A major purpose of the Technical Information Center is to provide the broadest dissemination possible of information contained in DOE's Research and Development Reports to business, industry, the academic community, and federal, state and local governments.

Although a small portion of this report is not reproducible, it is being made available to expedite the availability of information on the research discussed herein.

1

LA-UR -86-1175

RECEIVED MAY 1 2 1986

Los Alamos National Laboratory is operated by the University of California for the United States Department of Energy under contract W-7405-ENG-36

LA-UR--86-1175

DE86 010168

TITLE EFFECT OF DISSIPATION ON DYNAMICAL FUSION THRESHOLDS

AUTHOR(S) Arnold J. Sierk, T-9

SUBMITTED TO Argonne Symposium on "The Many Facets of Heavy-Ion Fusion Reactions"

MASTER

Los Alamos Los Alamos National Laboratory Los Alamos, New Mexico 87545

FORM NO 836 R4
ST NO 2629 5/81

DISTRIBUTION OF THIS DOCUMENT IS UNLIMITED

EFFECT OF DISSIPATION ON DYNAMICAL FUSION THRESHOLDS

Arnold J. Sierk
Theoretical Division, Los Alamos National Laboratory
Los Alamos, New Mexico 87545

ABSTRACT

The existence of dynamical thresholds to fusion in heavy nuclei ($A \geq 200$) due to the nature of the potential-energy surface is shown. These thresholds exist even in the absence of dissipative forces, due to the coupling between the various collective deformation degrees of freedom. Using a macroscopic model of nuclear shape dynamics, I show how three different suggested dissipation mechanisms increase by varying amounts the excitation energy over the one-dimensional barrier required to cause compound-nucleus formation. The recently introduced surface-plus-window dissipation may give a reasonable representation of experimental data on fusion thresholds, in addition to properly describing fission-fragment kinetic energies and isoscalar giant multipole widths. Scaling of threshold results to asymmetric systems is discussed.

I. INTRODUCTION

In order to form a compound nucleus in a heavy-ion collision, the composite system must pass through a configuration more compact than the fission saddle point[1-5]. In addition, of course, the energy initially in the fission degree of freedom must be shared relatively rapidly with other degrees of freedom, so the system does not "bounce apart" before energy equilibration can occur. For nuclei with mass numbers less than about 200, the fission saddle point is more elongated than the point of first contact of colliding ions, which lies inside the one-dimensional interaction barrier. Therefore fusion occurs naturally for lighter systems which are collided with an energy slightly above the interaction barrier height. In contrast, heavier systems will be driven apart by the disruptive Coulomb forces before passing inside the fission saddle point, since this point corresponds to a more compact configuration than the point of first contact. Thus, for nuclear systems heavier than some threshold value, an additional kinetic energy ΔE above the one-dimensional barrier energy is required to drive the system inside its saddle point and thus to fusion. This phenomenon was qualitatively discussed at least as early as 1969[1], and was quantitatively modeled with non-viscous liquid drops in 1973[3]. Calculations in 1976[4] and 1977[5] used an improved model of the nuclear force which included finite-range effects, still with no dissipative forces. Beginning in about 1980, Swiatecki and collaborators[7-11] developed a dynamical model which utilized a form of one-body dissipation. In this work the extra energy ΔE needed to cause fusion was given the colorful misnomer "extra push". Around the same time began appearing several experimental measurements of the existence of a dynamical threshold energy[12-21]. The conjunction of these latter two developments has led some to the conclusion that the existence of dynamical thresholds was evidence for the validity of this particular form of one-body dissipation. However, the mere existence of a dynamical threshold is predicted by a spectrum of models. More precise measurements and better models offer the possibility of restricting the types of dissipation which are consistent with the data. In fact, later measurements, although not conclusive, seem to indicate that the amount of dissipation predicted in Swiatecki's wall-and-window one-body dissipation is too high[22-26].

I will discuss in this paper the effects of three different models of nuclear dissipation on the dynamical thresholds to fusion in a single unified model, with the eventual aim of learning something about which dissipation models may be appropriate to describe real nuclei. This type of comparison within the confines of a single model is crucial because of the many details which differ among the various models[6-11,27,28]. Distinctions between results caused mainly by potential, inertial or dissipative effects will help to address the true nature of nuclear dissipation. In Section II, I discuss the dynamical model, in Section III, I show some calculated results, and in Section IV I discuss some effects of varying the mass asymmetry of the colliding ions. Also in this section I make a comparison of the calculations to the limited data available. Finally, in Section V I summarize and discuss the results.

II. DYNAMICAL MODEL

The details of the dynamical model I use have been considered elsewhere[29,6]. The basic idea is to specify the shape of an incompressible nucleus with a small number of parameters, then calculate a potential energy, a kinetic energy, and a Rayleigh dissipation function in terms of these parameters and their time derivatives. A set of modified Hamilton's equations with the appropriate initial conditions is numerically integrated to find the trajectories in parameter space of the time evolution of the nuclear shape.

A. Shape Coordinates

For the initial pre-contact stages of a collision, I describe the system by two spheres of the appropriate size. When the centers of the spheres reach a separation such that the matter redistributed into the neck from the region of overlap of the spheres reaches a radius of 3 fm, I then allow the nuclear shapes to evolve according to the modified Hamilton's equations appropriate to a deforming system[6].

For this second stage of the collision, I express the closed, axially symmetric surface specifying the shape in cylindrical coordinates (ρ,ϕ,z) as[30]

$$\rho_s^2(z,\phi) = \rho_s^2(z) = \sum_{i=0}^{N+1} a_i P_i(x) \quad , \tag{1}$$

where P_i is a Legendre Polynomial, a_i are the generalized coordinates, and

$$x = (z-\bar{z})/z_0 \quad , \tag{2}$$

where $2z_0$ is the length of the shape measured along the symmetry axis, and $\bar{z}$ is the value of the coordinate halfway between the end points. Imposing volume conservation and requiring the center of mass to lie at $z = 0$ leads to:

$$a_0 = -\sum_{i=2,4}^{N} a_i \quad , \tag{2a}$$

$$a_1 = -\sum_{i=3,5}^{N+1} a_i \quad , \tag{2b}$$

$$z_0 = 2/3a_0 \quad , \tag{2c}$$

$$\bar{z} = -a_1 z_0^2/2 \quad , \tag{2d}$$

and to the N independent coordinates $\{a_2, a_3, \ldots, a_{N+1}\}$. In this paper I will use $N = 10$. For purely symmetric shapes, all odd coefficients are zero, with 5 degrees of freedom remaining.

Instead of expressing results in terms of these coordinates $\{a_i\}$, which have little intuitive value by themselves, I will project the results onto a two-dimensional subspace of coordinate-independent mass moments. It would be very useful if all authors working in this field were to adopt this or a similar type of approach for presenting their results, as very few outsiders will have any feeling for the meaning of the coordinates describing a particular choice of shape parametrization. I first define a plane perpendicular to the symmetry axis through the neck (if the shape has a neck) or through x = 0 (if there is no neck), which separates the shape into two portions. For the right- and left-hand portions, I define $\langle z^n \rangle_{R \atop L}$ as the mass-weighted moments of z^n of the right- (left-) hand portion of the body[29]. The coordinates of the two-dimensional projected subspace are the center-of-mass separation

$$r = \langle z \rangle_R - \langle z \rangle_L \tag{3a}$$

and the fragment elongation

$$\sigma = \langle (z - \langle z \rangle_R)^2 \rangle_R^{1/2} + \langle (z - \langle z \rangle_L)^2 \rangle_L^{1/2} \quad . \tag{3b}$$

I also define a mass asymmetry coordinate

$$\alpha = (\langle z^0 \rangle_L - \langle z^0 \rangle_R)/(\langle z^0 \rangle_L + \langle z^0 \rangle_R) = (A_L - A_R)/A \quad . \tag{3c}$$

The moment r is the familiar separation of the centers of mass of two colliding ions, and σ is a measure of their relative oblateness or prolateness (before contact) or necking (after contact). For a spherical system, motions along r and σ are orthogonal linear combinations of the quadrupole and hexadecapole normal modes. The mass asymmetry α is 0 for symmetric fission or collisions and is nearly 1 for particle evaporation or for a single nucleon colliding with a heavy nucleus.

B. Potential Energy

From the surface function defined by two spheres or the parameters $\{a_i\}$ I calculate the Coulomb energy of a charge distribution made diffuse by folding a Yukawa function over a uniform density inside the surface[31], and the nuclear energy by twice folding a Yukawa-plus-exponential effective two-body interaction potential over the same uniform density distribution[32,33]. This form of the macroscopic energy, with four parameters determined from elastic electron scattering, from fission barrier heights and from heavy-ion elastic scattering[32-34], describes experimental fusion barriers of light nuclei without further parameters. This method also gives similar results to the proximity potential, but in contrast to the proximity potential is useful for arbitrary shapes. Since I will be concerned in this paper only with head-on collisions I will not discuss angular momentum effects.

C. Kinetic Energy

Since it has been shown that the nuclear inertia for large deformations approaches the incompressible, irrotational value[35], I will model the inertia by means of the Werner-Wheeler approximation to irrotational, incompressible flow, in which the fluid is assumed to move in circular layers[29]. The kinetic energy is expressed in terms of the coordinates and their time derivatives as

$$T = \frac{1}{2} M_{ij}(a)\, \dot{a}_i \dot{a}_j = \frac{1}{2} (M^{-1})_{ij} p_i p_j \quad , \tag{4}$$

where M_{ij} is the deformation-dependent inertia tensor, $p_i = M_{ij}\dot{a}_j$, and I use the convention of summing over repeated indices.

D. Dissipation

The coupling between the collective and internal degrees of freedom gives rise to a dissipative force (also to a fluctuating force which I ignore here) whose component along the coordinate a_i is

$$F_i = - \eta_{ij}(a)\dot{a}_j \quad . \tag{5}$$

For the preliminary stages of a collision, when the ions remain spherical, their radial motion is damped by Randrup's one-body proximity window dissipation[36]. After the dynamical shape evolution is started, either no dissipation or one of three different models is assumed. For historical reasons, I consider hydrodynamical viscosity with a viscosity coefficient of 0.02 terapoise, which reproduces fission-fragment kinetic energy data[6,37], even though there is no theoretical reason to suppose that this model should be appropriate to low energy nuclear processes. For the second model, I use the one-body wall-and-window dissipation originally suggested by Swiatecki[37-39].

Before describing the third dissipation model, I observe that neither of these historical mechanisms reproduces the observed widths of isoscalar giant multipole resonances[40], and both require physical assumptions that are not appropriate for the type of nuclear motion occuring in low energy collisions and fission. Furthermore, two-body collisions, which are assumed to occur uniformly throughout the nuclear volume in the first model, and totally neglected in the second will occur with greatest probability in the nuclear surface region. Various more detailed models of nucleon dynamics inside nuclei imply that one-body damping is significantly reduced from the value predicted by the classical wall formula[41,42].

By abandoning some of the questionable assumptions of the moving-wall model, but still realizing that a combination of one-body and two-body dissipative collisions will be concentrated near the nuclear surface, we are led (in lowest order in the surface diffuseness) to a dissipation rate which has the same functional form as that of the wall model, but a different strength. Specifically, the time rate of change of the collective Hamiltonian is

$$\frac{dH}{dt} = - k_s \rho \bar{v} \int (\dot{n}-D)^2 dS \quad , \qquad (6)$$

where $\dot{n}$ is the normal velocity of the surface element dS, D is the normal component of the average drift velocity of nucleons about to collide with dS, ρ is the nuclear density, $\bar{v}$ is the average speed of nucleons inside the nucleus (3/4 V_F for a Fermi gas), and k_s is a dimensionless strength parameter, measuring the relative probability of an energy-changing collision either between two nucleons or between a single nucleon and the moving one-body potential. The Swiatecki wall formula would correspond to the maximal value k_s=1. This coefficient has been determined by adjustment to isoscalar giant

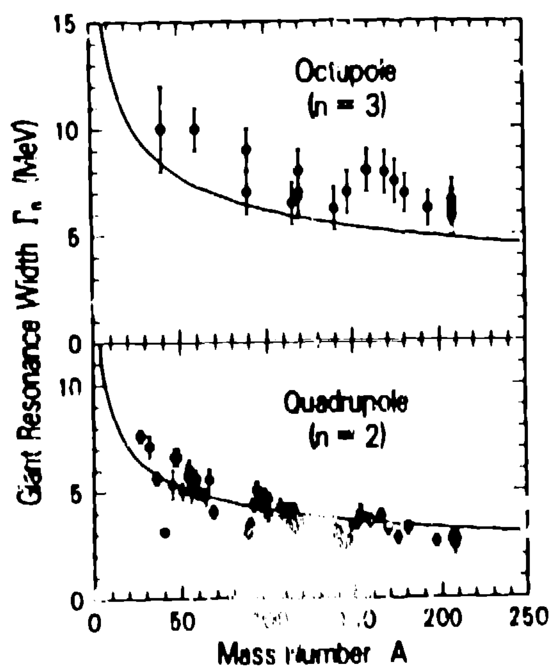

Fig. 1. Simultaneous reproduction of experimental isoscalar giant quadrupole and giant octupole widths by surface dissipation with k_s=0.27.

quadrupole and octupole resonance widths[43] to be k_s=0.27[44].

For necked-in or dumbbell-like shapes, the transfer of nucleons between the end regions leads to an additional dissipation similar to the classical window formula[38,39], but which also includes effects on this transfer due to the deforming of the end regions themselves. This modification causes a small but noticeable change in the results from using the classical window formula. In Figure 2 I show the degree to which data on the mean fission-fragment kinetic energies for nuclei throughout the Periodic Table are reproduced by the previously determined value of k_s = 0.27.

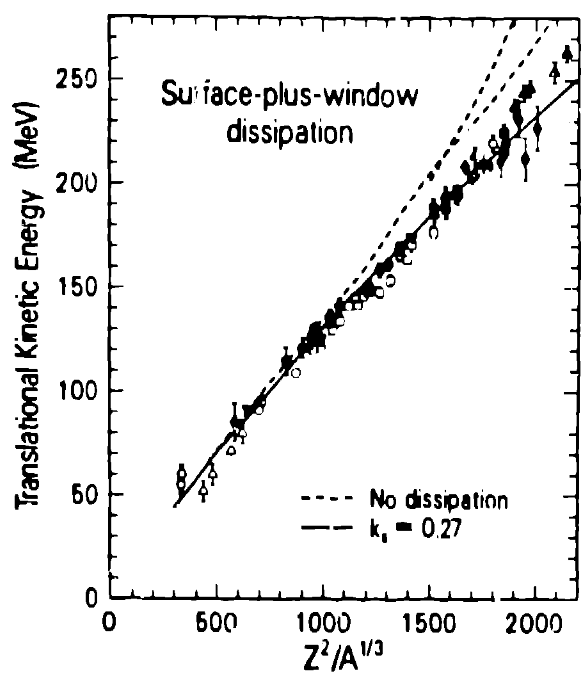

Fig. 2. Reduction of average fission-fragment kinetic energies by surface-plus-window dissipation, compared to experimental values.

A similar degree of reproduction is achieved with hydrodynamic viscosity with a strength of 0.02 TP[37], while the wall-and-window model significantly underestimates the energies for heavier nuclei[37], although it works well for $Z^2/A^{1/3}$ < 1300.

III. CALCULATED RESULTS

Through the remainder of this paper I shall consider either specific nuclear systems with integral Z and A or idealized systems in which the charge and mass numbers of each of the colliding nuclei are related by Greens approximation to the line of beta stability[45]

$$\bar{Z}_i = \frac{A_i}{2}\left(1 - \frac{0.4\,A_i}{A_i+200}\right) \quad . \qquad (7)$$

This choice, which is nearer to experimental conditions than is a beta-stable compound system leads to significantly different results for thresholds[27,28].

In Figure 3, taken from Ref. 46 I show calculated macroscopic potential-energy contours as a function of the mass moments r and σ for mass symmetric deformations of ^{220}U. Although quantitatively slightly different from our present potential-energy calculations, this Figure illustrates certain qualitative features relevant to dynamical fusion thresholds.

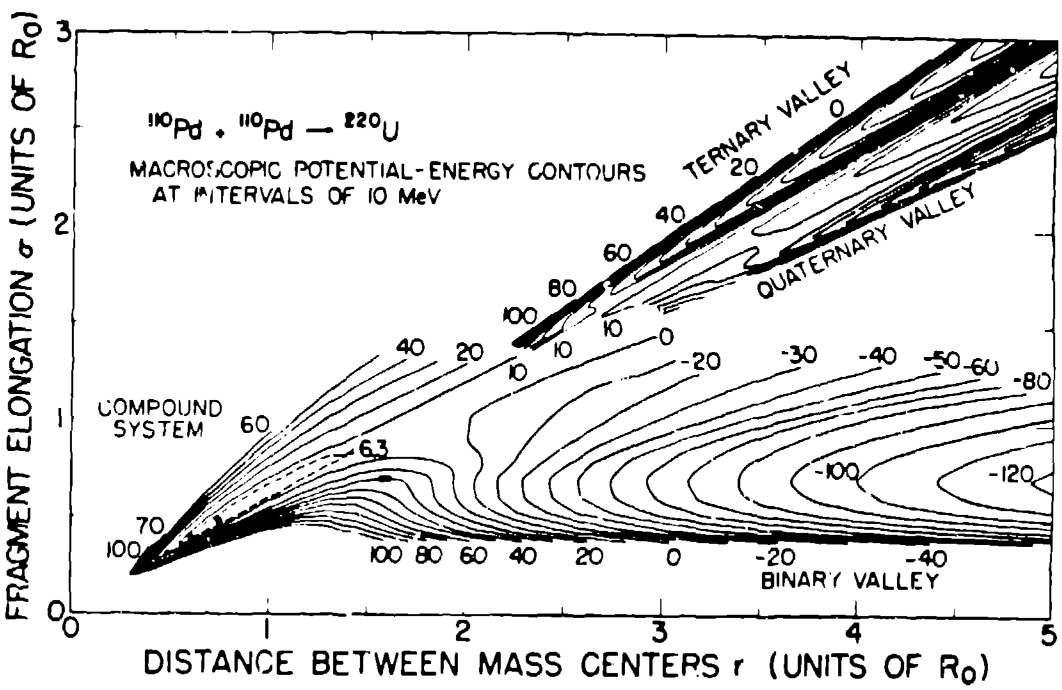

Fig. 3. Potential-energy contours, in units of MeV, for the reaction $^{110}Pd + ^{110}Pd \rightarrow ^{220}U$, calculated with a single-Yukawa macroscopic model. The location of the sphere is given by the solid point, the location of two touching spheres by the two adjacent solid points, and the fission saddle point by the intersecting dashed contours.

In a collision of two ^{110}Pd nuclei, the system moves from right to left up the binary valley with $\sigma \sim 0.7$. The calculated maximum in the one-dimensional interaction-energy barrier occurs just outside the point of tangency, which occurs at $r \sim 1.6$, $\sigma \sim 0.7$. The spherical macroscopic ground state (single point) occurs at $r = 0.75$, $\sigma \sim 0.5$. The fission saddle point ($r \sim 1.4$, $\sigma \sim 0.6$) is located where the dashed contours intersect. An important qualitative fact is that the maximum in the interaction barrier lies above the saddle point in energy, and is displaced from it[2]. If this system is started at rest at the top of the barrier (near the two touching points), it appears that the forces on the system will tend to push it on a trajectory which would not pass inside the saddle point. This expectation is borne out by dynamical calculations[5,6]. As one moves to heavier (more fissile) systems, the location of the saddle point becomes closer to the spherical ground state, as shown in Figure 4, and the height of the barrier decreases. The existence of a dynamical threshold is thus due to this qualitative nature of the potential-energy surface, while its exact nature depends on details of the collective dynamics.

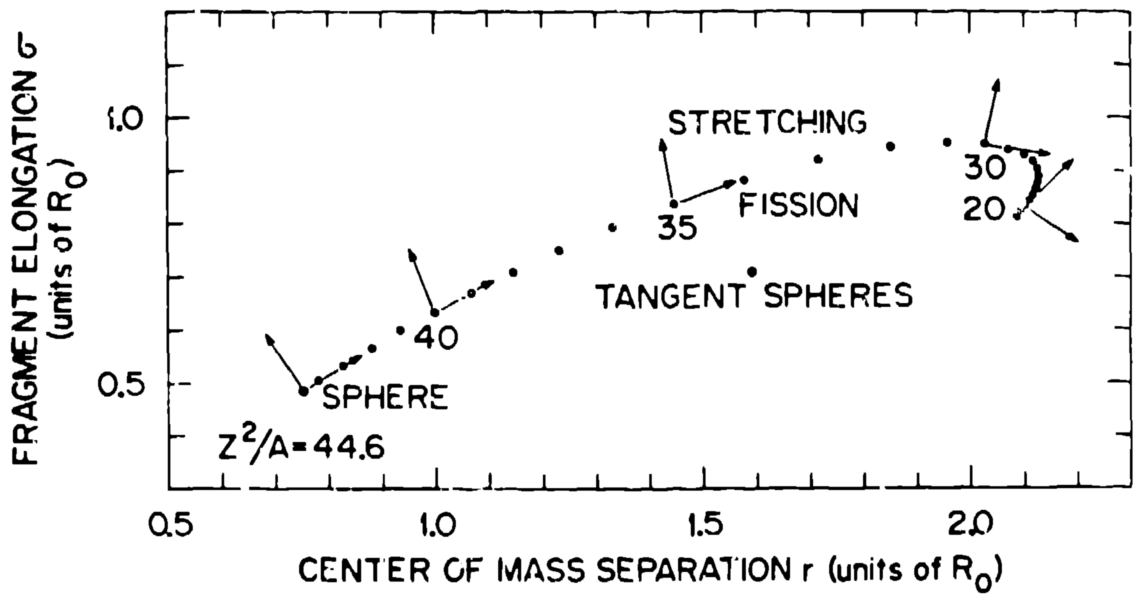

Fig. 4. Locations in r-σ space of the saddle points of beta-stable nuclei with Z^2/A from 18 to 44.6. The arrows denote the directions of the two lowest symmetric normal modes (quadrupole and hexadecapole for a sphere).

In Refs. 6-11 it was shown for a simplified system that if the dynamical trajectory passes through the saddle point for the system with $(Z^2/A)_{thresh}$ when started at rest at the top of the barrier, the extra energy above the barrier needed to drive a more fissile system with $Z^2/A = (Z^2/A)_{thresh} + \Delta(Z^2/A)$ through its saddle point is

$$\Delta E = \gamma[\Delta(Z^2/A)]^2 + \ldots \quad . \tag{8}$$

The assumptions used in this paper, namely that the colliding nuclei are constrained to be spherical until a 3 fm radius neck is formed lead to

$$\Delta E = \beta\,[\Delta(Z^2/A)] + \gamma'\,[\Delta(Z^2/A)]^2 + \ldots \quad . \tag{9}$$

These functional forms of ΔE *vs*. Z^2/A occur independently of the details of the mechanism of dissipation. Therefore, the existence of a dynamical threshold to fusion says nothing about the character of nuclear dissipation. However, the values of $(Z^2/A)_{thresh}$, β, γ', *etc*. will depend on such details.

As an illustration of the results of dynamical calculations of fusion, I show in Figure 5 the shapes as a function of time, and in Figure 6 the trajectories in r-σ space calculated both for no dissipation and for the three dissipation models introduced above. The calculations show the trajectories for the mass-symmetric fusion of beta-stable nuclei to form a system with

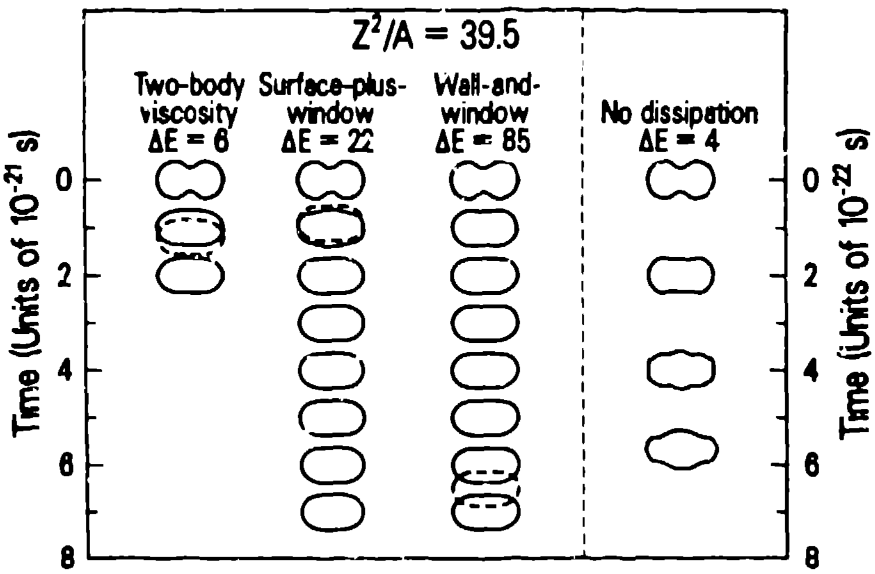

Fig. 5. Nuclear shapes as a function of time for the symmetric collision of beta-stable nuclei with just enough energy to form a compound nucleus with $Z^2/A = 39.5$, for no dissipation, two-body viscosity, wall-and-window one-body dissipation, and for surface-plus-window (one-body and two-body) dissipation. The dashed shapes are the fission saddle-point shape, shown at the times when the trajectories pass closest to it.

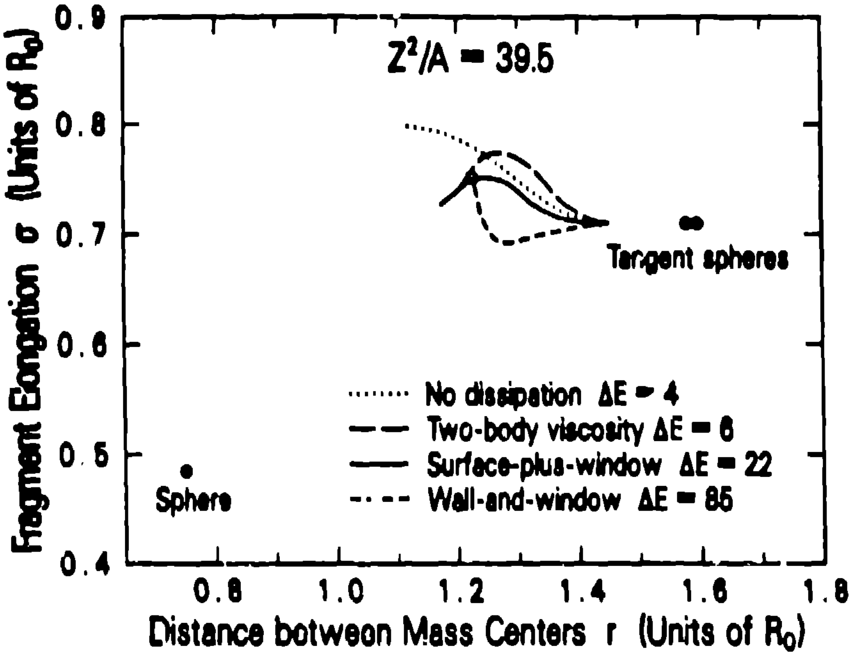

Fig. 6. Trajectories in r-σ space for the collisions pictured in Fig. 5. The location of the saddle-point shape is shown by the open point. Other points are as in Fig. 3.

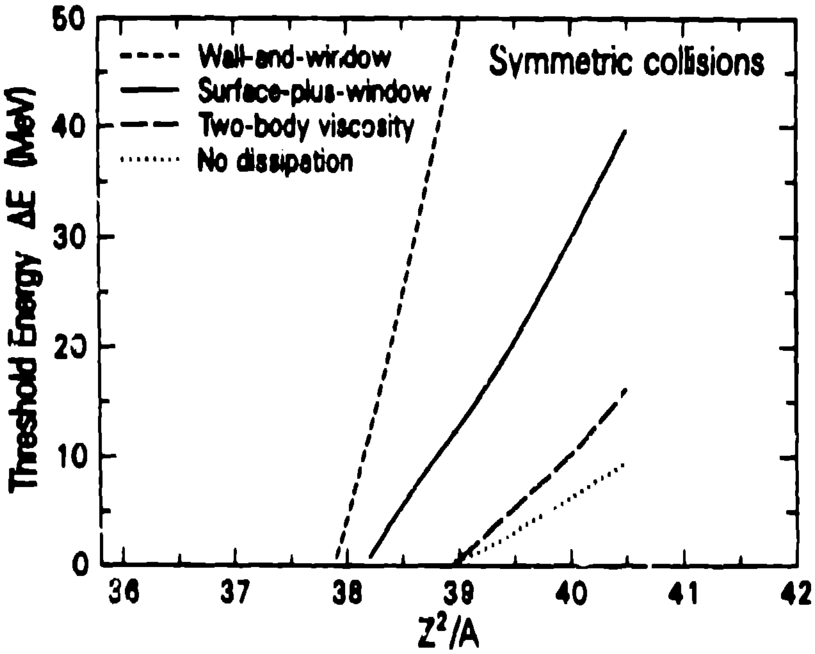

Fig. 7. Fusion threshold energy ΔE vs. Z^2/A of the composite system for symmetric collisions of beta-stable nuclei for no dissipation, two-body viscosity of 0.02 TP, wall-and-window dissipation, and for surface-plus-window dissipation.

total $Z^2/A = 39.5$. In each case the system has been given the specified ΔE in the center-of-mass frame sufficient to just cause fusion. In only the case of no dissipation, a trajectory which passes outside the saddle point initially may eventually pass inside, although the dynamical model breaks down before it is possible to unambiguously determine whether fusion is likely to occur. In all dissipative calculations, the threshold can be determined unambiguously.

Results from several analogous calculations are compiled in Figure 7, where I show ΔE vs. Z^2/A for symmetric collisions using the four types of dissipation. Because of the poor approximation of not allowing the colliding nuclei to deform before contact, the thresholds in this model occur at too high a value of Z^2/A. The most significant point here is the extremely different behavior of ΔE vs. Z^2/A for the different modes of dissipation.

IV MASS-ASYMMETRIC COLLISIONS AND SCALING

Up to this point I have discussed only mass-symmetric collisions. By considering the extreme example of a proton or alpha particle fusing with a heavy nucleus, it is easy to see that increasing the mass asymmetry while keeping fixed the overall mass and charge of the composite system leads to a reduction of ΔE[7-11]. This is because an asymmetric system of two tangent spheres is loosely speaking more compact than a symmetric one, while the same symmetric saddle point must be reached in each case to cause fusion. Swiatecki[7-11] has suggested than an effective Z^2/A for asymmetric systems be defined as

$$(Z^2/A)_{eff} = 4Z_1Z_2/[A_1^{1/3}A_2^{1/3}(A_1^{1/3}+A_2^{1/3})] \quad . \tag{10}$$

However, it is obvious that the fusion behavior cannot be entirely determined by the incident channel, but also is effected by the actual Z^2/A of the composite system. This observation and some experimental evidence[26] has led to the suggestion of using[6]

$$(Z^2/A)_{mean} = [(Z^2/A)(Z^2/A)_{eff}]^{1/2} \tag{11}$$

as the appropriate scaling variable for asymmetric systems.

I have calculated fusion thresholds using surface-plus-window dissipation for systems composed of beta-stable pairs with mass asymmetries $\alpha = 0.0$, 0.2,

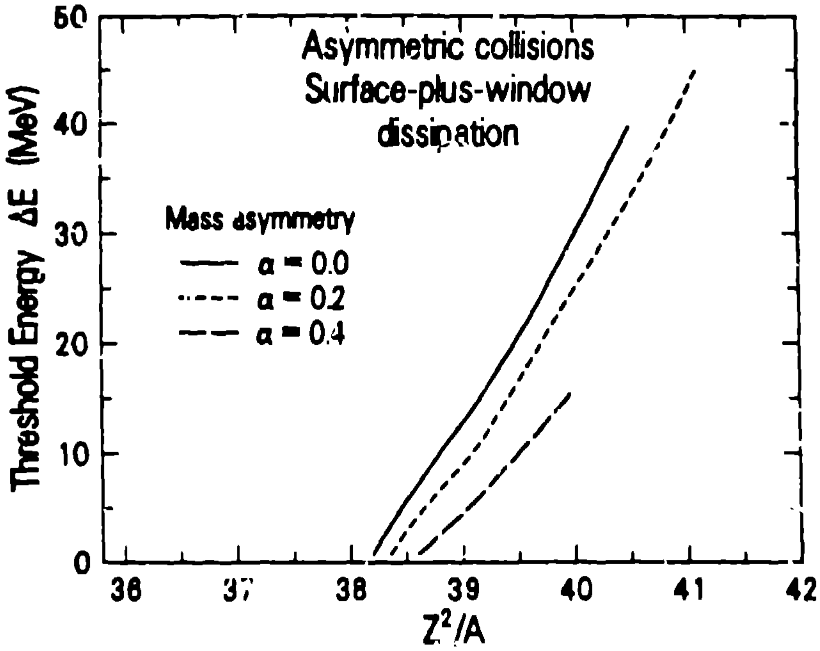

Figure 8. Fusion threshold energy ΔE vs. Z^2/A for fusion of beta-stable nuclei with mass asymmetry α = 0.0, 0.2, and 0.4.

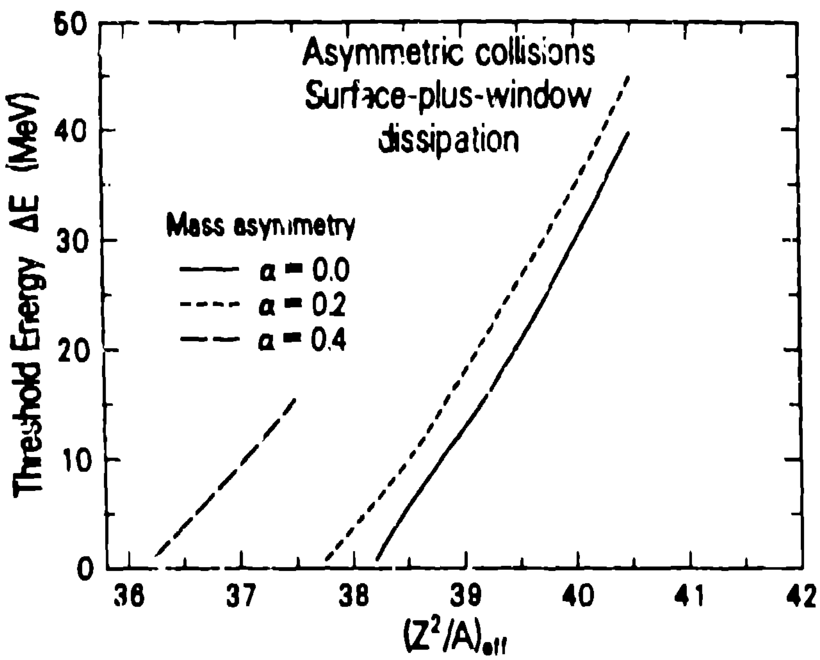

Fig. 9. Same as Fig. 8 plotted vs. $(Z^2/A)_{eff}$.

and 0.4 (technical difficulties with the calculations have so far precluded finding results for $\alpha > 0.4$). In Figures 8, 9, and 10, I show these calculated thresholds as a function of Z^2/A, $(Z^2/A)_{eff}$, and $(Z^2/A)_{mean}$, respectively. The Z^2/A and $(Z^2/A)_{eff}$ scalings are not appropriate for this particular model, while the $(Z^2/A)_{mean}$ scaling works very well for $\alpha = 0.2$, but less well for $\alpha = 0.4$. This deviation is at least reasonable, as there is no *a priori* reason to expect the exact weighting of incoming and outgoing channels predicted by Eq. 11 to be valid.

In Figure 10, I also show the limited data available from evaporation-residue measurements in this mass region[25,26].

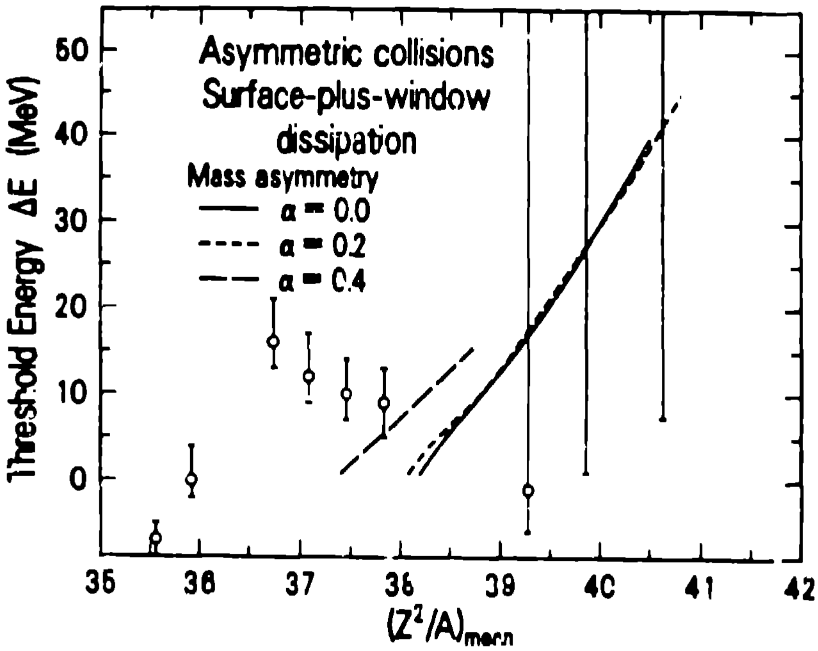

Fig. 10. Same as Fig. 8 plotted *vs.* $(Z^2/A)_{mean}$. Experimental points[25,26] from evaporation-residue measurements are also shown.

V. DISCUSSION AND SUMMARY

There are several problems inherent in trying to use experimental fusion data to unambiguously rule out one or more mechanisms of dissipation. In the first place, there is some uncertainty as to what to use for the interaction-barrier height; *i.e.*, what do you subtract from an experimentally measured fusion barrier height to arrive at a ΔE? In plotting the data on Figure 10 I have used the barriers calculated from the Yukawa-plus-exponential potential[32-34]. In the original references[25,26], barriers found in a Bass model were used; the differences between these two approaches have a deviation which increases with Z^2/A, reaching 17 MeV for the right-most point in Fig. 10. This sort of deviation of theoretical barriers can mask the true dynamical threshold behavior. Also, it appears that simple scaling formulas are not sufficiently accurate for less simple models[27,28]. These factors lead me to recommend that unless and until a model with some other simple scaling appears that results of calculations should be expressed as the amount of total energy necessary to cause fusion.

It is also clear from Figure 10 that thresholds that are only lower limits[26] cannot be used to distinguish between some models of dissipation. However, realizing that an improved model of the incoming channel will shift the curves in Figs. 7-10 to the left, and that a symmetric system has the largest threshold for a given $(Z^2/A)_{mean}$, it is clear that a system somewhat lighter than ^{244}Fm[26] would offer a much better chance of determine a threshold with a non-zero lower and a non-infinite upper limit. Specifically, I would recommend serious consideration of measuuring the dynamical thresholds for reactions forming evaporation residues in the immediate neighborhood of the reaction

$$^{110}_{46}Pd + ^{110}_{46}Pd \rightarrow ^{220}_{92}U \quad . \qquad (12)$$

which has a composite system Z^2/A = 38.5,

$$^{103}_{45}Rh + ^{110}_{46}Pd \rightarrow ^{213}_{91}Pa \qquad (13)$$

with

Z^2/A = 38.7, or

$$^{104}_{44}Ru + ^{104}_{44}Ru \rightarrow ^{208}_{88}Ra \quad . \qquad (14)$$

with Z^2/A = 37.2. This choice is not motivated by the fact that reaction (12) was extensively studied in Refs. 4 and 6, but by the considerations that.

1) For a given compound nucleus, the symmetric entrance channel has the largest predicted ΔE, with the greatest separation between the predictions of the various dissipation models.

2) A system heavy enough to have a significant ΔE (15 to 20 MeV), but light enough to have some probability of surviving as an evaporation residue is needed.

The converse of the preceeding discussion is that in order to form very heavy systems with the minimum amount of excitation energy, the most asymmetric target-projectile configurations are favored. However, the best (still crude) estimates using the surface-plus-window dissipation model for the system

$$^{48}Ca + ^{248}Cm \rightarrow ^{296}116 \qquad (15)$$

(α = 0.68) indicate that a ΔE of at least 30 MeV is needed to reach a nearly spherical configuration, which when added to the potential energy acquired in passing to the sphere, leads to a minimum excitation energy of the compound nucleus of roughly 50 MeV[48]; at this excitation energy there probably would be no shell stabilization energy. This is consistent with the lack of observation of superheavy elements formed in fusion reactions, despite the many heroic efforts that have been applied to the search.

ACKNOWLEDGEMENT

Many of the ideas expressed in this paper are the result of collaboration with J. R. Nix and K. T. R. Davies. This work was supported by the U. S. Department of Energy.

REFERENCES

1. W. J. Swiatecki, in *Proceedings of the International Conference on Nuclear Reactions Induced by Heavy Ions, Heidelberg, Germany, 1969*, edited by R. Bock and W. R. Hering (North Holland, Amsterdam/American Elsevier, New York, 1970), p. 729.

2. W. J. Swiatecki and S. Bjørnholm, Phys. Rep. 4C, 325 (1972).

3. A. J. Sierk and J. R. Nix, in *Proceedings of the Third International Atomic Energy Agency Symposium on the Physics and Chemistry of Fission, Rochester, New York, 1973* (IAEA, Vienna, 1974), Vol. II, p. 273.

4. A. J. Sierk and J. R. Nix, in *Proceedings of the Symposium on Macroscopic Features of Heavy-Ion Collisions, Argonne, Illinois, 1976*, Argonne National Laboratory Report No. ANL-PHY-76-2, 1976 (unpublished), Vol. I, p. 407.

5. J. R. Nix and A. J. Sierk, Phys. Rev. C 15, 2072 (1977).

6. K. T. R. Davies, A. J. Sierk, and J. R. Nix, Phys. Rev. C 28, 679 (1983)

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

8. W. J. Swiatecki, Phys. Scr. 24, 113 (1981).

9. W. J Swiatecki, in *Proceedings of the 4th International Conference on Nuclei Far from Stability, Helsingør, Denmark, 1981*, CERN Report No. CERN-81-09, 1981 (unpublished), p. 781.

10. W. J. Swiatecki, Nucl. Phys. A376, 275 (1982).

11. S. Bjørnholm and W. J. Swiatecki, Nucl. Phys. A391, 471 (1982).

12. H. Sann, R. Bock, Y. T. Chu, A. Gobbi, A. Olmi, U. Lynen, W. Müller, S. Bjørnholm, and H. Esbensen, Phys. Rev. Lett. 47, 1248 (1981).

13. R. Bock, Y. T. Chu, M. Dakowski, A. Gobbi, E. Grosse, A. Olmi, H. Sann, D. Schwalm, U. Lynen, W. Müller, S. Bjørnholm, H. Esbensen, W. Wölfi, and E. Morenzoni, Nucl. Phys. A388, 334 (1982).

14. S. Bjørnholm, in *Proceedings of the International Workshop X on Gross Properties of Nuclei and Nuclear Excitations, Hirschegg, Austria, 1982*, edited by H. Feldmeier, Technische Hochschule Darmstadt DReport ISSN-0720-8715, 1982, p. 15.

15. S. Bjørnholm in *Proceedings of the Nuclear Physics Workshop, Trieste, Italy, 1981*, edited by C. H. Dasso, R. A. Broglia, and A. Winther (North-Holland, Amsterdam, 1982), p. 521.

16. S. Bjørnholm, Nucl. Phys. A387, 51c (1982).

17. S. Bjørnholm, Comments Nucl. Part. Phys. 11, 9 (1982).

18. J. R. Huizenga, J. R. Birkelund, W. U. Schröder, W. W. Wilcke, and J. J. Wollersheim, in *Proceedings of the 3rd Adriatic Europhysics Study Conference on the Dynamics of Heavy-Ion Collisions, Hvar, Yugoslavia, 1981*, edited by N. Cindro, R. A. Ricci, and W. Greiner (North-Holland, Amsterdam, 1981), p. 15.

19. H. Gäggeler, W. Brüchle, J. V. Kratz, M. Schädel, K. Sümmerer, G. Wirth, and T. Sikkeland, in *Proceedings of the International Workshop X on Gross Properties of Nuclei and Nuclear Excitations, Hirschegg, Austria, 1982*, edited by H. Feldmeier, Technische Hochschule Darmstadt Report ISSN-0720-8715, 1982, p. 40.

20. W. Westmeier, R. A. Esterlund, A. Rox, and P. Patzelt, Phys. Lett. 117B, 163 (1982).

21. B. Sikora, J. Bisplinghoff, M. Blann, W. Scobel, M. Beckerman, F. Plasil, R. L. Ferguson, J. Birkelund, and W. Wilcke, Phys. Rev. C 25, 686 (1982).

22. G. Münzenberg, W. Reisdorf, S. Hofmann, Y. K. Agarwal, F. P. Hessberger, K. Poppensieker, J. R. H. Schneider, W. F. W. Schneider, K.-H. Schmidt, H.-J. Schött, P. Armbruster, C.-C. Sahm, D. Vermeulen, Z. Phys. A315, 145 (1984)

23. C.-C. Sahm, H.-G. Clerc, K.-H. Schmidt, W. Reisdorf, P. Armbruster, F. P. Hessberger, J. G. Keller, G. Münzenberg, D. Vermeulen, Z. Phys. A319, 113 (1984).

24. J. G. Keller, K.-H. Schmidt, H. Stelzer, W. Reisdorf, Y. K. Agarwal, F. P. Hessberger, G. Münzenberg, H.-G. Clerc, C.-C. Sahm, Phys. Rev. C 29, 1569 (1984).

25. C.-C. Sahm, H.-G. Clerc, K.-H. Schmidt, W. Reisdorf, P. Armbruster, F. P. Hessburger, J. G. Keller, G. Münzenberg, and D. Vermeulen, preprint IKDA 85/2, 1985 (unpublished).

26. H. Gäggeler, T. Sikkeland, G. Wirth, W. Brüchle, W. Bögl, G. Franz, G. Hermann, J. V. Kratz, M. Schädel, K. Sümmerer, W. Weber, Z. Phys. A316, 291 (1984).

27. H. Feldmeier, in *Proceedings of the International Workshop X on Gross Properties of Nuclei and Nuclear Excitations, Hirschegg, Austria, 1982*, edited by H. Feldmeier, Technische Hochschule Darmstadt Report ISSN-0720-8715, 1982. p. 26.

28. H. Feldmeier, Argonne National Laboratory Report ANL-PHY-85-2, 1985 (unpublished).

29. K. T. R. Davies, A. J. Sierk, and J. R. Nix, Phys. Rev. C13, 2385 (1976).

30. S. Trentalange, S. E. Koonin, and A. J. Sierk, Phys. Rev. C22, 1159 (1980).

31. K. T. R. Davies and J. R. Nix, Phys. Rev. C 14, 1977 (1976).

32. H. J. Krappe, J. R. Nix, and A. J. Sierk, Phys. Rev. Lett. 42, 215 (1979).

33. H. J. Krappe, J. R. Nix, and A. J. Sierk, Phys. Rev. C 20, 992 (1979).

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

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

36. J. Randrup, Ann. Phys. (N.Y.) 112, 356 (1978).

37. J. R. Nix and A. J. Sierk in *Proc. Int. Conf. on Nuclear Physics, Bombay, 1984* (World Scientific, Singapore, 1985), p. 365.

38. J. Błocki, Y. Boneh, J. R. Nix, J. Randrup, M. Robel, A. J. Sierk, and W. J. Swiatecki, Ann. Phys. (N.Y.) 113, 330 (1978).

39. J. Randrup and W. J. Swiatecki, Ann. Phys. (N.Y.) 125, 193 (1980).

40. J. R. Nix and A. J. Sierk, Phys. Rev. C 21, 396 (1980); 25, 1068 (1982).

41. J. J. Griffin and M. Dworzecka, Phys. Lett., to be published.

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

43. F. E. Bertrand, private communication (1985).

44. J. R. Nix and A. J. Sierk, in *Proc. of the Winter Workshop on Nuclear Dynamics IV, Copper Mountain, Colorado, 1986*, to be published.

45. A. E. S. Green, *Nuclear Physics* (McGraw-Hill, New York, 1955), pp. 185, 250.

46. P. Möller and J. R. Nix, Nucl. Phys. A272, 502 (1976).

47. K. T. R. Davies, J. R. Nix and A. J. Sierk, Phys. Rev. C 28, 1181 (1983).

48. A. J. Sierk, in *Proceedings of the International Symposium on Superheavy Elements, Lubbock, Texas, 1978*, edited by M. A. K. Lodhi (Pergamon, New York, 1978), p. 479.