DIABATICITY IN HEAVY-ION COLLISIONS AND ITS ROLE IN NON-EQUILIBRIUM NEUTRON EMISSION

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Slow collective nuclear motions such as heavy-ion collisions slightly above the Coulomb barrier and induced nuclear fission imply collective velocities $\dot{q}$ which are small compared to the Fermi velocity ($\dot{q} \ll v_F$). In the limit $\dot{q} \to 0$ the individual nucleons can completely adjust to the collective motion and the residual interactions ensure that the collective motion proceeds through constrained correlated ground-states in a time-reversible way. For realistic collective velocities, however, adiabaticity is not fulfilled. There is an interesting interval of collective kinetic energies per nucleon[1)]

$$1/40 \text{ MeV} \leq E_{coll}/A \ll 40 \text{ MeV}, \qquad (1)$$

where the motion of the individual nucleons is nearly adiabatic and the concept of single-particle states still remains valid. Only in the vicinity of quasi-crossings of the adiabatic levels the nuclear s.p. motion changes completely. According to the Landau-Zener formula[2)] a nucleon in the lower level before the crossing will find itself in the upper level after the crossing if this has been unoccupied before, or in other words, the nucleon will stay in the *diabatic* level preserving its nodal structure. By replacing the usual adiabatic s.p. basis by the diabatic states we take full account of the non-perturbative jump probabilities (close to 1). Conserving the occupation probabilities of these states the motion along the diabatic levels leads to the creation of particle-hole excitations whenever a diabatic level crosses the Fermi energy. With respect to the collective motion this excitation mechanism gives rise to a conservative potential which partially decays in time due to residual two-body collisions[3)].

The calculation of this diabatic potential with respect to the relative motion of two heavy-ions is based on the adiabatic single-particle states $\chi_\nu(q)$ generated in a two-center shell-model[4)]. The construction of the diabatic states $\phi_\nu(q)$ proceeds in a two-step iterative procedure. If the nuclei are far apart adiabatic and diabatic states coincide, i.e. $\chi_\nu(q) = \phi_\nu(q)$. The diabatic state $\phi_\nu(q')$ at a somewhat smaller relative distance q' is now generated by states $\psi_\mu(q')$ which are a linear combination of the adiabatic states $\chi_\nu(q')$,

$$\psi_\mu(q') = \sum_\nu \langle\chi_\nu(q')|\phi_\mu(q)\rangle|\chi_\nu(q')\rangle \quad . \qquad (2)$$

The expansion coefficients correspond to the overlap between the diabatic states at q and the adiabatic state at $q' = q-\Delta q$. The summation is restricted to states with maximum overlap. By orthogonalization of the set $\{\psi_\mu(q')\}$ the diabatic states $\phi_\mu(q')$ are obtained which allow for a calculation of $\phi_\nu(q-2\Delta q)$ in the next step and so forth.

The procedure described above was applied to the system $^{86}\mathrm{Kr} \rightarrow {}^{166}\mathrm{Er}$ for spherical nuclei including $\vec{l}\cdot\vec{s}$ and $\vec{l}^2$ coupling in the adiabatic s.p. Hamiltonian. The results for the lower s.p. energy branch in the $\Omega = 1/2$ shell are displayed in fig. 1 for the adiabatic s.p. levels (l.h.s.) and the constructed diabatic levels (r.h.s.):

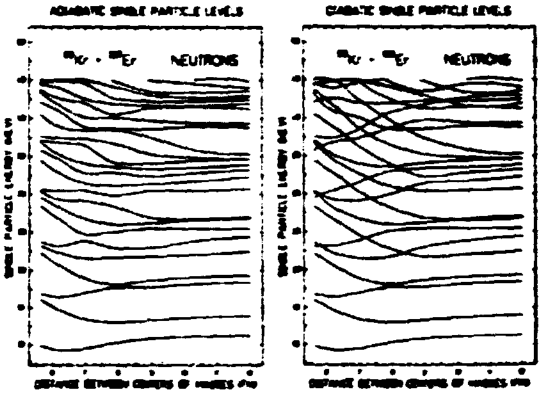

Fig. 1

The levels nicely demonstrate the changes at the pseudo-crossings of the adiabatic levels which become real crossings in the diabatic case.

The diabatic potential now corresponds to the extra energy pumped into the system due to the diabatic production of particle hole-states. All s.p. levels which contribute to this type of excitation are depicted in fig. 2a) in case of neutrons for $^{86}\mathrm{Kr} \rightarrow {}^{166}\mathrm{Er}$. The levels occupied at 14 fm (i.e. full lines for ^{166}Er and dashed lines for ^{86}Kr below the Fermi-level, which is indicated by open circles) remain occupied in the diabatic case while the Fermi-level corresponds to the highest adiabatic level. By taking the difference between the sums of diabatic and adiabatic s.p. energies, which are occupied as described before, the diabatic extra energy ΔV_{diab} for $^{86}\mathrm{Kr} \rightarrow {}^{166}\mathrm{Er}$ is obtained. The numerical result is displayed in fig. 2b) for the neutrons (dashed line) and protons (dotted line) separately as well as for their sum (full line).

The graphs show that a considerable amount of collective kinetic energy can be stored in the diabatic s.p. motion in the initial phase of a heavy-ion collision. In this respect the additional potential ΔV_{diab} yields an alternative explanation of the 'extra push'[5)], which due to friction forces alone explains the experimental observation[6)], that the barrier for the mutual capture of two heavy nuclei is considerably larger than the bare Coulomb barrier.

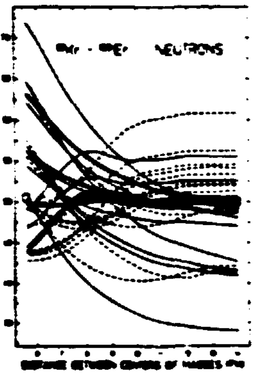

Fig. 2a)

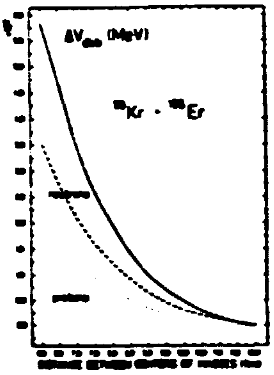

Fig. 2b)

Another interesting aspect of diabatic s.p. states is their direct connection with non-equilibrium neutron or proton emission in heavy-ion collisions. Several neutrons at least should be emitted from the colliding nuclei due to the direct coupling of the s.p. states to the collective motion, which produces highly excited s.p. states (cp. fig. 2a).

The excitation may be viewed as a squeezing of the nuclei in the radial direction leading to a considerable shift in energy of those states with several nodes in this direction. The excited states have a non-zero decay width and can be directly emitted in the continuum (*diabatic* emission of neutrons). As a first and rough approximation we assume the decay-width to be zero if the excitation energy above the local Fermi-level is less than the experimental separation energy for the corresponding neutron and to be infinite in the other case (kick-off condition). This yields an upper limit for the diabatic emission of neutrons.

The diabatic emission of neutrons (i) has to be distinguished from neutron evaporation from (ii) fully accelerated fragments and the neutron emission (iii) due to single nucleon-nucleon collisions. The process (ii) is described by isotropic emission from two sources moving with the velocities of the projectile- and target-like fragments respectively, while the mechanism (iii) is characterized by an emission (isotropic in case of s-wave scattering) from a single source moving with half of the beam velocity. The diabatic emission of neutrons is an alternative to the Fermi-jet assumption of Bondorf et al.[7] essentially describing non-equilibrium patterns.

Dynamical calculations for the collision of $^{86}Kr + ^{166}Er$ based on the diabatic potential, which is given by the sum of the energy excess ΔV_{diab} depicted in fig. 2b and the adiabatic potential of Möller and Nix[8] are performed for various energies per nucleon. The decay of the diabatic force due to residual two-body collisions is taken into account within the dissipative diabatic dynamics (DDD) described in ref. 3. The cross-section for the diabatic neutron emission is displayed in fig. 3 for various values of the local equilibration time τ_{loc} which is related to the mean free path λ of nucleons by $\tau_{loc} = 3\lambda/v_F$. The threshold for diabatic neutron emission as well as the magnitude of the cross-section only slightly depends on the local equilibration time because the time between contact and the classical turning point is small compared with reasonable values of τ_{loc}.

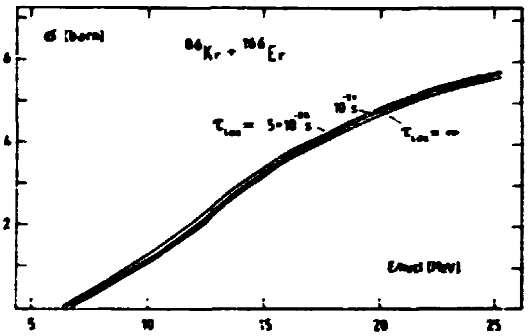

Fig. 3

The predictions up to now (within the error bars) are consistent with the experimental results for 5.7, 7.0 and 7.9 MeV/n where no significant cross-section for non-equilibrium neutrons was found and the whole yield could be explained by statistical evaporation from fully accelerated fragments[9]. In contrast to these measurements a considerable amount of non-equilibrium neutrons was detected at 11.9 MeV/n[10] which are also predicted by the calculations (fig. 3).

In order to obtain preliminary information about the angular distribution for the diabatic emission of neutrons we calculated the angular dependence (w.r.t. the center of mass of the dinuclear complex) of those states which are highly shifted in energy across the Fermi-level (cp. fig. 2a), i.e.

$$|\phi_\nu(q;\theta_{CM},\varphi)|^2 = \int_0^\infty dr\, |\phi_\nu(q;\theta_{CM},\varphi,r)|^2 \; . \quad (3)$$

In this initial study the distance between the center of masses of the two nuclei q was fixed by the kick-off condition described above. The numerical results for $|\phi_\nu(\theta_{CM})|^2$ (displayed in fig. 4) show a very pronounced forward- and backward peaking; the decrease from $\theta = 0^0$ to $\theta = 90^0$ is about one order of magnitude. Experimental measurements of Plasil et al. for $^{12}C + ^{158}Gd$ at

150 - 160 MeV show a similar behaviour[11]. They find a decrease in the angular distribution of fast neutrons from $\theta \simeq 15^0$ to $\theta \simeq 113^0$ by a factor 6.

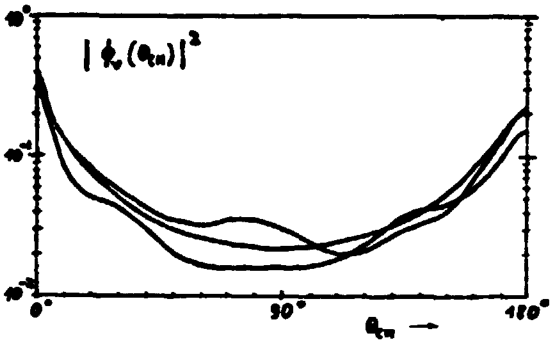

Fig. 4

We should keep in mind that the distribution shown in fig. 4 can not directly be compared with experimental results because the strong peaking of the s.p. density in z direction (fig. 5, $|\phi_\nu(q;\ \rho,z)|^2$ for three values of the radial distance q) is smoothed out due to the different orientation of the dinuclear complex with respect to the beam axis throughout the collision ($q = R_{cm}$).

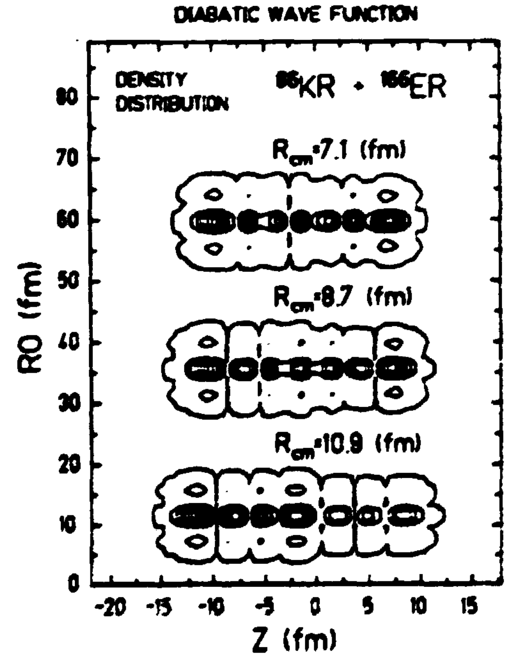

Fig. 5

Calculations along this line are in progress as well as for the energy distribution of the fast neutrons. Experimental efforts in this direction could possibly clarify, if the concept of diabaticity is valid for large amplitude but moderately slow collective nuclear motion.

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