Lawrence Berkeley Laboratory

UNIVERSITY OF CALIFORNIA

Presented at the International Conference on Heavy Ion Physics and Nuclear Physics, Catania, Italy, March 21-26, 1983

THE PHENOMENOLOGY OF DEEP-INELASTIC PROCESSES

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March 1983

Prepared for the U.S. Department of Energy under Contract DE-AC03-76SF00098

LBL-16372

THE PHENOMENOLOGY OF DEEP-INELASTIC PROCESSES+

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Abstract: The field of heavy-ion deep-inelastic reactions is reviewed with particular attention to the experimental picture. The most important degrees of freedom involved in the process are identified and illustrated with relevant experiments. Energy dissipation and mass transfer are discussed in terms of particles and/or phonons exchanged in the process. The equilibration of the fragment neutron-to-proton ratios is inspected for evidence of giant isovector resonances. The angular momentum effects are observed in the fragment angular distributions and the angular momentum transfer is inferred from the magnitude and alignment of the fragments spins. The possible sources of light particles accompanying the deep-inelastic reactions are discussed. The use of the sequentially emitted particles as angular momentum probes is illustrated. The significance and uses of a thermalized component emitted by the dinucleus is reviewed. The possible presence of Fermi jets in the prompt component is shown to be critical to the justification of the one-body theories.

Introduction

Let me begin by commenting upon the title of this talk. I did not choose it. Rather it was warmly "suggested" to me by the organizers and accompanied by a gentle but firm arm-twisting. So I shall make a vertue out of necessity, and shall give a view, if not a review, of the field of heavy-ion deep-inelastic reactions keeping my feet on the allegedly strong ground of experimental evidence. Therefore, no theory: this is of course a promise, not a fact. And if my "Phenomenology" does not turn out to be as firmly established and uncontroversial as that of Hegel, I hope that this will contribute to the livelihood of the conference and of the field of heavy ions.

As a general impression, I am pleasantly surprised in finding that the all too frequently haphazard and anecdotal nature of our own investigations can in fact be reconciled with a reasonably organic structure[1]). Historically, the early work was mainly dedicated to an exploration of the new macroscopic degrees of freedom made accessible by the dinuclear transient. These degrees of freedom were embrionically present in the fission process, but could not attain their full development due to our lack of control in the initial conditions. The ample choice of targets, projectiles and bombarding energies allowed us to study in some detail degrees of freedom such as relative distance, mass asymmetry, neutron-to-proton ratios, and a variety of angular-momentum-bearing modes. The picture that resulted from the early investigations was that of a system on its way to equilibrium with a variety of relaxation times, most of which were comparable with the nuclear interaction time and thus amenable to experimental investigation.

The extensive and pervasive presence of "conditional" thermalization, visible for instance in the relaxed component in the kinetic energy spectrum, in the ther-

* Alexander von Humboldt US senior scientist Award.

+ This work was supported in part by the Director, Office of Energy Research, Division of Nuclear Physics of the Office of High Energy and Nuclear Physics of the US Department of Energy under contract DE-AC03-76SF00098.

mal partition of the energy between the two fragments, and in the rigid rotation limit prompted the development of thermal models[2]) that were immediately effective in predicting first and second moments of a host of variables. The realization that the variance in a diffusive process is already almost at its equilibrium value after one relaxation time while the first moment may still be far away from equilibrium[3]) prompted the use of the equilibrium models to predict fluctuations for systems not completely equilibrated. The equilibrium statistical models are interesting because they do not require the knowledge of the reaction mechanism but they are satisfied with a rather sketchy Hamiltonian.

The observation of slowly evolving and broadening distributions suggested the introduction of diffusion models based upon the Master equation[4]) for the Fokker Planck equation[5]) in order to describe the final stages of approach to equilibrium. Most theories on the market at present are still based on such equations. The clarification process that has occurred since has been in the direction of specifying the reaction mechanism.

Much confusion exists at present in the theoretical field. Time-Dependent-Hartree-Fock claims the price for "microscopicity". Most likely its difficulties arise from "microscopic poverty" as the model is essentially prevented from using most of the available phase-space of which nature seems to make lavish use. Furthermore while some claim to similarity to experiment are being made, the very professionals seem at loss when requested to give an explanation as to "why" the model behaves as it does. With these limitations it is uncertain for some people whether the model "explains" anything even if it "reproduces" everything.

The fashionable exponents of the "new dynamics" are the one-body theories[6]) which seem to do an excellent job in fitting data. However, competing diffusion theories seem to do just as well[7]). It seems to me that, in most cases, when agreement with the data is reached, such an agreement is due to the common features of the two theories. For instance mass exchange does dissipate energy and transfer angular momentum. The remaining features, like long or short mean-free-path etc., characterizing and differentiating the models have little or no effect on the observables.

Similar difficulties arise when the particle transfer models[6]) and the collective vibration models[7]) are compared on the processes of energy and angular momentum dissipation. One has the impression that each model contains ingredients that tend to minimize the relevance of the opposing theory. So perhaps it is not a bad idea to stick to experiment.

Deep-inelastic processes should be viewed as a part of the physics of fusion and reseparation[9]). In this framework, with increasing angular momentum and/or effective fissionability, one moves from the mononuclear regime, consisting of fusion leading to compount nucleus, and fast fission, to the dinuclear regime consisting of deep-inelastic, quasi-elastic, and direct reactions. Such a broad picture is beyond the scope of this presentation, therefore I shall limit myself to the reseparation aspect of the process. We proceed with a presentation of the great chapters of this field, usually associated with specific degrees of freedom, taking the liberty to comment on the various interpretations of the data as it seems appropriate.

The energy dissipation

The dynamical inability of the system to coalesce and progress towards fusion is accompanied by a rather successful attempt to dispose of all of the entrance channel kinetic energy. The experimental picture is consistent with a complete dissipation of the radial energy and with a partial dissipation of the tangential energy[1]). A Wilczynski diagram gives at times a very vivid picture of the process of energy dissipation, hinting heavily towards a correlation between energy loss and impact parameter, the more deeply inelastic collisions being associated with smaller impact parameters. Nevertheless, large angular fluctuations, observed at any given energy loss, preclude an absolute correlation between energy dissipation and impact parameter. Normally a peak is observed in the Q-value spectrum at the largest energy losses. This peak appears from all evidence to be totally relaxed, corresponding to sub-Coulomb energies on one hand, and being associated with complete angular momentum relaxation or rigid rotation on the other[1]). The great tendency towards relaxation in these reactions can also be seen when an inquiry is made about the fate of the dissipated energy. For sufficiently low energies there is an overwhelming evidence that the dissipated energy reappears in form of fragment excitation energy. Even at higher bombarding energies where the presence of fast particles is undeniable, most of the energy ends up in the fragments. This has been demonstrated[10]) for instance in the study of the reaction $^{nat}Cu + {}^{20}Ne$ at bombarding energies of 158, 252 and 343 MeV. As shown in fig. 1, the overall missing charge, which varies dramatically with bombarding energy, is a linear function of the total energy loss, the slope corresponding to about 12.5 MeV/nucleon (assuming an equal number of neutrons being lost as protons). This slope is consistent with the overall missing energy being dumped in the two fragments.

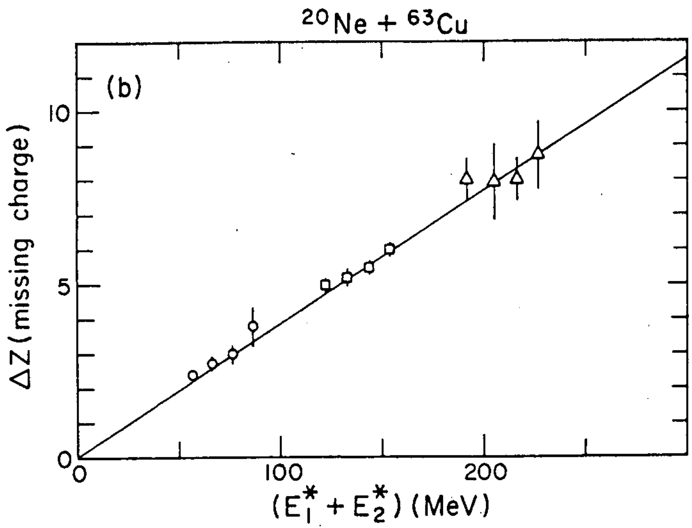

Fig. 1 Missing charge against excitation energy in the reaction $^{20}Ne + {}^{63}Cu$[10]).

The next problem to be solved is the partition of the dissipated energy between the two fragments. It is easy to verify that the thermal equilibrium between the two fragments requires an energy partition proportional to the fragment masses[11]). If the level densities of the two nuclei are ρ_1, ρ_2, then the probability of a partition of an amount of energy E into x and E-x is given by:

$$P(x) \propto \rho_1(x)\rho_2(E-x)$$

The maximum probability is given by:

$$\frac{d \ln P}{dx} = 0 = \frac{d \ln \rho_1}{dx} + \frac{d \ln \rho_2}{dx} = \frac{1}{T_1} - \frac{1}{T_2}$$

For a Fermi gas, $E = aT^2$ where $a \cong \frac{A}{8}$. Thus the result $\frac{x}{E-x} = \frac{A_1}{A_2}$.

Two techniques have been used to measure the energy partition. The first is a kinematic technique based upon the measurement of energy and angle of both fragments plus the atomic number of one fragment[12,13]). It is observed that the mean number of emitted neutrons by each fragment is proportional the fragment mass (or charge). The consistency of this result with an energy partition proportional to the masses can be verified more accurately by means of an evaporation code. Indeed, the results obtained in the deep-inelastic region are totally consistent with thermal equilibration between fragments. An example of this approach is given in fig. 2.

The second technique, which is the direct measurement of the neutrons[14-16]) emitted by each fragment, proved at one time the equal temperature of the two sources (from neutron spectra) and the thermal energy partition (from the neutron multiplicities), as shown in fig. 3. The most astounding result is the dependence of the neutron multiplicity ratio upon energy loss. The constancy of this ratio, as seen in fig. 3, seems to indicate that the two fragments reach thermal equilibrium even at the shortest interaction times. This poses a problem: since the mechanisms likely to contribute to energy dissipation tend to deposit equal energy in each fragment, how can the system rearrange (thermalize) the energy in such a short time? One can always speculate about the presence of a subdolous mechanism

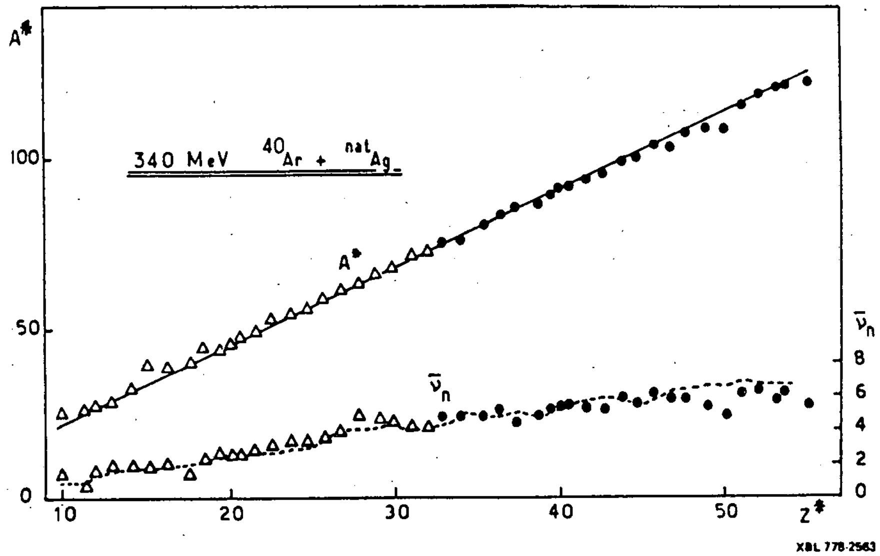

Fig. 2 Pre-evaporation masses A* (upper points) and average number of evaporated neutrons $\bar{\nu}_n$ (lower points). The dashed line is an evaporation calculation assuming an energy partition proportional to the masses[12]).

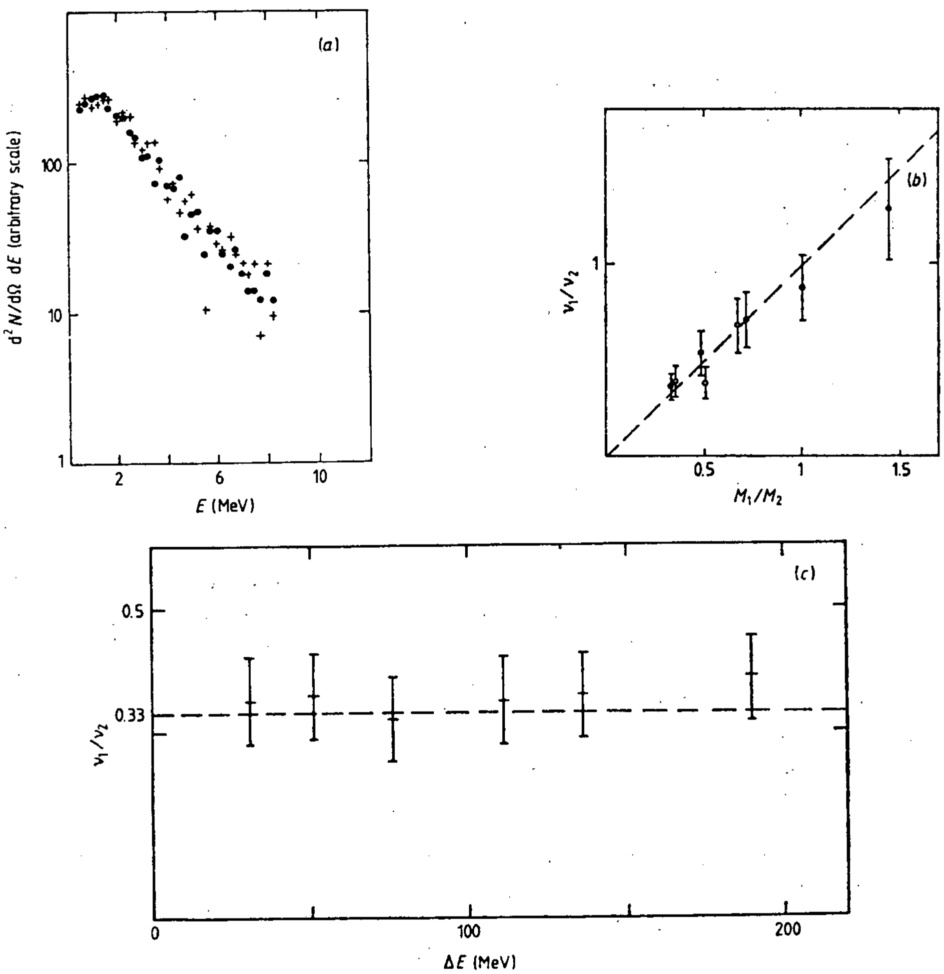

Fig. 3 Neutron emission in the reaction 400 MeV Cu + Au[14]). Top left: neutron kinetic energy spectra from the two fragments; top right: neutron multiplicity ratio as a function of mass asymmetry; bottom: neutron multiplicity ratio at fixed mass asymmetry (dashed line) as a function of energy loss.

able to partition the energy directly in the ratio of the masses. On the other hand one can check on the variance of the energy partition and verify whether it is in agreement with the thermal limit[17]. The variance is given by:

$$\sigma_E^2 = 2 \frac{a_1 a_2}{a_1 + a_2} T^3$$

Unfortunately, there is no direct measurement of this quantity as yet. Some indirect evidence can be obtained from the high energy tail present in the proton spectra associated with the deep-inelastic reaction 252 MeV Ne+Cu[18]. The energy fluctuation between the two fragments in contact at constant temperature, can create strong fluctuations in the temperature of a Ne-like fragment after separation, thus allowing for high energy tails in the particle evaporation spectra. In fact the use of thermal energy partition plus energy fluctuations is sufficient to reproduce the experimental spectra, as shown in fig. 4.

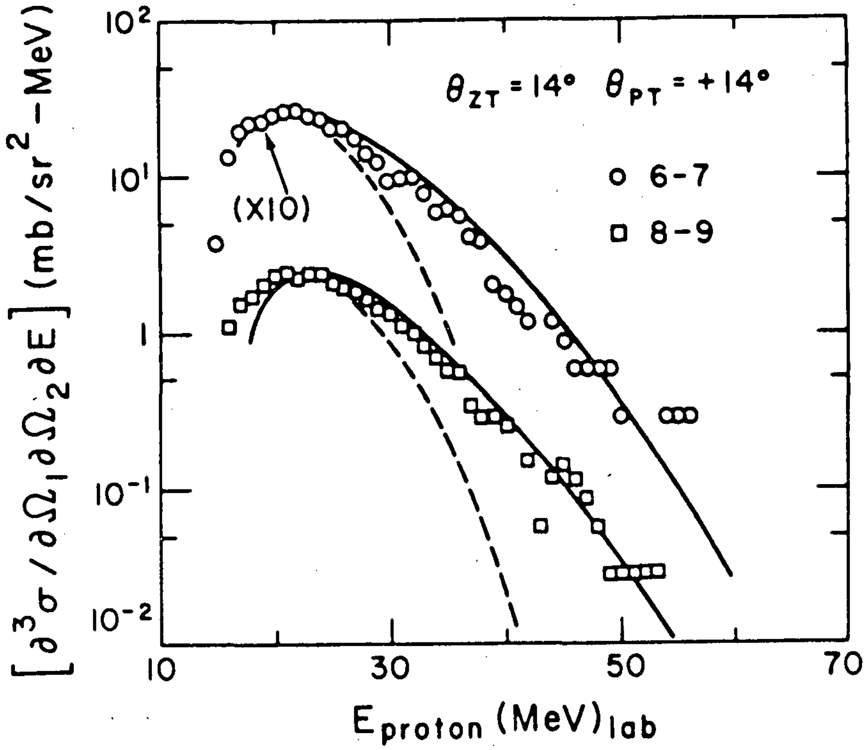

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Fig. 4 Proton kinetic energy spectra in coincidence with a light fragment in the reaction Cu + ^{20}Ne at 252 MeV bombarding energy[18]. The dashed lines are calculations performed using a fixed energy sharing proportional to the masses, while the solid lines allow for thermal fluctuations in the energy partition.

While waiting for experiments to give a direct answer to the question of energy fluctuations, let us consider the mechanisms that are likely to preside to the energy transfer process. The mass transfer is an obvious and unavoidable mean of energy dissipation. There is good evidence that nucleon transfer can account for a large fraction if not for the total energy dissipation in heavy systems. Incoherent particle-hole excitation produced by the diabatic single particle motion in a time-dependent mean field is another possibility. Finally the excitation of low and high-frequency vibrations in the two fragments is also able in principle to dissipate large amounts of energy[8]. This last possibility has been invoked to explain structures in the Q-value of some reactions[19]. Whether such structures can be identified with the excitation of giant oscillations in the two fragments is not clear at the moment.

The mass distribution

Despite the fact that the mass asymmetry degree of freedom was one of the first to be treated in the framework of a diffusion theory, Nature seems to have taken an Olympic disregard towards at least one of the most direct predictions of this theory, like the mass drift.

For reasons which are still not understood, the population of the mass asymmetry coordinate appears to be extremely simple. On one hand a projectile and a target-like peaks are observed, which broaden progressively with increasing energy loss. Only at the largest energy losses does one (if at all) observe a drift in the direction prescribed by the potential energy. This deep-inelastic component is readily identified by its angular distribution which is either forward-peaked or side-peaked. On the other hand a broad, fission-like component is observed near symmetry. Such a component has fission-like kinetic energies and an angular distribution symmetric about 90°. Only its inconsistence with the compound nucleus predictions allows one to distinguish it from true fission. These components are readily seen in fig. 5 and the energy dependence of the mass distribution is shown in fig. 6.

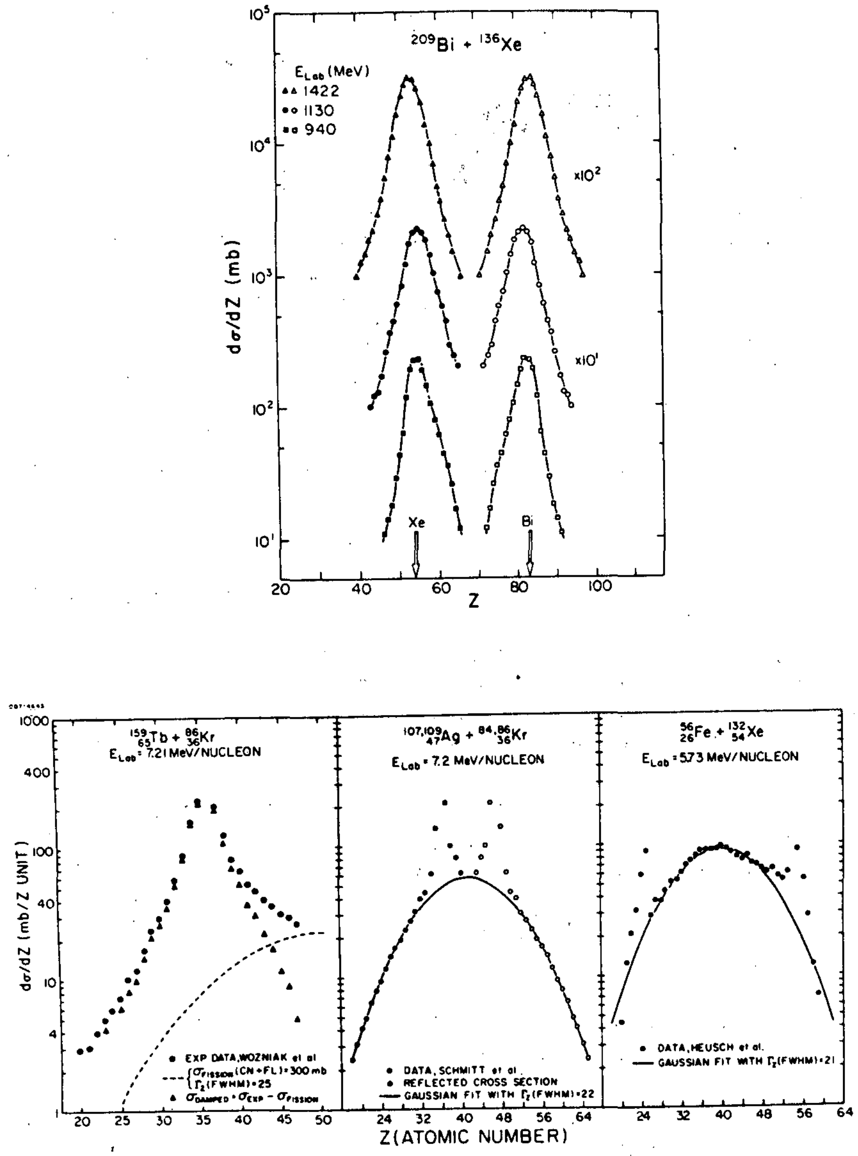

Fig. 5 A compilation of charge distributions for a series of deep-inelastic reactions[20]).

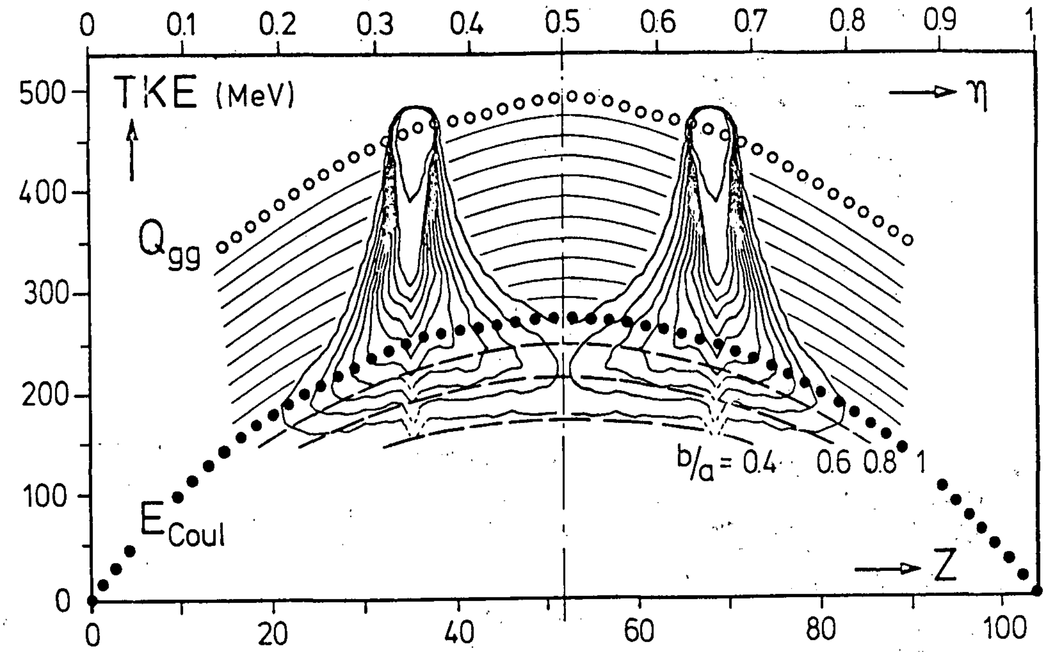

Fig. 6 Dependence of the charge distribution upon energy loss for the reaction Kr + Er[21]).

The explanation of these two components is not easy. Swiatecki in his model of coalescense and reseparation[9]) suggests that a dinuclear regime is prevailing for the deep inelastic processes while fast fission, the symmetric component is associated with the formation of a more compact shape, (mononucleus) which having grown a fat neck, but being otherwise unable to fuse, can easily rearrange its mass asymmetry (not necessarily by particle exchange, but also by a simple shape evolution) and decay out along the fission valley.

Early analyses of the dependence of the mass width upon angle and energy loss demonstrated the apparent diffusive nature of the time-evolution of the mass asymmetry degree of freedom and established a remarkably long relaxation time[1,4,5]). However, the frustrating unwillingness of the system to follow the potential energy casts an additional question-mark on a problem whose complexity has just begun to untangle.

Correlations between energy dissipation and mass distribution

The hypothesis of the energy dissipation being mediated by particle exchange as postulated in one-body theories can be put to test by observing the correlation between the energy loss and the variance of the mass distribution at different Q-values[6]). The energy loss associated with a mass exchange dm is

$$dE = E\frac{dm}{\mu} \quad \text{or} \quad \frac{dE}{dt} = E\,\frac{\Phi(t)}{\mu} \quad \text{and} \quad E = E_0\, e^{-\frac{\Phi\cdot t}{\mu}} \quad (\text{if } \Phi = \text{constant})$$

In these expressions Φ is the mass flux. These equations do not contain the mass of the particle that has been exchanged. On the other hand, from random walk one has:

$$\sigma^2 = Nm\cdot m = \Phi tm$$

where N is the number of exchanges and m is the mass of the particle. Thus:

$$E = E_0 \exp - \frac{\sigma^2}{\mu m}$$

A classical calculation along these lines substantially underpredicts the energy loss for any given value of σ^2 if nucleons are assumed to be the exchanged partic-

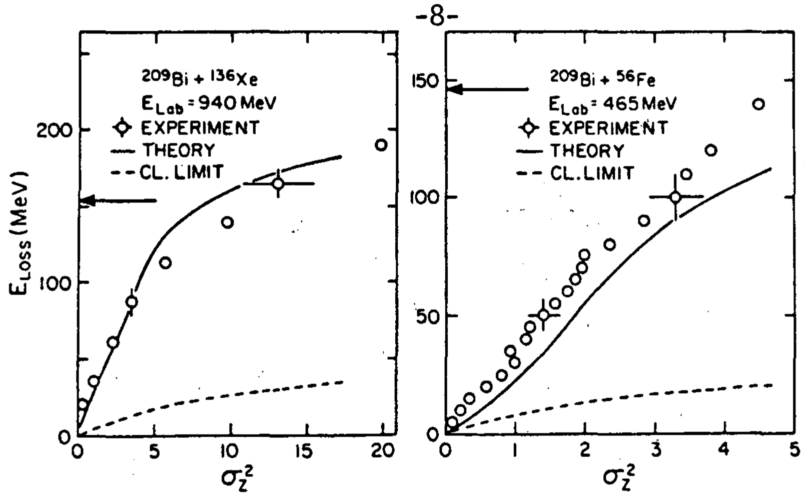

Fig. 7 Correlation between energy loss and charge variance for two reactions [22]. The dashed line and the solid line correspond to calculations without and with the Pauli principle respectively in terms of a one-body theory.

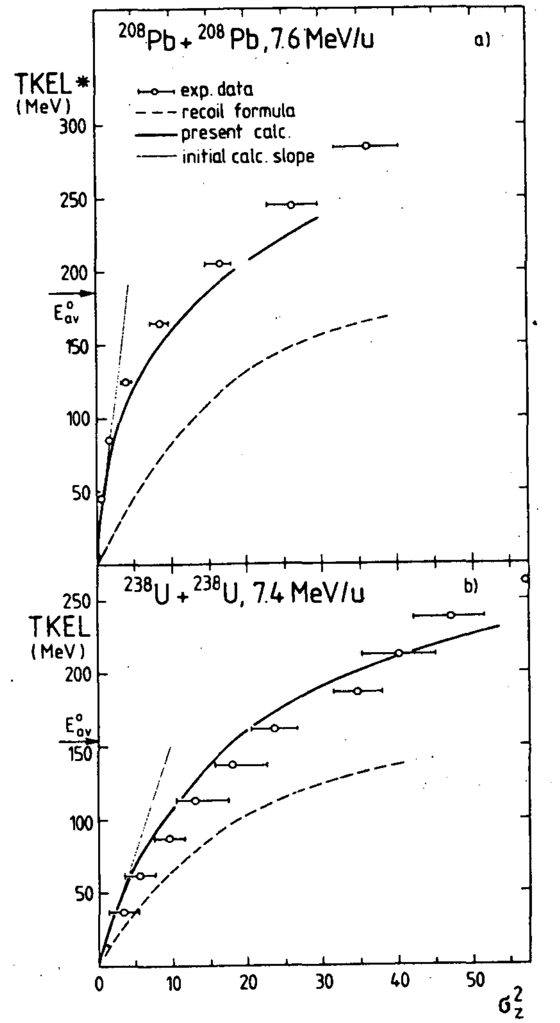

Fig. 8 Correlation between energy loss and charge variance for two reactions. The dashed line does not incorporate the Pauli principle, while the solid line does. The model assumes that particle transfer as well as particle-hole excitations are responsible for energy dissipation[7].

les[22]. The inclusion of the Pauli blocking alters σ^2 but not the energy loss, as required by the fact that the former quantity is the expectation value of a two-body operator while the latter is the expectation value of a one-body operator[6]. This inclusion seems to reproduce the data satisfactorily[22]. This agreement, illustrated in fig. 7, seems to prove that the particles exchanged are nucleons and may manifest one of the very rare macroscopic quantal effects in the field of deep-inelastic processes.

This nice picture is obscured somewhat by the fact that other theories which assume two simultaneous energy transfer mechanisms, namely particle transfer and particle-hole excitation[7]) seem to reproduce the experimental data equally well, as shown in fig. 8. In these theories as much as 50% of the energy is transferred by means of particle-hole excitation.

Also it is definitely worth considering that a transfer of a mass 4 object (alpha particle) seems to be able to reproduce the data just as well. This should certainly draw the attention of the experts in direct reactions who are not unfamiliar with alpha particle transfer.

It may be worth pointing out, at this particular stage, that the successful reproduction of the energy loss-mass variance correlation by the so-called one-body theories does not prove the validity of the new dynamics. The long mean-free-path is not required, rather the mean-free-path needs be only longer than the window region in order to prevent collisions within the window, leading to a particle coming back into the nucleus it started from. What the experiment seems to prove is that: a) there is mass transfer; b) that the mass transfer can be associated with nucleon exchange; c) that the mass exchange (not necessarily nucleons) accounts for a good fraction of the dissipated energy.

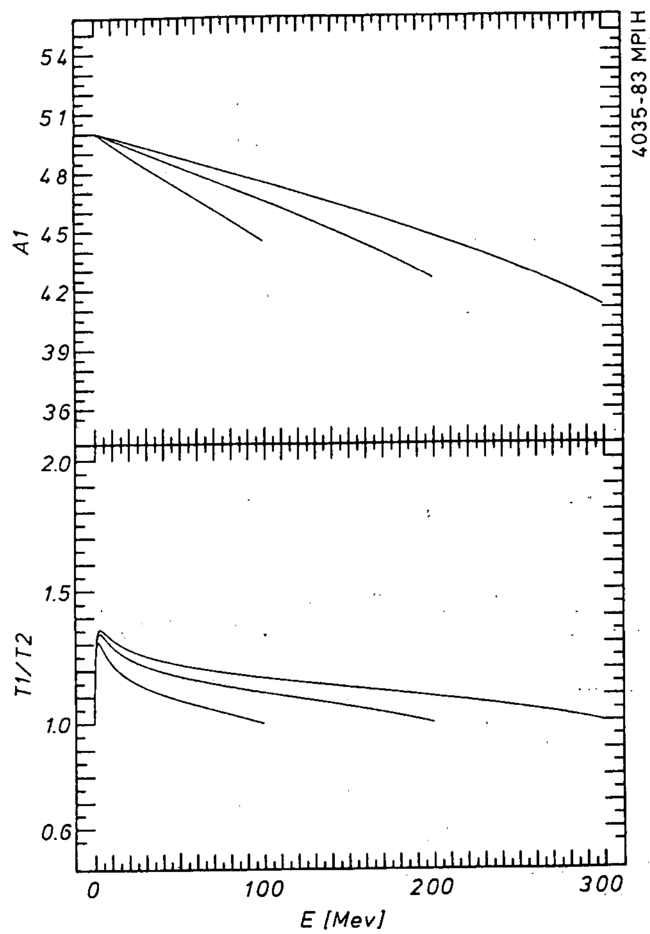

Fig. 9 Mass of the light fragment vs. energy loss for three different center of mass energies (300, 200, 100 MeV) for a system with initial masses 50, 150 (top). Light fragment/heavy fragment temperature ratio vs. energy loss for the same center of mass energies (bottom).

A second point worth considering at this stage is the apparently unrelated combination of two experimental facts noted above: 1) the lack of drift towards symmetry of the centroid of the mass distribution; and 2) the thermal equilibrium between fragments as shown by the energy partition. It may be possible to explain these two features in terms of the particle exchange picture as follows. At the beginning of the collision the energy deposition is equal in both fragments. In an asymmetric system the light fragment will grow hotter. If, as it is the case during the relevant part of the collision, the particle fluxes are strongly temperature-dependent, then the particle flux from light to heavy fragment will be greater than the flux from heavy to light fragment. Since the energy deposition into a given fragment is proportional to the number of particles landing into it, it follows that, due to the temperature gradient, the energy deposition will be redirected in favor of the cooler heavy fragment. This is a fast feedback process because it acts not on the energy already deposited but on the energy being deposited. In this way the light fragment tries to minimize the temperature gradient, but it does so at the expense of its mass! In other words a dynamical drift is created towards greater asymmetries. This drift is generated by the relative motion and it may be quite effective in counteracting the mass drift towards symmetry dictated by the potential energy.

In fig. 9 the mass of the light fragment is shown vs. Q-value for three bombarding energies. The initial mass ratio is 50:150. The same figure shows the temperature ratio. It can be seen that when no conservative force is acting on the mass asymmetry, the system tends to grow more asymmetric. The feedback process is also seen to control and contain the initial temperature gradient. In fig. 10 the light fragment mass as a function of Q-value is calculated for progressively stronger forces driving the system towards symmetry. The minor initial drift towards symmetry is rapidly controlled so that only after most of the energy is lost the system can drift again towards symmetry. In fig. 11 some recent pre-evaporation mass measurements[23]) as a function of energy loss show a pattern not unlike those observed in fig. 10.

Equilibration of the neutron-to-proton ratio

It is experimentally observed[1]) that, when two nuclei with different neutron-to-proton ratio are allowed to interact, they tend to evolve towards a conditional equilibrium given by the equation:

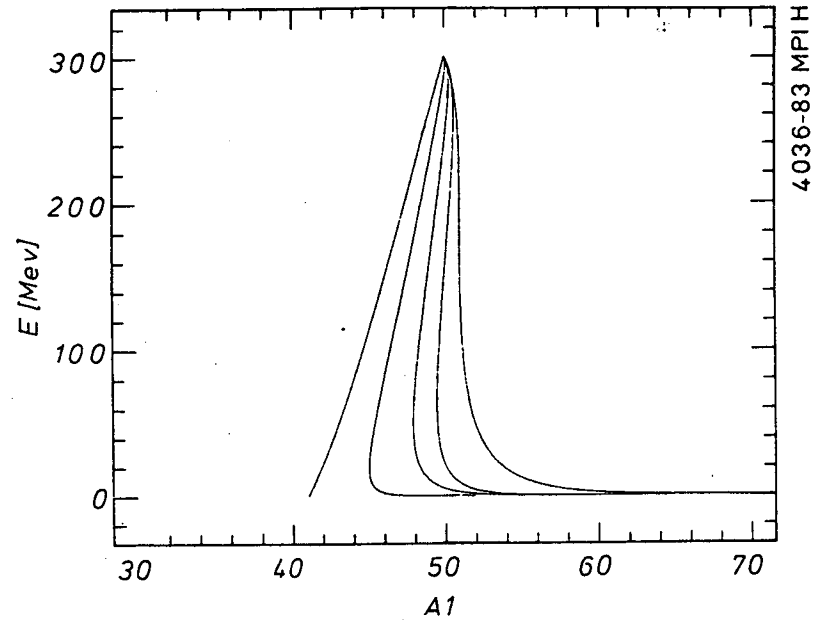

Fig. 10 Light fragment mass vs. energy loss for different driving forces. The driving force increases from left to right in steps of 0.5 MeV/amu.

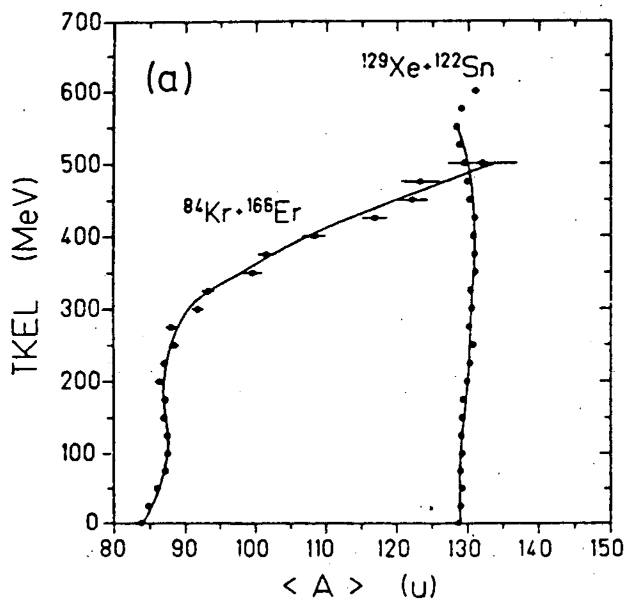

Fig. 11 Pre-evaporation masses as a function of energy loss for two reactions[23]).

$$\left.\frac{\partial E(A,A_1,Z_1)}{\partial Z_1}\right|_{A_1} = 0$$

where $E(A,A_1,Z_1)$ is the total energy of a dinucleus of total mass A, with one of the two fragments with mass A_1 and charge Z_1. In other words it appears that the system likes to move towards the conditional energy minimum at fixed mass asymmetry. This result comes naturally out of any diffusion calculation performed in the N-Z plane because of the steep and narrow valley represented by the function $E = E(A,A_1,Z_1)$ at fixed A_1. Along the same line, one would expect that the variances be reproduced by the diffusion models. However, one open problem appears to exist still on this subject. This is the problem of the variance σ_Z^2 (constant A). Some experiments[24]) have produced data whose Q-value dependence is completely consistent with diffusion models, or, even better with the equilibrium thermal model,

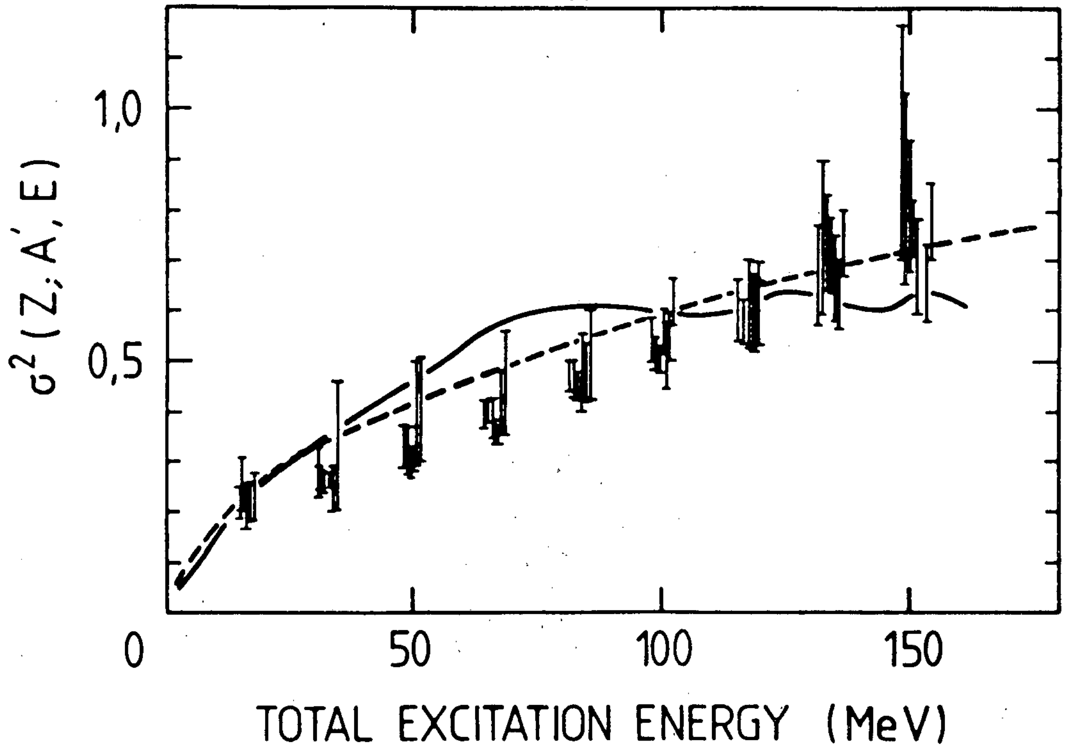

Fig. 12 Isobaric charge distribution variances as a function of excitation energy for the system ^{132}Xe + ^{197}Au[24]).

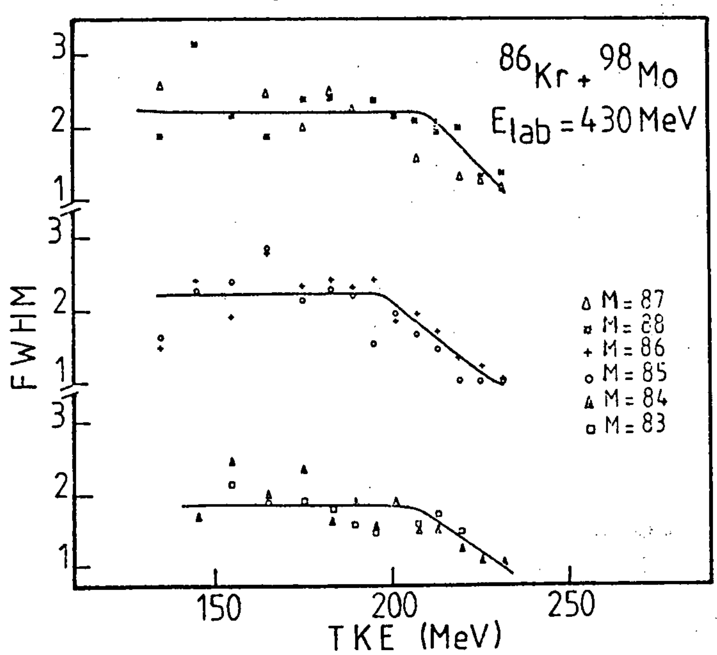

Fig. 13 Width of the isobaric charge distribution as a function of exit channel kinetic energy[25]).

namely

$$\sigma_Z^2 \text{ (constant A)} = \frac{T}{c}$$

where $c = \left.\frac{\partial^2 E}{\partial Z_1^2}\right|_{\tilde{Z}_1}$, showing that the system has reached a conditional thermal equilibrium. In fig. 12 an excellent example of this situation is shown.

On the other hand, other data (shown in fig. 13) give a completely different picture[25]). At small energy losses the variance increases rapidly to values substantially exceeding the thermal value and saturates for the rest of the Q-value region. A natural explanation has been given in terms of a collective isovector oscillation[1]) leading to a periodic enrichment and depletion of charge in each of the fragments at fixed mass asymmetry. The lowest multipole associated with such a vibration is of course E1. It is tempting then to suggest that the zero-point motion along such a degree of freedom is responsible for the fluctuation in Z. If the phonon $\hbar\omega >> T$ then $\sigma^2 = \hbar\omega/2c$. In this way one explains both the large value of σ^2 and its lack of Q-value dependence. The difficulty with this explanation is twofold: on one hand it is not obvious how to calculate $\hbar\omega$; on the other one does not know how to account for higher multipole isovector modes that may contribute to the same effect[26]). Be as it may, the difficutly at this stage appears to be mainly experimental.

The angular distributions

The angular distributions observed in deep-inelastic processes vary from side-peaked to forward peaked, frequently in the same reaction, as one considers greater energy dissipations and greater distances in mass or charge from the projectile. This great variation in pattern must correspond to an equivalent variation in the interaction time on one hand, and in the width of the orbital angular momentum distribution on the other. In particular the behavior of the angular distribution with mass may carry an important and yet untapped information on the differential spreading of the various l-waves along the mass asymmetry coordinate (angular momentum fractionation).

Perhaps one of the most puzzling features of the side-peaked angular distributions is their rapid, almost instantaneous spread in angle even for small energy losses. This feature is difficult to explain in term of certain thermal diffusion models since a sizeable amount of energy loss is needed by them in order to generate a sufficient temperature and thus a sizeable angular momentum width. However, it is clear, that an extremely narrow l-window must be accompanied by a large diffraction width in angle: $\Delta\theta=1/(2\Delta l)$. To be sure, the incorporation of such a quantal effect seems to fix-up the problem quite neatly[27]), as it is beautifully shown in fig. 14. However, calculations in terms of one-body models do not seem to need any such quantal fix[28]): they just reproduce the Wilczynski plot as they are, as shown in fig. 15. Even more interesting is the fact that the incorporation of quantal effects does not make such distributions wider, but it actually bearly changes them[29]). In my opinion the explanation is obvious. One-body theories explicitly take advantage of the mis-match of the Fermi spheres due to relative motion to drive diffusion, while other theories need to wait for a temperature build up.

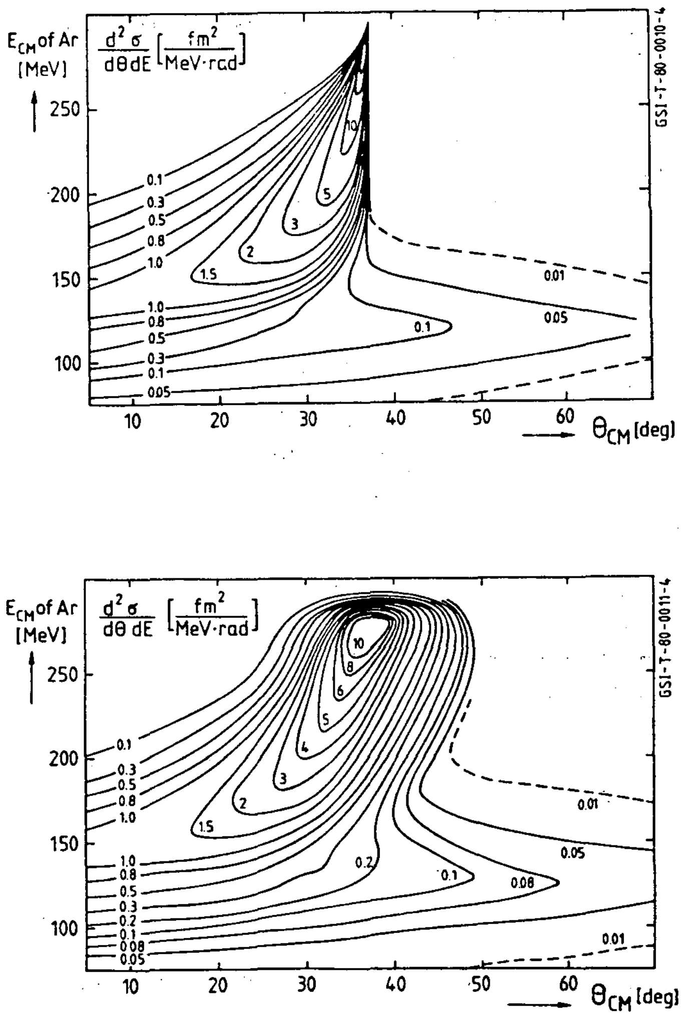

Fig. 14 Wilczynski plot calculated without (top) and with (bottom) quantal fluctuations[27]).

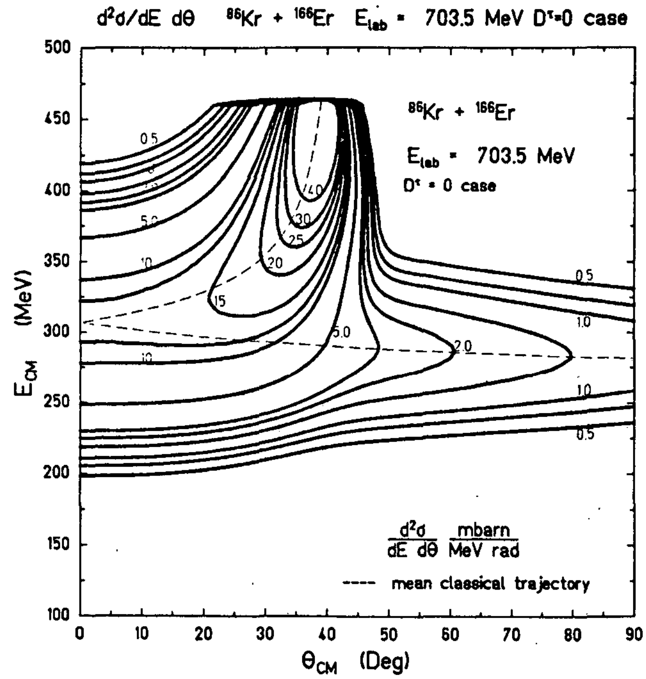

Fig. 15 Wilczynski plot calculated from a one-body theory and without quantal fluctuations[28].

Consequently the l-window σ_l available to one-body theories is quite large even at the beginning of the collisions. As a consequence we have[30])

$$\sigma_\theta^2 = \frac{1}{4\sigma_l^2} + \left(\frac{d\theta}{dl}\right)^2 \sigma_l^2$$

The first term is due to diffraction, the second is the spread in θ due to a classical spread in l. One-body theories have a fairly large σ_l since the beginning of the collision: thus the diffraction term is small and the fluctuation in θ is dominated by the second term.

The above equation is a very useful one because in principle one can use it to extract the fluctuations in orbital angular momentum and compare them with theory, for instance with the equilibrium thermal fluctuations[2]). The quantity σ_l^2 can be considered as

$$\sigma_l^2 = \sigma_l^{E^2} + \sigma_l^{F^2}$$

where the first term is the l-window (entrance channel) associated with the energy bin and the second is the exit channel fluctuation arising from the trade-in between orbital and intrinsic angular momenta. By using the empirical deflection function one can obtain σ_l^2 from experiment for energy bins of various magnitude. Then the limit:

$$\lim_{\Delta E \to 0} \sigma_l^2 = \sigma_l^{F^2}$$

allows one to estimate the fluctuations (thermal?) that can be inferred in a less direct way from the fluctuations in magnitude and orientation of the fragment spin[34]).

Angular momentum transfer

The angular momentum transferred to the fragments spins can be measured in various ways. The measurement of the gamma ray multiplicity gives the sum of the spins of the two fragments, while by measuring the angular distribution of sequentially emitted alpha particles or fission fragments one can obtain the spin of an individual fragment[1]).

The dependence of the transferred spin upon Q-value is characterized by an early rapid rise with increasing energy loss followed by a saturation or perhaps even by a slight decrease at the deepest inelasticities. The rising part of the curve seems to portray the progressive action of frictional forces towards the relaxation of the angular momentum. The question then arises whether complete relaxation for the angular momentum (rigid rotation) even occurs. An early test of rigid rotation showed the predicted rise of the transferred spin with increasing mass asymmetry[31]). More recent measurements of the spin of an individual fragment vs. charge[32]) proved the limit of rigid rotation to a good and convincing degree (fig. 16). For many systems however the expected rise of the total spin (gamma ray multiplicity) with increasing mass asymmetry fails to materialize[1]). The standard explanation is based upon the concept of angular momentum fractionation. If a broad l-window is available for the process one should <u>not</u> in principle expect that each l-wave spreads uniformely along the mass asymmetry coordinate. Rather, the high l-waves associated with a short interaction time will spread only a little around the entrance channel mass asymmetry, while the low l-waves being associated with longer interaction times can actually spread much farther and populate asymmetries far from the entrance channel value. As a consequence the average l associated with a given mass asymmetry decreases as one moves farther away from the initial asymmetry[1]). Assuming rigid rotation to transform the average l-value into the sum of the spins one finds that the decrease in l with increasing mass asymmetry approximately compensates the rising ratio of the total angular momentum going into spin, thus producing an approximately flat dependence of the gamma ray multiplicity upon mass asymmetry. One should appreciate that rigid rotation is not disproved by such experiments in favor, say, of rolling. Rather the experiment is explained in terms of rigid rotation plus angular momentum fractionation.

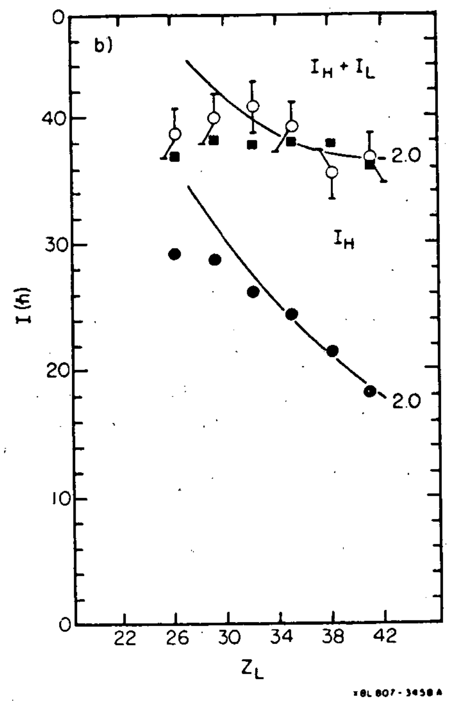

Fig. 16 Spin of the heavy fragment as a function of the atomic number of the coincident fragment determined from alpha particle angular distributions (dots). The sum of the spins inferred from the previous measurement is also shown (squares) as well as that determined from γ-ray multiplicities (open circles). The lines correspond to the rigid rotation limit for two touching spheroids with the ratio of axes equal to 2:1[32]).

Spin alignment and polarization

The angular momentum transferred from orbital to intrinsic motion should approximately retain its orientation perpendicular to the reaction plane. Thus the expectation of a strong polarization and alignment of the fragments spins. Strong spin polarization measured from the circular polarization of the emitted γ-rays has been observed[33]), varying from ∿-100% to ∿+100%, associated with the quasi-elastic and deep-inelastic component in the reaction ^{58}Ni + ^{16}O. This inversion in the spin polarization has been used to confirm the near-side scattering in the quasi-elastic region and the far-side scattering in the deep-inelastic region.

Relatively strong alignment has been observed from the out-of-plane angular distributions of alpha and sequential fission fragments. However, there is strong evidence for some angular momentum misalignment which of course arises from the presence of in-plane components of the fragment angular momentum. The origin of these components may be found either in the angular momentum transferring process, or in the thermal excitation of the angular momentum-bearing modes of the dinucleus[2]). The latter possibility leads to a simple result.

The angular momentum components of one of the fragments can be given in terms of a probability distribution[34]):

$$P(I_x,I_y,I_z) \propto \exp - \{\frac{I_x^2}{2\sigma_x^2} + \frac{I_y^2}{2\sigma_y^2} + \frac{(I_z-\langle I_z\rangle)^2}{2\sigma_z^2}\}$$

The cartesian reference system is chosen in such a way, that the y-axis coincides with the line between centers and the z-axis is parallel to the total angular momentum of the system. The quantity $\langle I_z\rangle$ is the expected rigid rotation contribution to the spin of one fragment. For two touching equal spheres the variances are approximately equal and have the following values[2]):

$$\sigma_x^2 = \frac{6}{5}\mathcal{I}T;\ \sigma_y^2 = \frac{6}{7}\mathcal{I}T,\ \sigma_z^2 = \frac{6}{7}\mathcal{I}T$$

where $\mathcal{I}$ is the moment of inertia of one of the two spheres and T is the temperature.

A particularly thorough study of the spin alignment has been made by studying the angular distribution of the continuum γ-ray spectrum associated with the fragment de-excitation. If the fragments are good rotational nuclei, the angular momentum is removed mainly by stretched quadrupole gamma ray transitions. For a completely aligned system, the angular distribution is $W_2(\theta) = 5/4(1-\cos^4\theta)$. By folding this distribution with the distribution of the angular momentum components one can obtain expressions for the γ-ray angular distributions that depend on the degree of alignment of the spin[34]).

The study of the angular distribution of the quadrupole component has provided us with information on the alignment of the fragments spins[35]). As an example let us consider the reactions: 1400 MeV ^{165}Ho + (^{176}Yb, ^{148}Sn, ^{nat}Ag)[36]). The fragments spins obtained from the multiplicities are given on fig. 17 as a func-

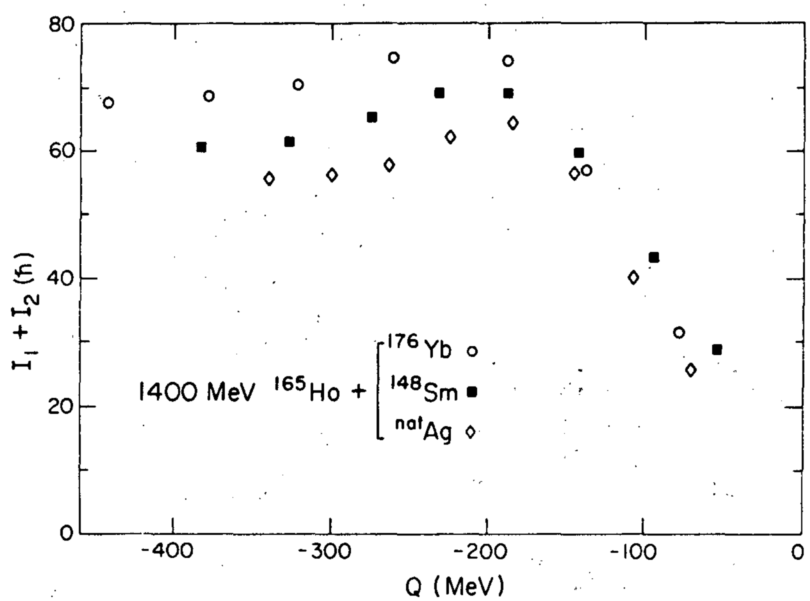

Fig. 17 Sum of the fragments spins vs. Q-value for three reactions[36]).

tion of Q-value. The dependence of the anisotropy on Q-value is given on fig. 18. A most remarkable picture of a rising and falling anisotropy is seen as the Q-value increases. The explanation in terms of the statistical model is quite straightforward. At small Q-values, spin is rapidly fed to the fragments. This angular momentum is aligned. The small Q-value implies small temperatures and thus little excitation of the angular momentum-bearing modes. Therefore there are small in plane components and the spin becomes rapidly aligned. When the Q-value exceeds a critical value, the fragments angular momentum does not increase any longer, actually the z-component decreases, while the temperature and thus the fluctuating angular momentum components keep increasing. As a consequence the system becomes more and more misaligned and the anisotropy falls. The figure shows also the anisotropies calculated on the basis of the statistical model, including the contribution of dipole transitions and the effect of neutron evaporation on the alignment of the system.

Light particles accompanying heavy ion collisions

The great majority of the light particles emitted in deep-inelastic collisions originates from the fully accelerated fragments as evaporation products. As

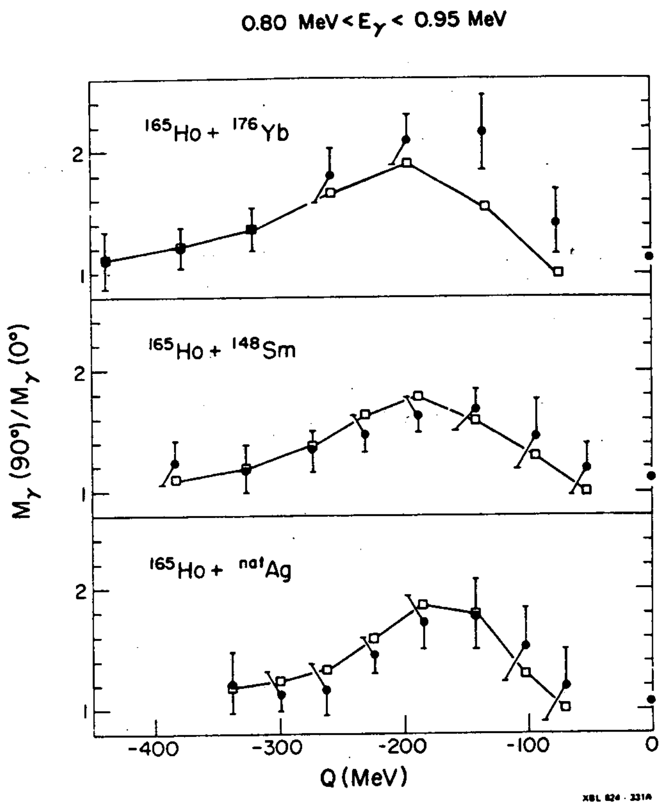

Fig. 18 Gamma-ray anisotropies vs. Q-value for three reactions. The open squares connected with a line represent a calculation based upon the statistical model for the angular momentum misalignment[36].

such they have been used to infer various properties associated with the collision itself[1]). For instance, the emitted neutrons have been used to determine the partition of the dissipated energy between the two fragments. Similarily the emitted alpha particles have been used to determine the angular momentum deposition through their angular distributions and gamma-rays have provided information about the magnitude and alignment of the transferred angular momentum.

In the same class of processes one should place sequencial fission. Being an ordinary compound nucleus fission of an excited deep-inelastic fragment, it has been put to good use in determining the fragment spin and its alignment[1]).

Light fragments apparently being thermally emitted by the dinuclear system have also been reported[37]). Such emission is of the utmost interest. On one hand it demonstrates the possibility of "thermal" evaporation from a system which is not equilibrated in its collective modes, like evaporation from running water, on the other hand it may contain information on the shape and on the collective dynamics of an extremely short-lived system.

With progressively increasing bombarding energy, non-equilibrium components begin appearing. A great deal of hope has been placed in these non-equilibrium components because they may contain information regarding the primary dissipation processes. Two opposing theories, the "Fermi jets"[38,39]) and the "hot spot" theories[40]) have been proposed to explain such an emission. The first theory is based upon the assumption of long nucleonic mean-free-path. The vectorial combination of the relative velocity with the nucleonic (Fermi) velocity may lead to a situation whereby fast nucleons are emitted opposite to the collision point. The second theory is based upon the assumption of short mean-free-path leading to local heating on the surface of the nuclei followed by evaporation. General theoretical considerations favor the "Fermi jet" assumption of long mean-free-path. Unfortunately, the experiment does not seem to favor either theory. In fact the picture of prompt particle emission is rather disheartening in its complexity. Nevertheless it may very well be from this kind of evidence that the most important information on the microscopic nature of dissipation may be retrieved.

Experimental evidence of the "new" one-body dynamics

Swiatecki[9]) has proposed that, as one can understand a great deal of the static properties of nuclei in terms of a one-body Hamiltonian (shell model), so one could also understand the dynamical nuclear properties in terms of one-body features, like one-body viscosity as portrayed by the wall formula and, to a lesser extend by the window formula. Such an approach has been used to describe the physics of coalescence and reseparation of two colliding nuclei, with some success. The truly novel aspect of this approach is the one-body-friction, while the potential energy aspect is based on the traditional liquid drop model. It is interesting to critically analyze the experimental data to see whether there is any evidence for this one-body picture which arises from the long nucleon mean-free-path.

The good success of the one-body theory in reproducing energy, angular momentum, particle, and charge transfer[6,28]) seems at first sight a strong support of the new dynamics. However, in my opinion, this is not the case yet. A great deal of the success of these models is due to the fact that they assume nucleon transfer. This automatically defines the energy and angular momentum transfer. Yet particle transfer itself does not imply long mean-free-path, but a mean-free-path somewhat larger than the window "thickness". This is especially true of the "window formula" whose derivation does not require anywhere the specific condition of a long mean-free-path. As a counterproof one can present other theories which are not "one-body" theories, which allow, of course, for particle transfer, and which seem to do a reasonably good job in fitting data[7]).

This is by no means a criticism of the new dynamics. It is a statement of the simple fact that the above mentioned observables depend only weakly upon the nuclear mean-free-path.

Is there some other observable that is specifically sensitive to the long mean-free-path? The answer is yes. The theory of Fermi jets or Pep jets specifically uses the concept of the long mean-free-path which allows a nucleon, with a velocity equal to the Fermi velocity plus relative velocity, to cross the dinuclear system and escape to the other side. The energy and angular distribution of these nucleons is predicted to be very characteristic. The observation of Fermi jets would be a practically unequivocal proof of the long mean-free-path and of the new dynamics. Unfortunately, the various searches for such jets have not been successful. While the "prompt" non-thermal emission of particles is well established at sufficiently high bombarding energy, it does not conform to the Fermi jet picture. Still it is obvious that in the observed great confusion of prompt light particle emission may lie the key to the microscopic make-up of the macroscopic dynamics.

Conclusion

There are several lessons that can be learned from this review. Perhaps the most important is the fact that, while there are many experiments and many theories, a serious attempt to put the two together with a critical analysis is still lacking. In a way it is perhaps unfortunate that the great expanse of experiments gives room to any theory to find a comfortable spot. On the other hand the richness of the phenomena encompassed by the field of heavy-ion reactions should be taken as a challenge for more pointed experimental work and more critical theoretical development.

Acknowledgements

It is a pleasure to thank prof. H. A. Weidenmüller for his kind hospitality at the Max-Planck-Institute. This work was supported in part by the Dir., Off. of Energy Research, Div. Nucl. Physics of the Off. of High Energy and Nucl. Physics of the U.S. DOE under contract DE-AC03-76SF00098.

References

1) The interested reader can benefit from the following review articles:
W. U. Schröder, J. R. Huizenga, Ann. Rev. Nucl. Sci. 27 (1977) 465;
M. Lefort and C. Ngô, Ann. Phys. (NY) 3 (1978) 5;
L. G. Moretto and R. P. Schmitt, Rep. Prog. Phys. 44 (1981) 543.
2) L. G. Moretto and R. P. Schmitt, Phys. Rev. C21 (1980) 204.
3) L. G. Moretto and L. G. Sobotka, Z. Phys. A303 (1981) 299.
4) L. G. Moretto and J. S. Sventek, Phys. Lett. 58B (1975) 26.
5) W. Nörenberg, Phys. Lett. 52B (1974) 289.
6) J. Randrup, Nucl. Phys. A307 (1978) 319; ibid A327 (1979) 490.

7) M. Dakowski, A. Gobbi and W. Nörenberg, Nucl. Phys. A378 (1982) 189, and references therein.
8) R. A. Broglia, C. M. Dasso and A. Winther, Phys. Lett. 53B (1974) 301; ibid 61B (1976) 113.
9) W. J. Swiatecki, Nucl. Phys. A376 (1982) 275.
10) R. P. Schmitt, G. Bizard, G. J. Wozniak and L. G. Moretto, Phys. Rev. Lett. 41 (1978) 1152.
11) D. J. Morrissey and L. G. Moretto, Phys. Rev. C23 (1981) 1835.
12) B. Cauvin, R. C. Jared, P. Russo, R. P. Schmitt, R. Babinet and L. G. Moretto, Nucl. Phys. A301 (1978) 511.
13) R. Babinet, B. Cauvin, J. Girard, H. Nifenecker, G. Gatty, D. Guerreau, M. Lefort and X. Tarrago, Nucl. Phys. A296 (1978) 160.
14) B. Tamain, R. Chechik, H. Fuchs, F. Hanappe, M. Morjean, C. Ngô, J. Pēter, Nucl. Phys. A330 (1979) 253.
15) Y. Eyal, A. Gavron, I. Tserruya, Z. Fraenkel, Y. Isen, S. Wald, R. Bass, J. R. Gould, G. Kreyling, R. Renfort, K. Stelzer, R. Zitzmann, A. Gobbi, U. Lynen, H. Stelzer, I. Rode, R. Bock, Phys. Rev. Lett. 41 (1978) 625.
16) D. Hilscher, J. R. Birkelund, A. D. Hoover, W. U. Schröder, W. W. Wilke, J. R. Huizenga, A. Mignerey, K. C. Wolf, H. F. Breuer, V. I. Viola, Phys. Rev. C20 (1979) 576.
17) D. J. Morrissey and L. G. Moretto, Phys. Rev. C23 (1981) 1835.
18) R. P. Schmitt, G. J. Wozniak, G. U. Rattazzi, G. Mathews, R. Regimbart and L. G. Moretto, Phys. Rev. Lett. 46 (1981) 522.
19) N. Frascaria, P. Colombani, A. Gamp, J. P. Garron, M. Riou, J. C. Roynette, C. Stephan, A. Ameaume, G. Bizard, J. L. Laville, M. Louvel, Z. Phys. A294 (1980) 167.
20) J. R. Huizenga, W. U. Schröder, J. R. Birkelund and W. W. Wilcke, Nucl. Phys. A387 (1982) 257.
21) G. Rudolf, A. Gobbi, H. Stelzer, U. Lynen, A. Olmi, N. Sann, R. G. Stokstad and D. Pelte, Nucl. Phys. A330 (1979) 245.
22) W. U. Schröder, J. R. Birkelund, J. R. Huizenga, W. W. Wilcke and J. Randrup, Phys. Rev. Lett. 44 (1980) 300.
23) P. Glässel, D. V. Harrach, L. Grodzins, H. J. Specht, Z. Phys. A310 (1983) 189.
24) J. V. Kratz, J. Poitou, W. Brüchle, H. Gäggeler, M. Schädel, G. Wirth, R. Lucas, Nucl. Phys. A357 (1981) 437.
25) M. Berlanger, A. Gobbi, F. Hanappe, U. Lynen, C. Ngô, A. Olmi, H. Sann, H. Stelzer, H. Richel, M. F. Rivet, Z. Phys. A291 (1979) 133.
26) L. G. Moretto, C. R. Albiston and G. Mantzouranis, Phys. Rev. Lett. 44 (1980) 924.
27) W. Cassing and H. Friedrich, Z. Phys. A298 (1980) 129.
28) H. Feldmeier, International School of Physics "Enrico Fermi", Varenna 1982.
29) H. Spangenberger, F. Beck and H. Feldmeier, GSI Scientific Report 1982 ISSNO 174-0814.
30) V. M. Strutinski, Phys. Lett. 44B (1973) 245.
31) P. Glässel, R. S. Simon, R. M. Diamond, R. C. Jared, I. Y. Lee, L. G. Moretto, J. O. Newton, R. P. Schmitt and F. S. Stephens, Phys. Rev. Lett. 38 (1977) 331.
32) L. G. Sobotka, C. C. Hsu, G. J. Wozniak, D. J. Morrisey and L. G. Moretto, Nucl. Phys. A371 (1981) 510.
33) C. Lauterbach, W. Dünnweber, G. Graw, W. Hering, H. Puchta and W. Trautmann, Phys. Rev. Lett. 41 (1978) 1774.
34) L. G. Moretto, S. K. Blau and A. J. Pacheco, Nucl. Phys. A364 (1981) 125.
35) R. J. McDonald, A. J. Pacheco, G. J. Wozniak, H. B. Bolotin, C. Schuck, S. Shih, R. M. Diamond, F. Stephens and L. G. Moretto, Nucl. Phys. A373 (1982) 54.
36) A. J. Pacheco, G. J. Wozniak, R. J. McDonald, R. M. Diamond, C. C. Hsu, L. G. Moretto, D. J. Morrisey, L. C. Sobotka and F. S. Stephens, LBL14091, Nucl. Phys., in press.
37) D. Logan, M. Rajagopalan, M. S. Zisman, J. M. Alexander, M. Kaplan and L. Kowalski, Phys. Rev. C22 (1980) 104; D. Logan, M. Delagrange, M. F. Rivet, M. Rajagopalan, J. M. Alexander, M. Kaplan, M. S. Zisman and E. Duek, Phys. Rev. C22 (1980) 1080.

38) M. Robel, PhD Thesis LBL 8181, 1979.
39) J. P. Bondorf, J. N. De, G. Fai, A. O. T. Karvinen, B. Jakobson and J. Randrup, Nucl. Phys. A333 (1980) 285.
40) R. Weiner and M. Weström, Phys. Rev. Lett. 34 (1975) 1523.

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