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Gesellschaft für Schwerionenforschung mbH
Planckstr. 1 · Postfach 110541 · D-6100 Darmstadt 11 · Germany

Influence of single-particle shell structure on the energy dissipation in heavy-ion collisions

M. Dakowski¹, A. Gobbi and W. Nörenberg

Gesellschaft für Schwerionenforschung (GSI),

D-6100 Darmstadt, W. Germany

Abstract

Shell effects in the mean energy dissipated per transferred nucleon are analysed within a statistical model of random transfers and excitations. Quantitative agreement for all 21 analysed reactions is obtained without fitting any parameter to the data.

¹) on leave from Institute for Nuclear Research, 05-400 Swierk
k.Otwocka, Poland

1. Introduction

In collisions between very heavy nuclei the microscopic mechanism of energy dissipation has not been revealed completely. There exist several models for calculating the average energy dissipation per transferred nucleon. The classical window formula of Blocki et al.¹⁾ as applied by Beck et al.²⁾ and Schröder et al.³⁾ severely underestimates the energy loss rate. "Pauli blocking" as introduced by Randrup⁴⁾ improves the calculations⁵⁾. However, all these models as well as the more sophisticated transport theories⁶⁾ do not take the shell structure of the single-particle levels into account. Therefore they are not able to account for the large experimental differences in the energy loss rates for closed-shell systems like $^{208}\text{Pb}+^{208}\text{Pb}$ ⁷⁾ or $^{136}\text{Xe}+^{209}\text{Bi}$ ^{8,9)} as compared to systems like $^{238}\text{U}+^{238}\text{U}$ ¹⁰⁾ or $^{154}\text{Sm}+^{154}\text{Sm}$ ¹¹⁾ with half-filled shells. Although the importance of shell effects was demonstrated on the basis of experimental systematics ^{12,13)}, no well founded theoretical explanation has been proposed.

It is the purpose of this paper to examine in detail the shell-structure effects on the mean energy dissipated per transferred nucleon. We introduce in chapter 2 a statistical model which treats simultaneously all possible one-particle transitions within and between the interacting nuclei. We start from a rate equation for the occupation probabilities of the single-particle states taking full account of the Pauli principle. The energy transfer in a single transition is determined from the allowed momentum (angular momentum) transfer between collective (relative) and intrinsic motion. Using realistic single-particle levels the energy loss rate is obtained without fitting any parameter. Chapter 3 compares the calculated results with the experimental

data from 21 reactions.

2. Model of random transitions

In a statistical model the probabilities $F_i(t)$ for the occupation of single-particle states i is determined by the rate equation

$$dF_i(t)/dt = \sum_j w_{ij}(t) \{ F_j(t) [1 - F_i(t)] - F_i(t) [1 - F_j(t)] \}. \quad (2.1)$$

The first term on the r.h.s. describes the gain of occupation probability of F_i due to transitions from all states j to the state i , whereas the second term represents the total loss of probability from i . The transition probability $w_{ij}(t)$ is symmetric with respect to the exchange of i and j and vanishes whenever i and j have different isospin. This guarantees conservation of the total probability for protons and neutrons separately. The rate equation (2.1) takes full account of the Pauli exclusion principle.

The energy-loss rate $\dot{E}$ is obtained from (2.1) by

$$\dot{E} = \sum_i \varepsilon_i \frac{dF_i(t)}{dt} = \sum_{ij} (\varepsilon_i - \varepsilon_j) w_{ij}(t) F_j(t) [1 - F_i(t)] \quad (2.2)$$

where ε_i denotes the energy of the single-particle state i . It should be noted that this expression includes both the excitations with and without transfer. For the rate of exchanged nucleons we can readily write

$$\dot{N} = \sum_{ij} w_{ij}(t) F_j(t) [1 - F_i(t)] \quad (2.3)$$

where the prime indicates that the sum is taken only over states i, j belonging to different fragments. Combining (2.2) and (2.3) the mean

energy loss per transferred nucleon is given by

$$\frac{dE}{dN} = \frac{\sum_{ij} (\varepsilon_i - \varepsilon_j) w_{ij}(t) F_j(t) [1 - F_i(t)]}{\sum_{ij} w_{ij}(t) F_j(t) [1 - F_i(t)]}. \quad (2.4)$$

This expression can be discussed qualitatively by referring to fig. 1. Figure 1a shows in a schematic way all types of transitions (excitations with and without transfer) which are taken into account in (2.4). The effect of single-particle shell structure is most easily understood by regarding two magic fragments (cf. fig. 1b). It is seen that with each transition an extra energy roughly equal to the gap energy $\Delta \varepsilon_{\text{shell}}$ is needed. Therefore the rate (2.4) for magic fragments should be roughly given by dE/dN for non-magic fragments plus $\Delta \varepsilon_{\text{shell}}$.

2.1. The energy loss rate for grazing collisions

In the further evaluation of (2.4) we first restrict to grazing collisions where the situation is most clean for the discussion of shell-structure effects. In this case the allowed energy transfer $\Delta_{gr} = \hbar^2 l_{gr} \Delta j / I_{rel}$ is determined by the uncertainty Δj of angular momentum in a single-particle transition via angular momentum and energy conservation for intrinsic and relative motion. Here l_{gr} and I_{rel} denote the grazing angular momentum and the relative moment of inertia, respectively. An estimate of Δj from the overlap of the fragments results in (ref. 13, p. 253)

$$\Delta_{gr}/\text{MeV} = \frac{13.5}{R/\text{fm}} \left[\frac{E_{cm} - V(R)}{A_{red} \cdot \text{MeV}} \right]^{1/2} \cdot (A_1 + A_2)^{1/6} \quad (2.5)$$

for grazing collisions. In accordance with ¹⁴⁾ we write the transition probabilities in the form

$$w_{ij}(t) = \chi_{ij}(t) G_{ij} \quad (2.6)$$

where the time-independent Gauss function

$$G_{ij} = \frac{1}{\sqrt{2\pi} \cdot \Delta q} \exp \left[-\frac{(\epsilon_i - \epsilon_j)^2}{2 \Delta q^2} \right] \quad (2.7)$$

takes care of the range of allowed energy change in each transition. Replacing the magnitudes $\chi_{ij}(t)$ of the transition probabilities by the mean value $\bar{\chi}(t)$, thereby assuming that the fluctuating part $\chi_{ij}(t) - \bar{\chi}(t)$ gives a negligible contribution in the sums of (2.4), we find

$$\frac{dE}{dN} = \frac{\sum_{ij} (\epsilon_i - \epsilon_j) G_{ij} F_j(t) [1 - F_i(t)]}{\sum_{ij} G_{ij} F_j(t) [1 - F_i(t)]} \quad (2.8)$$

We note that $\bar{\chi}(t)$ does not enter the expression for dE/dN and hence, may have an arbitrary time dependence which is determined by the trajectory of relative motion. Since the mean matrix elements for excitations and transfers are not very different we have taken them to be the same (cf. refs. 13, 15). A generalisation for $\chi_{ex} \neq \chi_{tr}$ is straightforward but is not pursued here.

The occupation probabilities are assumed to be Fermi functions

$$F_i(t) = \left[1 + \exp \left(\frac{\epsilon_i - \bar{\epsilon}_F^{(v)}}{T_v} \right) \right]^{-1} \quad (2.9)$$

where v denotes both the isospin and the fragment according to whether i is a proton or neutron state in fragment 1 or 2. Thus we have four different Fermi energies $\bar{\epsilon}_F^{(v)}$ and temperatures T_v . For symmetric

systems the differences between the Fermi energies of the fragments vanish. For asymmetric systems these differences are less than about 1 MeV (only for Fe+Bi about 3 MeV) and are expected to become negligible after the transfer of one or two nucleons (charge equilibration). We

have checked that a difference of 3 MeV produces an error of less than 15% in the ratio (2.8).

The interaction between the fragments gives rise to shifts of the single-particle levels. This results in particle-hole excitations and hence, to an occupation probability which deviates in general from the $T = 0$ Fermi distribution. In order to take this smearing of the Fermi surface into account we introduce an effective temperature T_{eff} which is determined from the equivalent shifts of levels due to an increase of the deformation parameter β by 0.15. The parameter β describes the quadrupole deformation of the nuclear surface by $R_{surface} = R_0 [1 + \beta Y_{20}(\theta, \varphi)]$. In all cases where no level crossings occur, $T_{eff} = 0$, i.e. essentially for closed shells. In case of level crossings we find $T_{eff}^{(v)} \approx 1.5$ MeV with large fluctuations for different nuclei. A unique value of 1.5 MeV is used, however, in the calculations for all these situations.

With the specifications discussed so far the initial energy loss per transferred nucleon can be calculated according to (2.8). In evaluating the sums over the single-particle states i of the incident nuclei realistic levels ϵ_i are used ¹⁵⁾.

2.2. Extension to more central collisions

In extending the preceding discussion to more central collisions we

have to consider three effects which all reduce dE/dN :

- (i) The temperature becomes time dependent. For a given total kinetic energy loss TKEL we introduce a mean temperature by

$$T(\text{TKEL}) = \sqrt{\frac{\text{TKEL}}{2a} + T_{\text{eff}}^2} \quad (2.10)$$

with the level density parameter $a = (A_1 + A_2)/8$.

- (ii) The allowed energy change Δ decreases with decreasing angular momentum. Defining the mean angular momentum by $\bar{l} = (l_i + l_f)/2$, where l_i and l_f denote the initial and final angular momentum, we obtain (cf. eq. (4.54) of ref. 13).

$$\Delta = \Delta_{\text{gr}} \bar{l}/l_{\text{gr}} \quad (2.11)$$

with Δ_{gr} given by (2.5).

- (iii) With these modified expressions for T and Δ we are able to calculate a mean energy loss rate $\langle dE/dN \rangle_{\text{TKEL}}$ for a given total kinetic energy loss. However, this quantity has been derived under the assumption that the reservoir of kinetic energy is infinite in the relative motion. In order to take into account that the energy loss vanishes if all energy is dissipated we introduce an explicit decay by

$$dE/dN = \langle dE/dN \rangle_{\text{TKEL}} \cdot \exp(-t/\tau). \quad (2.12)$$

The observed energy loss rate is the average over the interaction time τ_{int}

$$\langle dE/dN \rangle_{\tau_{\text{int}}} = \langle dE/dN \rangle_{\text{TKEL}} \left[1 - \exp(-\tau_{\text{int}}/\tau) \right] \tau / \tau_{\text{int}} \quad (2.13)$$

where the decay time τ is identified with the decay time of the angular momentum dissipation τ_{ang} .

For the evaluation of the expressions (2.11) and (2.13) we need the interaction times τ_{int} and the relaxation times of energy and angular momentum. The phenomenological model of Riedel et al.¹⁷⁾, as described in ¹³⁾, p.222-223, is used to obtain these quantities. We note that the use of this model which does not contain any shell effects is not consistent with the model discussed here. Moreover, our treatment is well justified only for the more peripheral collisions. However, we feel that (2.13) still contains the major shell effect through the factor $\langle dE/dN \rangle_{\text{TKEL}}$ also for the more central collisions.

3. Comparison to the experimental results

The calculations were performed for 21 systems, for which the experimental data on σ^2_A or σ^2_Z as functions of TKEL were available. For a comparison of the calculated dE/dN functions to the experimental data the number of exchanged nucleons is obtained for each energy loss bin from the measured variance of the mass distribution by taking

$$N = \sigma^2_A. \quad (3.1)$$

This relation is valid under the assumption of a random walk process for the nucleon transfer. A Monte Carlo calculation ¹⁸⁾ shows that this equation is valid even for a narrow valley along the stability line in the (N,Z)-plane.

A comparison of the calculations with the experimental results

becomes more ambiguous for reactions where only charge distributions were measured. In such cases the relation $\sigma_A^2 = (A/Z)^2 \sigma_Z^2$ was assumed. This can be considered as an acceptable approximation, since we are interested only in the gross properties of the energy loss rate. The calculations were performed only for those systems, where the initial c.m. kinetic energy lies at least 15% above the Coulomb barrier. This condition eliminates those reactions, where the neutron flow (which could dominate at incident energies just above the Coulomb barrier) could make the $(A/Z)^2$ assumption unjustified. In fact for low energy losses the experimental σ_A^2/σ_Z^2 values have been seen to lie, for some systems below the $(A/Z)^2$ line ^{19,20}, for other systems above it ^{21,22}. Generally for large energy-losses the σ_A^2/σ_Z^2 values tend to reach a plateau, usually near the $(A/Z)^2$ -value. As a consequence the correlations presented in fig. 2 are affected by the $(A/Z)^2$ approximation which can cause some scattering of points in the case of the σ_Z^2 data.

3.1. Initial energy loss rate

Figure 2 presents a comparison between the theoretical values of $\langle dE/dN \rangle^0$ (abscissa) and the experimental values of $\langle dE/d(\sigma_A^2) \rangle^0$ extracted for the first 25 MeV of TKEL. The agreement of the calculations with the experimental $\langle dE/d(\sigma_A^2) \rangle^0$ values, emerging on an absolute scale in fig. 2, is quite striking (differences $\leq 20\%$) and can be considered as a valuable result in determining the nature of the microscopic mechanism responsible for the energy loss of dissipative collisions. This result demonstrates that the single-particle level

structure plays a dominant role in the energy dissipation process.

3.2. Comparisons with σ_A^2 vs. TKEL correlations

Using relation (3.1) and the calculated dE/dN dependence on TKEL one obtains

$$\sigma_A^2 = \int_0^{TKEL} \frac{dE'}{(dE/dN)} \quad (3.2)$$

A comparison between the σ_Z^2 vs. TKEL correlation calculated on the basis of (3.2) and the experimental data is presented in fig. 3 for the systems Pb+Pb as well as U+U (σ_Z^2 data from refs. 7 and 10), and on

fig. 4 for two Sm+Sm systems (σ_A^2 data from refs. 11 and 20). The ¹⁴⁴Sm+¹⁴⁴Sm system has a magic number of neutrons, while in the ¹⁵⁴Sm+¹⁵⁴Sm system both protons and neutron shells are half-filled. In fig. 5 the calculations are compared with the experimental data obtained for the bombardement of three targets (Fe, Ho, Bi) with 8.3 MeV/u Fe-ions ¹⁹. Similar comparisons are made in fig. 6 for ¹³⁹La+⁸⁶Kr ²⁴) and in fig. 7 for ⁵⁶Fe+¹³⁶Xe and ¹²⁴Sn+¹²⁹Xe (data from ref. 21).

For most of the investigated systems the ratio between the calculated energy loss rates for protons $(dE/dN)_p$ and the one for neutrons $(dE/dN)_n$ is very near or equal to one (e.g. for Pb+Pb, Pb+U, U+U, Er+Kr, Fe+Bi, Fe+Ho, Fe+Fe etc.). However those systems, where either the proton or the neutron shells are closed, are characterised by large differences in the energy loss rates for protons and neutrons, respectively. For the system ¹⁴⁴Sm+¹⁴⁴Sm e.g. the calculated energy loss rate

for neutrons is, in the quasi-elastic region, about three times higher than that for protons. For the $^{12}\text{Sn}+^{12}\text{Sn}$ system the closed proton shell reverses the situation. Figure 8 summarises calculations for several such reactions. The striking similarity of these calculated curves with the σ_A^2/σ_Z^2 values measured for the systems $^{144}\text{Sn}+^{144}\text{Sm}$ 20 and $^{124}\text{Sn}+^{129}\text{Xe}$ 21 suggests that the σ_A^2/σ_Z^2 values are also strongly influenced by the shell structure.

3.3. "Recoil" and "structure" energy

In a systematic comparison among the available experimental TKEL vs. σ_Z^2 correlations it was found empirically 12 that the average energy loss rate expressed in terms of $dE/d(\sigma_A^2)$ can be generally approximated by the expression:

$$-dE/d(\sigma_A^2) = E_{av}/\mu + DE(\text{TKEL}). \quad (3.3)$$

In such an ansatz the first term E_{av}/μ (denominated recoil energy) depends only on the relative energy, while the structure energy $DE(\text{TKEL})$ is independent of the bombarding energy. In fig. 9a calculated values for dE/dN are presented for the Kr+Fr reaction measured over a wide range of incident energies namely at 6.0, 7.2, 8.2 and 12.1 MeV/u 12,25 . In fig. 9b the theoretical DE values deduced according to (3.3) are compared with the corresponding experimental values. The agreement gives a theoretical justification to the phenomenological ansatz (3.3).

A comparison between the calculated and measured DE values is presented in fig. 10a for the Pb+Pb and U+U reactions 7,19 , and in fig. 10b for

the Bi+Xe reaction 8,9 . Fig. 11 presents again as in fig. 2 the initial energy loss rates but after subtraction of the recoil term. In this way the dependence on the incident energy is largely eliminated and the structure dependence is isolated: the systems ($^{208}\text{Pb}+^{208}\text{Pb}$, $^{136}\text{Xe}+^{208}\text{Bi}$) with all four shells closed (or almost closed) display large energy-loss rates, while systems with only half-closed shells (U+U , $^{154}\text{Sm}+^{154}\text{Sm}$) are grouped in the low-energy-loss region.

4. Conclusions

Within a statistical model we have examined the shell-structure effects on the mean energy dissipated per transferred nucleon. The observed effects are quantitatively related to the shell structure of the interacting nuclei. The excellent agreement between experimental and calculated values supports also the underlying statistical picture of energy dissipation:

(i) Pure excitations as well as single-particle transfers contribute about equal amounts to the observed energy loss rate.

(ii) The energy loss in a given transition is determined by the occupation probabilities of initial and final states and by the uncertainty in the momentum (angular momentum) which is exchanged between relative and intrinsic motion.

Both points are in conflict with the assumption of the one-body dissipation mechanism where only the recoil energy of the transferred nucleons contributes to the energy loss.

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Figure captions

Figure 1. Schematic picture of possible transitions for two interacting nuclei (a) without and (b) with shell structure.

Figure 2. Correlation between the theoretical $\langle dE/dN \rangle^0$ values (abscissa) and the $\langle dE/d(\sigma_A^2) \rangle^0$ values extracted from experiments.

Both values are obtained from an averaging over the first 25 MeV energy bin of TKEL. The straight line indicates complete agreement between experimental and calculated values. The squares represent systems for which the experimental results were extracted using mass widths. The rest of the data represents those systems, where only charge distributions were experimentally available.

Figure 3. Comparison of the calculated and measured correlations

between σ_Z^2 and TKEL: (a) for PbPb⁷), (b) for U+U¹⁰).

Figure 4. Comparison of the calculated and measured correlations between σ_A^2 and TKEL for two Samarium systems:

(a) ¹⁴⁴Sm+¹⁴⁴Sm (beam energy 6.95 MeV/u, ²⁰)

(b) ¹⁵⁴Sm+¹⁵⁴Sm (beam energy 6.3 MeV/u, ¹¹)

Figure 5. Comparison of the calculated and measured correlations between σ_A^2 and TKEL for targets of Fe, Ho, Bi bombarded with Fe ions of 8.3 MeV/u¹⁹).

Figure 6. Comparison of the calculated and measured correlations between σ_Z^2 and TKEL for the ¹³⁹La+⁸⁶Kr reaction²⁴)

Figure 7. Comparison of the calculated and measured correlations between σ_Z^2 and TKEL for $^{56}\text{Fe}+^{136}\text{Xe}$ and $^{124}\text{Sn}+^{129}\text{Xe}$ (data from ref.21).

Figure 8. Calculated ratio of $(dE/dN)_p / (dE/dN)_n$ for selected systems. In brackets beam energies in MeV/u.

Figure 9. a) Calculated dE/dN as function of TKEL for different beam energies in the case of Kr+Er.

b) Comparison of the calculated "structure term" DE with the experimental ones 12,25 . The much weaker beam-energy dependence of calculated DE function in comparison with dE/dN (Fig.9a) should be noted. The experimental σ_Z^2 vs. TKEL correlation for 12.1 MeV/u agrees well 25 with the assumption of independence of DE from the beam energy.

Figure 10. Comparison of DE curves calculated and extracted from the measurements: (a) for Pb+Pb reaction and for U+U reaction 7,19 , (b) for Bi+Xe reactions 8,9 .

Figure 11. Correlation between theoretical and experimental $(DE)^0$ values (from the initial energy loss rates the recoil term is subtracted). Both values are the averaged values for the first 25 MeV energy bin in TKEL. The straight line indicates complete agreement between experimental and calculated values.

Fig. 1

Fig. 4

Fig. 5

Fig. 6

Fig. 7

Fig. 8

Fig. 9

Fig. 10

Fig. 11