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DYNAMICS OF DAMPED NUCLEAR REACTIONS

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Abstract

The nucleon-exchange model for damped nuclear reactions is briefly reviewed in the framework of macroscopic transport theory. Recent developments are discussed and some future directions indicated.

1. Introduction

During the last decade a new arena for nuclear physics has developed: that of damped nuclear reactions. This class of reactions occurs when heavy nuclei collide with kinetic energies of several MeV per nucleon. Such collisions introduce large-scale distortions and thus enable us to study novel aspects of nuclear structure. More generally we are provided with a unique possibility for studying the dynamics of relatively small quantal systems far from equilibrium.

Damped nuclear reactions can be briefly characterized as follows: In the interval between the relatively distant quasi-elastic reactions and the nearly head-on collisions leading to mononuclear configurations there is a regime where the two reactants engage in an intimate interaction but eventually reseparate. The associated kinetic energy loss is substantial, often amounting to hundreds of MeV. This loss of translational energy is counterbalanced by a corresponding high degree of excitation of the two emerging fragments. It is this characteristic feature of large energy loss that has motivated the names for this type of reaction: strongly damped or deep inelastic.

In spite of their strong interaction, as evidenced by the large energy loss, the two emerging fragments are usually not very different from the initial nuclei. This feature indicates that a binary configuration is maintained throughout the process. However, some deviation from the initial nucleides does occur, and it is noteworthy that the dispersions in nucleon number, and in most other observables, increase steadily as functions of the energy loss suffered. The appreciable dispersons in particle number suggest that many nucleons are exchanged in the course of a damped nuclear reaction.

In the following, we discuss the development of a dynamical theory for damped nuclear reactions based on the idealization that the exchange of nucleons between the two reaction partners constitutes the sole dissipative mechanism. The general derivation of the theory was given in ref. [1]; more recently, the application to angular momentum variables has been given special attention in ref. [2], which also contains a description of the standard implementation of the model for actual numerical calculations. The interested reader may find more detailed discussions in those papers.

2. Nuclear macrodynamics

There are different types of theoretical approach to the nuclear many-body problem. One is to start from a phenomenological few-body interaction and then try to derive the nuclear properties by solving the corresponding Dirac equation for the many-particle density matrix, in some suitable approximation scheme, such as the mean-field approximation. This task poses formidable intellectual challenges, and, although good progress has been made with respect to the formal development, we are still far from able to address processes as complicated as damped nuclear reactions. Therefore, for the time being, if we want to treat these processes, it is necessary to develop a different, less microscopic approach.

However, this purely practical motivation is not the most important one. Indeed, even if a tractable microscopic theory were available, we would still like to understand the microscopic description of the complicated many-body system in terms of simple mechanicms and elementary dynamical modes. Before this has been accomplished, we have not fully "understood" the physics of the problem. [In fact, one might say that the two approaches are mutually complementary.] In my discussion I shall denote this second type of approach

by the 'macroscopic' approach, just to distinguish it from the microscopic one. Before addressing the specific topic of damped nuclear reactions, it is useful to recall some of the characteristic features of such an approach.

When discussing the dynamics of a physical system, a natural first step is to specify the relevant degrees of fredom. In doing so we must take guidance from experiment, and a good deal of intuition is also helpful. A major drawback of the microscopic approach is that it offers no prescription for selecting the important variables. So it is with the aid of our physical feeling that we select a relatively small number of variables whose dynamical evolution we wish to consider explicitly. Let them be denoted by $\mathbb{C} = \{\mathbb{C}_1, \mathbb{C}_2, \ldots\}$ and let them be referred to as the macroscopic variables. [For a single nucleus these variables might be $\{NZ\vec{R}\vec{P}\vec{S}\alpha\}$, where N and Z are the neutron and proton numbers, $\vec{R}$ and $\vec{P}$ the nuclear position and momentum, $\vec{S}$ the total spin, and α represents a set of variables describing the nuclear shape; if the nucleus is isolated only α is of dynamical interest.]

Since the remaining degrees of freedom remain unspecified, or are only specified statistically (for example in terms of a temperature), a particular set of values for the macrovariables $\mathbb{C}$ is compatible with many microscopically different dynamical states of the system, whose only common feature is that they happen to have the same values for the selected macroscopic variables. Thus, a given value set $\mathbb{C}$ actually characterizes an entire ensemble of states of the system.

As a consequence of this inherent feature, the way in which the unspecified part of the system interacts with the retained variables is not totally predictable. This uncontrolled interaction introduces a stochastic element in the evolution of the variables considered, much like the unknown motion of individual molecules produces a random force on a macroscopic body

immersed in a fluid. In fairly small many-body systems, such as the atomic nucleus, this stochastic feature can be relatively important and need be included in the dynamical description.

Therefore, it is useful to consider the distribution function $f(\mathbb{C};t)$, which gives the probability that the retained variables have the value $\mathbb{C}$ at the time t. The introduction of the positive definite probability distribution f presumes that the macroscopic variables $\mathbb{C}$ exhibit classical dynamics.

Rather generally, the rate of change of the distribution function $f(\mathbb{C};t)$ can be written as a gain resulting from transitions into the group of states characterized by the macroscopic variables $\mathbb{C}$ minus a loss due to transitions from the current state into states characterized by a different value $\mathbb{C}'$ of the macroscopic variable. These terms generally depend on the entire history of the system. However, when the memory time is short in comparison with the time scale characterizing the evolution of the macroscopic variable, the Markov approximation can be made, resulting in a time-local master equation:

$$\frac{d}{dt} f(\mathbb{C};t) = \sum_{\mathbb{C}'} [W(\mathbb{C}' \to \mathbb{C};t)\, f(\mathbb{C}';t) - W(\mathbb{C} \to \mathbb{C}';t)\, f(\mathbb{C};t)] \tag{1}$$

A major task in nuclear dynamics is to calculate the transition probabilities W and subsequently solve the associated master equation (1) for f. It is useful to introduce the following moments of W, usually denoted the drift and diffusion coefficients, respectively,

$$\begin{aligned} V_i(\mathbb{C}) &= \sum_{\mathbb{C}'} (\zeta'_i - \zeta_i)\, W(\mathbb{C} \to \mathbb{C}') \\ D_{ij}(\mathbb{C}) &= \sum_{\mathbb{C}'} \frac{1}{2} (\zeta'_i - \zeta_i)(\zeta'_j - \zeta_j)\, W(\mathbb{C} \to \mathbb{C}') \end{aligned} \tag{2}$$

In terms of these transport coefficients, it is possible to show that the following moments of the distribution function $f(\mathbb{C};t)$

$$\bar{\mathfrak{C}}_i(t) \equiv \sum_{\mathfrak{C}} \mathfrak{C}_i \, f(\mathfrak{C};t) = \langle \mathfrak{C}_i \rangle$$

$$\sigma_{ij}(t) \equiv \sum_{\mathfrak{C}} (\mathfrak{C}_i - \bar{\mathfrak{C}}_i)(\mathfrak{C}_j - \bar{\mathfrak{C}}_j) \, f(\mathfrak{C};t) = \langle (\mathfrak{C}_i - \bar{\mathfrak{C}}_i)(\mathfrak{C}_j - \bar{\mathfrak{C}}_j) \rangle \tag{3}$$

evolve according to

$$\dot{\bar{\mathfrak{C}}}_i(t) = \langle V_i \rangle \approx V_i(\mathfrak{C} = \bar{\mathfrak{C}}(t))$$

$$\dot{\sigma}_{ij}(t) = \langle 2D_{ij} + (\mathfrak{C}_i - \bar{\mathfrak{C}}_i) \, V_j + V_i \, (\mathfrak{C}_j - \bar{\mathfrak{C}}_j) \rangle \tag{4}$$

$$\approx [2D_{ij} + \sum_k (\sigma_{ik} \frac{\partial V_j}{\partial C_k} + \frac{\partial V_i}{\partial C_k} \sigma_{kj})]_{\mathfrak{C}=\mathfrak{C}(t)}$$

Here the last, approximate relations follow when f is sufficiently narrow to permit a first-order expansion of the transport coefficients around the mean value of $\mathfrak{C}$.

When W varies smoothly the master equation (1) may be replaced by a differential equation. It is often conceptually convenient to use a representation where the transitions described by W are only those induced by a specific dissipative coupling. That is to say, in the absence of that coupling the time evolution of f is considered trivial. This can be accomplished by employing a conservative term to govern the nondissipative part of the evolution. The equation of motion may then be written in Liouville-Fokker-Planck form,

$$\dot{f} = \{\mathcal{H}, f\} - \sum_i \frac{\partial}{\partial \mathfrak{C}_i} V_i f + \sum_{ij} \frac{\partial}{\partial \mathfrak{C}_i} \frac{\partial}{\partial \mathfrak{C}_j} D_{ij} \, f \tag{5}$$

where $\mathcal{H}$ is the macroscopic hamiltonian. Although technically more convenient, it is still a formidable task to solve this equation due to the multi-dimensionality of the macroscopic variable $\mathfrak{C}$. However, often one is only interested in determining the moments (3) of the distribution function. In that case (4) forms a convenient closed set of coupled differential

equations, which is much easier to solve than the original multi-dimensional equation (1) or (5) for f. This is often called the mean trajectory method.

3. Theory of Nucleon Exchange

In damped nuclear reactions the collision system maintains a binary character throughout the contact phase. It is therefore natural to discuss the process in terms of the degrees of freedom associated with a dinucleus. The most important dynamical variables are then the following. 1) The binary partition of mass and charge (as specified, for example, by the number of neutrons, N, and protons, Z, associated with the projectile-like nucleide). 2) The relative nuclear motion (as given by the relative position $\vec{R} = \vec{R}_A - \vec{R}_B$ of the two reacting nucleides A and B together with the conjugate momentum $\vec{P}$). 3) The angular momenta carried by the individual nucleides, $\vec{S}_A$ and $\vec{S}_B$. [There is here some redundancy due to the conservation of the total angular momentum $\vec{J} = \vec{S}_A + \vec{S}_B + \vec{R} \times \vec{P}$.] 4) In addition, variables describing the interaction zone are necessary (such as the size of a "neck" between the two nucleides); distortions of the individual nucleides may also be considered (they are of particular importance in very damped reactions where large fragment deformations have time to develop). The generated nuclear excitation of the unretained, microscopic degrees of freedom, the heat Q, is also an important, though redundant, dynamical variable.

In terms of these macroscopic variables, we wish to discuss the dynamics of the system, employing the framework of time-local transport theory described in the preceding section. As noted in the introduction, the large dispersions in particle number incurred by the reacting nuclei suggest that many nucleons are exchanged in the course of a damped reaction. Elementary kinematical considerations indicate that a transferred nucleon deposits an

appreciable amount of energy and (angular) momentum. [The transfer of a nucleon with momentum $\vec{p}$ between two nuclei with a slow relative velocity $\vec{U}$ leads to an energy dissipation of approximately $\Delta Q \approx \vec{U}\cdot\vec{p}$ and an angular momentum dissipation of approximately $\vec{\Delta L} \approx \vec{R} \times \vec{p}$. Using the typical values $p^2/2m \approx T_F \approx 37$ MeV, $\frac{1}{2} mU^2 \approx 2$ MeV, and $R \approx 12$ fm, we find $\Delta Q \approx \frac{1}{2} UP_F \approx 8$ MeV and $\Delta L \approx \frac{1}{2} RP_F \approx 8\ \hbar$.] Therefore, it appears that multiple nucleon transfer may be an important reaction mechanism and its role must be examined.

It follows from the discussion in the preceding section that the basic quantity is the probability rate $\nu(i \to j)$ for a nucleon transfer from a given orbital i in one nucleide to the orbital j in the other nucleide. This quantity depends on the entire dinuclear dynamical state and in particular on the details of the interaction zone. Until now, no reliable microscopic calculation has been made of the elementary transition rates, and we rely in the following on simple estimates based on semi-classical considerations. Thus, if the donor orbital i is represented by the point $(\vec{r}_i,\vec{p}_i)$ in phase space and the receptor orbital j analogously by $(\vec{r}_j,\vec{p}_j)$, we employ the form

$$\nu(i \to j) = h^3\delta(\vec{p}_i - \vec{p}_j)\ \delta(\vec{r}_i - \vec{r}_j)\ \delta(n)\ \frac{|v_n|}{2}\ f_i\ \bar{f}_j \tag{6}$$

which expresses the assumption of quasi-free transfer of nucleons located at the contact surface where $n = \vec{r}\cdot\hat{n} = 0$. The transition rate is proportional to the speed $|v_n| = |\vec{v}\cdot\hat{n}|$ in the direction normal to the contact surface, $\hat{n}$. Finally, the Fermi-Dirac nature of nucleons dictates that the transition only take place if the donor orbital i is initially occupied (which has the probability f_i), and the receptor orbital j is available (which has the probability $\bar{f}_j \equiv 1 - f_j$).

Once the microscopic transition rate has been specified, for example as by (6), it is possible to calculate the corresponding macroscopic transition rates $W(\mathbb{C} \to \mathbb{C}')$ and the associated transport coefficients (2). Let us in the following consider macroscopic variables of the form $\mathcal{A} = \sum_i \mathcal{A}_i$ where $\mathcal{A}_i$ refers to the orbital i in the nucleide A. Examples of such variables are the neutron and proton numbers N and Z associated with A, the momentum $\vec{P}_A$ of A and its spin $\vec{S}_A$. For such additive variables we find

$$V_{\mathcal{A}} = \sum_{ij} (\nu(j \to i) - \nu(i \to j))\mathcal{A}_i$$

$$= \int \frac{d\vec{r}d\vec{p}}{h^3} \delta(n) \frac{|v_n|}{2} (f^B - f^A)\mathcal{A} \tag{8a}$$

$$\approx N'(\varepsilon_F) \langle \omega \mathcal{A} \rangle_F$$

$$2D_{\mathcal{A}\mathcal{A}'} = \sum_{ij} (\nu(j \to i) + \nu(i \to j)\mathcal{A}_i \mathcal{A}_i'$$

$$= \int \frac{d\vec{r}d\vec{p}}{h^3} \delta(n) \frac{|v_n|}{2} (f^A f^B + f^{\overline{A}\overline{f}^B})\mathcal{A}\mathcal{A}' \tag{8b}$$

$$\approx 2N'(\varepsilon_F) \langle \frac{\omega}{2} \coth (\frac{\omega}{2})\mathcal{A}\mathcal{A}' \rangle_F$$

Here the second relations follow from insertion of the specific semi-classical form (6). The third, approximate, relations emerge in the idealized case when the relative nuclear velocity is small, $U \ll V_F$, the nuclear Fermi levels are nearly equal, $F \equiv \varepsilon_B - \varepsilon_A \ll T_F$, and the microscopic degrees of freedom are in statistical equilibrium with a small temperature $\tau \ll T_F$. The amount of energy dissipated in a transfer of a nucleon with momentum $\vec{p}$ from B to A is given by $\omega = F - \vec{U}\cdot\vec{p} \ll T_F$ and the brackets $\langle\cdot\rangle_F$ denote a directional average over orbitals at the Fermi surface, which are the only ones participating in this limit. The overall common form factor $N'(\varepsilon_F)$, the differential nucleon current at the Fermi surface, depends sensitively on

the nuclear geometry. For a circular contact surface of area σ it is usually estimated as $N' \approx \frac{3}{4}\frac{\sigma}{P_F}\rho$ where ρ is the nuclear density.

Before proceeding, it may be in order to examine the justification for the approach. The idealization of the collision system as a dinucleus is sensical when the dinuclear coupling, as provided by the nucleon exchange mechanism, is sufficiently weak. After having derived the results (8), we are now in a position to address this important question. The rate of nucleon exchange is given approximately by the diffusion coefficient for the nucleon number, D_{AA}. Using a typical contact area of $\sigma = 30\ \mathrm{fm}^2$ (corresponding to a neck radius $c \approx 3$ fm), we find a rate of approximately 1.5 transfers per 10^{-22} sec. The coherence time, or duration, of a given transfer can be estimated as $t \approx \lambda/V_F \approx 10^{-22}$ sec where λ is the mean free path of a transferred nucleon. Thus there are around $\Delta N = 1.5$ transfers in progress simultaneously. This finding might seem to invalidate the basic assumption of statistically independent transfers. However, it must be remembered that the domain over which the transfers can occur is so large that several transfers could progress simultaneously without any appreciable interference. The markovian assumption is therefore reasonably well justified.

The transfer process, while in progress, introduces some uncertainty in the value of the macroscopic variables affected by the transfer. For the mass partition this effect amounts to $\Delta A \approx \sqrt{\Delta N} \approx 1.2$, which is small in comparison with typical nucleon numbers (≈ 100). For the dissipated energy we find ΔQ $\omega_{ave}\sqrt{\Delta N} \approx 3$ MeV, using $\omega_{ave} = 2$ MeV; this quantity is also small in comparison with typical energy losses, which amount to hundreds of MeV. Finally, for the angular momentum we have $\Delta L \approx (\Delta j)_{ave}\sqrt{\Delta N} \approx 10\ \hbar$. While this quantity is reasonably small in comparison with typical orbital angular momenta (which are $\approx 200\ \hbar$), it is not entirely negligible in comparison with

typical fragment spins (which are $\approx 30\ \hbar$). On the whole, though, it seems reasonably justified to consider the transfer mechanism as a weak random coupling between the two dinuclear partners, thus allowing the present type of theoretical description.

4. Results

In the preceding section we have briefly outlined the theoretical treatment of the nucleon exchange mechanism. By exploiting the stochastic and perturbative character of this coupling, it is possible to derive a consistent transport theory for the dynamics of the dinucleus. The theory provides fundamental relations between the observed mass, charge, (linear and angular) momentum, and energy. This is valuable when confrontation of the theory with experiment is made. Below we discuss briefly some of the attempts to test the theory.

4.1 Relation between energy loss and particle-number variance

The model implies a simple approximate reaction between the dissipated energy and the variance in the nucleon partition. This is easily understood qualitatively, since each transfer dissipates a kinematics-dependent amount of energy and the variance in mass partition provides a measure of the number of transfers that have occurred. Quantitatively, the rate of energy loss is given by (use $\mathcal{A} = \omega$ in (8a))

$$\dot{Q} = V_Q \approx N' \langle\omega^2\rangle_F = N' \omega_{ave}^2 \tag{9a}$$

Here $\omega_{ave}^2 \equiv \langle\omega^2\rangle_F \approx \langle(\vec{U}\cdot\vec{p})^2\rangle_F \approx \frac{1}{4} U^2 P_F^2 = \frac{1}{2} mU^2\ T_F$ (see ref. 5). Furthermore, the initial growth rate of the nucleon-number variance is given by (use $\mathcal{A} = 1$ in (8b))

$$\dot{\sigma}_{AA} \approx 2D_{AA} \approx 2N' \left\langle \frac{\omega}{2} \coth \frac{\omega}{2} \right\rangle_F \approx 2N' \left\langle \frac{|\omega|}{2} \right\rangle_F \approx N' \omega_{ave} \tag{9b}$$

Here we have assumed that $\tau << \omega$ as is justified as long as the energy loss is not too large. By taking the ratio of the above two relations the somewhat uncertain form factor N' is eliminated and we are left with the relation

$$\frac{dQ}{d\sigma_A^2} \approx \omega_{ave} \approx \frac{1}{2} UP_F = \left(\frac{1}{2} mU^2 \; T_F \right)^{1/2} \tag{10}$$

This quantity can be interpreted as the energy dissipated per nucleon exchange.

The above relation (10) makes the general prediction that the increase of energy loss with mass variance should be proportional to the square root of the available macroscopic energy per nucleon, equal to the kinetic energy of the relative dinuclear motion (above the barrier). This feature has been verified by examining experimentally obtained values of $dQ/d\sigma_A 2$ for the system Kr + La at two different bombarding energies.[3] It was found that the data do indeed quantitatively follow the simple form (10).

More extensive tests of the theoretical relation between energy loss and mass variance have been made by Schröder et al.[4] by comparing undifferentiated data directly with solutions of the dynamical moment equations (4). This is a harder, but less specific test, since the agreement with data also depends on the dynamical evolution of the form factor N'.

In order to appreciate the importance of such tests it must be recognized that alternative nucleon exchange models without proper inclusion of the Pauli blocking would lead to substantial discrepancies with this kind of data. One can therefore conclude that the picture that the major part of the damping is caused by random exchange of nucleons that are subject to the restrictions imposed by Fermi-Dirac statistics is basically a sound one.

4.2 Simultaneous transport of charge and mass

In the nucleon exchange model, the evolution of the mass and charge partition is considered as a two-dimensional random walk in the NZ-plane. The average rate of change of N and Z are given by the associated drift coefficients V_N and V_Z, which are proportional to the corresponding driving forces F_N and F_Z, as derived from the dinuclear energy function:

$$\begin{aligned} \dot{N} &= V_N = N'_N F_N \\ \dot{Z} &= V_Z = N'_Z F_Z \end{aligned} \tag{11}$$

The evolution of the associated covariance tensor is governed by the three coupled equations

$$\begin{aligned} \dot{\sigma}_{NN} &= 2D_{NN} + \sigma_{NN} \frac{\partial V_N}{\partial_N} + \sigma_{NZ} \frac{\partial V_N}{\partial_Z} \\ \dot{\sigma}_{ZZ} &= 2D_{ZZ} + \sigma_{ZZ} \frac{\partial V_Z}{\partial_Z} + \sigma_{NZ} \frac{\partial V_Z}{\partial_N} \\ \dot{\sigma}_{NZ} &= \sigma_{NN} \frac{\partial V_Z}{\partial N} + \sigma_{ZZ} \frac{\partial V_N}{\partial Z} + \sigma_{NZ}\left(\frac{\partial V_N}{\partial N} + \frac{\partial V_Z}{\partial Z}\right) \end{aligned} \tag{12}$$

The appearance of the variances on the RHS of (12) guarantees that the distribution approaches its proper equilibrium form. Thus, due to this saturation effect, it is not generally possible to deduce the number of elementary transfers from the observed variances.

The basic statistical independence of the individual transfers implies that the mixed diffusion coefficient vanishes, $D_{NZ} = 0$. However, this fact does <u>not</u> imply that the distribution function f(N,Z) remains uncorrelated; in the course of time it will adjust to the shape of the potential energy surface, which is elongated in the A-direction. Conversely, the presence of a covariance in the final distribution does not imply that the basic steps

consist of correlated (cluster) transfers. Unfortunately, much confusion has arisen in the literature about this elementary point.

The NZ moment equations (11) and (12) are solved in conjunction with the dynamical model for the mean trajectory of the dinucleus, in order to determine the time development of the form factors in the transport coefficients. The results can be directly confronted with data (apart from possible corrections due to subsequent evaporation processes, which may distort the function f(N,Z)).

Comparisons with isobaric and isotopic distributions have been made by Schröder et al.[5], especially for the case of Fe + Ho. Britt et al.[6] have made comparisons with similar data for Fe on different Ni isotopes, where the dinucleus is considerably lighter. These studies have yielded good quantitative agreement with the data. Although it is not yet fully clarified how specific this kind of test is, the overall good success lends strong support to the model, especially considering the fact that no adjustments were made prior to the confrontations.

4.3 Fragment spins

An important testing ground for theories of damped nuclear reactions is provided by the angular momentum carrying variables. The directions and magnitudes of the spins of the emerging fragments can be probed by measurements of multiplicities and angular distributions of the subsequent decay products such as γ-rays, α-particles, and fission fragments. Such techniques have produced a wealth of data, which theory must confront.

The exchange of a single nucleon transfers an appreciable amount of angular momentum from the relative orbital motion to the individual fragments. Because of the large random component of the nucleon momentum the angular momentum transfer is associated with large fluctuations. The

nucleon-exchange model makes definite statements about the coefficients characterizing the angular momentum transport. There are six independent spin variables that interplay dynamically with one another and with the additional variables. The detailed consideration of angular momentum transport has only begun recently and the quantitative confrontation with data must await the proper treatment of the subsequent deexcitation processes, which is presently in progress. It is therefore only possible at this point to state that the preliminary comparisons of the theory appear very promising. The behavior of the transferred angular momentum, its total magnitude as well as its aligned component, exhibit the same behavior as observed experimentally: a quick accumulation of aligned spin through the moderately damped region followed by fluctuation dominance in the strongly damped region.

Of special interest is the correlation between the two fragment spins, a quantity that may be probed in a double fission experiment, for example. The nucleon exchange model predicts a substantial correlation between the two spins of the emerging fragments, resulting from the preferentially parallel recoil spins deposited when a nucleon is transferred.[7] This feature is in sharp contrast with the small and negative correlation derived in a statistical model, such as that advocated by Moretto et al.[8] and also contrary to the results calculated by Wolschin et al. in a dynamical model.[9] An experiemental determination of the spin-spin correlations is therefore of great interest and would help discriminate between various current models for damped nuclear reactions.

5. Concluding remarks

The field of damped nuclear statistics is now about a decade old. In the course of this period a wealth of data has been accumulated and our qualitative insight has improved steadily. Still we are far from a comprehensive understanding of these complicated processes, which continue to pose a major challenge in nuclear theory. It is now possible to carry out rather complicated experiments and our ability to obtain more detailed information has thus improved considerably. This fact calls on theory to make more specific predictions so that decisive tests of our theoretical ideas can be made.

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