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A quantum chaotic clock and damping of the coherent nuclear rotation in the ^{28}Si +^{64} Ni dissipative collision

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Abstract. We employ the statistical reactions with memory approach to study oscillating excitation functions in the $^{28}Si(E_{lab}$ = 120 − 126.75 MeV)+^{64}Ni strongly dissipative reaction and the time evolution of the collision process. The nonself-averaging of the oscillations in the excitation functions is interpreted as indication of quantum chaos and damping of the coherent nuclear rotation in dissipative heavy-ion collisions.

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1 Introduction

The measurement of excitation function oscillations in dissipative heavy-ion collisions (DHIC) [1-16] is a valuable tool in studying the time volution of the interaction process. The advantage of analysing excitation functions oscillations is that the characteristic scales of the observable energy structure provides direct access to the associated time scales of DHIC. On the contrary, energy-averaged cross sections are time-integrated characteristics of the reaction, and therefore can reveal only indirect information on the time evolution of the collision process.

There have been already many works devoted to interpreting excitation function oscillations by studing the energy coherence lengths and the shape of the cross section energy autocorrelation functions in DHIC [2, 17-28]. However these studies are not completely consistent since the central question as to why the oscillations are not washed out in spite of the high excitations of the intermediate system, typical for DHIC, and the fact that there are a very large number of final microchannels has not been answered. A possible solution of this “intriguing puzzle” of DHIC [29] has been presented recently in refs. [30-34]. It was suggested that the nonself-averaging of the excitation function oscillations is due to spontaneous symmetry breaking resulting in the onset of coherent nuclear rotation and quantum chaos. Quantum chaos causes damping of the coherent rotation of the highly excited intermediate system with strongly overlapping resonances. In this paper we employ this approach [30-34] to analyse both the excitation function oscillations [1, 2] and the angular distributions [2, 35] of the Z = 12 and 13 reaction fragments from the $^{28}Si(E_{lab}$ = 120 − 126.75 MeV)+^{64}Ni strongly dissipative collision. This allows us to determine the angular velocity of the coherent dinuclear rotation and thereby to reconstruct the time power spectrum of the decaying intermediate system in absolute time units. It is found that the decay width β of the S-matrix spin and parity correlations is very small: β = 3.7 keV, which indicates that the coherent nuclear rotation is an extremely stable nuclear mode. This is consistent with recent analysis [16, 33, 34] of the ^{19}F+^{89}Y system [3], where it was found that β = 3.5 keV.

2 Energy-averaged cross section and time power spectrum of the collision

In order to interpret the excitation function oscillations in the ^{28}Si+^{64}Ni dissipative collision [1, 2], we take the measurable, summed over very large number $N_{\bar{b}} \to \infty$ of final microchannels ($\bar{b}$), cross section of the binary reaction in the form

$$d\sigma(E,\theta)/d\theta \equiv \sigma(E,\theta) = \sum_{\bar{b}} \sigma_{\bar{a}\bar{b}}(E,\theta), \tag{1}$$

where ($\bar{a}$) specifies the microstates of the colliding nuclei, E is the total energy and θ is the detection angle. The partial cross section can be written as [31]

$$\sigma_{\bar{a}\bar{b}}(E,\theta) = |F_{\bar{a}\bar{b}}(E,\theta)|^2 = \sigma^{(+)}_{\bar{a}\bar{b}}(E,\theta) + \sigma^{(-)}_{\bar{a}\bar{b}}(E,\theta) + 2\mathrm{Re}[F^{(+)}_{\bar{a}\bar{b}}(E,\theta)F^{(-)}_{\bar{a}\bar{b}}(E,\theta)^*], \tag{2}$$

where

$$F_{\bar{a}\bar{b}}(E,\theta) = F^{(+)}_{\bar{a}\bar{b}}(E,\theta) + F^{(-)}_{\bar{a}\bar{b}}(E,\theta), \tag{3}$$

$$F^{(\pm)}_{\bar{a}\bar{b}}(E,\theta) = \exp(\pm i(\theta/2-\pi/4)) \cdot \sum_J (W(J,I))^{1/2} \exp(iJ(\Phi \pm \theta))\bar{S}^J_{\bar{a}\bar{b}}(E), \tag{4}$$

and

$$\sigma^{(\pm)}_{\bar{a}\bar{b}}(E,\theta) = |F^{(\pm)}_{\bar{a}\bar{b}}(E,\theta)|^2. \tag{5}$$

In (2-5), $F_{\bar{a}\bar{b}}(E,\theta)$ is the total reaction amplitude, $F^{(-)}_{\bar{a}\bar{b}}(E,\theta)$ is the near-side amplitude and $F^{(+)}_{\bar{a}\bar{b}}(E,\theta)$ is the far-side amplitude. In (4), J is the total spin, Φ is the deflection angle due to the J-dependence of the potential phase shifts in the entrance and exit channels, $\bar{S}^{J}_{\bar{a}\bar{b}}(E)$ are the normalized S-matrix elements [33, 34] which do not depend on the orbital momenta or channel spins due to continuum-continuum correlation-coupling [30, 31, 46]. Equations (2-5) are obtained by using the asymptotic form of Legendre polynomials for $J \gg 1$ and $\theta \geq 1/J, \pi - \theta \geq 1/J$. Although (2-5) are formally the same as for the spinless reaction fragments, they still account for the spins of outgoing fragments [36] provided the channel spin in the exit channel is approximately perpendicular to the reaction plane, *i.e.* the reaction is planar. Therefore the argumentation of [36] and thereby (2-5) are consistent with the macroscopic picture of DHIC.

Since the colliding ions are spinless, the J-value determines the total parity uniquely. Accordingly, the parity labels are omitted. The average partial reaction probability is taken in the J-window form

$$W(J,I) \sim \exp[-(J-I)^2/d^2], \tag{6}$$

where $I \gg 1$ is the average spin of the intermediate system and d is the J-window width.

The excitation function data [1, 2] was taken with finite angular resolution of $\Delta\theta \simeq 3^\circ$ c.m. due to integration over the whole energy range of the dissipative yield. Consequently, in our theoretical analysis, we study the statistical properties of the angle-averaged cross section

$$\begin{aligned}\overline{\sigma(E,\theta)}^{\Delta\theta} &= (1/\Delta\theta)\int_{\theta-\Delta\theta/2}^{\theta+\Delta\theta/2} d\theta'\sigma(E,\theta') \\ &= \sum_{\bar{b}} \overline{\sigma_{\bar{a}\bar{b}}(E,\theta)}^{\Delta\theta},\end{aligned} \tag{7}$$

where

$$\begin{aligned}\overline{\sigma_{\bar{a}\bar{b}}(E,\theta)}^{\Delta\theta} &= \sum_{J_1 J_2}(W(J_1,I)W(J_2,I))^{1/2} \quad \cdot \\ &\bar{S}^{J_1}_{\bar{a}\bar{b}}(E)\bar{S}^{J_2}_{\bar{a}\bar{b}}(E)^* Q^{J_1J_2}(\Phi,\theta,\Delta\theta)\end{aligned} \tag{8}$$

with

$$\begin{aligned}&Q^{J_1\neq J_2}(\Phi,\theta,\Delta\theta) = \exp(i\Phi(J_1-J_2))\{(1/((J_1-J_2)\Delta\theta)) \\ &\times \sin(\Delta\theta(J_1-J_2)/2)\cos(\theta(J_1-J_2)) + (1/(J_1+J_2+1)\Delta\theta) \\ &\times \sin((J_1+J_2+1)\Delta\theta/2)\sin((J_1+J_2+1)\theta)\},\end{aligned} \tag{9}$$

and

$$\begin{aligned}&Q^{J_1=J_2}(\Phi,\theta,\Delta\theta) = 1/2 + (1/(2J_1+1)\Delta\theta) \\ &\times \sin((2J_1+1)\Delta\theta/2)\sin((2J_1+1)\theta).\end{aligned} \tag{10}$$

Using (8) we obtain the energy-averaged cross section in the form

$$\begin{aligned}&< \overline{\sigma_{\bar{a}\bar{b}}(E,\theta)}^{\Delta\theta} >_E = \sum_{J_1J_2}(W(J_1,I)W(J_2,I))^{1/2} \\ &\times < \bar{S}^{J_1}_{\bar{a}\bar{b}}(E)\bar{S}^{J_2}_{\bar{a}\bar{b}}(E)^* >_E Q^{J_1J_2}(\Phi,\theta,\Delta\theta)\end{aligned} \tag{11}$$

with [31]

$$< \bar{S}^{J_1}_{\bar{a}\bar{b}}(E)\bar{S}^{J_2}_{\bar{a}\bar{b}}(E)^* >_E = \Gamma/(\Gamma+\beta|J_1-J_2|+i\hbar\omega(J_1-J_2)). \tag{12}$$

Here, Γ is the total decay width of the intermediate dinucleus, β is the S-matrix spin and parity decoherence width and ω is the angular velocity of the coherent nuclear rotation [31].

The time evolution of the collision process is studied in terms of the time power spectrum $P(t,\theta)$ [38]. It is given by the Fourier component of the amplitude energy autocorrelation function:

$$\begin{aligned}P(t,\theta) &= (1/2\pi\hbar)\int_{-\infty}^{\infty} d\varepsilon \exp(-i\varepsilon t/\hbar) < F_{\bar{a}\bar{b}}(E+\frac{\varepsilon}{2},\theta) \\ &\quad \times F_{\bar{a}\bar{b}}(E-\frac{\varepsilon}{2},\theta)^* >_E \\ &= \sum_{J_1J_2}(W(J_1,I)W(J_2,I))^{1/2}Q^{J_1J_2}(\Phi,\theta,\Delta\theta) \\ &\quad \times (1/2\pi\hbar)\int_{-\infty}^{\infty} d\varepsilon \exp(-i\varepsilon t/\hbar) < \bar{S}^{J_1}_{\bar{a}\bar{b}}(E+\frac{\varepsilon}{2}) \\ &\quad \times \bar{S}^{J_2}_{\bar{a}\bar{b}}(E-\frac{\varepsilon}{2})^* >_E \\ &= (\Gamma/\hbar)\exp(-\Gamma t/\hbar)\sum_{J_1J_2}(W(J_1,I)W(J_2,I))^{1/2} \\ &\quad \times Q^{J_1J_2}(\Phi,\theta,\Delta\theta) \times \\ &\quad \exp(-i\omega t(J_1-J_2))\exp(-|J_1-J_2|\beta t/\hbar),\end{aligned} \tag{13}$$

where we have used the relation [31]

$$\begin{aligned}&< \bar{S}^{J_1}_{\bar{a}\bar{b}}(E+\frac{\varepsilon}{2})\bar{S}^{J_2}_{\bar{a}\bar{b}}(E-\frac{\varepsilon}{2})^* >_E \\ &= \Gamma/(\Gamma+\beta|J_1-J_2|+i\hbar\omega(J_1-J_2)-i\varepsilon).\end{aligned} \tag{14}$$

3 The cross section energy autocorrelation function

The energy autocorrelation function of the angle-averaged oscillating cross section is given by

$$\begin{aligned}C(\varepsilon,\theta) &= (< \overline{\sigma(E+\frac{\varepsilon}{2},\theta)}^{\Delta\theta}\,\overline{\sigma(E-\frac{\varepsilon}{2},\theta)}^{\Delta\theta} >_E - \\ &\quad < \overline{\sigma(E,\theta)}^{\Delta\theta} >^2_E)/ < \overline{\sigma(E,\theta)}^{\Delta\theta} >^2_E .\end{aligned} \tag{15}$$

It is easy to show that for a very large number $N_{\bar{b}} \to \infty$ of final microchannels contributing in the sum (1),

$$C(\varepsilon,\theta) = \rho(\varepsilon,\theta,\Delta\theta), \tag{16}$$

where

$$\begin{aligned}\rho(\varepsilon,\theta,\Delta\theta) &= (<\overline{\sigma_{\bar{a}\bar{b}}(E+\frac{\varepsilon}{2},\theta)}^{\Delta\theta}\,\overline{\sigma_{\bar{a}\bar{b}'}(E-\frac{\varepsilon}{2},\theta)}^{\Delta\theta} >_E \\ &\quad - <\overline{\sigma_{\bar{a}\bar{b}}(E,\theta)}^{\Delta\theta}>_E<\overline{\sigma_{\bar{a}\bar{b}'}(E,\theta)}^{\Delta\theta}>_E) \\ &\quad /(<\overline{\sigma_{\bar{a}\bar{b}}(E,\theta)}^{\Delta\theta}>_E<\overline{\sigma_{\bar{a}\bar{b}'}(E,\theta)}^{\Delta\theta}>_E)\end{aligned} \tag{17}$$

is the generalized cross-channel energy autocorrelation coefficient which does not depend on a particular pair ($\bar{b} \neq \bar{b}'$) of the exit microchannels [30, 31, 33, 34]. The numerator in the r.h.s. of (15) can be represented as

$$\begin{aligned}&\sum_{J_1J_2J_3J_4}(W(J_1,I)W(J_2,I)W(J_3,I)W(J_4,I))^{1/2} \\ &\times Q^{J_1J_2}(\Phi,\theta,\Delta\theta)Q^{J_3J_4}(\Phi,\theta,\Delta\theta) \\ &\times [<\bar{S}^{J_1}_{\bar{a}\bar{b}}(E+\frac{\varepsilon}{2})\bar{S}^{J_2}_{\bar{a}\bar{b}}(E+\frac{\varepsilon}{2})^*\bar{S}^{J_3}_{\bar{a}\bar{b}'}(E-\frac{\varepsilon}{2})\bar{S}^{J_4}_{\bar{a}\bar{b}'}(E-\frac{\varepsilon}{2})^*>_E \\ &- < \bar{S}^{J_1}_{\bar{a}\bar{b}}(E)\bar{S}^{J_2}_{\bar{a}\bar{b}}(E)^* >_E< \bar{S}^{J_3}_{\bar{a}\bar{b}'}(E)\bar{S}^{J_4}_{\bar{a}\bar{b}'}(E)^* >_E].\end{aligned} \tag{18}$$

The expression in square brackets in (14) is calculated by changing from energy averaging to ensemble averaging [33, 34]. Then, considering the S-matrix elements to be Gaussian stationary stochastic processes [48-50], we obtain [33, 34]

$$\begin{aligned}
&< \bar{S}^{J_1}_{\bar{a}b}(E+\frac{\varepsilon}{2})\bar{S}^{J_2}_{\bar{a}b}(E+\frac{\varepsilon}{2})^* \bar{S}^{J_3}_{\bar{a}b'}(E-\frac{\varepsilon}{2})\bar{S}^{J_4}_{\bar{a}b'}(E-\frac{\varepsilon}{2})^* >_E \\
&\quad - < \bar{S}^{J_1}_{\bar{a}b}(E)\bar{S}^{J_2}_{\bar{a}b}(E)^* >_E < \bar{S}^{J_3}_{\bar{a}b'}(E)\bar{S}^{J_4}_{\bar{a}b'}(E)^* >_E \\
&= \overline{\bar{S}^{J_1}_{\bar{a}b}(E+\frac{\varepsilon}{2})\bar{S}^{J_2}_{\bar{a}b}(E+\frac{\varepsilon}{2})^* \bar{S}^{J_3}_{\bar{a}b'}(E-\frac{\varepsilon}{2})\bar{S}^{J_4}_{\bar{a}b'}(E-\frac{\varepsilon}{2})^*} \\
&\quad - \overline{\bar{S}^{J_1}_{\bar{a}b}(E)\bar{S}^{J_2}_{\bar{a}b}(E)^*}\ \overline{\bar{S}^{J_3}_{\bar{a}b'}(E)\bar{S}^{J_4}_{\bar{a}b'}(E)^*} \\
&= \overline{\bar{S}^{J_1}_{\bar{a}b}(E+\frac{\varepsilon}{2})\bar{S}^{J_4}_{\bar{a}b'}(E-\frac{\varepsilon}{2})^*}\ \overline{\bar{S}^{J_2}_{\bar{a}b}(E+\frac{\varepsilon}{2})^* \bar{S}^{J_3}_{\bar{a}b'}(E-\frac{\varepsilon}{2})},
\end{aligned} \tag{19}$$

where

$$\begin{aligned}
&\overline{\bar{S}^{J}_{\bar{a}b}(E+\frac{\varepsilon}{2})\bar{S}^{J'}_{\bar{a}b'}(E-\frac{\varepsilon}{2})^*} \\
&= \Gamma/(\Gamma+\beta|J-J'|+i\hbar\omega(J-J')-i\varepsilon) \\
&\quad -\Gamma/(\Gamma+\beta|J-I|+\beta|J'-I|+i\hbar\omega(J-J')-i\varepsilon).
\end{aligned} \tag{20}$$

In (19, 20) the overbars denote the ensemble averaging. Basic argument in favour of a Gaussian distribution of S-matrix elements in DHIC is the following. DHIC are complicated processes in which an extremely large number of strongly overlapping resonance states is excited. Accordingly S-matrix elements are the sums over a very large number of partial amplitudes corresponding to different resonance levels. Then assuming that the phases of the partial amplitudes with fixed total spin and parity values quickly become randomazed [48-50] and employing the central limit theorem of probability theory [41] one arrives to a Gaussian distribution for S-matrix elements in DHIC.

4 Analysis of experimental data

We consider first the energy-averaged angular distributions

$$d\sigma(E,\theta)/d\Omega = < \overline{\sigma(E,\theta)}^{\Delta\theta} >_E / \sin\theta, \tag{21}$$

where $< \overline{\sigma(E,\theta)}^{\Delta\theta} >_E$ is given by (11). The experimental angular distributions [2, 35] (see Fig. 1) demonstrate strong forward peaking. The average deflection angle $\Phi = 45°$ is indicated by the region of sharpest variation in $d\sigma(E,\theta)/d\Omega$. In order to reproduce this sharp increase around $\Phi = 45°$ we take $d = 15/(2)^{1/2} \simeq 11$. Finally, from the ratio of experimental cross sections at $\theta \simeq 30° - 40°$ and $\theta \simeq 60° - 70°$, we find $\Gamma/\hbar\omega = 2.06$. Note that $d\sigma(E,\theta)/d\Omega$ is insensitive to the average spin I on the studied angular interval provided $I \geq d$.

While the analysis of the energy-averaged angular distributions enabled us to determine the $\Gamma/\hbar\omega$-ratio, it can not provide us with the absolute values of Γ and $\hbar\omega$. Accordingly, analysis of $d\sigma(E,\theta)/d\Omega$ does not allow one to find the absolute time scales of the coherent rotation of the intermediate dinucleus and its average life-time $\hbar/\Gamma$. The analysis of angular distributions is also powerless in the detection of quantum chaotic phenomena in DHIC, whose indication is a nonvanishing β-width [30-34], since $d\sigma(E,\theta)/d\Omega$ is insensitive to the imaginary part $\beta/\hbar$ of the angular velocity of the coherent nuclear rotation provided $\beta \ll \hbar\omega$.

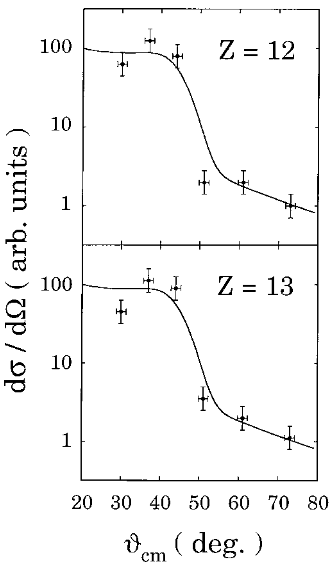

Fig. 1. Angular distributions of the $Z = 12$ and 13 fragments from the ^{28}Si+^{64}Ni dissipative collision. Data is from [2, 35]. *Solid lines* are obtained from (11) (see text)

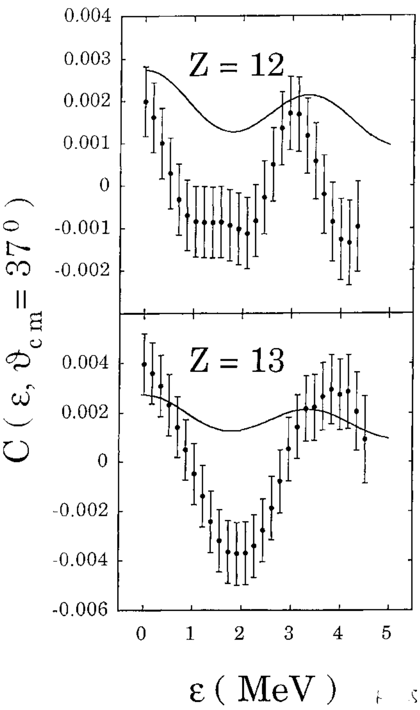

Fig. 2. Cross section energy autocorrelation functions for $Z = 12$ and 13 fragments from the ^{28}Si +^{64} Ni dissipative collision. Data is from [1, 2]. *Solid lines* are obtained from (16-20) (see text)

In order to find Γ, $\hbar\omega$ and β we must turn to the analysis of the cross section energy autocorrelation function $C(\varepsilon,\theta)$ obtained from measurements of excitation function oscillations [1, 2]. For the $Z = 12$ and 13 reaction fragments the experimental $C(\varepsilon, \theta = 37°)$'s are similar and demonstrate quasiperiodic structure (see Fig. 2). They were obtained by extracting the smooth linear background in the measured excitation functions. This linear background was deduced from a least square fit applied to the data. Such a procedure corresponds to the linear approximation in the Pappalardo method [39] of separation of gross structure against the Ericson compound nucleus fluctuations [40]. Although this quasiperiodicity is considerably overshadowed by large statistical uncertainties due to the finite data range, one can still deduce approximately a quasiperiod of $\Delta\varepsilon \simeq 3 - 3.5$ MeV. We calculate $C(\varepsilon, \theta < \Phi)$ and find that, independent on $I \geq d$, it is quasiperiodic with quasiperiod $\Delta\varepsilon = 2\pi\hbar/(t_{f.s} - t_{n.s.}) = \pi\hbar\omega/\theta$, where $t_{f.s.} = (\Phi + \theta)/\omega$ and $t_{n.s.} = (\Phi - \theta)/\omega$ are the time delays for the decay of the far-side and near-side rotating wave packets, respectively [19-21, 38]. Taking from the experiment $\Delta\varepsilon = 3.3$ MeV we uniquely obtain $\hbar\omega = 0.68$ MeV. This corresponds to $I \simeq 30$ calculated with the moment of inertia of two touched nuclei. Having determined the $\Gamma/\hbar\omega$-ratio from the analysis of the angular distributions we also obtain $\Gamma = 1.4$ MeV.

Finally, from fitting $C(\varepsilon = 0, \theta = 37°)$, we uniquely find $\beta = 3.7$ keV. The comparison between the experimental $C(\varepsilon, \theta = 37°)$'s and those calculated from (16-20) is presented in Fig. 2. One can see that although the oscillating behaviour of $C(\varepsilon, \theta = 37°)$ with quasiperiod $\Delta\varepsilon \simeq 3.3$ MeV is reproduced, there is still clear disagreement between the experimental and theoretical energy autocorrelations. Indeed, while the experimental $C(\varepsilon, \theta = 37°)$'s are negative in between the maxima at $\varepsilon = 0$ and $\varepsilon \simeq 3 - 3.5$ MeV, the theoretical ones are non-negative for the entire ε-interval. The reason for this is the following. Our present analysis is based on the results [30, 31, 33, 34], where the theoretical energy autocorrelation was obtained by performing ensemble averaging, rather than energy averaging. In order to ensure the equality of the two averaging, which is referred to as ergodicity of the underlying stochastic process, one has to perform the averaging on the energy interval $\mathcal{I}$, such that [41, 42]

$$C(\varepsilon \ll \mathcal{I}, \theta)/C(\varepsilon = 0, \theta) \to 0. \tag{22}$$

This is not the case for the experimental data [1, 2], which was obtained on the energy interval 4.7 MeV in c.m. Therefore, in order to satisfy the above condition (22) one should considerably increase the energy interval, which would also reduce statistical uncertainties in the experimental $C(\varepsilon, \theta)$'s. Alternatively, one should try to develop further the approach [30, 31, 33, 34] and calculate $C(\varepsilon, \theta)$ performing the energy averaging on the finite energy interval. The fact, that one indeed can obtain negative values of $C(\varepsilon, \theta)$ by carrying out energy averaging is illustrated in [43, 44, 53] for the case of one or several final micro-channels. However it is not obvious at present stage how to generalize these calculations [43, 44] to the case of a very large number of final micro-channels and to extend the approach [30, 31, 33, 34] to explicitely performing an energy averaging on a finite energy interval. The stochastic modelling of the oscillating excitation functions in the DHIC, which we intend to develop and apply in the future, may serve as useful tool in studying deviations from ergodicity in DHIC and, in general, in statistical processes with memory [30-34, 37, 43-47, 53, 54].

5 Quantum chaos and the dinuclear clock in the $^{28}Si + ^{64}Ni$ dissipative collision

Using (13) we calculate the time power spectra of the collision at different angles (see Fig. 3). One observes two distinct maxima which clearly determine the time delays for the near-side and far-side decays of the rotating intermediate system. For comparison we also depict exponential time power spectra for the compound nucleus decay.

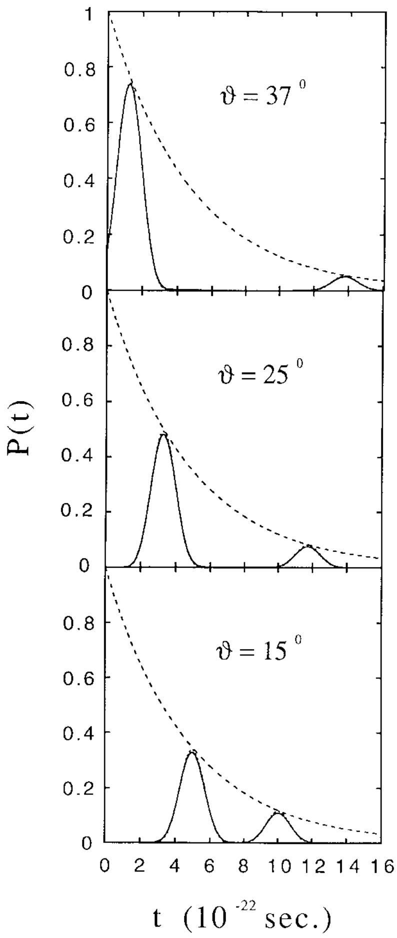

Fig. 3. Time power spectra reconstructed for the $^{28}Si + ^{64}Ni$ dissipative collision at different angles. *Solid lines* are obtained with $\beta = 3.7$ keV, and *dotted lines* correspond to $\beta/\Gamma = 0$ (see text). *Dashed lines* correspond to exponential compound nucleus decay

The fact that we are able to make an absolute normalization of the dinuclear clock is due to the analysis of excitation function oscillations which provided us with the absolute value of the angular velocity ω. The presence of any structure in the excitation functions, *i.e.* the nonself-averaging of the excitation function oscillations in spite of the strong overlap of the dinuclear resonances and enormous number of final micro-channels, indicates non-vanishing of the β-width [30-34]. The $\beta/\hbar$-rate has a physical meaning as the imaginary part of the angular velocity of the coherent nuclear rotation [31]. The origin of the damping of the coherent nuclear rotation is the S-matrix spin and parity decoherence, which is a typical quantum chaotic phenomenon [30, 31] occurring because of the statistical relaxation in the discrete spectrum on finite time intervals [51, 52]. Therefore the quantum chaos and damping of the coherent nuclear rotation enable us to obtain an absolute time normalization of the time power spectra and thereby supply us with the dinuclear clock.

Yet, although being non-vanishing, the damping width of the coherent rotation is extremely small: $\beta = 3.7$ keV. This implies a very slow time-space delocalization of the highly excited dinucleus. Indeed, the difference between the time power spectra with β = 3.7 keV and β/Γ = 0 is practically invisible in Fig. 3. However the signal from the finite non-vanishing β-width is clearly seen due to oscillating, non smooth and structureless, excitation functions in the ^{28}Si +^{64} Ni dissipative collision [1, 2].

Reduction of β results in the suppression of the time-space delocalization of the dinucleus which improves the precision of the dinuclear clock. In the limit $\beta/\Gamma \to 0$, *i.e.* in the limit of absence of quantum chaos resulting in undamped regular coherent rotation, any structures in the excitation functions are washed out. Accordingly, information about the value of ω is lost and absolute normalization of time is no longer possible. Thus a dinuclear clock can never be precise, otherwise it becomes invisible. However, since $\beta/\hbar\omega \simeq$ 0.005 rad $\simeq$ 0.3°, coherent rotation is an extremely stable nuclear mode and thus a dinuclear clock is still quite reliable.

6 Conclusions

We have analyzed the angular distributions and excitation function oscillations of the Z = 12 and 13 reaction fragments in the $^{28}Si(E_{lab}$ = 120 − 126.75 MeV) +^{64} Ni dissipative collision. This enabled us to reconstruct the time power spectrum of the collision process in absolute time units. We have interpreted the presence of oscillating structure in the excitation functions as an indication of quantum chaos and damping of the coherent nuclear rotation in DHIC [30-34]. However the damping width of the coherent nuclear rotation was found to be 3.7 keV. This is consistent with recent analysis [16, 33, 34] of the ^{19}F +^{89} Y dissipative collision [3], which resulted in β=3.5 keV. This indicates that coherent nuclear rotation is an extremely stable mode in spite of the high excitations typical for DHIC. On the contrary, collective rotational bands and, in general, nuclear collective states acquire substantial spreading widths already at several MeV excitation energy [50]. Accordingly, since the spreading widths of the highly excited collective states in DHIC are from about 2-3 MeV to 10-20 MeV [48, 49], these collective modes lose their individuality as excited states very quickly. Therefore, one observes that, unlike the extremely stable nuclear coherent motion, the collective nuclear motion in DHIC is strongly overdamped and thus can not contribute to the nuclear dynamics on time intervals longer than 10^{-22} sec.

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