$\lesssim 30°$), it appears that the mechanism of momentum transfer and energy dissipation at projectile energies of about 20 MeV/nucleon is closely connected with the fast emission of a "jet" of particles preferentially into the forward direction. The measurement of the projectile-energy dependence of this process should yield additional important information on the change in the reaction mechanism from deeply inelastic and quasielastic reactions at low energies to fragmentation reactions at higher energies.

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[6] Qualitative insight may be gained by noting that $\theta_A \approx \theta_B \approx 90°$; then $p_R^{\parallel} = p_A \cos\theta_A + p_B \cos\theta_B$ is the sum of two small terms and $p_R^{\perp} = p_A \sin\theta_A - p_B \sin\theta_B$ is the difference of two large terms. Therefore, $p_R^{\parallel}$ is better determined than $p_R^{\perp}$.

Giant *E*1 Mode and Its Energy Broadening from the Charge Distributions in Heavy-Ion Reactions

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The narrow experimental width of the Z distribution at fixed mass asymmetry in heavy-ion reactions is interpreted in terms of a giant dipole mode whose damping increases with excitation energy. Further theoretical predictions and relevant experiments are discussed.

The gaint $E1$ mode in cold nuclei is best known through its photoexcitation which is manifested in a peak at an energy $E = 78A^{-1/3}$ with a width of typically 4–6 MeV. While it would be of extreme interest to study it also in *hot systems,* this is not feasible by means of either photoexcitation or photon decay. We want to illustrate here a possible way to study the $E1$ mode in hot systems. The $E1$ degree of freedom is involved in the charge distribution at fixed mass asymmetry in the binary heavy-ion reactions (and in fission). This can be seen if one thinks of the intermediate complex in heavy-ion reactions as a single deformed nucleus. Because of the deformations, the giant $E1$ resonance is split into two components: a singly degenerate longitudinal component (parallel to the deformation axis) and a doubly degenerate transversal component perpendicular to the deformation axis. The former component is clearly responsible for the left-right charge fluctuations and controls the fragment charge distribution at fixed mass asymmetry after the intermediate complex breaks up. The equilibration of the $E1$ mode in heavy-ion reactions, or the equilibration of the neutron-to-proton ratio of the two fragments, occurs quite fast, faster in fact than the mass asymmetry degree of freedom, possibly as a result of exchange effects which allow for charge transfer without mass transfer. Accordingly, the most probable charges can be obtained by minimizing the potential energy of the two fragments in contact with respect to the charge of one of the fragments at constant fragment mass. This well-documented feature of heavy-ion reactions only provides information about the potential energy term of the collective $E1$ Hamiltonian. In principle one could obtain information for the whole Hamiltonian by a measurement of the charge distribution at fixed mass. Since in the great majority of cases the

$E1$ phonon energy is expected to be much larger than the temperature, the $E1$ mode is expected to be in its ground state. As an example, let us consider the reaction Ni + Ar at 280-MeV bombarding energy, whose mass and charge distribution has been studied in detail.[1] From the maximum linear dimension of the intermediate complex one obtains the relevant $E1$ phonon energy:

$$\hbar\omega = 94/d = 8\text{–}10 \text{ MeV},$$

where d is the semimajor axis of the intermediate complex. The effect of the neck between the two fragments is not clear at the present time. This may shed some doubt on the estimate of the phonon energy. From the internal excitation energy of the complex one obtains

$$T = (E_x/a)^{1/2} = 2 \text{ MeV}.$$

Since $\hbar\omega/T \simeq 4\text{–}5 \gg 1$, the collective $E1$ mode should be mainly in its ground state. Therefore the Z distribution at fixed mass asymmetry should be given by the square of the modulus of the ground-state wave function and the second moment of the distribution is expected to be

$$\sigma_z^2 = \hbar\omega/2c \simeq 0.6\text{–}0.8 \text{ (charge units)}^2,$$

where c is the stiffness constant associated with the $E1$ mode, or

$$V_{(E1)} = \tfrac{1}{2}c(z - z_0)^2.$$

The analysis of the experimental charge and mass distribution shows that mass and charge are strongly correlated as expected, with a correlation coefficient $r = 0.97$. However, the intriguing result for the second moment of the Z distribution at constant A is

$$\sigma_z^2 = 0.3 \text{ (charge units)}^2$$

substantially *smaller* than expected. The correction of σ_z^2 for particle emission is expected to be minor because of the stabilizing effect of the particle binding energies on the probability of particle emission. (If a proton is emitted, the next proton will be more strongly bound and less likely to be emitted.) Even more surprising is the fact that the experimental value of σ_z^2 is well reproduced if one assumes just a classical statistical distribution in Z, namely,

$$\sigma_z^2 = T/c \simeq 0.3 \text{ (charge units)}^2.$$

The outstanding problem is then to understand why the distribution in Z is classical rather than quantal, as expected.

Some hint of what might be happening comes from the energy widths of the giant $E1$ resonances. These widths are nonzero, thus indicating that the $E1$ mode does not give rise to a pure state but to a state coupled to some doorway states. In other words, the collective $E1$ mode is damped. In photoexcitation, the giant resonance is mainly a one-particle, one-hole (1p-1h) state and presumably owes its width to the coupling into the 2p-2h states. In the present case, at relatively high excitation energy (60 MeV), the collective mode is an (np-nh) state which may couple into $[(n+1)$p-$(n+1)$h$]$, or (np-nh), or again, $[(n-1)$p-$(n-1)$h$]$ states. The transition may be given by[2]

$$\lambda_+ = \frac{2\pi}{\hbar} v^2 \frac{g^3U^2}{(p+h+1)},$$

$$\lambda_0 = \frac{2\pi}{\hbar} v^2 g^2 U \frac{3(p+h)-2}{4},$$

$$\lambda_- = \frac{2\pi}{\hbar} v^2 g[ph(p+h)-2],$$

where p and h are the particle and hole number, respectively, U is the excitation energy, and v^2 is the average matrix element connecting the giant mode to the doorway states. These three transition rates become approximately equal for p-h states which are near the equilibrium configuration. The resulting damping is energy dependent and due mainly to the increasing density of the doorway states with increasing energy. The energy width is given by $\Gamma = \hbar(\lambda_+ + \lambda_0 + \lambda_-)$. It is interesting to see the consequence of this coupling to the Z distribution. Following Bohr and Mottelson[3] with a simple generalization, we can describe the coupling of the collective state $|a\rangle$ to the doorway states $|\alpha\rangle$. The exact state $|i\rangle$ is given by

$$|i\rangle = |a\rangle + \frac{P}{E_i - H_0 - V} V|a\rangle,$$

where $P = \sum_\alpha |\alpha\rangle\langle\alpha|$, H_0 is the unperturbed Hamiltonian, and V is the perturbation. The overall normalization is determined by the condition $\langle i|i\rangle = 1$. [This, together with $(E_i - H_0 - V)^{-1} \simeq (E_i - H_0)^{-1}$, in zeroth-order perturbation theory, leads to the known results.]

The relevant charge distribution is given by $p_i(z) = \int dx\,|\psi_i(z,x)|^2$, where $\psi_i(z,x) = \langle z,x|i\rangle$ and x denotes all other variables which must be projected out. In order to compare theory with experiment we have to consider the average of the distribution over an energy interval around E_i.

We can write

$$p_i(z)_{\rm av} = \int dx\left[|\{\psi_i(z,x)\}_{\rm av}|^2 + \{|\psi_i^{\rm fl}(z,x)|^2\}_{\rm av}\right]$$

with $\psi_i^{\rm fl} = \psi_i - \{\psi_i\}_{\rm av}$ the "fluctuating" wave function. The fluctuating part can be shown to be responsible for the *broadening* of the distribution. It leads to a statistical distribution for Z. We want to show that the first term can lead to a *narrowing* of the distribution. For this purpose we have to consider the averaged Green's function $\{1/(E_i - H_0 - V)\}_{\rm av}$. This average has been considered extensively in the literature.[4] For large systems and high excitation energies only the average diagonal matrix elements of the resolvent have to be considered and it can be shown that

$$\left\{\left\langle \alpha \left| \frac{1}{E_i - H_0 - V} \right| \alpha \right\rangle\right\}_{\rm av} = \frac{1}{E_i - E_\alpha - i\Gamma},$$

where Γ is the imaginary part of the "equivalent optical potential" describing the dissipation of the state $|a\rangle$ into the states $|\alpha\rangle$. Since we have to deal with particle-hole configurations near equilibrium, the previous width is equal to the width which is responsible for the coupling of the doorway α with other states. The amplitude of the state $|a\rangle$ contained in the average eigenstate $|i\rangle$ is given by

$$c_a(i) = \left(1 + \sum \frac{V_{\alpha a}^2}{(E_i - E_\alpha - i\Gamma)^2}\right)^{-1/2}.$$

Correspondingly,

$$c_\alpha(i) = c_a(i)\,\frac{v_{\alpha a}}{E_\alpha - E_i - i\Gamma}.$$

In summary, and omitting for simplicity the bracket of the average,

$$|i\rangle = c_a(i)|a\rangle + \sum_\alpha c_\alpha(i)|\alpha\rangle.$$

The next step is to establish that the sum over α in the above equation is a *coherent* one and thus the corresponding term describes a *wave packet*, i.e., it leads to a *narrowing* of the distribution. Although one can simply see that the bracket combination $|a\rangle\langle\alpha|$ is invariant under phase transformation of the vectors $|\alpha\rangle$, the following formal proof demonstrates it clearly. Assume, namely, that $V_{a\alpha}$ are random numbers, $\sum_\alpha V_{a\alpha} = 0$, i.e., their phases $\chi_{a\alpha}$ [$V_{a\alpha} = \exp(i\chi_{a\alpha}) \times |V_{a\alpha}|$], are random. Using the identity $\sum_\alpha |\alpha\rangle\langle\alpha| = 1 - |a\rangle\langle a|$ and the fact that $(E_\alpha - E_i - i\Gamma)^{-1}$ is a smooth function, we can simply show that $\sum_\alpha c_\alpha(i)|\alpha\rangle \neq 0$, i.e.,

$$\sum_\alpha \frac{|V_{a\alpha}|\exp(i\chi_{a\alpha})}{E_i - E_a - i\Gamma}|\alpha\rangle \neq 0.$$

This proves that if $V_{a\alpha}$ is random, the vectors $|\alpha\rangle$ contain phases which destroy the random property of $V_{a\alpha}$. Having established this point from first principles, we are entitled to use as first guess a simple-as-possible model. First we skip the irrelevant variables χ and we write the average wave function associated with the charge asymmetry coordinate as

$$\varphi_i(z) = c_a(i)\psi_a(z) + D^{-1}\int_0^\infty dE_\alpha\, c_\alpha(i)\psi_\alpha(z), \qquad (1)$$

where D is the level spacing of the available doorway states and $\psi_a(z)$ is the ground-state wave function of the $E1$ mode: $\psi_a(z) = (2\pi\hbar\omega/c)^{1/2} \times[-cz^2/2\hbar\omega]$. Qualitatively one sees already that, as the coupling increases, the integral in (1) becomes progressively dominant and the more $|\alpha\rangle$ states that are called into play by the strength of the coupling, the narrower $\varphi_i(z)$ becomes. As a qualitative first guess on the $\psi_\alpha(z)$ we can use the plane-wave expression

$$\psi_a(z) = (2\pi\hbar\omega/c)^{1/2}\exp[iz(c/2\hbar\omega)^{1/2}(E_\alpha/D_1)^{1/2}],$$

where the plane waves are normalized to unity in a z box of volume corresponding to that of the harmonic oscillator. The term $(E_\alpha/D_1)^{1/2}$ scales the energies in terms of the average single-particle spacing D_1. We further assume the matrix element $V_{a\alpha}$ can be replaced by an overall average value v. The integral in (1) can be evaluated by integration in the complex plane and gives a result

$$\varphi_i(z) = c_a(i)\left\{\psi_a(z) + \frac{2\pi V_{a\alpha}}{D}\left(\frac{2\pi\hbar\omega}{c}\right)^{1/2}\exp\left[-iz(c/2\hbar\omega)^{1/2}(1/D_1)^{1/2}(E_i - i\Gamma)^{1/2}\right]\right\}. \qquad (2)$$

The second moment of the z distribution, σ_z^2, can then be obtained from the z distributions given by the square of the modulus of Eq. (2).

A simple calculation can be performed in order to obtain the energy dependence of σ_z^2. The matrix element v^2 can be estimated by requiring the energy E and the width Γ of the giant resonance for the corresponding spherical nucleus to be 15 and 4 MeV, respectively. Furthermore, we take the most probable number of quasiparticles in the excited intermediate complex to be $1 + p + h \approx (2gE)^{1/2}$, where $g = D_1^{-1}$ is the single-particle level density and has been taken to be equal to g

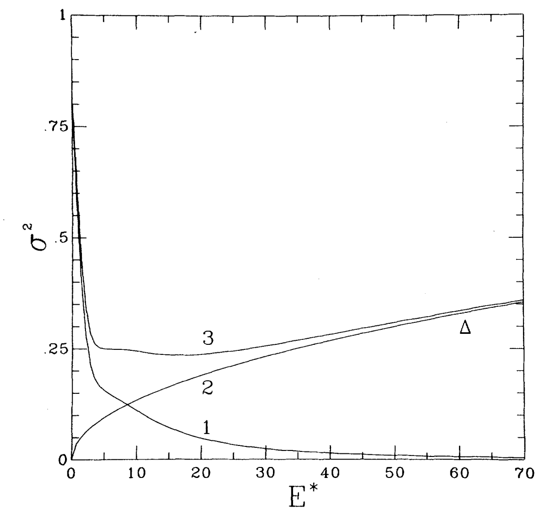

FIG. 1. Energy dependence of the quantal (curve 1) and of the classical statistical width (curve 2). Curve 3 represents the sum of both widths and the point indicates the experimental value.

$=(6/\pi^2)\times A/10\ \mathrm{MeV}^{-1}$.

The second moment of the quantal distribution σ_z^2 versus excitation energy is given in Fig. 1. The narrowing of the distribution with increasing energy is quite evident. Since this calculation does not include thermal fluctuations, which correspond to the fluctuating part of the wave function, they are introduced in the simplest way,

$$\sigma_z^2=\sigma_{z,Q}^2+\sigma_{z,T}^2,$$

where the labels Q and T stand for quantal and thermal. The thermal width can be rigorously estimated by the same techniques as for the fluctuating cross sections in the statistical theory. It depends on the level densities only. The estimate which we gave for the thermal fluctuations corresponds to classical Boltzmann statistics.

The possibility of experimentally observing the minimum of σ_z^2 and its rapid rise with decreasing energy is of extreme importance because it would provide us with information on the damping of a giant resonance in a hot nucleus. This is particularly true in view of the extremely difficult alternatives, like γ decay from highly excited nuclei, etc.

The only experimental result shown in the figure is a heavy-ion example. Similar data are available in fission. Of course they do not prove our point. Until we can be assured that our guess for $\hbar\omega$ is a reasonable one (within a factor of 2), the experimental data should be considered as circumstantial evidence in favor of the present theory. Further theoretical work and experimental measurements at energies closer to the barrier will eventually tell the rest of the story.

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Origin of Gross Structure in $^{12}C+^{12}C$ and $^{16}O+^{16}O$ Inelastic Scattering

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A nonresonant diffraction-model calculation is found adequate to describe the gross-structure behavior thus far observed in the $^{12}C+^{12}C$ and $^{16}O+^{16}O$ inelastic scattering excitation functions.

The prominent gross structures observed[1-3] in the single and mutual inelastic excitation yield of the 2^+ first excited state of ^{12}C in $^{12}C+^{12}C$ collisions have been discussed in terms of carbon-carbon molecular resonances.[1,4] Similar gross-structure behavior has been found in the cross