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Wolfgang Nörenberg

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Gesellschaft für Schwerionenforschung mbH
Planckstr. 1 · Postfach 1105 41 · D-6100 Darmstadt 11 · Germany

DISSIPATIVE DIABATIC DYNAMICS

Wolfgang NÖRENBERG

Gesellschaft für Schwerionenforschung (GSI)
D-6100 Darmstadt, Germany

Abstract: A time-dependent shell-model approach to large-amplitude collective nuclear motion is introduced which simultaneously treats the coherent (diabatic) single-particle motion in the time-dependent mean field and the incoherent two-body collisions which lead to statistical equilibrium within the intrinsic degrees of freedom. An important feature of the combined effects from diabatic single-particle motion and dissipative two-body collisions is the retardation of dissipation for realistic values of the mean free path of nucleons. This elastoplastic behaviour represents an interesting dynamical property of nuclear matter. A diabatic two-center shell-model is introduced and results on diabatic single-particle levels and potentials are given and compared with corresponding quantities from TDHF calculations. Effects on cross-sections of heavy-ion collisions are discussed and preliminary results on the diabatic pre-equilibrium emission of nucleons in central collisions are presented.

1. Introduction

In low-energy heavy-ion collisions the velocity of the mean field is small as compared to the velocities of individual nucleons within the mean field. This natural separation of collective and intrinsic degrees of freedom is the basis for any description of such large-amplitude collective nuclear motions. However, different assumptions about the role of two-body collisions between nucleons lead to contradictory theoretical concepts. Current transport theories, for a review cf. ref.¹⁾, like the linear response theory²⁾, the one-body dissipation theory³⁾ and the (non-perturbative) random matrix theories^{1,4-6)} imply statistical equilibrium within the intrinsic degrees of freedom ('local' equilibrium) throughout the process. As a consequence the incoherent coupling between collective and intrinsic degrees of freedom is treated by time-irreversible transport equations. Whereas the instantaneous local equilibrium requires a very short mean free path of nucleons, the time-dependent Hartree-Fock theory (TDHF) is based on the assumption of a very large mean free path⁷⁾. Here, the coupling between the mean field and the single-particle motion is treated quantum-mechanically and hence, in a coherent time-reversible way.

The intention for describing both the coherent and incoherent aspects have recently initiated strong activities⁸⁻¹⁸⁾. The inclusion of residual two-body collisions in extended time-dependent Hartree-Fock theories (ETDHF) is a natural step towards a more realistic description⁸⁾. However, because of numerical difficulties, realistic applications of these approaches are impossible with even the fastest available computers. An alternative model has been introduced by the concept of dissipative diabatic dynamics (DDD)¹²⁻¹⁸⁾. Giving up self-consistency the main coherent aspects of the coupling between the collective and the single-particle motion is treated within the diabatic approximation for a time-dependent shell model. Simultaneously, the two-body collisions are treated as the main incoherent processes which result in dissipation.

The aim of this paper is to describe the basic concept of DDD (cf. §2) and to draw the attention to the elastoplastic behaviour of nuclear matter (cf. §3) which is due to the interplay of the coherent coupling between collective and intrinsic degrees of freedom. Model calculations, comparison with TDHF results and applications are discussed in §4.

2. Basic concept

The basic idea of DDD is to consider simultaneously the coherent coupling of the single-particle motion to the time-dependent mean field (§2.1) and the incoherence introduced by the residual two-body collisions (§2.2). As a consequence of the interplay between these coherent and incoherent aspects of the coupling between collective and intrinsic degrees of freedom, the effective equation of motion for the collective degrees contains a time-retardation in the dissipative term (retarded friction, §2.3).

2.1 DIABATIC SINGLE-PARTICLE MOTION

We consider slow collective motions such that the Fermi velocity v_F is large compared to the collective velocity v_{coll} . In this limit the individual nucleons can essentially adjust their wave functions to the change of the mean field in a time-reversible manner. A particular problem, studied by Landau, Zener and Stueckelberg^{19,20} is encountered at quasi-crossings of adiabatic levels.

The essential points are illustrated in fig. 1. Consider a single-particle Hamiltonian $H_0(q)$ with some symmetries (e.g. axial and reflection symmetry). Levels with different quantum numbers may cross as function of q . These levels are called *diabatic* (greek: crossing) levels. The total Hamiltonian $H(q)$ differs in general from $H_0(q)$ by some symmetry-breaking term H' which leads to a coupling between the diabatic states. The diagonalization of the total Hamiltonian results in

Fig. 1: Diabatic and adiabatic levels in the Landau-Zener-Stueckelberg (quasi-) crossing. Consider for example the quadrupole deformation q of a nucleus: Then, the diabatic level rising (descending) with increasing q corresponds to a wave function with a larger (smaller) number of nodes along that axis which is squeezed. In the diabatic motion the nucleon keeps the nodal structure of the wave function.

adiabatic (i.e. non-crossing) levels with wave functions changing drastically their (approximate) symmetry properties. For large enough collective velocities the nucleons are not able to adjust their wave functions adiabatically and follow the adiabatic level. Instead, a nucleon occupying the lower level before the crossing with an unoccupied level, will stay (while preserving the symmetry of its wave function) on the diabatic level and finds itself in the upper level after the crossing. This diabatic motion leads to the excitation of a particle-hole state. The jump probability (i.e. the degree of diabaticity) is determined by

$$J = \exp \left\{ -\frac{1}{\Delta} \right\} \quad (2.1)$$

with the diabaticity parameter

$$\Delta = \frac{\hbar |\partial(\epsilon_1 - \epsilon_2)/\partial q| |\dot{q}|}{2\pi |H_{12}^1|^2} \quad (2.2)$$

which is proportional to the difference in the slopes of the diabatic levels and to the collective velocity, and inverse proportional to the absolute square of the coupling between the diabatic states. For $\Delta \gg 1$ we have $J=1$, i.e. the diabatic limit; and for $\Delta \ll 1$, $J=0$ the adiabatic limit. In order to have a jump probability $J > 0.5$ the relative kinetic energy per particle E_{rel}/A_{red} (with $A_{red} = A_1 A_2 / (A_1 + A_2)$) needs to be larger than 0.04 MeV. In obtaining this estimate, the typical values $|H_{12}^1| = 0.5$ MeV and $\partial \epsilon_{12}/\partial q = 1$ MeV/fm have been used¹⁸). Combining this estimate with the condition for slow collective motion ($v_{rel} \ll v_F$) we find the energy range

$$0.04 \text{ MeV} \ll E_{rel}/A_{red} \ll 40 \text{ MeV} \quad (2.3)$$

where the single-particle motion is diabatic. The construction of diabatic states is discussed in § 4.

2.2 LOCAL EQUILIBRATION BY TWO-BODY COLLISIONS

Residual two-body collisions limit the single-particle picture discussed in the preceding § 2.1. The *mean free path* of a nucleon (a hole for $\epsilon < \epsilon_F$) with energy ϵ close to the Fermi energy ϵ_F and for low nuclear temperatures $T \ll \epsilon_F$, is approximately given by²¹⁾

$$\lambda = \alpha \frac{1 + \exp(-|\epsilon - \epsilon_F|/T)}{(\epsilon - \epsilon_F)^2 + \pi^2 T^2} \quad (2.4)$$

The factor α is ≈ 540 fm MeV² when estimated from the Fermi gas model²²⁾ and ≈ 850 fm MeV² when estimated from the imaginary part of the optical potential (cf. fig. 23 of ref. ²³). With an average value of $\alpha \approx 700$ fm MeV² we expect (cf. fig. 2) the mean free path λ to vary from large values at the initial stage of nucleus-nucleus collisions to ≈ 15 fm at the completely damped stage ($T \approx 2.5$ MeV).

Assuming *three* subsequent collisions to be sufficient for establishing intrinsic equilibrium for fixed values of the collective variables (*local* or *intrinsic* equilibrium), we have

$$\tau_{loc} \approx 3\lambda/v_F, \quad (2.5)$$

for the *local equilibration time* (v_F = Fermi velocity). A direct estimate for the local equilibration time has been given by Bertsch²⁴⁾. He considers the equilibration of the momentum distribution within a Fermi gas at finite temperatures with an initial quadrupole deformation and finds numerically

$$\tau_{loc} \approx 3 \cdot 10^{-22} \text{ s} / \epsilon^* \quad (2.6)$$

Here, ϵ^* denotes the total excitation energy per particle in MeV stored as heat and deformation energy of the Fermi sphere. The constant $3 \cdot 10^{-22}$ s corresponds to

Fig. 2 Mean free paths of nucleons and local equilibration times as calculated from eqs. (2.4) and (2.5). The hatched region indicates the mean free paths and local equilibration times during a nucleus-nucleus collision. The dashed curve is an estimate by Bertsch²⁴⁾ and is discussed in the text.

$\alpha = 700 \text{ fm MeV}^2$ in eq. (2.4). Using $T = (8\epsilon^*)^{1/2}$ the relation (2.6) is illustrated by the dashed curve in fig. 2 which compares quite well with the estimate from the mean free path.

For typical temperatures around 2.5 MeV in the completely damped stage we find $\tau_{\text{loc}} \approx 0.5 \cdot 10^{-21} \text{ s}$. In the completely damped stage of the process the collective velocities might be small enough such that the system is always close to local equilibrium and hence, that current transport theories are applicable. However, at the initial stage of heavy-ion collisions and for fission at small and intermediate energies the local equilibration time is considerably larger (of the order of a few 10^{-21} s). Therefore we expect large *memory effects* in the approach phase of nucleus-nucleus collisions (duration about $0.3 \cdot 10^{-21} \text{ s}$).

2.3 RETARDED FRICTION

We have discussed in the preceding §§ 2.1 and 2.2 the main coherent and incoherent processes in large-amplitude collective nuclear motion: the diabatic single-particle motion in the time-dependent mean field and the residual two-body collisions. Starting with some equilibrium distribution for the single-particle occupation probabilities (e.g. corresponding to the ground state if we consider the initial stage of a nucleus-nucleus collision) the diabatic excitation of particle-hole states (cf. fig. 3) will produce a repulsive force on the collective motion. Thus, kinetic energy is temporarily stored as a conservative potential. However, the two-body collisions try to establish a new equilibrium distribution for the occupation probabilities and thereby destroy the diabatic part of the potential. The local equilibration via two-body collisions is a time-irreversible process which leads to the dissipation of collective energy. We have

Fig. 3: Illustration of diabatic single-particle motion and dissipative two-body collisions as the basic elements of dissipative diabatic dynamics (DDD).

called this dynamical process *dissipative diabatic dynamics* (DDD) because it is characterized simultaneously by diabatic single-particle motion and dissipative two-body collisions.

The temporary storage of collective kinetic energy as a conservative potential and the subsequent dissipation by (one- and) two-body collisions gives rise to memory effects. These memory effects can be illustrated by approximating the two-body collision term by a simple relaxation term as

$$\frac{\partial n_\alpha}{\partial t} = \frac{1}{\tau_{loc}(t)} (n_\alpha - \bar{n}_\alpha) \quad (2.7)$$

which describes the approach of the occupation probabilities n_α of diabatic states α towards the local equilibrium distribution $\bar{n}_\alpha$. According to eq. (2.6) the local equilibration time depends on time via the time dependence of the total excitation energy. Temperature and chemical potential of $\bar{n}_\alpha$ are determined by the conservation of total excitation energy and total mass. Essentially due to the q -dependence of the single-particle energies, $\bar{n}_\alpha$ is strongly varying with $q(t)$.

In the single-particle picture the force on the collective variable is determined by

$$F(q, t) = - \sum_\alpha \frac{d\epsilon_\alpha(q)}{dq} n_\alpha(t) . \quad (2.8)$$

With the relaxation ansatz (2.7) for $n_\alpha(t)$ we obtain

$$\frac{d(F - \bar{F})}{dt} = - \frac{1}{\tau_{loc}(t)} (F - \bar{F}) - C\dot{q} \quad (2.9)$$

where $d^2\epsilon_\alpha/d^2q$ has been neglected for convenience. The equilibrium force $\bar{F}$ is determined by eq. (2.8) with n_α replaced by $\bar{n}_\alpha$. The quantity.

$$C = - \sum_{\alpha} \frac{d\epsilon_{\alpha}}{dq} \frac{d\bar{n}_{\alpha}}{dq} = \sum_{\alpha} \left[\frac{d(\epsilon_{\alpha} - \mu)}{dq} \right]^2 \left| \frac{d\bar{n}_{\alpha}}{d(\epsilon_{\alpha} - \mu)} \right| \quad (2.10)$$

is positive definite whenever the single-particle energies depend on q . Note that $(d/dq)\sum_{\alpha} n_{\alpha} = 0$. The first term on the r.h.s. of eq. (2.9) describes the decay rate of the diabatic force $F - \bar{F}$. The second term is a source term resulting from the diabatic excitation of particle-hole states. It is always repulsive since $C > 0$. Integration of eq. (2.9) from t_0 to t with $F(t_0) = \bar{F}(t_0)$ yields the force

$$F(t) = \bar{F}(t) + \int_{t_0}^t dt' \exp\left[-\int_{t'}^t dt''/\tau_{loc}(t'')\right] C(t') \dot{q}(t') \quad (2.11)$$

which deviates from the equilibrium force $\bar{F}$ by a *retarded friction* term. The collective equation of motion is given by

$$B\ddot{q} = F \quad (2.12)$$

with the mass parameter B which can be calculated in the cranking approximation for the diabatic single-particle motion and therefore is expected to be close to the irrotational flow value²⁰⁾.

Two limits can be considered for the retarded friction term:

(i) *Elastic limit* ($\tau_{loc} \rightarrow \infty$): In this limit the exponential in eq. (2.11) is 1. Neglecting the time dependence of C , eq. (2.12) describes for small amplitudes a harmonic oscillation,

$$B\ddot{q} - C(q - q_0) = 0 \quad (2.13)$$

with frequency $\omega_0 = (C/B)^{1/2} \approx 70 A^{-1/3} \text{ MeV}/\hbar$ for the quadrupole motion of nuclei. This pure diabatic vibrational mode is readily identified with the isoscalar giant quadrupole vibration^{12,25)}. Therefore the use of the giant vibrations of the colliding nuclei by Broglia et al²⁶⁾ as agents for the dissipation of energy in heavy-ion collisions can be regarded as an example of the dissipation mechanism described here. However, our concept goes beyond²⁶⁾ because it is not restricted to oscillations around the equilibrium shapes of the colliding nuclei.

(ii) *Dissipative limit* ($\tau_{loc} \ll \tau_{coll}$): This is the limit of very slow collective motion (τ_{coll} characteristic time of mean field changes) where $\dot{q}(t')$ and $C(t')$ can be replaced by $\dot{q}(t)$ and $C(t)$. Consequently we obtain

$$B\ddot{q} = \bar{F}(q) + F_{friction}(q, t) \quad (2.14)$$

with a *conventional friction force* given by

$$F_{friction}(q, t) = -C\dot{q} \int_{t_0}^t dt' \exp\left[-\int_{t'}^t dt''/\tau_{loc}(t'')\right] \quad (2.15)$$

$$\approx -C\tau_{loc} \dot{q} \quad (2.15a)$$

the last approximation being satisfied whenever τ_{loc} does not change appreciably during time intervals of the order τ_{loc} . Since the friction coefficient is proportional to the local equilibration time, i.e. proportional to the mean free path (2.4) its temperature dependence becomes proportional to T^{-2} for n_{α} close to the equilibrium distribution. Thus the friction coefficient has the same temperature dependence as the viscosity coefficient for large systems. From the study of an interacting Fermi gas within moving walls it has been shown¹²⁾ that the friction coefficient $\xi = C\tau_{loc}$ becomes equal to the wall friction coefficient ξ_w ³⁾ if τ_{loc} is replaced by the time of flight τ_F of the nucleons through the container. In general our friction coefficient is larger than ξ_w because $\tau_{loc} > \tau_F$.

Fig. 4: Time dependence of a quadrupole vibration (collective variable δ) within an oscillator model with $\tau_{loc} = 10^{-21}s$. The fast giant vibrations are damped on the time scale of τ_{loc} whereas the system relaxes very slowly to the equilibrium point ($\delta = 0$). From ref. ¹⁵⁾ (1bs = $10^{-23}s$).

The concept of DDD as discussed so far has been discussed in detail for the interacting Fermi gas within moving walls¹²⁾ and for the oscillator model¹⁵⁾. The qualitative features of the dissipation process as obtained within the oscillator model¹⁵⁾ are discussed in fig. 4. A simple classical analog of DDD is illustrated in fig. 5. Also in ref. ¹⁵⁾ we have treated the fluctuations in the collective velocities. The main difference to Brownian motion is that in DDD the fluctuations are produced by the stochastic changes in the occupation probabilities and hence, in the forces or accelerations and not in the velocities.

A general transport theory based on the diabatic single-particle motion has been formulated recently ¹⁶⁾. In the derivation of transport equations time averaging has been introduced as a natural procedure to apply on the macroscopic or collective motion which by definition should be the slow modes of the system. In connection with the time-dependence of the single-particle levels, the time average introduces irreversibility in the equations of macroscopic motions and justifies the use of the weak-coupling limit and the Markov approximation in the basic collision terms. These collision terms conserve not only the total occupation probability (i.e. conservation of total mass and charge) but also the total energy. This time-dependent shell-model theory with residual interactions can be regarded as a simplified version of extended TDHF theories. However, apart from two-body collisions, we also describe some couplings of the one-body type by irreversible collision terms. At the cost of selfconsistency and with the

Fig. 5: Classical analog of the dissipation mechanism: The local equilibration time is determined by the time the acid needs for dissolving the spring. Whenever the interaction time (between spring and acid) is small compared to τ_{loc} the system moves elastically, whereas it becomes dissipative for $\tau_{loc} \lesssim \tau_{int}$. The energy of elastic deformation of the spring is transformed into heat.

constraint of the collective motion to a single or a few degrees of freedom, the formulation looks more transparent and can easily supply transport coefficients which may be used in transport equations. Thus this shell-model approach to the transport theory of large amplitude collective motion of nuclei seems to be promising for future applications.

3. Elastoplastic behaviour of nuclei and phase diagram of nuclear dynamics

As discussed in §2 collective nuclear dynamics has a coherent and an incoherent aspect. We characterize the coherent coupling between the mean field and the single-particle motion by the mean diabaticity coefficient $\langle \Delta \rangle$ (cf. the definition by eq. (2.2)) and the rate of dissipation by the ratio τ_{loc}/τ_{coll} of the local equilibration time over the collective time which is typical for changes in the mean field (collective variables). The two parameters, i.e. the mean diabaticity and the degree τ_{loc}/τ_{coll} of dissipation allow to define a *phase diagram of nuclear dynamics* shown in fig. 6. As discussed in §2.1 the single-particle motion is essentially diabatic for $\langle \Delta \rangle \gg 1$ and adiabatic for $\langle \Delta \rangle \ll 1$. For $\tau_{loc}/\tau_{coll} \gg 1$ the dissipation is negligible and hence, the motion is elastic, i.e. either diabatic or adiabatic. In between these two limits there is a transition region (hatched area) where the single-particle motion is neither diabatic nor adiabatic. For $\tau_{loc}/\tau_{coll} \approx 1$ dissipation is essential, leading to frictional adiabatic motion for $\langle \Delta \rangle \ll 1$ and dissipative diabatic motion for $\langle \Delta \rangle \gg 1$.

Fig. 6: Phase diagram of nuclear dynamics with regimes of qualitatively different dynamical properties which depend on the mean diabaticity $\langle \Delta \rangle$ of single-particle motion and the degree τ_{loc}/τ_{coll} of dissipation. Note in particular the elasto-plastic regime.

Finally for $\tau_{loc}/\tau_{coll} \ll 1$ the system is chaotic, i.e. always close to local equilibrium. In the chaotic regime the motion is determined by the available phase space and thus, it will hardly make sense to distinguish here between diabatic and adiabatic single-particle motion.

The phase diagram also demonstrates the richness of nuclear dynamics in analogy to properties of macroscopic systems. Depending on the dynamical situation, nuclear matter behaves like a super-fluid, a viscous fluid or a creepy fluid on one hand, and like an elastic solid, an *elasto-plastic* (Silly Putty) or a plastic substance on the other hand.

The richness of nuclear dynamics as shown in fig. 6 has been explored only partially. Induced and spontaneous fission is characterized by very slow collective motion such that the nucleus can adjust their wave functions adiabatically. Therefore, only the lower part of the phase diagram has been explored by fission processes. Giant vibrations lie in the interesting elasto-plastic regime. They are, however, limited to small amplitudes. Large-amplitude collective motions in this regime are supplied by nucleus-nucleus collisions. In general, a deeply inelastic collision passes from the upper right-hand corner of the phase diagram through the dissipative diabatic (elasto-plastic) regime to the chaotic region. In order to explore the interesting elasto-plastic features of nuclear matter we have to study thoroughly such nucleus-nucleus collisions.

Large-amplitude nuclear motion has been studied theoretically within various approximations. In TDHF⁷⁾ the residual two-body collisions are neglected. This is the limit of $\tau_{loc}/\tau_{coll} \gg 1$ and hence, TDHF is restricted to the right-hand part of the phase diagram (fig. 6). Linear response theory²⁾ starts from adiabatic single-particle motion which limits its applicability to the lower strip. The (non-perturbative) random-matrix theories^{1,4-6)} imply local equilibrium throughout the process and therefore, are restricted to the left strip ($\tau_{loc}/\tau_{coll} \ll 1$). Finally, our DDD (dissipative diabatic dynamics) is applicable in the upper part of the phase diagram i.e. in particular to the elasto-plastic behaviour of nuclei. We regard this region as a big challenge to experimentalists and theorists. In §§ 4.3 and 4.4 some measurable quantities of nucleus-nucleus collisions are discussed which may render important information on specific aspects of dissipative diabatic (elasto-plastic) dynamics.

4. Applications

In order to study the elasto-plastic properties of nuclei in more detail we have started numerical calculations in the framework of DDD. The construction and properties of diabatic single-particle states within a two-center shell-model is discussed in §4.1. We compare in §4.2 properties of these diabatic states with "stationary" states occurring in TDHF calculations. Effects of DDD on cross-sections of deeply inelastic collisions have been calculated (cf. §4.3). First results on the emission of neutrons by diabatic single-particle motion are given in §4.4.

Fig. 7: Adiabatic and diabatic levels with $\Omega = 1/2$ for $^{86}\text{Kr} + ^{166}\text{Er}$ as functions of the relative distance. The potentials include spin-orbit coupling. The diabatic states are constructed from the adiabatic states with stepwidth $\Delta q = 1$ fm starting from the right-hand side.

4.1 DIABATIC TWO-CENTER SHELL-MODEL¹⁸⁾

We use a modified version of the two-center shell-model of Maruhn and Greiner²⁷⁾ with two oscillator potentials joined smoothly by an inverted oscillator potential. Within this model we have studied two methods for the construction of diabatic states:

(i) The method of *maximum overlap* makes use of the large overlap of diabatic states at different values of the collective variable. The procedure is based on the adiabatic single-particle states generated in the two-center shell-model. The construction of the diabatic states as functions of the relative distance starts from separated nuclei where adiabatic and diabatic states coincide. The diabatic state at a somewhat smaller distance is obtained from a linear combination of a few adiabatic states maximizing the overlap with the original diabatic state. Using the new states one is able to construct step by step the diabatic states as functions of the distance. This procedure has been applied to the system $^{86}\text{Kr} + ^{166}\text{Er}$ and results are shown in fig. 7. The comparison of the adiabatic levels with the diabatic levels nicely demonstrates the change of quasicrossings to real crossings.

(ii) The method of *maximum symmetry* makes use of diabatic states which are constructed from a potential with maximum symmetry compatible with the collective motion. For the relative motion we expand the single-particle states for the separated nuclei in terms of such maximum symmetry states. Keeping the expansion coefficients fixed as functions of the distance we obtain diabatic states with the desired property that their wave functions adjust to the shape of the potential by a minimum change in their structure. Diabatic levels obtained with this method agree with those from the maximum overlap construction (for sufficiently large steps). In the following calculation we have used the method of maximum symmetry because it turns out to be numerically easier to handle.

Fig. 8: Diabatic neutron levels near the Fermi level (denoted by 0) and the diabatic part ΔV_{diab} of the nucleus-nucleus potential for $^{86}\text{Kr} + ^{166}\text{Er}$. The solid levels correspond to Er asymptotically, the dashed levels to Kr. The total diabatic potential (solid line) consists of contributions from neutrons (dashed line) and protons (dotted line). Only levels which cross the Fermi level are shown.

Figure 8 shows the diabatic levels in the vicinity of the Fermi energy and the corresponding *diabatic potential* ΔV_{diab} which is produced by keeping the occupation probabilities of the diabatic levels fixed. The graphs show that a considerable amount of collective kinetic energy can be stored in the diabatic excitation of particle-hole states leading to a strong repulsive force during the approach of the nuclei (cf. §4.3). Another interesting aspect of the diabatic single-particle motion is the lifting of the levels into the continuum such that nucleons can be emitted (cf. 4.4).

4.2 DIABATIC FEATURES IN TDHF RESULTS²⁸⁾

The main argument behind the concept of DDD is the expectation that the single-particle dynamics is essentially taken account by the diabatic basis states. Therefore, it is interesting to ask whether TDHF results for nucleus-nucleus collisions show diabatic features. If TDHF dynamics would be diabatic, the corresponding single-particle wave functions ψ_v should differ from stationary diabatic states only by a phase factor of collective motion²⁹⁾. We have extracted single-particle energies from the single-particle wave functions by $\epsilon_v(t) = i\hbar \langle \psi_v | \partial/\partial t | \psi_v \rangle$. The "stationary" single-particle energies are obtained by subtracting from $\epsilon_v(t)$ the collective kinetic energy of relative motion (corresponding to the additional collective phase factor). Here, the required relative velocity $\dot{R}(t)$ is obtained from the change of the density distribution. The resulting single-particle levels are shown in fig. 9 for two initial relative kinetic energies $^{86}\text{Kr} + ^{166}\text{Er}$ at 15 and 20 MeV/u. The levels coincide for two energies in general within less than 0.5 MeV. A complete coincidence would be expected for true diabatic motion if the subtraction of the collective energy is exact.

The diagram of fig. 9, however, should be considered with some care. As is

Fig. 9: "Diabatic" neutron levels for $^{86}\text{Kr} + ^{166}\text{Er}$ at 15 MeV/u (dash-dotted lines) and 20 MeV/u (solid lines) from TDHF results.

well known, the individual single-particle states and hence, also the corresponding levels, are physically not relevant, because the total wave function (Slater determinant) is invariant with respect to any unitary transformation within the occupied states. Thus we have to consider properties which are independent of such unitary transformations.

(i) The qualitative feature of *level crossings* of occupied levels with unoccupied levels is such a property. This, however, cannot be seen directly from fig. 9, because TDHF only gives the occupied levels. Here we refer to the two-center shell-model calculation of fig. 8.

(ii) The crossing with unoccupied levels will give rise to a *repulsive* (adiabatic) potential which is also an invariant quantity. The occurrence of such a repulsive potential is indicated in fig. 10 by the sum of single-particle energies and by the difference of the initial kinetic energy and the kinetic energy as function of the relative distance during the approach. Again we find practically the same repulsive potentials for both initial energies, a property expected for a diabatic potential.

From this study it is strongly suggested that TDHF is indeed close to the diabatic limit. This conclusion is further supported by a direct comparison of calculations within the two-center shell-model (§4.1). Using some adjustments to the TDHF calculations (no spin-orbit coupling, larger values of nuclear radii) give results for the diabatic levels and potential which are close to the values extracted from TDHF.

Fig. 10: TDHF potentials for $^{86}\text{Kr} + ^{166}\text{Er}$ at 15 MeV/u from the sum of single-particle energies (cf. fig. 9) multiplied by 2/3 (—) and the variation of the relative kinetic energy (:.....), and at 20 MeV/u (----- sum of single-particle energies, -.-.-.-relative kinetic energies).

4.3 EFFECTS ON CROSS-SECTIONS OF DEEPLY INELASTIC COLLISIONS

Fig. 11: Three different trajectories in the (r, α) plane for Ar (400 MeV) + Pb. Here r corresponds to the distance between the fragments and α describes the mass asymmetry. The trajectory with the smallest l -value ($l=51$) leads to the formation of a compound nucleus. The largest l -value ($l=104$) corresponds to a fast process known as dissipative (or deeply inelastic) collision. The trajectory with $l=102$, although similar to that with $l=104$ for small times, is captured. In contrast to the trajectory with $l=51$ it does not lead to a compound nucleus because $l > l(\text{crit}) \approx 70$ (largest l -value for the existence of the compound nucleus). Instead, the system develops towards mass symmetry ($\alpha=0$) and splits into two fragments if deformation of the fragments are allowed. This long-living component of dissipative collisions rotates several times and hence, exhibits essentially a $1/\sin \theta$ angular distribution like the fragments from compound-nucleus fission. From ref. ¹³⁾

In the light of the preceding discussion (§§4.1 and 4.2) a reasonable approximation to the friction kernel in nucleus-nucleus collisions can be considered by assuming that the approach of the nuclei is largely diabatic. The fast approach gives rise to a *diabatic potential* which then decays in time. Since this diabatic potential arises from the diabatic single-particle motion and not from nuclear compression it is conceptually different from other sudden potentials³⁰⁾. Such a decaying diabatic potential has been introduced in ref. ¹³⁾ in the description of heavy-ion collisions (cf. fig. 11) including fusion and a long-living or slow

dissipative component, the latter being known today as fast fission³¹).

The concept of a decaying potential has been applied by several other authors^{31,32}). However, in these calculations (like in ref.¹³) no attention has been paid to a consistent modification of the transport coefficients. Such a more consistent description has been performed recently¹⁴). In the numerical treatment of the Fokker-Planck equation (for six collective variables) we have introduced a moment expansion up to second order. Figure 12 gives the angular distribution for $^{86}\text{Kr}(8.18 \text{ MeV/u}) + ^{166}\text{Er}$, showing a remarkable dependence on the size of the local equilibration time. Also other cross-sections, like the energy distribution¹⁴), show a significant dependence on the local equilibration time. Therefore, one may hope to decide by detailed analyses of experimental results whether memory effects from diabatic single-particle motion are present in nucleus-nucleus collisions or not. However, this will not be easy because the cross-section depends also on many other quantities (transport coefficients, mass parameters and potentials).

Fig. 12: Angular distributions for the reaction $^{86}\text{Kr}(8.18 \text{ MeV/u}) + ^{166}\text{Er}$ using three different local equilibration times, $\tau_{\text{loc}} \cdot 10^{21} \text{ s} = 2$ (---), 1 (—) and 0.5 (-.-). $\tau_{\text{loc}} = 1.6 \cdot 10^{-21} \text{ s}$ corresponds to the value consistent with eq. (2.6). The experimental points are from ref.³³).

4.4 DIABATIC PRE-EQUILIBRIUM EMISSION OF NUCLEONS^{18,34)}

A direct confirmation of diabatic single-particle motion may be obtained from measurements of pre-equilibrium particles in central nucleus-nucleus collisions. As shown in fig. 8 the diabatic motion lifts nucleons to energies which lie in the continuum. These highly excited states decay by residual two-body collisions down to bound states or emit particles due to the coupling to the continuum. The latter process is referred to as the *diabatic emission of nucleons* (because of the correspondence of diabatic motion to giant vibrations, cf. §2.3, we may also call this process *giant emission*).

We estimate the energy distribution (multiplicity M per energy unit) by

$$\frac{dM}{dE} = \sum_{\alpha} \int dt' \frac{\Gamma_{\alpha}^{\dagger}(t') n_{\alpha}(t')}{[E - \epsilon_{\alpha}(t')]^2 + \Gamma_{\alpha}^2(t')/4} \quad (4.1)$$

Here $\Gamma_{\alpha}^{\dagger}$ is the time-dependent emission width of a nucleon from the diabatic state α with energy ϵ_{α} . This width has been calculated as function of the energy above threshold for a two-dimensional box. The total width Γ_{α} includes the decay width $\Gamma_{\alpha}^{\dagger}$ determined by the mean free path (2.4) with a temperature calculated from the total excitation energy ϵ^* (cf. eq. (2.6)). The occupation probability $n_{\alpha}(t)$ is determined from the decays into the continuum and to other states.

Results for pre-equilibrium neutrons are shown in fig. 13. for the central collisions of $^{86}\text{Kr}(15 \text{ MeV/u}) + ^{166}\text{Er}$. The spectra show an almost exponential tail with a slope parameter $T_{\text{eff}} \approx 6.2 \text{ MeV}$ (dashed line) for head-on collisions ($b=0$). The spectrum is very close to the high-energy tail in the reaction $^{20}\text{Ne}(14.6 \text{ MeV/u}) + ^{165}\text{Ho}$, cf. ref.³⁵⁾. Further investigations, (also measurements!) e.g. for angular distributions¹⁸⁾ are necessary, however, to confirm the "giant" emission of nucleons in nucleus-nucleus collisions.

Fig. 13: Spectrum of diabolically emitted neutrons in the reaction $^{86}\text{Kr}(15 \text{ MeV/u}) + ^{166}\text{Er}$ for three impact parameters $b=0, 2$ and 4 fm .

5. Concluding remarks

I have discussed a time-dependent shell-model approach to nucleus-nucleus collisions. This approach treats simultaneously the coherent (adiabatic) single-particle motion in the time-dependent mean field and the incoherent two-body collisions which lead to statistical equilibrium within the intrinsic degrees of freedom. An important feature of the combined effects from adiabatic single-particle motion and dissipative two-body collisions is the retardation of dissipation (memory effect) for realistic values of the mean free path of nucleons.

I have tried to focus your attention to this elasto-plastic behaviour which is located in the upper part of fig. 6. It is this part which is described by the dissipative adiabatic dynamics (DDD) and which can be studied in nucleus-nucleus collisions up to 20 MeV/u. At present we have essentially understood the chaotic regime (at least mass diffusion and angular-momentum dissipation are successfully described in this limit). The challenge now is to establish also the elasto-plastic regime. We have found strong indications that TDHF describes essentially adiabatic motions. Another step would be to confirm the adiabatic (giant) emission of nucleons. Other phenomena like extra push and capture processes (leading to fast fission), may also contain conclusive informations about the behaviour of nuclear matter in this region. Adiabatic shell-effects may be used as finger prints in such investigations.

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