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Presented at the International Symposium on Nuclear Physics at Large Tandem Accelerators, Legnaro-Padova, Italy, March 15-17, 1983

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L.G. Moretto

March 1983

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

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NEW STATISTICAL AND NON-STATISTICAL EFFECTS IN DEEP INELASTIC HEAVY ION COLLISIONS

by

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Abstract: The physics of deep inelastic heavy ion collisions is briefly reviewed in the light of the experimental capabilities offered by large Tandem Accelerators. Three aspects are selected for illustration. The first is the problem posed by the angular distributions. The role of quantal fluctuations (diffraction) and of dynamical fluctuations is discussed. The possibility of extracting angular momentum fluctuations from the angular distribution width is pointed out, as well as the way to compare such results with those obtained from gamma-ray multiplicities or angular distributions. The likely sources of angular momentum fluctuations are discussed. The second problem is that of energy dissipation and partition. The experimental evidence for thermal equilibrium is reviewed and experiments to test thermal equilibrium through the measurement of the variance in the energy distribution are suggested. The correlation between energy loss and the mass variance is reviewed in terms of the mass transfer mechanism. Possible ways to determine the mass of the transferred particle are illustrated. The third problem considered is the experimental lack of mass drift in the mass distributions vs Q-value when it is expected on the basis of potential energy considerations. This problem is associated with the fast fragment thermalization and explained in terms of a feed-back effect associated with the temperature gradient that developes between two colliding nuclei of different mass.

Introduction The last decade has seen a rather extensive exploration of the heavy ion deep inelastic reaction field. The boundaries of this field with the direct reactions on one hand and with the compound nucleus reactions and perhaps fast-fission on the other, have been more or less established. The degrees of freedom associated with deep inelastic reactions have been identified and the dynamical regimes prevailing in their time-evolution have been demonstrated. In particular the dominant role of non-equilibrium and equilibrium statistical mechanics in almost all the important degrees of freedom has been stressed by a vast amount of experimental work[1].

The evolution of the field in the next decade will occurr, in my opinion, in two directions. On one hand there will be an exploration along the energy coordinate which will show a new boundary most likely related to the ability of the mean field to contain fast-moving nucleons and to the establishment of a shorter mean-free-path regime. On the other hand the "classical" field of deep inelastic reactions will be reworked with more pointed experiments in order to answer specific questions brought forth by competing theories. It is mainly in the latter direction that machines like large Tandem Accelerators can give important contributions.

In what follows I would like to outline three themes that are in great need

*Alexander von Humboldt Senior U.S.Scientist Award.

This work was partially supported by the Director, Office of Energy Research, Division of Nuclear Physics of the Office of High Energy and Nuclear Physics of the U.S. Department of Energy under Contract DE-AC03-76SF00098

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of clarification and that could be readily accessible from an experimental point of view, namely the angular distributions, the energy dissipation and partition, and finally the mass distribution.

Angular Distributions: Quantal and Statistical Fluctuations Experimentally the angular distributions observed in deep inelastic processes vary from side-peaked to forward-peaked, frequently in the same reactions. The side-peaking is found at relatively low Q-values and near the entrance-channel mass asymmetry, while the forward-peaking is apparent with increasing inelasticity and for mass asymmetries far from that of the entrance channel. Such a remarkable evolution must find its counterpart in a great variation in the interaction time on one hand, and in a rapid change in the width of the orbital angular momentum distribution on the other.

When a side-peaked angular distribution is present, it can be parameterized in terms of a center angle $\hat{\theta}$ and of its width σ_θ. What information can be obtained from these two quantities? The center angle $\hat{\theta}$ can be of use in constructing the deflection function if it can be associated with some average angular momentum. This can be tried by assuming a monotonic dependence[1,2] of the energy on angular momentum and by finding the l-space bin corresponding to the energy bin under consideration by means of the associated cross section. The width of the distribution arises from two contributions, as pointed out by Strutinski[3], namely a "diffractive" contribution, associated with the width of the l-packet, and a "dynamical" contribution associated with the classical deflection function. More precisely:

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

The first term can be readily derived from the indetermination principle, while the second propagates the width of the l-packet in angle through the deflection function $\theta = \theta(l)$. An interesting application of the above equation is in providing a simple explanation to some complex calculations. A Wilczinski diagram generated with a diffusion calculation[4] is unable to reproduce the width of the angular distribution at small energy losses as shown in fig. 1a. This is due to the fact that the fluctuations in l depend on the temperature, namely $\sigma_l^2 \propto T$. The introduction of quantal fluctuations results in a greatly broadened and more realistic diagram as illustrated in fig. 1b. The quantal broadening is of course arising from the narrow thermal l-distribution. On the other hand, a Wilczinski diagram generated with a one-body diffusion model (fig.2) is quite realistic even at small Q-values without quantal fluctuations. In fact, the introduction of quantal fluctuations does not broaden the angular distributions appreciably[6]. The explanation lies in the fact that the one-body theory predicts a large "dynamical" σ_l even for low Q-values, which is of course responsible for the broad angular distributions. But a large σ_l implies a small quantal fluctuation. Therefore the introduction of quantal fluctuations hardly changes the results. This illustrates how from true experimental data one can extract σ_l and compare it directly with theoretical predictions. Of course one trivial component of σ_l is due to the energy binning or

$$\sigma_l^2 = \left(\frac{dl}{dE}\right)^2 \sigma_E^2 .$$

The remainder of the width arises from true orbital angular momentum fluctuations.

An obvious source of l-fluctuations is the coupling with a "wriggling" mode[7]. In this coupling, orbital angular momentum can be traded for spin and vice-versa. The thermal limit is easy to calculate, especially for two touching spheres[8].

If I is the total angular momentum, the total rotational energy is:

$$E_R = \frac{l^2}{2\mathcal{J}^*} + \frac{I^2}{4\mathcal{J}} - \frac{Il}{2\mathcal{J}} \quad \text{where } \frac{1}{\mathcal{J}^*} = \frac{1}{\mu R^2} + \frac{1}{2\mathcal{J}}$$

and $\mathcal{J}$ is the moment of inertia of one of the spheres. The orbital angular momentum distribution at fixed I is:

$$P(l)dl = (2\pi\mathcal{J}T)^{-1/2}\exp-\left(\frac{l^2}{2\mathcal{J}^*T} - \frac{Il}{2\mathcal{J}T} + I^2\frac{\mathcal{J}^*}{8\mathcal{J}^2T}\right)$$

Introducing a 2IdI weight, remembering that the exit channel kinetic energy above the Coulomb energy is $E = l^2/2\mu R^2$ and introducing the reduced variables $\varepsilon = E/T$, $\lambda = I/(\mathcal{J}T)^{1/2}$, one obtains:

$$P(\varepsilon,\lambda) \propto \frac{\lambda}{\varepsilon}\exp-\left[\frac{7}{2}\varepsilon - \frac{5}{2}\lambda\varepsilon + \frac{5}{28}\lambda^2\right]$$

A map of this function is shown in fig. 3. At constant energy ε the most probable value of the angular momentum is:

$$\hat{\lambda} = \frac{7}{\sqrt{10}}\sqrt{\varepsilon}\left(1+\sqrt{1+\frac{4}{7\varepsilon}}\right)$$

while the variance, in which we are mainly interested, is

$$\sigma_\lambda^2 = \frac{14}{5} \quad \text{or} \quad \sigma_l^2 = \frac{14}{5}\mathcal{J}T$$

A further source of l-fluctuations is due to shape fluctuations in the dinucleus[8]. Calculations in the thermal limit are shown in fig. 4 for the system Fe + Fe. It can be seen that very substantial fluctuations in orbital angular momentum can be expected at fixed energy if shape fluctuations should occur.

In both cases considered above we have an interesting correlation: the leading edge of the angular distribution corresponds to $l > \hat{l}$ while the trailing edge corresponds to $l < \hat{l}$. This discussion suggests some possible experimental tests. For instance, the fluctuations in l measured from the angular distribution widths should find their counterpart in the spin fluctuations in the same energy bin. The spin fluctuations can be obtained for instance from a measurement of the spin misalignment[7,9].

Perhaps even more interestingly one could measure the spins (e.g. gamma-ray multiplicities) to the left and to the right of the angular distribution peak and verify that the gamma multiplicity to the right is greater than the gamma multiplicity to the left.

Energy Dissipation and Partition We have learned early in the study of heavy ion reactions that at low bombarding energies all of the dissipated energy goes into the fragments. Even at rather high energy (15-20 MeV/nucleon) most of the energy still goes into the fragments and only a minor part is associated with prompt particle emission[1].

The dissipated energy is partitioned between the two fragments proportionally to their masses even at very low energy losses[10-12]. This seems to signify that the two fragments reach thermal equilibrium even after very short interaction times which are associated with small energy losses.

These observations raise some serious problems which are still unanswered.
1) Is the energy partition truly thermal? In order to verify that such a partition is not due to some unexpected property of the primary energy deposition mechanism, one should check whether the higher moments of the distribution are thermal. The statistical model gives a straight-forward prediction for the energy fluctuations[13], namely

$$\sigma_E^2 = \frac{Cv_1Cv_2}{Cv_1+Cv_2}T^2 = 2\frac{a_1a_2}{a_1+a_2}T^3$$

So far no direct check exists for this thermal limit of the variance. Any measurement along this line would help a great deal to clarify this still misteri-

ous problem.

2) If the fragments are truly thermalized, how can they have done it so fast (in 2 or 3 x 10^{-22} sec.)? This question becomes even more pressing as we consider that the likely mechanisms for energy transfer tend to deposit approximately the same amount of energy in each fragment irrespective of the mass asymmetry. We will try to give an answer to this question later on.

3) How is energy transferred from relative motion to the internal degrees of freedom? This a question that is still looking for an elusive answer. Various obvious possibilities have been considered, like mass transfer, particle-hole excitation, and phonon excitation[1]. Let us consider the mass transfer in some detail since it is amenable to a series of experimental verifications.

Mass exchange is inevitably associated with energy transfer. In particular an amount of transferred mass dm will transfer an amount of energy given by:

$$dE = E\frac{dm}{\mu}$$

where μ is the reduced mass of the system. If Φ is the (time independent) mass flux, we have

$$\frac{dE}{dT} = E\,\frac{\Phi}{\mu} \quad \text{or} \quad E = E_o\, e^{-\Phi t/\mu}$$

Notice how the mass of the transferred particle does not appear explicitly in the above equations. This means that the energy loss due to mass transfer does not contain any information "per se" on the mass of the particles exchanged. If we assume a stochastic exchange of N particles of mass m (random walk), the variance of the mass of a fragment is:

$$\sigma^2 = Nmxm = \Phi txm$$

With this the energy expression becomes

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

This expression now depends on the mass of the transferred particle and suggests that such an information is contained in the correlation between energy loss and mass variance. An attempt to fit the experimental data using for m the nucleon mass fails altogether[2,14,15]. However we must remember that we have derived just a classical expression. The inclusion of the Pauli principle[15] requires the following modifications:

$$\sigma^2 = \psi tm$$

$$E = E_o \exp - \frac{\sigma^2\Phi}{\mu m\Psi}$$

where Ψ is the Pauli blocked flux. This means that

$$\sigma^2_{true} = \sigma^2_{class}\,\frac{\psi}{\Phi}.$$

In this way a good fit is in fact obtained, as shown in fig. 5. Can we conclude that the mass of the transferred object is m = 1 ? Not necessarily. In fact, if we assume that an alpha particle (m = 4) is being exchanged we do not need to include any Pauli principle and we obtain an equally good fit. There may be additional evidence of the alpha particle transfer in the even-odd staggering of the Z distributions observed at low Z. It should be possible to solve the dilemma by studying the correlation E_{loss} vs σ^2 at different bombarding energies.

<u>Mass Distribution</u> The open problem associated with the mass distribution is the lack of mass drift towards symmetry which we expect on the basis of the potential energy (liquid drop + rotational) as a function of mass asymmetry. Such a drift seems to appear, if at all, only after all the available energy has been dissi-

pated. This puzzle is as great as that associated with the thermalization of the energy between the two fragments and may be related to it. Here is a qualitative sketch for a possible solution.

Let us assume that particle transfer alone is responsible for energy dissition. While the excitation of collective vibrations may be quite important in the process of energy dissipation, it does not provide us with an easy coupling to the mass drift. Therefore we do not consider it for the moment.

Let us also assume that the particle fluxes are temperature-dependent. This can be readily verified for the regime in which a barrier is present between the two potential wells, which is, incidentally, the regime prevailing in the quasi elastic region where energy is being rapidly dissipated. Since particle exchange deposits nearly the same amount of energy in each fragment, if the system is asymmetric the lighter fragment must grow hotter. This temperature gradient now creates an imbalance in the particle fluxes, the large fragment receiving more nucleons than it gives away. Since the energy deposition on each fragment is proportional to the number of nucleons that land on it, more energy will now be deposited on the heavy fragment in an attempt to correct the temperature gradient. Therefore we see here a feedback effect that leads to a selfregulation in the energy partition. More specifically, the initially generated thermal imbalance redirects the energy deposition towards the larger fragment and controls the thermal gradient at the expense of the mass of the light fragment. Notice that the advocation of a feedback effect satisfactorily answers the requirement of speed because it acts not on the energy already dissipated, but on the energy that is being dissipated, deciding where it ought to go.

Notice also that, when the relative motion is damped, the effect that tends to make the small fragment smaller disappears altogether, allowing the system to drift down the static potential energy slope.

Let us consider two objects moving tangent to one another with a window open between them which allows for particle transfer. Let $\Phi(T)$ be the particle flux and $\Phi_E(T)$ be the energy flux exclusive of the relative motion. Also let the energy-temperature relation be given by $E = Nf(T)$.
Particle conservation gives the equation:

$$\frac{dN_1}{dt} = \Phi(T_2) - \Phi(T_1)$$

Notice that $\frac{dN_1}{dt} < 0$ if $T_2 > T_1$

From momentum conservation we have:

$$\frac{dp_1}{dt} = -V_1\Phi(T_1) - V_2\Phi(T_2) = N_1\frac{dV_1}{dt} + V_1\frac{dN_1}{dt}$$

$$N_1\frac{dV_1}{dt} = -(V_1+V_2)\Phi(T_2); \quad N_2\frac{dV_2}{dt} = -(V_1+V_2)\Phi(T_1)$$

Finally $\frac{d\ln V}{dt} = -\left(\frac{\Phi(T_2)}{N_1} + \frac{\Phi(T_1)}{N_2}\right); \quad V = V_1 + V_2$

From energy conservation we have:

$$\frac{dE_1}{dt} = \Phi E(T_2) - \Phi_E(T_1) + \frac{1}{2}\Phi(T_2)V^2$$

In order to obtain an equation for the temperature we can use:

$$E = Nf(T); \quad \frac{dE}{dt} = f(T)\frac{dN}{dL} + N\frac{df(T)}{dT}\frac{dT}{dt} = f(t)\frac{dN}{dt} + Cv\frac{dT}{dt}$$

where C_v is the heat capacity.
Solving for dT/dt we have

$$\frac{dT_1}{dt} = \frac{1}{C_{v,1}}\left[\Phi_E(T_2)-\Phi_E(T_1) + \frac{1}{2}\Phi(T_2)V^2 - f(T_1)\left(\Phi(T_2)-\Phi(T_1)\right)\right]$$

$$\frac{dT_2}{dt} = \frac{1}{C_{v,2}}\left[\Phi_E(T_1)-\Phi_E(T_2) + \frac{1}{2}\Phi(T_1)V^2 - f(T_2)\left(\Phi(T_1)-\Phi(T_2)\right)\right]$$

The last two equations show the feedback quite clearly. If for instance $T_1 > T_2$, the last term of the last equation will increase the value of dT_2/dt until there is a non-zero relative velocity.

We need now to calculate the Fermi gas fluxes for any given barrier B. For instance, the particle flux can be written as

$$I = \frac{2x2\pi}{(2\pi h)^3} \int \frac{v\cos\theta\sin\theta d\theta p^2 dp}{1+\exp(\varepsilon-\mu)/T} = K \int_B^\infty \frac{(\varepsilon-B)d\varepsilon}{1+\exp(\varepsilon-\mu)/T}$$

Three limits are readily calculable:

1) $\frac{B-\mu}{T}$ large, positive: $I \approx K\, T^2 \exp-(\frac{B-\mu}{T})$

2) $\frac{B-\mu}{T}$ large, negative:

$$I = K\left[\frac{\pi^2}{6} T^2 + T(\mu-B)\ln(1+\exp\frac{\mu-B}{T}) - \frac{1}{2}(\mu-B)^2 - T^2 e^{-\frac{\mu-B}{T}} (\frac{\mu-B}{T}+1)\right]$$

3) $\frac{B-\mu}{T}$ small:

$$I = K\left[\frac{\pi^2}{12} T^2 + T(\mu-B)\ln(1+\exp\frac{\mu-B}{T}) - \frac{1}{4}(B-\mu)^2 + \frac{1}{12}\frac{(B-\mu)^3}{T} - \frac{1}{240}\frac{(B-\mu)^5}{T^3} \ldots\right]$$

With these fluxes one can perform some simple calculations. Let us consider a dinucleus with mass 50-150. In fig.6a we see the dependence of the mass of the light fragment upon energy loss at three bombarding energies. Since in this calculation we did not include any driving force, the light fragment tends to lose mass, as expected. In fig.6b the temperature ratio of the two fragments is shown vs energy loss. We see that the initial temperature gradient which is established between the two fragments is readily controlled and reduced by the feedback mechanism.

In fig. 7 the dependence of the light fragment mass upon energy loss is shown for various driving forces. The general behaviour of these curves is quite interesting. At small energy losses the system tries to follow the driving force until the temperature gradient starts pushing mass in the opposite direction. In this way the mass asymmetry is more or less stabilized until the energy is almost completely dissipated. After this has occurred the system can finally drift in the direction of the driving force. Evidence for a structure in the dependence of the mass upon Q-value similar to that of fig. 7 has recently been found in a measurement where the pre-evaporation masses of the two fragments were determined kinematically[16]. This evidence is shown in fig. 8.

It is quite obvious that detailed studies of the primary masses of the fragments as a function of Q-value for different bombarding energies and initial mass asymmetries are necessary to test our theories.

Conclusions Hopefully this brief and selective survey on the field of heavy ion deep inelastic collisions has shown that a great number of exciting problems is still open to our investigation and that a large amount of detailed work needs to be done in order to transform many educated guesses or qualitative inferences into solid science. It is my opinion that large Tandem machines are ideally suited for this kind of work.

Acknowledgment It is a pleasure to thank W.A. Weidenmüller for his kind hospitality at the Max-Planck-Institut.

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Figure Captions

Fig. 1a Calculated Wilczynski plot without quantal calculations[4].

Fig. 1b Same as in 1a but with quantal fluctuations.

Fig. 2 Calculated Wilczynski plot from a one-body diffusion model[5] without quantal fluctuations.

Fig. 3 Two-dimensional plot of the energy-angular momentum distribution function obtained with a thermally fluctuating wriggling mode[8].

Fig. 4a Two-dimensional plot of the energy-angular momentum distribution function obtained by including shape fluctuations[8].

Fig. 4b Constant energy cuts of the function in a.

Fig. 5 Correlation between energy loss and charge variance[14]. The dashed line represents the classical transfers of particles with m = 1 while the solid lines include the Pauli blocking.

Fig. 6a Dependence of the light fragment mass upon energy loss for 100,200,300 MeV center of mass energy. The initial masses are 50,150 (top).

Fig. 6b Temperature ratios vs energy loss for the same system and energies as in fig. 6a (bottom).

Fig. 7 Dependence of the light fragment mass upon energy loss for the same system as in fig. 6b with (left to right) 0.5,1.0,1.5,2.0 MeV/A driving force towards symmetry.

Fig. 8 Kinematically measured average masses as a function of energy loss for the reaction $Kr + Er^{16}$.

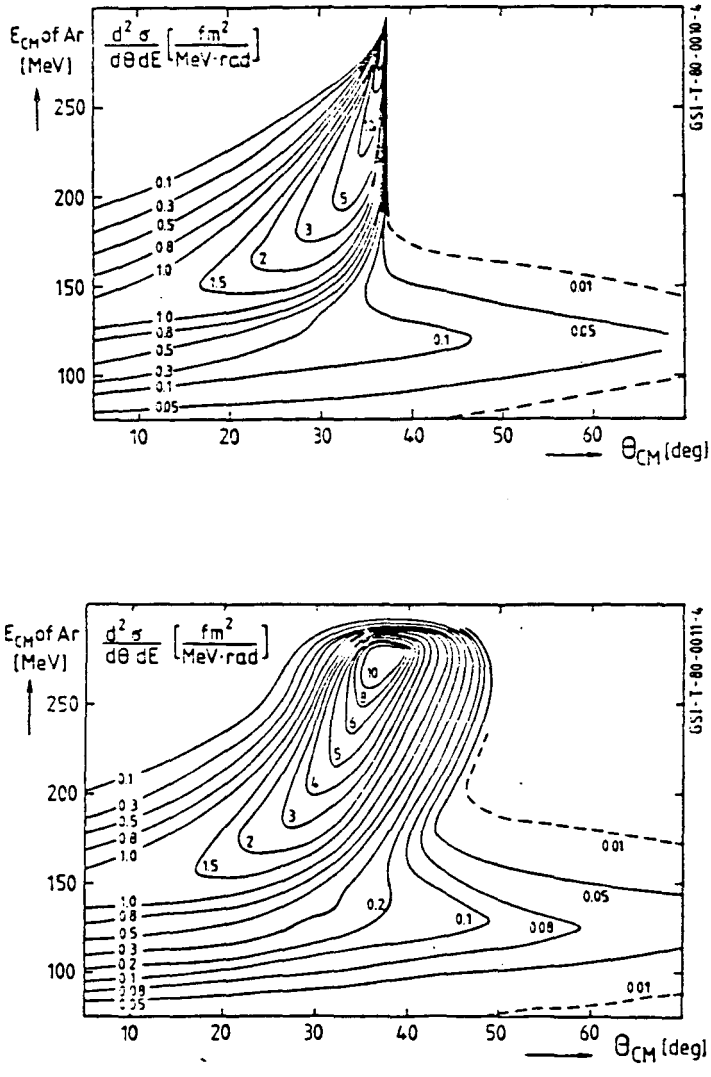
E_{CM} of Ar [MeV]
$\frac{d^2\sigma}{d\theta\, dE}\left[\frac{fm^2}{MeV \cdot rad}\right]$
θ_{CM} [deg]
GSI-T-80-0010-4
GSI-T-80-0011-4

fig. 1a, b

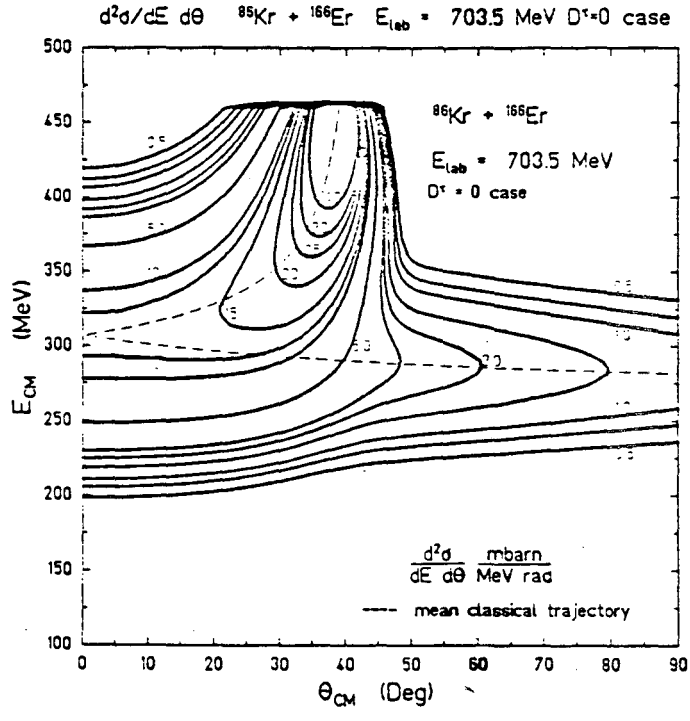
$d^2\sigma/dE\, d\theta$ $^{86}Kr + ^{166}Er$ E_{lab} = 703.5 MeV D^*=0 case
$^{86}Kr + ^{166}Er$
E_{lab} = 703.5 MeV
D^* = 0 case
E_{CM} (MeV)
θ_{CM} (Deg)
$\frac{d^2\sigma}{dE\, d\theta}$ $\frac{mbarn}{MeV\ rad}$
mean classical trajectory

fig. 2

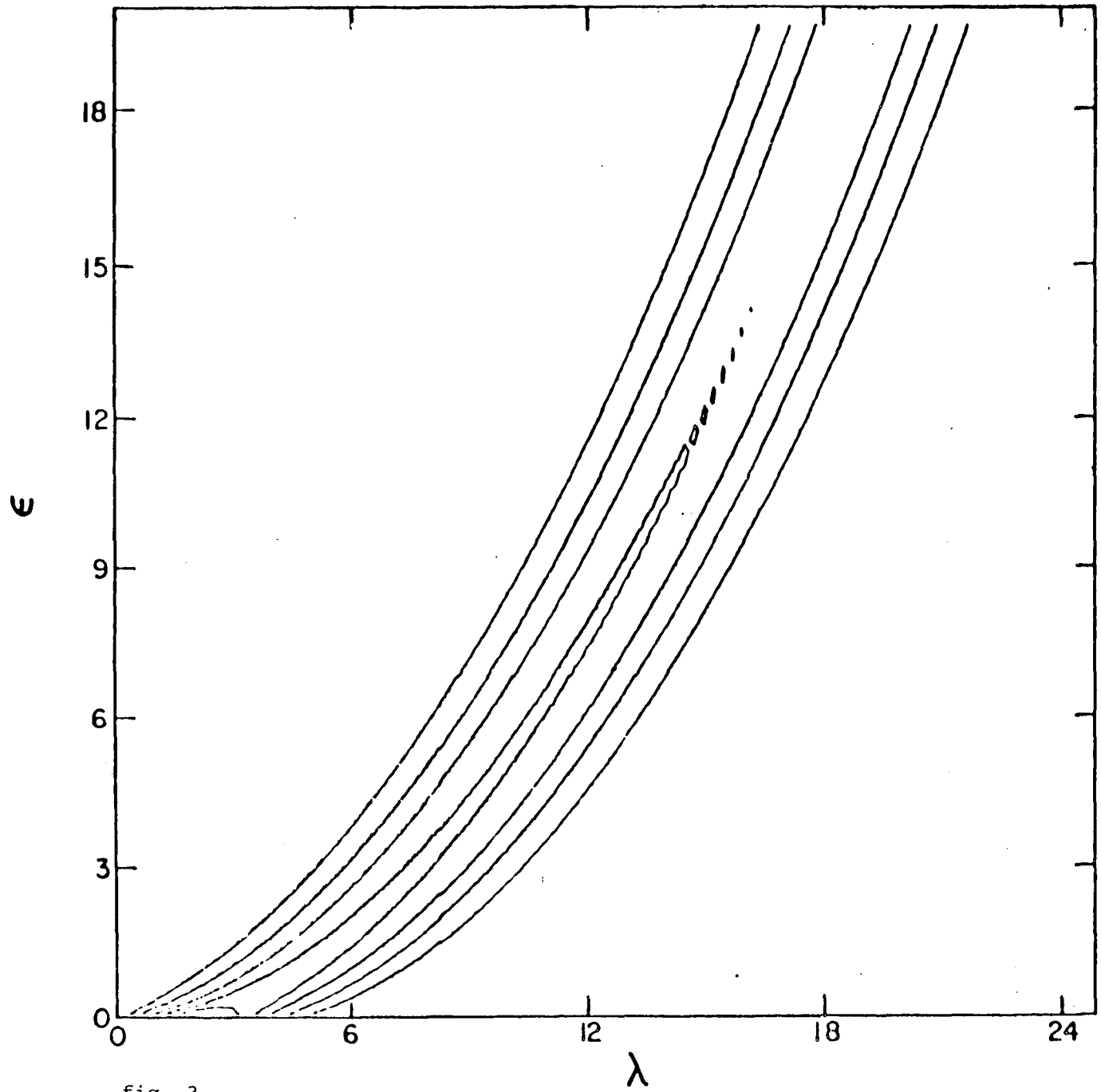

fig. 3

XBL815-2262

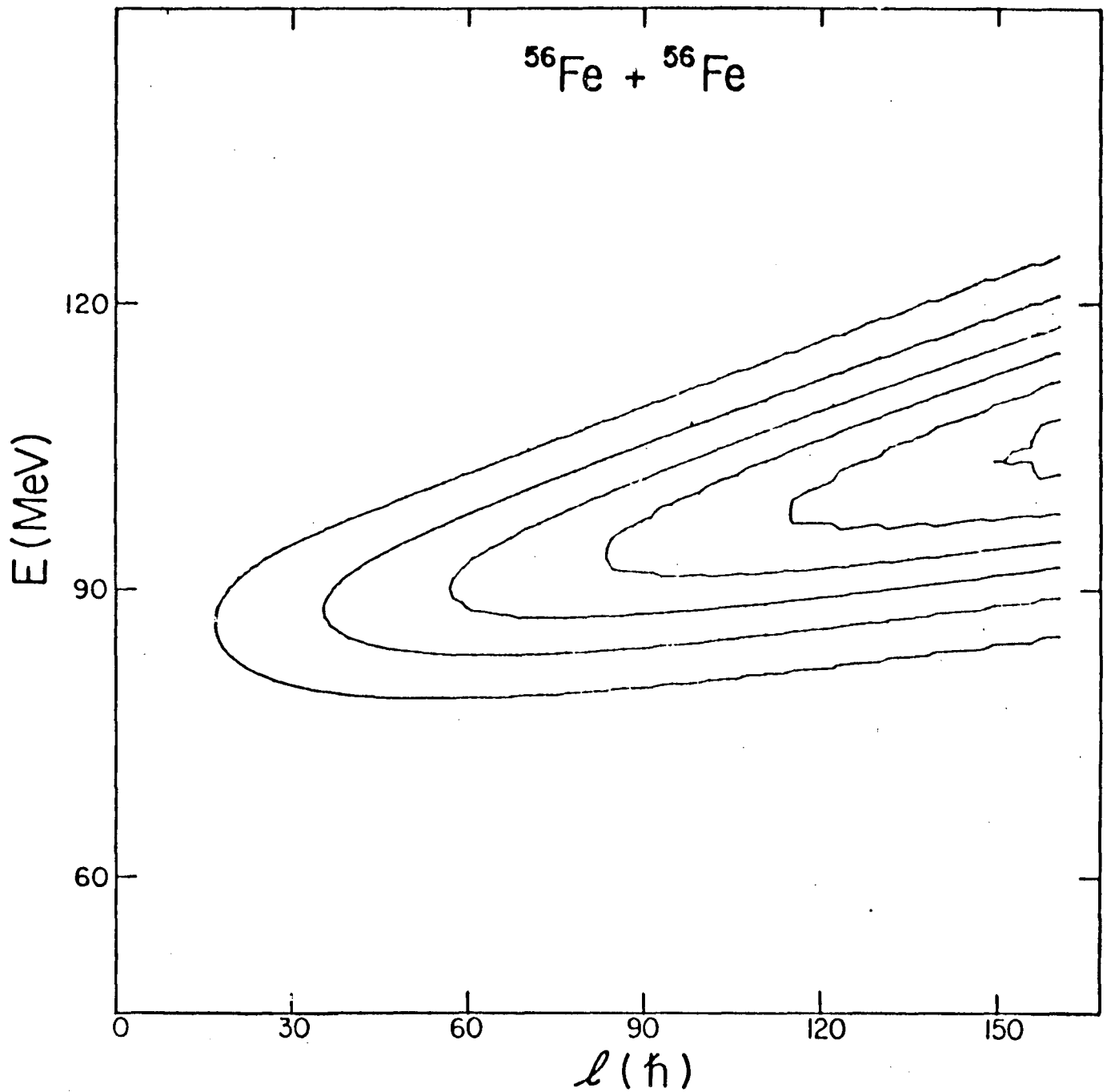

fig. 4a

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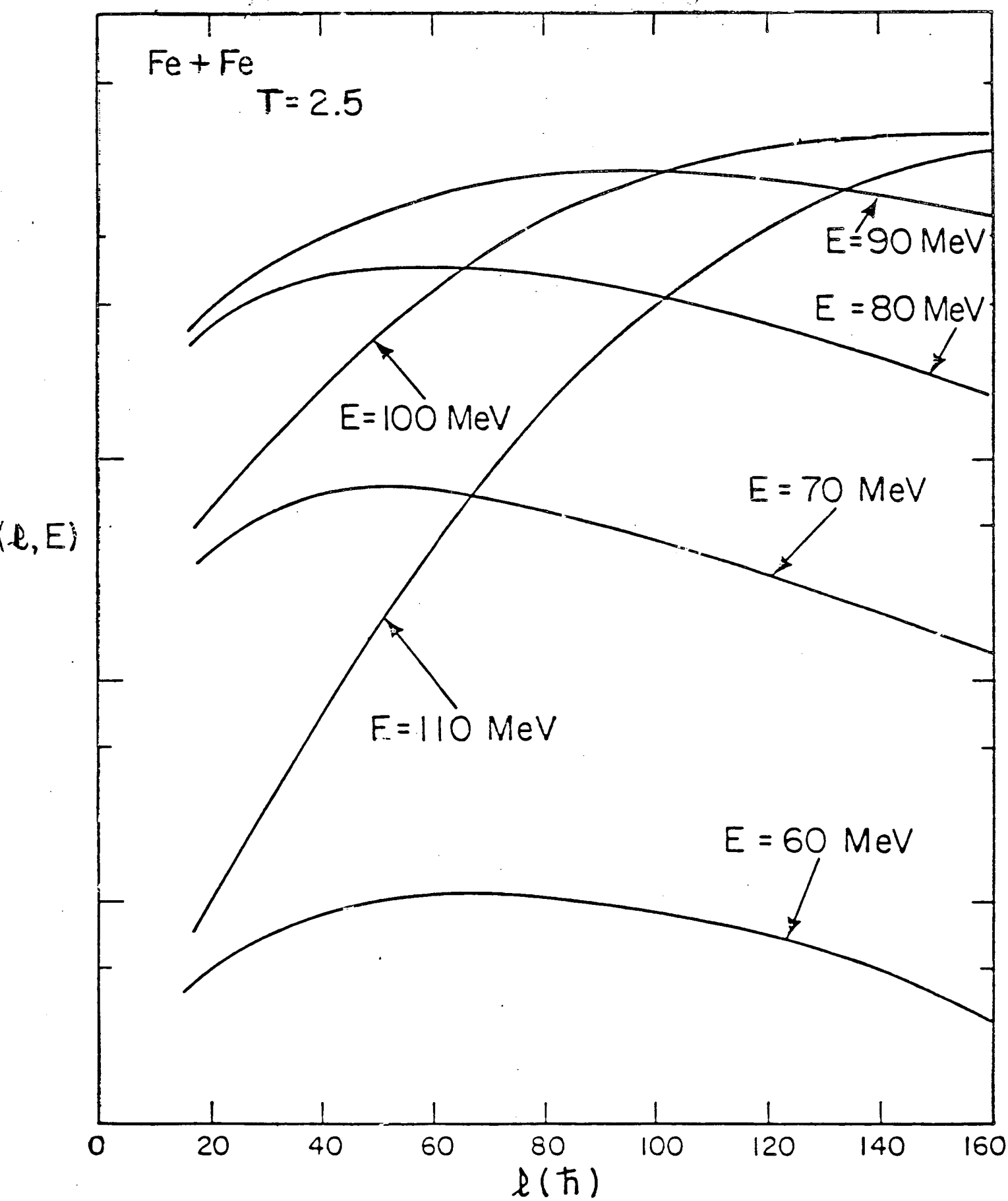
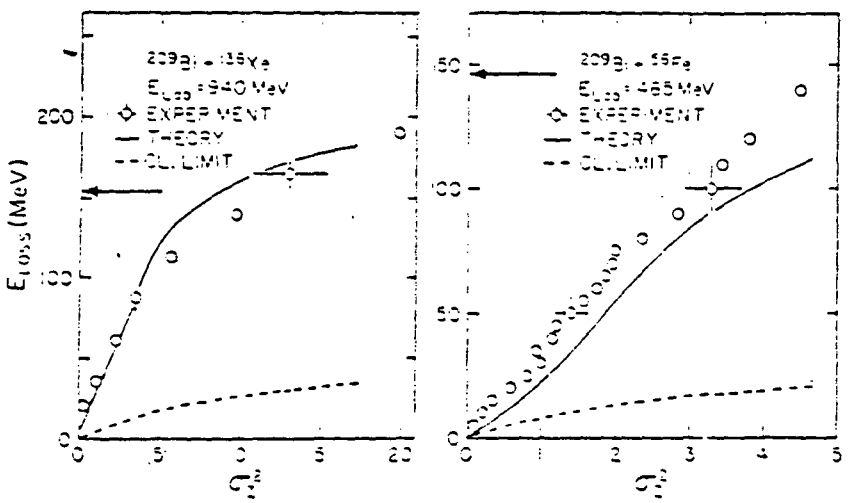


fig. 4b

XBL 814-768


EXPERIMENT
THEORY
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THEORY
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E_{loss} (MeV)
σ_2^2
σ_2^2

fig. 5

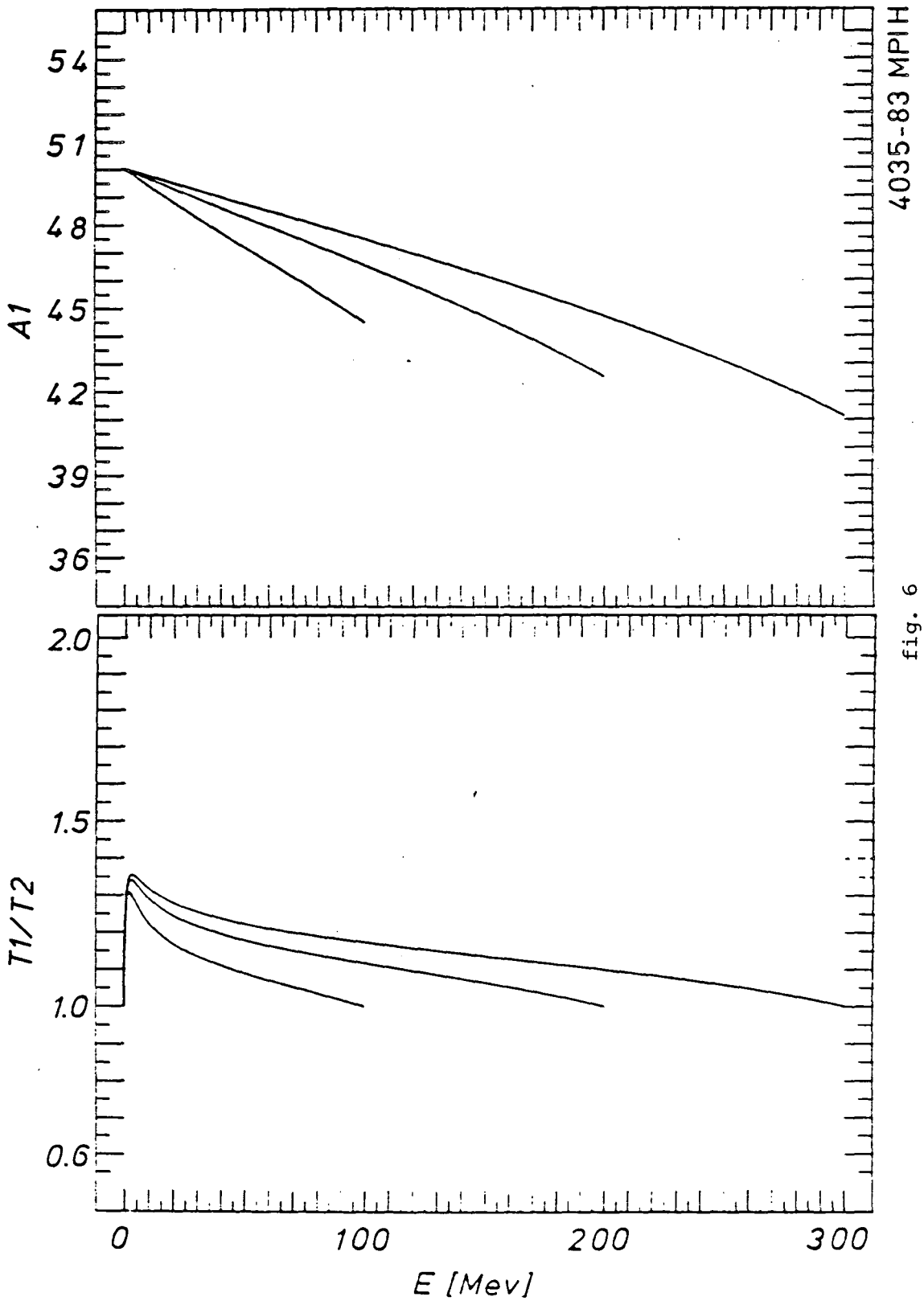
54
51
48
45
42
39
36
A1
2.0
1.5
1.0
0.6
T1/T2
0
100
200
300
E [Mev]
4035-83 MPIH

fig. 6

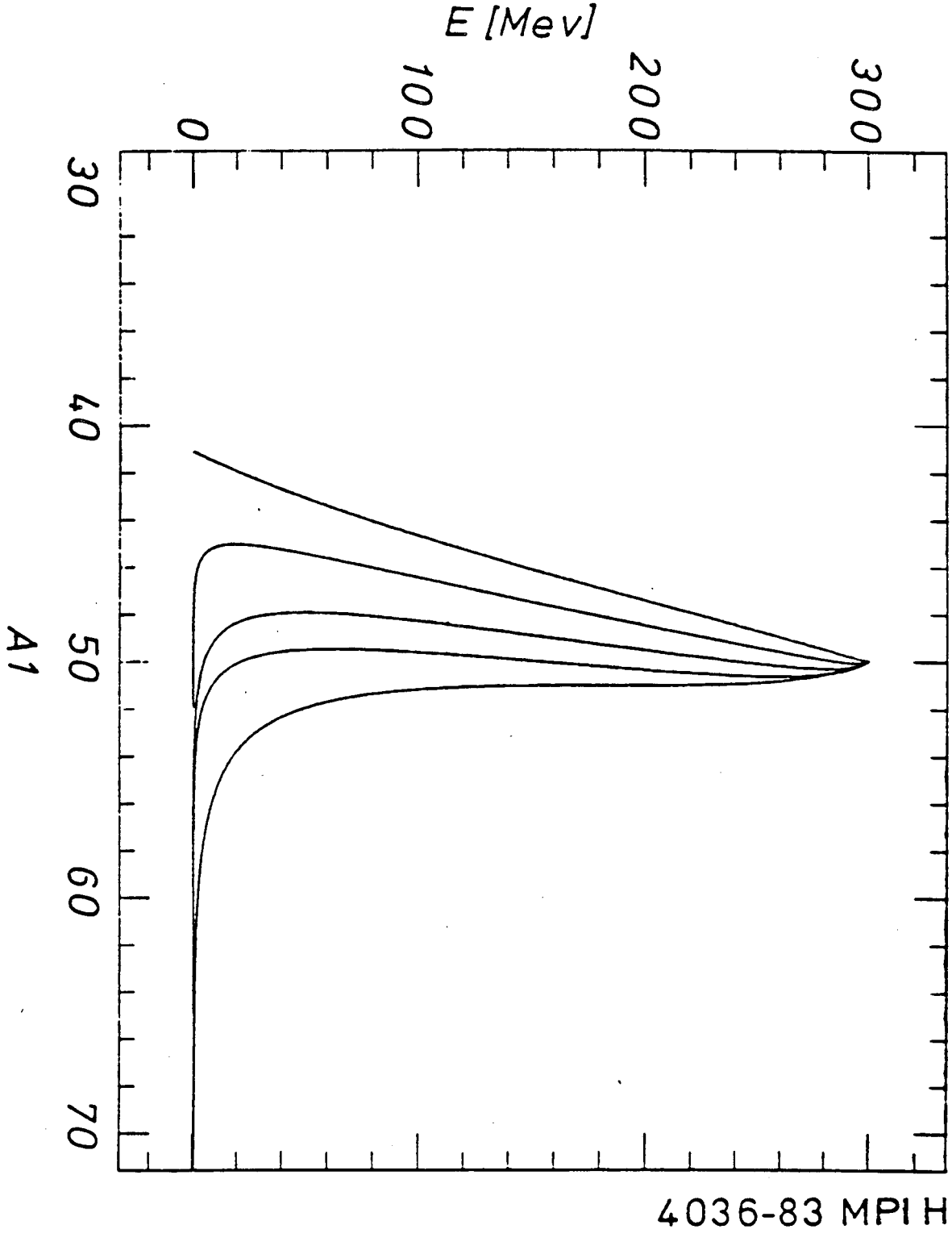

fig. 7

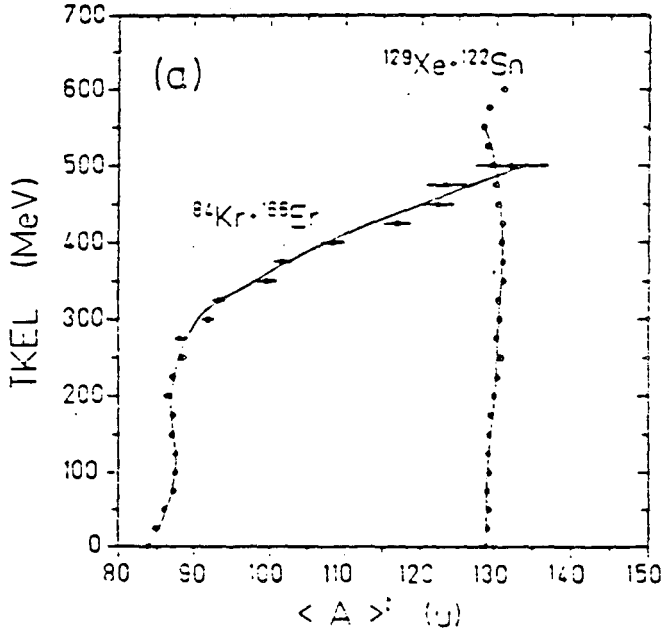
(a)
$^{129}Xe+^{122}Sn$
$^{84}Kr+^{166}Er$
TKEL (MeV)
< A >
700
600
500
400
300
200
100
0
80
90
100
110
120
130
140
150

fig. 8

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