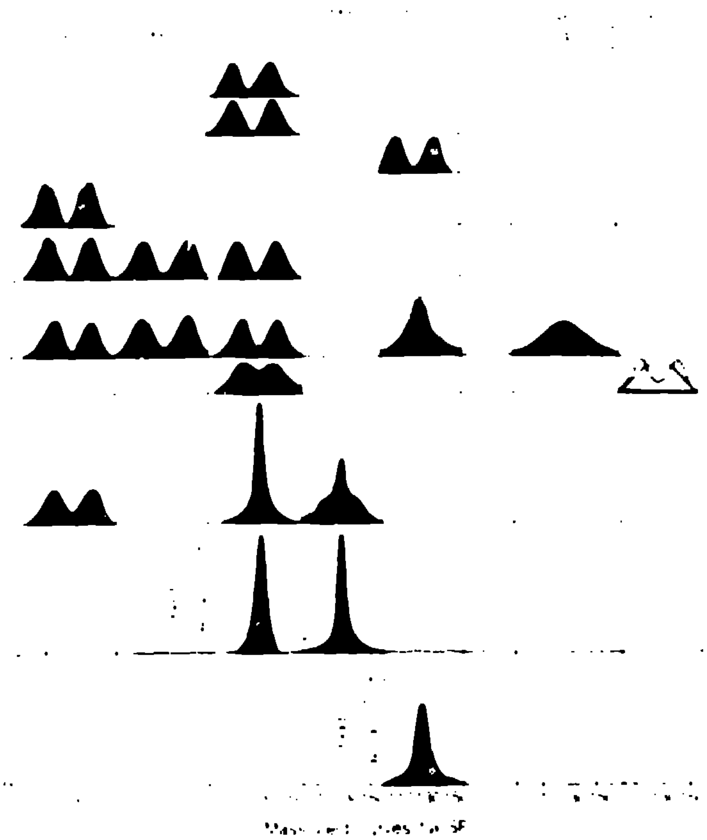

Fig. 6. Schematic illustration of experimental fission-fragment mass distribution for the spontaneous fission of heavy nuclei[45,46].

Theobald[34]. At the secondary minimum, some of the energy is in the form of deformation energy and is therefore unavailable for forming compound states. The distance between the groups of strong resonances is therefore much larger than the distance between the individual resonances, which occur when the excitation energy is close to a compound nuclear level for the nucleus situated at its ground state.

Experimental values are now available for the heights of the relevant saddle points and subsequent minima in the potential-energy surface for numerous nuclei. From previous comparisons[3,36,40-42] between experimental and calculated values, we conclude that the macroscopic-microscopic method is capable of reproducing the extrema in the potential-energy surface to within an accuracy of about 1 MeV. A detailed up-to-date comparison for nuclei throughout the periodic table would be extremely valuable.

We heard from several speakers about the important role that single-particle effects play in fission-fragment mass and kinetic-energy distributions. As discussed by Itkis[43], recent measurements show that at sufficiently low excitation energy the mass distributions for ^{201}Tl and other light nuclei are asymmetric. Wagemans[44] presented new results, shown in Fig. 5, indicating that in the spontaneous fission of plutonium isotopes, the addition of only two neutrons enhances the doubly magic fragment shell effects at 50 protons and 82 neutrons.

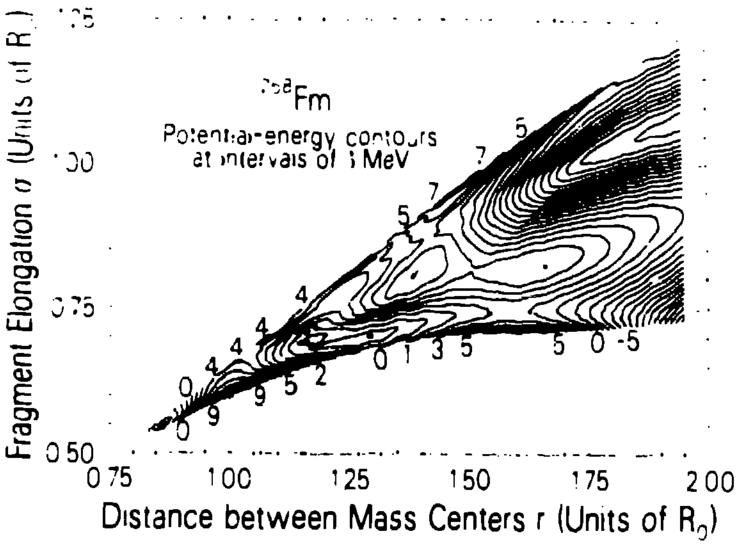

Fig. 7. Calculated potential-energy surface for ^{258}Fm, showing a new fission valley leading to compact scission shapes in the lower right-hand corner[49].

sufficiently to dramatically shift the peak in the mass distribution and increase the kinetic energy.

Hoffman[45] told us about the abrupt transition that occurs at ^{258}Fm in fission-fragment mass distributions, kinetic-energy distributions, and half-lives for the spontaneous fission of heavy nuclei[45-47]. As illustrated in Fig. 6, for ^{258}Fm and certain nuclei beyond, the mass distribution becomes very narrow, with a single peak at symmetry. The kinetic-energy distribution for some of these nuclei becomes skewed, with a peak at high energy and a broad low-energy tail. The spontaneous-fission half-life decreases by several orders of magnitude. All three of these phenomena are explained qualitatively by a new fission valley in the multidimensional potential-energy surface.

As discussed by Brosa[48] and others, this new valley is associated with doubly magic fragment shell effects at 50 protons and 82 neutrons. Since these shell effects are maximum for spherical shapes, it is essential for an accurate calculation that the shape parameterization be capable of describing touching spherical fragments and that the finite range of the nuclear force be taken into account in calculating the macroscopic energy. These items are included in the potential-energy surface[49] for ^{258}Fm shown in Fig. 7.

Because the saddle point leading into the new fission valley in the lower right-hand portion of Fig. 7 is lower than that leading into the old fission valley in the upper right-hand portion of the figure, spontaneous fission will proceed primarily into the new valley. However, there is also a switchback path, which is lowered by mass-asymmetric deformations, that branches off from the new path and leads into the old valley across another saddle point. This switchback path is probably responsible for the low-energy tail of the kinetic-energy distribution for this nucleus[49].

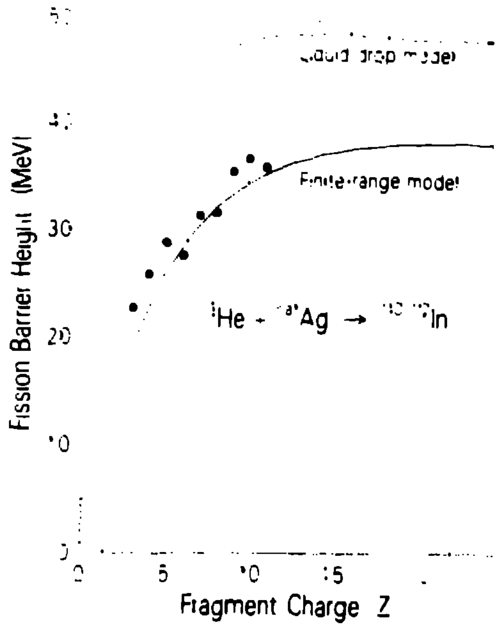

Fig. 8. Fission-barrier height as a function of the charge of the lighter fragment[50,51,53].

For compact dumbbell-like shapes, which are present both in the new fission valley and at the fission saddle points for medium-weight nuclei, the finite range of the nuclear force substantially lowers the macroscopic energy relative to that calculated in the liquid-drop model. This is illustrated in Fig. 8, which was shown by Schröder[50] in his talk. The solid curve, calculated by Sierk[51] with previously determined constants[42] in the Yukawa-plus-exponential model[52], satisfactorily reproduces the experimental[53] fission-barrier heights for asymmetric mass divisions. In contrast, the liquid-drop model yields barriers that are approximately 10 MeV too high.

Moretto[54] stressed in his talk that mass-asymmetric fission of the type shown in Fig. 8 evolves continuously with decreasing mass asymmetry into the statistical evaporation of complex fragments. The spontaneous emission of complex fragments was discussed by Price[55] and Poenaru[56]. As shown in Fig. 9, clusters ranging from ^{14}C to Si have been observed experimentally[55], with half-lives ranging from 10^{11} s to 10^{28} s. Since cluster emission is intermediate between fission and α decay, theoretical approaches based on both fission theory and the preformation of clusters have been used. The emitted clusters are neutron-rich because this leads to tightly bound daughter nuclei close to ^{208}Pb, with large energy release.

3. NUCLEAR INERTIA

3.1. Role in spontaneous fission

We turn now to the nuclear inertia, which at low excitation energy is experimentally tested primarily by spontaneous fission. Recent progress in selfconsistent microscopic treat-

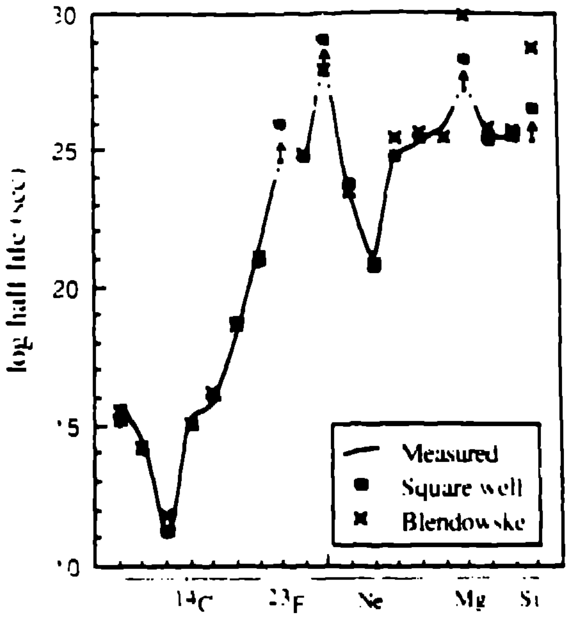

Fig. 9. Half-life for cluster emission[55].

ments of spontaneous fission was discussed by Negele[28], who used a Feynman path integral to formulate an imaginary-time mean-field theory of tunneling in many-body systems. This approach has the nice feature that the contributing paths in deformation space are determined automatically by the system's Hamiltonian. For a ^{32}S nucleus without spin-orbit interaction and with the charge artificially increased to cause fission to occur, the resulting dominant path is significantly different from that corresponding to constrained mean-field theory. However, because of computational difficulties, this approach has not yet been applied to the spontaneous fission of a realistic heavy nucleus.

As discussed by Sobiczewski[57], actual calculations of spontaneous-fission half-lives are usually performed by use of the semiclassical WKB approximation applied to a particular one-dimensional path through the multi-dimensional deformation space. This path is in practice determined either from considerations of statics alone or dynamically by maximizing the penetrability, which is related to an integral along the path involving the product of the inertia with respect to this path and the potential energy relative to the ground-state energy.

3.2. Microscopic and semi-empirical inertias

The nuclear inertia tensor for spontaneous fission is frequently calculated microscopically[57,58] by use of the Inglis cranking model. As illustrated by the solid and short-dashed curves[24] in Fig. 10, the microscopic inertia with respect to the distance between mass centers r is an oscillating function of deformation and for small deformations is several times the irrotational value. These properties arise from the rearrangement of nodal structure in single-particle wave functions which occurs at the near crossings of single-particle levels. The microscopic inertia oscillates about a semi-empirical inertia, which is related to the irrotational inertia in such a way that both approach the reduced mass for separated fragments.

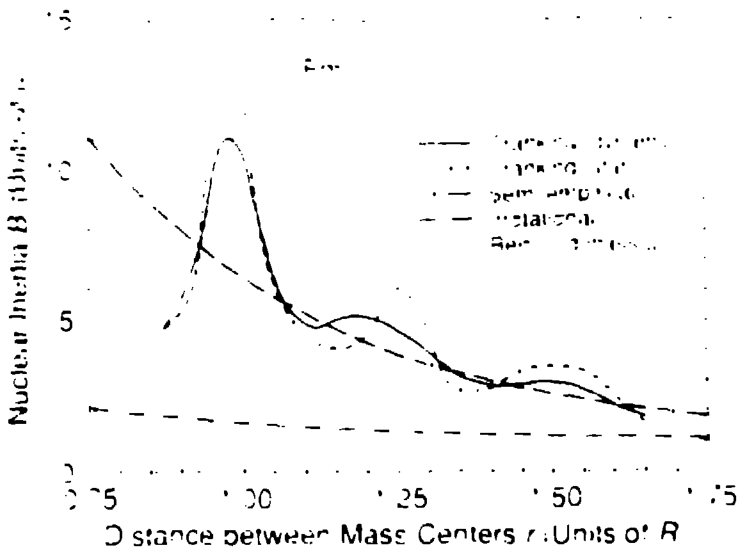

Fig. 10. Increase in the nuclear inertia for spontaneous fission (cranking and semi-empirical values) relative to the irrotational value[58].

with an adjustable constant extracted from spontaneous-fission half-lives.

The original Inglis cranking model has been generalized by Kunz[59,60] to include the velocity dependence of the single-particle potential and the reaction of the pairing field to the collective motion. The inertia calculated in this generalized cranking model approaches the correct value in the limit of large pairing, in contrast to that calculated in the original Inglis cranking model.

3.3. Spontaneous-fission half-lives

The important role that spontaneous fission plays in limiting the production of very heavy nuclei was stressed by Lazarev[61] in his discussion of recent attempts at the JINR in Dubna to produce element 110 through such reactions as ^{232}Th + ^{48}Ca → ^{280}110. He also described how pairing vibrations can increase the tunneling, an idea introduced in 1971 by Moretto and Babinet[62] and recently used by Pomorski and co-workers[63] in their half-life calculations.

Münzenberg[64] described the production of three atoms of ^{266}109 at the GSI in Darmstadt by use of ^{58}Fe projectiles. The production of other heavy elements with neutron-rich ^{48}Ca projectiles was discussed by Gäggeler[65].

When comparing experimental and calculated spontaneous-fission half-lives, Hoffman[66] noted the large discrepancy that previously existed for very heavy elements with neutron number $N > 158$. As illustrated in Fig. 11, this discrepancy is to a large extent resolved by considering tunneling into the new fission valley that exists in such cases[47,48]. The reduction in calculated half-life arises both because the potential energy is lower and because the inertia is less along the path leading into the new valley than the corresponding quantity along the path leading into the old valley. The agreement is even further improved when the influence of pairing vibrations on the tunneling is taken into account[63].

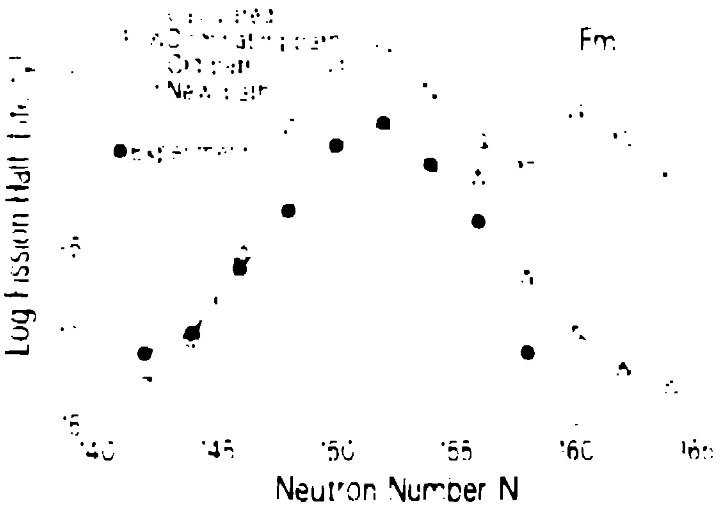

Fig. 11. Improved agreement between experimental and calculated spontaneous-fission half-lives when the new fission valley is considered[45,49].

3.4. Inertia for induced fission

At high excitation energies and large deformations, where pairing correlations have disappeared and near crossings of single-particle levels have become less frequent, the rotational moment of inertia is close to the rigid-body value and the vibrational inertia is close to the incompressible, irrotational value[60]. The inertia tensor can then be calculated for a superposition of rigid-body rotation and incompressible, nearly irrotational flow. For this purpose the Werner-Wheeler method, which determines the flow in terms of circular layers of fluid[61], can be used.

4. NUCLEAR DISSIPATION

We turn now to nuclear dissipation and its role in the dynamical evolution of a fissioning nucleus. Dissipation mechanisms were discussed at the conference by Weidenmüller[7], and Pashkevich[68] presented a survey of macroscopic approaches to fission dynamics.

Up until about 1974, it was generally believed that the dissipation mechanism in fission is two-body collisions, like that responsible for viscosity in fluids, and that the viscosity coefficient is relatively small. But then, through the work of Gross[11], Swiatecki[2], and others, it was realized that the long mean free path of nucleons inside a nucleus at low and moderate excitation energies, arising from the Pauli exclusion principle for fermions and the details of the nucleon-nucleon interaction, alters both the mechanism and magnitude

4.1. Long-mean-free-path dynamics

The effects of a long nucleon mean free path on fission dynamics have been incorporated by means of several different macroscopic and microscopic approximations, which have led to radically different pictures. By assuming that the velocity distribution of nucleons striking a moving container wall is completely random, Swiatecki and his colleagues derived a [illegible]

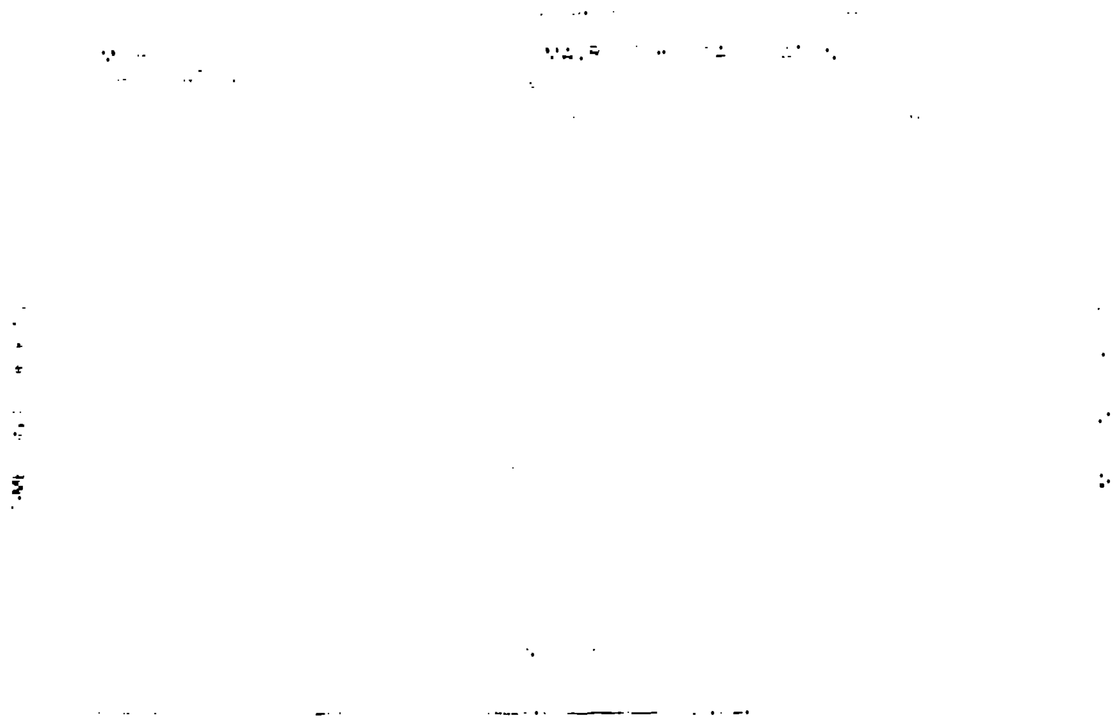

Fig. 12. Fission of ^{236}U, starting from rest 1 MeV beyond the saddle point[26].

wall formula for describing such one-body dissipation, in terms of which nuclei are predicted to be highly overdamped[20,69,70]. In contrast, by constraining the many-body wave function to at all times be a Slater determinant of single-particle wave functions, Negele *et al.* treated the dynamics by use of the time-dependent Hartree-Fock (TDHF) approximation, in terms of which nuclei dissipate much less energy[26,27].

These two contrasting extremes are illustrated[26] in the first and last columns of Fig. 12 for the fission of ^{236}U. In the TDHF calculation, which is performed with zero spin-orbit interaction and an effective pairing gap $\Delta = 2$ MeV, the time required to reach the point neck rupture is 3.4×10^{-21} s. This is only slightly longer than the 2.5×10^{-21} s required for the nonviscous descent shown in the second column, but is substantially less than the 12.9×10^{-21} s required for the descent with Swiatecki's wall formula shown in the fifth column.

The relatively small dissipation present in the TDHF solution is associated with the complete neglect of two-particle collisions in this approximation. Several steps have been taken in the important direction of incorporating two-particle collisions into the TDHF approximation[27]. Although promising, these attempts have not yet led to any definitive conclusions because of both conceptual and computational difficulties.

4.2 Surface-plus-window dissipation

Pashkevich[68] discussed another macroscopic approach, where long-mean free path dynamics and two-particle collisions are merged in a natural way[71]. In this approach, which is valid for intermediate excitation energies above which pairing has disappeared and below which the nucleon mean free path exceeds the nuclear diameter, the dissipation power primarily in the surface region from two distinct mechanisms.

The first mechanism is one-body dissipation, but with a magnitude that is substantially reduced relative to that of the wall formula. In calculations based on the random-phase approximation for spherical nuclei, Griffin, Dworzecka, and Yannouleas have shown that the effect of replacing three idealizations of the wall formula by more realistic features appropriate to real nuclei is to reduce the one-body dissipation coefficient to roughly 10% of the wall-formula value[72,73]. Alternatively, the reduction could arise because the nucleons retain some memory of their previous collisions with the wall, which invalidates the assumption of a random velocity distribution that was used to derive the wall formula.

The second mechanism is two-body collisions in the surface region. The Pauli exclusion principle, which suppresses two-body collisions in the nuclear interior, disappears as one passes through the nuclear surface to the exterior[74]. In addition, the [illegible] nucleon-nucleon cross section itself increases as one passes through the surface to the exterior because of its increase with decreasing kinetic energy. Since the density decreases to zero outside the nucleus, the probability for two-body collisions peaks in the nuclear surface.

Under the assumption that the surface dissipation is local, it can be calculated from the leading term in an expansion of the time rate of change of the collective Hamiltonian H in powers of the surface diffuseness divided by the nuclear radius[70]. This leading term can be written as

$$\left(\frac{dH}{dt}\right)_{\mathrm{surface}} = k_s \rho v \int (\dot{n} - D)^2 dS \ . \qquad (2)$$

where $\dot{n}$ is the velocity of a surface element dS, D is the normal drift velocity of nucleons about to strike the surface element dS, v is the average speed of the nucleons inside the nucleus, ρ is the nuclear mass density, and k_s is a dimensionless parameter that specifies the total strength of the interaction of either one or two nucleons with the moving nuclear surface. A value of $k_s = 1$ would correspond to the wall formula, but several types of experimental data indicate that for real nuclei its value is much less than unity. The value of k_s could depend upon both the excitation energy and type of collective motion involved.

For dumbbell-like shapes, the transfer of nucleons through the window separating the two portions of the system leads to an additional dissipation that is analogous to the classical window formula of Swiatecki[21,66,75]. The result is

$$\left(\frac{dH}{dt}\right)_{\mathrm{wall}} = \frac{1}{2}\rho v a \dot{r}^2 F(q,\dot{q}) \ . \qquad (3)$$

where a is the area of the window, $\dot{r}$ is the relative velocity of the centers of mass of the two portions of the system, q denotes the collective coordinates, $\dot{q}$ denotes the collective velocities, and $F(q,\dot{q})$ describes the effect of a nonuniform velocity as a function of position in the deforming fragments. There is no need to renormalize this part of the dissipation because nucleons that have passed through a small window have a low probability of returning through it while still retaining memory of their previous passage. The combination of these two mechanisms leads to surface-plus-window dissipation.

Pashkevich[68] and Schröder[50] both discussed dynamical calculations of fission that have been performed with surface-plus-window dissipation[71]. For values of the strength corresponding to dynamical motion that is somewhat overdamped, these calculations satisfactorily reproduce the average kinetic energies for the fission of nuclei throughout the periodic table at moderate excitation energies[71].

4.3. Generalized Fokker-Plank equation

Thus far we have been considering the effect of dissipation on mean quantities in fission. As discussed by Weidenmüller[67] and Pashkevich[68], the coupling between the collective and internal degrees of freedom also gives rise to residual fluctuating forces, which diffuse the dynamical paths in phase space. When these stochastic forces are treated under the Markovian assumption that they do not depend upon the system's previous history, one arrives at the generalized Fokker-Planck equation

$$\frac{\partial f}{\partial t} + (M^{-1})_{ij} p_j \frac{\partial f}{\partial q_i} - \left[\frac{\partial V}{\partial q_i} + \frac{1}{2}\frac{\partial (M^{-1})_{jk}}{\partial q_i} p_j p_k\right] \frac{\partial f}{\partial p_i} = \eta_{ij} (M^{-1})_{jk} \frac{\partial}{\partial p_i}(p_k f) + \tau \eta_{ij} \frac{\partial^2 f}{\partial p_i \partial p_j} \quad (4)$$

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 η and the nuclear temperature τ, which is measured here in energy units.

By solving a stationary Fokker-Planck equation in one dimension for the probability flow over the barrier, Kramers[75] showed in 1940 that dissipation increases the asymptotic value of the fission lifetime relative to the Bohr-Wheeler transition-state value[6]. As discussed by Weidenmüller[67] and Pashkevich[68], this important result has only recently been incorporated into studies of fission.

4.4. Neutron emission prior to fission

Solution of the Fokker-Planck equation (4) for two other situations has made it possible to extract information on fission time scales and nuclear dissipation from neutron emission prior to fission, as discussed by Hinde[76], Gavron[77], Dietrich[78], and Schröder[50]. First, an analytical solution for the mean saddle-to-scission time has been obtained from a one-dimensional stationary Fokker-Planck equation[79,80]. Second, a numerical solution for the transient time required to build up the quasi-stationary probability flow over the fission barrier has been obtained from a one-dimensional time-dependent Fokker-Planck equation[80,81].

During these saddle-to-scission and transient times, additional neutron emission can take place relative to that calculated from evaporation in a standard statistical model. Such enhanced neutron emission has been observed experimentally in heavy-ion-induced fission reactions[76-77]. Although some important differences in the experimental results still remain, the analysis of this enhancement in terms of neutron emission during the saddle-to-scission and transient times nevertheless suggests that fission is somewhat overdamped[81].

4.5. Additional experiments related to dissipation

Bocquet[82] described measurements of the odd-even effect in fission-fragment charge distributions for the neutron-induced fission of nuclei ranging from thorium to californium. For the fission of uranium isotopes, this odd-even effect decreases from about 23% for zero excitation energy at the saddle point to about 4% for 4 MeV excitation energy at the saddle point. This dramatic decrease in odd-even effect with such a small increase in excitation energy led Bocquet to conclude that the energy dissipation between the saddle and scission points is very small at low excitation energies. In particular, he found that the energy dissipated between saddle and scission ranges from about 3 MeV for thorium to about 10 MeV for californium.

Presenting a paper of Signarbieux[83], Gönnenwein showed highly resolved scatter plots of fission fragments versus their mass and kinetic energy. In the extreme high-energy tails of these distributions, essentially all of the energy released in fission goes into the kinetic energy of the fragments, with zero neutron emission. Although these high-energy events are only a tiny fraction of the total events, they nevertheless correspond to situations in which there is no dissipation of energy between the saddle and scission points.

Several other experiments were discussed at the conference whose proper analysis could in principle yield information on nuclear dissipation. These include fission induced by muons, antiprotons, and lambdas[34], light-particle-accompanied fission[85,86], scission neutrons[87], and the emission of charged particles and gammas prior to scission.

Although several issues still remain to be clarified, our present picture is that the energy dissipated between the saddle and scission points is small for low initial excitation energy at the saddle point and increases with increasing excitation energy. At moderate excitation energies, the dissipation of large-amplitude nuclear collective energy into internal single-particle excitation energy arises primarily from the interaction of nucleons with the mean field and with each other in the vicinity of the nuclear surface, as well as from the transfer of nucleons through the window separating the two portions of a dumbbell-like system. The magnitude of dissipation at moderate excitation energies corresponds to dynamical motion that is somewhat overdamped.

5. FUTURE CHALLENGES

Our 50-year odyssey with fission has led to a vast wealth of experimental data, a semiquantitative understanding of the process based on the macroscopic-microscopic method and important progress with selfconsistent microscopic methods. Now that we have turned to the birthplace of our odyssey and completed our assessment, we see that it is not yet over. Fission continues to surprise us, and important challenges lie ahead. These include devising experiments to unambiguously distinguish dissipative effects from analogous effects caused by collective degrees of freedom, further refining the predictions of the macroscopic-microscopic method, and computing fission directly from the underlying hadronic interaction.

In the last area, much remains to be done even within the restriction of an effective two-nucleon interaction treated in the nonrelativistic approximation. But it is also important to extend this work by using a more realistic interaction, as well as in another direction by using a relativistic approximation suitable for spin-$\frac{1}{2}$ nucleons. The ultimate challenge is of course to start at the level of quantum chromodynamics, with explicit quark and gluon degrees of freedom taken into account.

Large-amplitude collective nuclear motion, as exemplified by fission, should continue to provide an invaluable testing ground for nuclear many-body theories. The next 50 years could be even more exciting than the first!

ACKNOWLEDGEMENT

I am grateful to D. C. Hoffman, J. Machen, P. Möller, and R. Vandenbosch for stimulating discussions while writing this summary.

REFERENCES

1) O. Hahn and F. Straßmann, Naturwiss. 27 (1939) 11.

2) G. Herrmann, Discovery and confirmation of fission, this volume.

3) E. Fermi, Nature 133 (1934) 898.

4) I. Noddack, Angew. Chem. 47 (1934) 653.

5) L. Meitner and O. R. Frisch, Nature 143 (1939) 239.

6) N. Bohr and J. A. Wheeler, Phys. Rev. 56 (1939) 426.

7) R. Vandenbosch, Nuclear fission: what have we learned in 50 years?, this volume.

8) S. Frankel and N. Metropolis, Phys. Rev. 72 (1947) 914.

9) D. L. Hill and J. A. Wheeler, Phys. Rev. 89 (1953) 1102.

10) U. L Businaro and S. Gallone, Nuovo Cim. 5 (1957) 315.

11) S. Cohen and W. J. Swiatecki, Ann. of Phys. 22 (1963) 406.

12) V. M. Strutinsky, N. Ya. Lyashchenko, and N. A. Popov, Nucl. Phys. 46 (1963) 639.

13) I. Kelson, Phys. Rev. 136 (1964) B1667.

14) J. R. Nix and W. J. Swiatecki, Nucl. Phys. 71 (1965) 1.

15) J. R. Nix, Nucl. Phys. A130 (1969) 241

16) V. M. Strutinsky, Yad. Fiz. 3 (1966) 614; Sov. J. Nucl. Phys. 3 (1966) 449.

17) V. M. Strutinsky, Potential energy surfaces, this volume.

18) W. J. Swiatecki, in Proc. Int. Conf. on nuclidic masses, Vienna, 1963 (Springer-Verlag, Vienna, 1964) p. 58.

19) D. H. E. Gross, Nucl. Phys. A240 (1975) 472.

20) W. J. Swiatecki, in Proc. Int. School-Seminar on reactions of heavy ions with nuclei and synthesis of new elements, Dubna, 1975, Joint Institute for Nuclear Research report JINR-D7-9734 (1976) p. 89.

21) J. Bartel, P. Quentin, M. Brack, C. Guet, and H. B. Håkansson, Nucl. Phys. A386 (1982) 79.

22) J. F. Berger, M. Girod, and D. Gogny, Nucl. Phys. A428 (1984) 23c.

23) M. Brack, C. Guet, and H. B. Håkansson, Phys. Rep. 123 (1985) 275.

24) L. Bennour, J. Libert, and P. Quentin, in Contributed Papers, Third Int. Conf. on nucleus nucleus collisions, Saint-Malo, 1988, Université de Caen report, ISBN No. 2-905461-30-6 (1988) p. 137.

25) J. F. Berger, Constrained Hartree-Fock, this volume.

26) J. W. Negele, S. E. Koonin, P. Möller, J. R. Nix, and A. J. Sierk, Phys. Rev. C17 (1978) 1098.

27) J. W. Negele, Rev. Mod. Phys. 54 (1982) 913.

28) J. W. Negele, Microscopic theory of fission dynamics, this volume.

29) F. Wilczek, Ann. Rev. Nucl. Part. Sci. 32 (1982) 177.

30) S. Bjørnholm and J. E. Lynn, Rev. Mod. Phys. 52 (1980) 725.

31) J. R. Nix, Ann. Rev. Nucl. Sci. 22 (1972) 65.

32) P. Möller and J. R. Nix, in Proc. Third IAEA Symp. on the physics and chemistry of fission, Rochester, 1973, Vol. I (IAEA, Vienna, 1974) p. 103.

33) D. R. Habs, Spectroscopy in the secondary minimum, this volume.

34) S. M. Polikanov, V. A. Druin, V. A. Karnaukhov, V. L. Mikheev, A. A. Pleve, N. K. Skobelev, V. G. Subbotin, G. M. Ter-Akopjan, and V. A. Fomichev, Zh. Eksp. Teor. Fiz. 42 (1962) 1464; Sov. Phys. JETP 15 (1962) 1016.

35) G. D. James, J. E. Lynn, and L. G. Earwaker, Nucl. Phys. A189 (1972) 225.

36) J. Blons, A third minimum in the fission barrier, this volume.

37) F. M. Baumann, K. T. Brinkmann, H. Freiesleben, J. Kiesewetter, and H. Sohlbach, Angular distributions of ^{230,232}Th and the third minimum hypothesis, this volume.

38) H. Weigmann, Light particle and photo-induced fission, this volume.

39) E. Migneco and J. P. Theobald, Nucl. Phys. A112 (1968) 603.

40) H. C. Britt, in Proc. Fourth IAEA Symp. on physics and chemistry of fission, Julich, 1979, Vol. I (IAEA, Vienna, 1980) p. 3.

41) M. Brack, in Proc. Fourth IAEA Symp. on physics and chemistry of fission, Julich, 1979, Vol. I (IAEA, Vienna, 1980) p. 227.

42) P. Möller and J. R. Nix, Nucl. Phys. A361 (1981) 117.

43) M. G. Itkis, Symmetric and asymmetric fission of nuclei lighter than radium, this volume.

44) C. Wagemans, P. Schillebeeckx, and A. Deruytter, Investigation of neutron shell effects and fission channels in the spontaneous fission of the Pu-isotopes, this volume.

45) D. C. Hoffman, Spontaneous fission, lifetime systematics, this volume.

46) D. C. Hoffman and L. P. Somerville, in Particle Emission from Nuclei, Vol. III, Fission and Beta-Delayed Decay Modes (CRC Press, Boca Raton, 1988) p. 1.

47) E. K. Hulet, J. F. Wild, R. J. Dougan, R. W. Lougheed, J. H. Landrum, A. D. Dougan, M. Schädel, R. L. Hahn, P. A. Baisden, C. M. Henderson, R. J. Dupzyk, K. Sümmerer, and G. R. Bethune, Phys. Rev. Lett. 56 (1986) 313.

48) U. Brosa, Random neck rupture, this volume.

49) P. Möller, J. R. Nix, and W. J. Swiatecki, Nucl. Phys. A492 (1989) 349.

50) W. U. Schröder, Experimental status of heavy-ion-induced fission, this volume.

51) A. J. Sierk, Phys. Rev. Lett. 55 (1985) 582.

52) H. J. Krappe, J. R. Nix, and A. J. Sierk, Phys. Rev. C20 (1979) 992.

53) M. A. McMahan, L. G. Moretto, M. L. Padgett, G. J. Wozniak, L. G. Sobotka, and M. G. Mustafa, Phys. Rev. Lett. 54 (1985) 1995.

54) L. G. Moretto, Angular momentum bearing modes in nuclear fission, this volume.

55) P. B. Price, Complex radioactivities, this volume.

56) D. N. Poenaru, W. Greiner, and M. Ivascu, Predicted halflives for cluster radioactivities, this volume.

57) Z. Patyk, J. Skalski, A. Sobiczewski, and S. Ćwiok, Potential energy and spontaneous-fission half-lives for heavy and superheavy nuclei, this volume.

58) A. Baran, K. Pomorski, S. E. Larsson, P. Möller, S. G. Nilsson, J. Randrup, A. Lukasiak, and A. Sobiczewski, in Proc. Fourth IAEA Symp. on physics and chemistry of fission, Jülich, 1979, Vol. I (IAEA, Vienna, 1980) p. 143.

59) J. Kunz and J. R. Nix, Nucl. Phys. A426 (1984) 353.

60) J. Kunz and J. R. Nix, Z. Phys. A321 (1985) 455.

61) Yu. A. Lazarev, Very heavy element production, this volume.

62) L. G. Moretto and R. P. Babinet, Phys. Lett. 49B (1974) 147.

63) A. Staszczak, S. Pilat, and K. Pomorski, Influence of pairing vibrations on spontaneous fission probability, Nucl. Phys. A, submitted.

64) G. Münzenberg, Heavy element production and limits to fusion, this volume.

65) H. Gäggeler, D. Jost, A. Türler, P. Armbruster, W. Brüchle, H. Folger, F. P. Heßberger, S. Hofmann, G. Münzenberg, V. Ninov, M. Schädel, K. Sümmerer, J. V. Kratz, U. Scherer, and M. Leino, Cold fusion reactions with ^{48}Ca, this volume.

66) J. R. Nix, Nucl. Phys. A130 (1969) 241

67) H. A. Weidenmüller, Dissipation mechanisms, this volume.

68) V. Pashkevich, Theory of macroscopic fission dynamics, this volume.

69) J. Blocki, Y. Boneh, J. R. Nix, J. Randrup, M. Robel, A. J. Sierk, and W. J. Swiatecki, Ann. of Phys. 113 (1978) 330.

70) W. J. Swiatecki, Prog. Part. Nucl. Phys. 4 (1980) 383.

71) J. R. Nix and A. J. Sierk, in Proc. Int. School-Seminar on heavy ion physics, Dubna, 1986, Joint Institute for Nuclear Research report JINR-D7-87-68 (1987) p. 453.

72) J. J. Griffin and M. Dworzecka, Nucl. Phys. A455 (1986) 61.

73) C. Yannouleas, Nucl. Phys. A439 (1985) 336.

74) R. W. Hasse and P. Schuck, Nucl. Phys. A438 (1985) 157.

75) H. A. Kramers, Physica 7 (1940) 284.

76) D. J. Hinde, Fission time scales of excited nuclei, this volume.

77) A. Gavron, Heavy-ion fission—an inherently non-equilibrium process?, this volume.

78) K. Dietrich, E. Strumberger, and K. Pomorski, What can we learn on the fission process from the spectrum of prefission neutrons?, this volume.

79) J. R. Nix, A. J. Sierk, H. Hofmann, F. Scheuter, and D. Vautherin, Nucl. Phys. A424 (1984) 239.

80) P. Grangé, S. Hassani, H. A. Weidenmüller, A. Gavron, J. R. Nix, and A. J. Sierk, Phys. Rev. C34 (1986) 209.

81) P. Grangé, J. Q. Li, and H. A. Weidenmüller, Phys. Rev. C27 (1983) 2063.

82) J. P. Bocquet and R. Brissot, Mass, energy, and nuclear charge distribution of fission fragments, this volume.

83) C. Signarbieux, Primary fragment mass-charge-energy correlations from cold fission, this volume.

84) S. M. Polikanov, Meson, antiproton, and hyperon induced fission, this volume.

85) J. P. Theobald, Ternary fission, this volume.

86) V. E. Makarenko, Yu. D. Molchanov, G. A. Otroshenko, and G. B. Yankov, Ternary fission of neutron induced uranium fissioning isomers, this volume.

87) D. Seeliger, Modes of deexcitation of compound nuclei, this volume.