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MACROSCOPIC DESCRIPTION OF THE INTERACTION BETWEEN TWO COMPLEX NUCLEI

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MACROSCOPIC DESCRIPTION OF THE INTERACTION BETWEEN TWO COMPLEX NUCLEI

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1. THE MACROSCOPIC AND LEPTODERMOUS IDEALIZATIONS

The word macroscopic means that the particle number A is large, i.e. $A^{-1} \ll 1$, and that the discreteness of nucleons is disregarded. In addition I shall use the leptodermous idealization:

$$\frac{\text{Nuclear surface thickness}}{\text{Cube root of nuclear volume}} \ll 1,$$

in virtue of which a nuclear shape may be defined. The central problem in macroscopic nuclear dynamics is to give an account of the time evolution of the shape of the nuclear system.

2. THE SIXFOLD WAY--A MADMAN'S EXPERIMENT

To establish a broad framework for our thinking about this problem imagine the following experiment designed by a madman: a beam consisting of a mixture of all elements, with energies from zero to relativistic, impinges on a target also made of a mixture of all elements! What kinds of reactions would one expect in such a nightmare experiment?

Figure 1 is an attempt to classify the fate of the stream of particles into

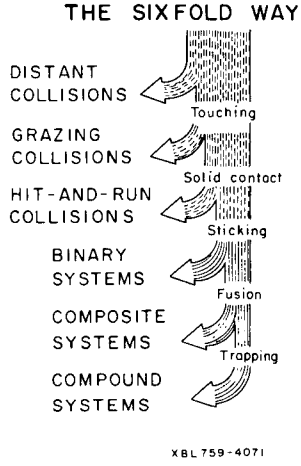

Fig. 1. An attempt to classify collisions between two nuclei into six main categories.

six categories: Nuclei that don't Touch undergo Distant Collisions, where only electromagnetic interactions, modified by the tail of the nuclear potential, come into play. (Touching could be defined approximately as touching of the Fermi levels of the two systems.) Nuclei that touch but do not make Solid Contact undergo Grazing Collisions. (Solid Contact might be defined approximately as beginning when the half-density contour acquires the topology of a single object.) Nuclei that make solid contact but do not Stick make Hit-and-Run Collisions. (Sticking might be defined approximately in terms of the sharing of linear momentum.) Nuclei that stick but do not Fuse form a Binary System. (Fusion might be defined in terms of a loss of identity of the pieces, associated with a filling-in of the neck.) Nuclei that fuse but are not Trapped form a Composite System. The remainder are trapped in a potential energy hollow, and form a Compound System.

The first problem in describing the above processes, whether quantally or classically, is to map the potential energy as a function of the shape of the system.

3. THE POTENTIAL ENERGY OF A LEPTODERMOUS SYSTEM

With certain qualifications it is possible to prove the following theorem about the functional form of the potential energy V(shape) of a thin-skinned system of volume $\mathscr{V}$ whose shape is specified by an (effective sharp) surface Σ.

$$V(\text{shape}) = V_{\text{Bulk}} + V_{\text{Surface Layer}} \tag{1}$$

where $V_{\text{Bulk}} = c_1 \mathscr{V}$, and

$$\begin{aligned} V_{\text{Surface Layer}} = {} & c_2 \mathscr{S} + c_3 \mathscr{K} + c_4 \mathscr{G} + c_4' \mathscr{Q} + \\ & + \text{corrections that tend to zero as } \mathscr{V} \text{ grows} \\ & + \text{Proximity Energy } V_p. \end{aligned} \tag{2}$$

Here $c_1 \cdots c_4'$ are, for a fixed $\mathscr{V}$, constants independent of shape. The shape dependence is given by the Proximity Energy V_p and by the following functionals of shape: $\mathscr{S}$ = area of surface Σ, $\mathscr{K}$ = integral over Σ of $\kappa (\equiv R_1^{-1} + R_2^{-1})$, $\mathscr{G}$ = integral over Σ of $(R_1 R_2)^{-1}$, $\mathscr{Q}$ = integral over Σ of κ^2. Here R_1, R_2 are the principal radii of curvature at a point on Σ. For a nucleus the terms in c_i are of order $A(A^{-1/3})^{i-1}$. The last term in Eq. (2) is less familiar and is the subject of the next section.

4. THE PROXIMITY ENERGY[1]

The expansion in $A^{-1/3}$ given above is an asymptotic[2] expansion and fails if there are violent contortions in Σ, for example such that elements of the surface face each other across a gap or crevice whose width is of the order of the surface thickness. If this width D is only a slowly varying function of position the additional energy may be approximated by

$$V_P \approx \int\int e(D)dx\ dy. \tag{3}$$

Here e(D) is the interaction energy per unit area between two flat parallel surfaces at separation D and the integral is over the gap or crevice, assumed to lie approximately in the x-y plane.

For the family of gap or crevice functions D(x,y) defined implicitly by

$$\sum_{n=0}^{N} c_n D^n = \alpha x^2 + \frac{1}{\alpha} y^2 \quad (\equiv r^2),$$

where α is an area-preserving transverse stretching parameter, Eq. (3) leads to the simple result (independent of α)

$$V_P = \pi \sum_{1}^{N} n\, c_n\, I_{n-1},$$

where

$$I_n = \int_{s \text{ or } 0}^{\infty} D^n\, e(D)dD,$$

are moments of the universal interaction function e(D). The lower limit s or 0, distinguishes between gaps (in which case D has a least value s and, for negative s, the two objects overlap) and crevices (in which case D is zero at a finite value of r, given by $\sqrt{c_0}$, and there is no density doubling.) The case N = 2, for which the relation between D and the "stretched" cylindrical radius vector r is a conic section, gives

$$V_P = \pi\, c_1 I_0 + 2\pi\, c_2 I_1 \tag{4}$$

where, for gaps, I_0, I_1 are two "universal" functions and, for crevices, two "universal" constants, characteristic of the material of which the surfaces are made. Gaps or crevices corresponding to the following shapes are covered by Eq.(4). (For crevices the overlapping portions are erased.) (i) For two juxtaposed coaxial paraboloids with radii of curvature at tips C_1, C_2 and tip distance s (which may be negative): $c_1 = 2\bar{R}$, $c_2 = 0$, where $\bar{R} = C_1C_2/(C_1 + C_2) \equiv \bar{C}$. (ii) For two equal coaxial spheroids with semi-axes C, B (C along the line of centers) and tip distance s: $c_1 = B^2(C + \frac{s}{2})/C^2$, $c_2 = -B^2/4C^2$. For spheres C = B. (iii) As above but with one spheroid infinitely large (i.e. a plane): $c_1 = 2B^2(C + s)/C^2$, $c_2 = -B^2/C^2$. (iv) For two equal juxtaposed hyperboloids with axes C, B, (C along the line of centers) and tip distance s: $c_1 = B^2(C - \frac{s}{2})/C^2$, $c_2 = B^2/4C^2$. (v) As above, with one hyperboloid degenerating into a plane: $c_1 = 2B^2(C-s)/C^2$, $c_2 = B^2/C^2$. (vi) Single hyperboloid of one or two sheets: $c_1 = 0$, $c_2 = B^2/4C^2$. (vii) Coaxial juxtaposed circular cones with semi-

opening angles α_1, α_2 and tip distance s: $c_2 = (\cot\ \alpha_1 + \cot\ \alpha_2)^{-2}$, $c_1 = -\ 2sc_2$.
(viii) In addition, for the case of a gap between two coaxial elliptic paraboloids with tip distance s, with radii of curvature A_1 and B_1 in the principal planes of curvature through the tip of paraboloid 1, and A_2, B_2 for paraboloid 2, and azimuthal angle ϕ between the principal planes of curvature of 1 and 2: $c_1 = 2\bar{R}$, $c_2 = 0$, where

$$\frac{1}{\bar{R}^2} = \frac{1}{A_1B_1} + \frac{1}{A_2B_2} + \left(\frac{1}{A_1A_2} + \frac{1}{B_1B_2}\right)\sin^2\phi + \left(\frac{1}{A_1B_2} + \frac{1}{A_2B_1}\right)\cos^2\phi.$$

By taking the negative derivative of V_p with respect to s in case (i) (which is a lowest order approximation to the general case of arbitrary gently curved surfaces with least separation s) we arrive at the following Proximity Force Theorem: "The force between two gently curved surfaces as a function of the separation degree of freedom is equal to the interaction potential per unit area between two flat surfaces, times 2π into the geometric mean of the principal radii of curvature of the gap width function." Thus

$$F(s) \equiv -\frac{\partial V_p}{\partial s} = 2\pi\ \bar{R}\ e(s). \tag{5}$$

The (universal) function e(s) vanishes for s beyond the range of the surface width, has a minimum of about minus twice the surface energy coefficient γ when the surfaces are in solid contact at $s \approx 0$ (and the two bodies have combined to make approximately uniform bulk matter), and then rises linearly to infinity as the bodies pile up, with s going to negative values. It follows that the maximum attraction between two gently curved bodies is $F_{max} \approx -\ 4\pi\bar{R}\gamma$.

Introducing a unit of length b proportional to the width (diffuseness) of the surface, the proximity potential may be written in terms of a dimensionless universal function $\Phi(\zeta)$:

$$V_p = 4\pi\gamma\bar{R}b\bullet\Phi(\zeta), \tag{6}$$

where

$$\Phi(\zeta) = (2\gamma b)^{-1}\int_s^\infty e(s)ds, \qquad \zeta = s/b.$$

J. Randrup, using W. D. Myers' version of the Seyler-Blanchard nuclear Thomas-Fermi method[3] has calculated the universal nuclear function Φ, to which an excellent analytic approximation is (see Fig. 2):

$$\Phi(\zeta < 1.2511) = -\frac{1}{2}(\zeta - 2.54)^2 - n\ (\zeta - 2.54)^3, \tag{7}$$

with $n = 0.0852 \approx 1/12$.

$$\Phi(\zeta > 1.2511) = -\ 3.437\ \exp(-\zeta/0.75). \tag{7a}$$

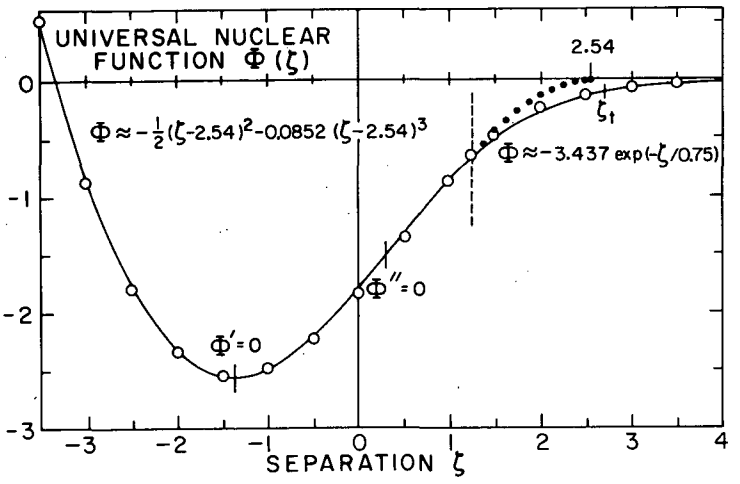

Fig. 2.

The universal nuclear function $\Phi(\zeta)$. The circles give the cubic and exponential approximations joined smoothly at the dashed line. The dots are the continuation of the cubic to where it touches the ζ axis at 2.54. The Fermi levels touch at $\zeta_t = 2.71$.

The following values of the parameters entering V_p would be reasonable for a nuclear system with total neutron, proton and mass numbers N, Z, A:

Surface Energy[4] $\gamma = 0.9517 \left[1 - 1.7826 \left(\frac{N-Z}{A}\right)^2\right] \approx 1 - 2\left(\frac{N-Z}{A}\right)^2$ MeV/fm^2,

Surface Width[5] $b \approx 1$ fm, Nuclear Central[5],[6] Radius $C = R - b^2/R$,

where R = Effective Sharp Radius = $(1.13 + 0.0002A)A^{1/3}$fm $\approx 1.15\ A^{1/3}$ fm.

5. ACCURACY OF THE LEPTODERMOUS THEOREM

First a comparison with the exactly soluble mathematical model of Krappe and Nix[7] of two sharp spheres (surface separation s), whose volume elements interact by a Yukawa interaction of range a and strength written as $V_0/4\pi a^3$. For a single such sphere of radius R the exact energy is

$V = V_0\ [-4\pi R^3/3 + 2\pi a R^2 + 0.R - 2\pi a^3 + 2\pi a(R + a)^2 \exp(-2R/a)]$. With

$V_0 = 2.2105$ MeV/fm^3, a = 1 fm, R = 1.2 $A^{1/3}$ fm this gives:

$V = -16A + 20A^{2/3} + 0.A^{1/3} - 13.89 + 13.89(1.2A^{1/3} + 1)^2 \exp(-2.4A^{1/3})$ MeV

$= -250 + 125 + 0 - 13.89 + 0.55$ MeV for A = 16.625. Equations (1) and (2) applied to this case give $V = 2.2105\ \text{MeV/fm}^3\ [\mathscr{V} + a\mathscr{S}/2 + 0.\mathscr{K} - 3a^3\mathscr{G}/32 - a^3\mathscr{Q}/8] = -250 + 125 + 0 - 13.89$ MeV. For two equal spheres separated by s the exact additional interaction energy is $V_{int}/4\pi a^2\gamma = -4(\rho \cosh \rho - \sinh \rho)^2 \cdot (\sigma + 2\rho)^{-1} \exp(-\sigma + 2\rho)$, where $\rho = R/a$, $\sigma = s/a$ and γ is the surface energy coefficient given by $aV_0/2$. This amounts to -9.35 MeV for $\sigma = 0$ and -2.95 MeV for $\sigma = 1$. The Proximity Formula applied to this case (using for the radius of a Krappe-Nix sharp sphere, generating a diffuse potential, the mean

between the central radii of the density and potential distributions) gives $V_P/4\pi a^2\gamma = -\frac{1}{2}(\rho - \rho^{-1})\exp(-\sigma - 2\rho)$, which amounts to -9.51 MeV for $\sigma = 0$ and -3.50 MeV for $\sigma = 1$.

The second comparison is with the Orsay Group calculations[8] of the interaction potential as a function of the separation degree of freedom for 17 pairs of model nuclei described by the "energy-density formalism." As required by the proximity theorem the 17 potentials are approximately reducible to a universal function whose appearance is quite similar to our Φ. The absolute magnitudes are also about the same. Figure 3 illustrates this in the case of $^{63}Cu + {}^{197}Au$.

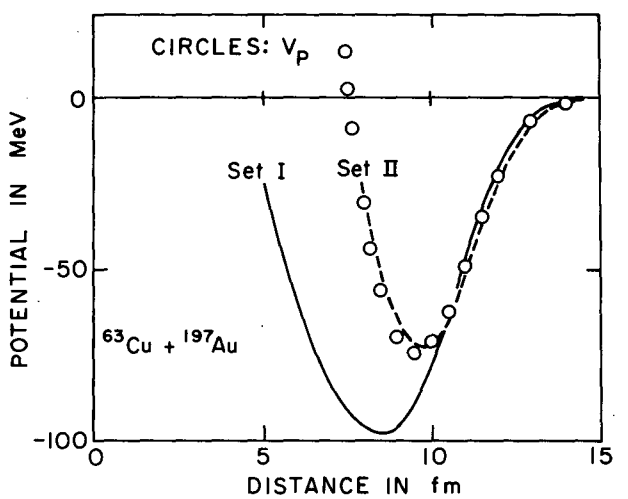

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Fig. 3.

The proximity potential Vp between ^{63}Cu and ^{197}Au (circles) compared with the Orsay Group calculations with two sets of parameters. Set II (dashed curve) is the preferred choice.

6. THE FROZEN IDEALIZATION

If all the degrees of freedom except the separation s are frozen one can have a lot of fun with the cubic or cubic-exponential approximation (7,7a) to the nucleus-nucleus interaction. Thus let us write the interaction, including Coulomb energy, as

$$\begin{aligned} V(r < r_1) &= V_P + V_{Coul} \\ &= 4\pi\gamma\bar{C}\, b(-\tfrac{1}{2}\delta^2 + n\,\delta^3) + (Z_T Z_P e^2/r_1)\,(1 + \nu\delta + \nu^2\delta^2 + \nu^3\delta^3 + \cdots) \\ &= (Z_T Z_P e^2/r_1)\,(1 + k_1\delta - k_2\,\delta^2 + k_3\,\delta^3) \equiv E_t k(\delta). \end{aligned} \tag{8}$$

Here $Z_T e$, $Z_P e$ are the target and projectile charges, r_1 (equal to $C_T + C_P + 2.54b$) is the distance at which the cubic nuclear interaction comes into play, δ is the dimensionless penetration variable $(r_1 - r)/b$, ν is b/r_1 (about 0.1 for many systems), and $k(\delta)$ stands for the cubic whose coefficients are

$k_1 = \nu$, $k_2 = \nu(\frac{1}{2}X^{-1} - \nu)$, $k_3 = \nu(nX^{-1} + \nu^2)$, where X is a dimensionless "Coulomb Parameter":

$$X \equiv \frac{Z_T Z_P e^2/r_1^2}{4\pi\gamma\,\bar{C}} = \frac{\text{Coulomb repulsion at } r_1}{\text{Nominal maximum nuclear attraction}}.$$

Finally E_t is $Z_T Z_P e^2/r_1$, a natural energy unit of the problem. The distance r_1 is a natural length--it is also close to where the Fermi levels of the nuclei would touch, and is thus approximately where nuclear reactions would begin to occur. Here is a list of analytical results that follow: Touching ($\approx$ reaction) cross section $\sigma_t \approx \pi r_1^2(1 - E_t/E)$. It starts from zero at the touching energy E_t (see Fig. 4). Barrier position δ_B and height B in

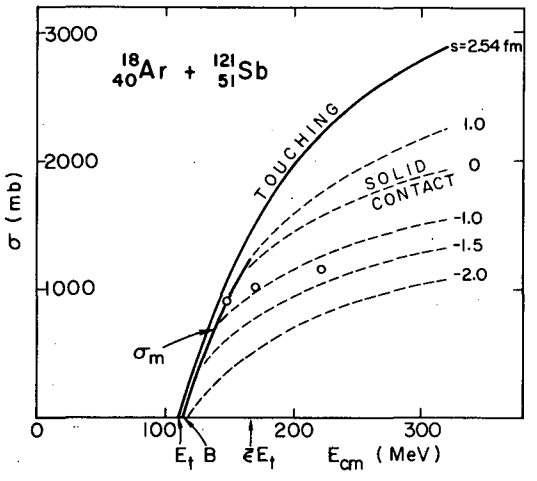

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Fig. 4. A scattering diagram for $^{18}_{40}$Ar + $^{121}_{51}$Sb calculated using analytic formulae. The cross sections for penetrating to various values indicated on the right are shown by the dashed curves, some of which, for lower energies, are intercepted by the curve σ_m. The circles are recent experimental measurements (by H. Gauvin et al., to be published in Physics Letters) which suggest penetrations to about 0.8 - 1.2 fm beyond solid contact.

$V(r)$: $\delta_B = \delta_0(1-\sqrt{1-\lambda})$, $B = E_t k(\delta_B)$, where $\delta_0 \equiv k_2/3k_3$ and λ is a "Pocket Parameter" defined by $\lambda \equiv 3k_1k_3/k_2^2 = 12nX \cdot (1 + \nu^2X/n)/(1 - 2\nu X)^2 \approx X/(1 - 2\nu X)^2$ for small ν, $\approx X$ for very small ν. The "pocket" in $V(r)$ spills out at $\lambda_{crit} = 1$, i.e. at a critical Coulomb parameter $X_{cr} = [\sqrt{(3n + \nu)^2 + 2\nu^2} - (3n + \nu)]/4\nu^2 \approx (1 + 4\nu)^{-1}$. The critical (limiting) value of δ_B is then $(\frac{1}{2} - \nu X_{cr})/3(n+\nu^2X_{cr}) \approx 2(1 + 2\nu)/(1 + 4\nu)$. The limiting value of B is $B_{cr} = E_t(1 + 2\nu/3(X_{cr}^{-1} - 2\nu)) \approx E_t(1 + 2\nu/3(1 + 2\nu))$. The depth of the pocket in $V(r)$ is $4k_2^3E_t(1-\lambda)^{3/2}/27\,k_3^2$.

The cross section for penetrating to below a pre-assigned distance r is $\pi r^2(1 - V(r)/E)$ (see Fig. 4), except that with decreasing energy some of these hyperbolae are intercepted by the curve σ_m, which then gives the cross section for penetrating to r. The equation for σ_m is $\pi r_m^2(1 - V(r_m)/E)$, where r_m locates the maximum in the effective energy $V_{eff}(r)$ when this maximum osculates the energy E (V_{eff} is $V(r) + L^2/2Mr^2$, where L is the angular momentum and M the reduced mass of the system). The value of r_m is $r_1 - b\delta_m$, where δ_m is the lowest root of the

cubic $K(\delta) = \varepsilon$, where $\varepsilon \equiv E/E_t$ and $K(\delta) = \frac{1}{2} + K_1\delta - K_2\delta^2 + K_3\delta^3$, with $K_1 = \frac{1}{2}(X^{-1} + \nu)$, $K_2 = -\frac{1}{2}\nu^2 + (\frac{3}{2}n + \nu)X^{-1}$, $K_3 = \frac{5}{2}\nu(nX^{-1} + \nu^2)$. The pocket in V_{eff} spills out when the maximum and minimum in V_{eff} coalesce, which occurs at a limiting value $\bar{\delta}$ given by $D - U$, where $D = K_2/3K_3$, $U = (K_2^2 - 3K_1K_3)^{1/2}/3K_3$. The limiting energy at that moment is given by $\bar{\varepsilon} = K(\bar{\delta})$. In terms of $\bar{\varepsilon}$ the solution for δ_m locating the maximum is given by

$$\delta_m = D - 2U \cos \frac{1}{3} \{\pi - \cos^{-1}[1-(\bar{\varepsilon}-\varepsilon)/2K_3U^3]\}.$$

The depth and width of the pocket in V_{eff} are given approximately by $\frac{2}{3} B_2^3/B_3^2$ and $3B_2/B_3$, where $B_2 = (2k_2 - 6k_3\delta_m)E_t/b^2 - 3L^2/Mr_m^4$, $B_3 = 6k_3E_t/b^3 + 12L^2/Mr_m^5$.

The above formulas bring the discussion of nucleus-nucleus scattering in the frozen idealization within the realm of an analytic treatment. In particular, attempts to correlate fusion cross sections in terms of a "critical distance" of penetration/9/, can be carried out analytically.

7. SCALING

In the above theory there are two energies (Coulomb and nuclear) and two lengths (r_1 and b); hence the two dimensionless parameters X and ν. It follows from dimensional analysis that for two pairs of target and projectile combinations with the same X and ν, their dimensionless scattering diagrams (with σ measured in units of πr_1^2 and E in units of $Z_TZ_Pe^2/r_1$) would be identical (if the frozen idealization were valid).

It turns out that in practice once X is matched for two pairs of nuclei, ν is approximately matched, so the situation is analogous to macroscopic fission theory, where properly scaled fission barriers, spontaneous half-lives, fragment energies etc. should be primarily universal functions of the fissility parameter x. The following table lists two systems which are "similar" in the above sense.

	λ	X	ν	Z_TZ_P	E_t(MeV)	πr_1^2(mb)
$Hf_{72}^{180} + Ar_{18}^{40}$	0.5542	0.4517	0.0796	1296	148.6	4954
$Kr_{36}^{84} + Kr_{36}^{84}$	0.5531	0.4486	0.0819	1296	152.8	4688

Note that λ (or X) are roughly proportional to the fraction $Z_TZ_P/A_T^{1/3}A_P^{1/3}(A_T^{1/3} + A_P^{1/3})$, of which the numerator is often the controlling factor. The ordering of experimental cross-section data according to λ (or X)

might serve a similar purpose as the ordering of fission data according to x: the assessment of the validity of the underlying idealizations and the isolation of the expected deviations.

8. OTHER DEGREES OF FREEDOM

The expected deviations from the frozen idealization are associated with unfreezing: (a) Other shape degrees of freedom, (b) Fragment-orientation degrees of freedom/10/, (c) Internal degrees of freedom (this leads to dissipation). In connection with (a) J. Błocki is preparing an extensive atlas of nuclear potential energy maps (primarily sums of Coulomb and surface energies). The parametrization is two spheres joined by a conic; each map is for a fixed mass asymmetry of the spheres and displays the deformation energy as function of elongation and neck degrees of freedom. Figure 5a shows such a family of shapes and the associated energy for a mass ratio 1:1, Fig. 5b for 1:4 and Fig. 5c for 1:19. The fissility parameter x for the whole system is 0.85 in each case ($Z \simeq 102$). Colliding nuclei would come in from the right, touch at the top of the ridge, and either be trapped near the sphere (left lower corner) or re-separate, making use--as in fission--of the neck degree of freedom. Błocki's atlas covers x's from 0 to 1.5 and mass ratios 1:1 to 1:19.

I have nothing new to say about (b). Concerning (c) I would like to pose a question:

9. IS NUCLEAR DYNAMICS, SUPERFLUIDITY ASIDE, DOMINATED BY ONE-BODY DISSIPATION?

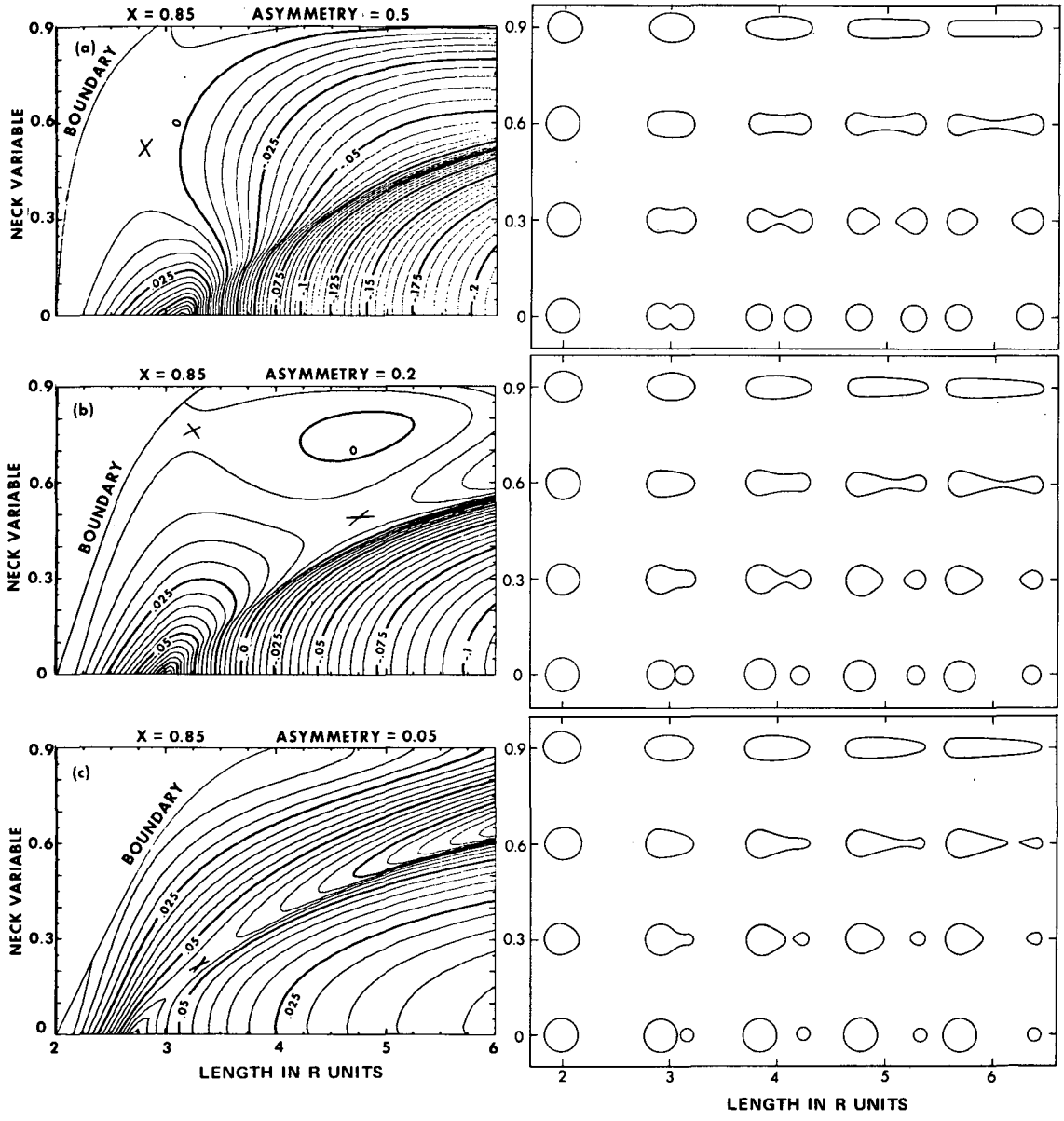
Ordinary viscosity in liquids is the result of <u>two</u>-body molecular forces which degrade nonuniform collective flows into heat. A different, <u>one</u>-body mechanism for exchanging energy between collective and internal degrees of freedom is the collision of particles with a moving container wall. This energy exchange between wall and particles, calculated to lowest order in the ratio of particle to wall velocities, is responsible for the Ideal Gas Law $pV = RT$! By going to the next order in the ratio of velocities, J. Błocki and M. Robel arrive at the following formula for the rate of energy flow $\dot{E}$ into a gas, (macroscopically at rest and with mass density ρ) composed of independent particles with average speed v, and contained in a vessel of fixed volume, whose walls deform with normal velocities specified by $\dot{n}$:

$$\dot{E} = \rho\, v \oint \dot{n}^2\, d\sigma\,. \tag{9}$$

[For a degenerate Fermi gas $v = (3/4)$(Fermi velocity).] This formula applies only for $|\dot{n}| \ll v$, and for time intervals during which the original velocity distribution of the particles has not been modified substantially by the wall motion. To illustrate the order of magnitude of the dissipation predicted by the

Fig. 5.

Potential energy maps corresponding to families of shapes shown on the right. (The maps stop at the line marked "Boundary.") The deformation energy is in units of the surface energy of the standard sphere.



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above formula Błocki and Robel have added such a damping to the equations of motion of an idealized nuclear droplet vibrating irrotationally around the spherical shape in its nth multipole mode. One may define a factor of overdamping μ, which gives the number of times the dissipation term exceeds the value for which critical damping would occur for the given mode. The formula for μ is

$$\mu = (3/8)\sqrt{3}(9\pi)^{1/3}\sqrt{(\hbar^2/mr_0^2)/(4\pi r_0^2\gamma)}\left\{\frac{n-1}{n}\left(n+2-\frac{20x}{2n+1}\right)\right\}^{-\frac{1}{2}}\cdot A^{1/6},$$

where m is the nucleon mass and γ the surface energy coefficient. The following table illustrates the degree of overdamping for 3 idealized nuclei and four harmonics (Lysekil[4] nuclear parameters were used):

	n=2	n=4	n=8	n=16
$^{20}_{10}Ne$	3.01	1.94	1.38	0.99
$^{120}_{50}Sn$	5.24	2.88	1.94	1.37
$^{238}_{92}U$	9.46	3.58	2.27	1.58

The damping appears to be huge (as with ordinary viscosity when applied to small systems[11]), but note that contrary to conventionally damped drops the damping is largest for long-wavelength modes (and large systems). A high nuclear viscosity is indeed implied by deep inelastic scattering, which suggests that nuclei often behave like two pieces of tar that get stuck but do not flow together. However, high two-body viscosity is inadmissible in fission[12] where it would predict stretched out scission shapes leading to much too low fission-fragment kinetic energies. This stretching out is due to the stronger damping of short-wavelength modes (like necking) compared to a general elongation. Since with one-body damping the situation is just the reverse, it is quite possible that in a fission process dominated by one-body damping the scission shape would be moderately compact and would reproduce the observed fragment energies. One difficulty that this might remove is the apparent lack of temperature dependence of the conventional viscosity that would be deduced from fission (the fragment energies are practically independent of the excitation of the nucleus in the range of a few to 100 MeV). This is difficult to explain with ordinary viscosity which should depend sensitively on temperature. For a Fermi liquid like He^3 viscosity goes like T^{-2}, which is associated with the need to find, for the two colliding particles,

unoccupied two-body scattering states in the diffuse region (of width proportional to T) in the Fermi momentum distribution. In one-body excitations this problem is not present, and so the one-body dissipation might only vary slowly with temperature. (The v in Eq. 9 would only change substantially when $T \simeq 20$ MeV).

It seems to me that in nuclear dynamics we may be at a stage analogous to the time around 1950 when, after some 15 years dominated by the prejudice that two-nucleon interactions were overwhelming, it was realized that, after all, a one-body approach, with independent nucleons put into a common container, was a fair starting point for discussing nuclear structure. The one-body, dissipation-dominated, nuclear dynamics should be explored in the spirit of the fifties, by taking seriously the usefulness of the simple concept of a common container, also when its shape is a function of time.

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