Quantum-classical correspondence in microscopic and mesoscopic complex collisions

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We propose a novel method to study quantum-classical correspondence in complex, e.g., heavy-ion, atomic, molecular, and atomic cluster collisions. The many-body rotating wave packets are formed spontaneously and their stability is due to the slow decoherence between highly excited strongly overlapping states of the intermediate complex. The phenomenon is illustrated by analyzing the $^{12}C+^{24}Mg$ collision. The significant deviations from the random-matrix theory are reproduced in terms of the macroscopically rotating molecule formed in this collision.

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The search for coherent-state wave packets started in the early days of quantum mechanics [1]. These studies addressed a single-particle problem: the formation of atomic-electron wave packets that were spatially localized and moving in classical Kepler orbits. Since the 1980's it has been shown that such states can be formed using a short laser pulse to excite a coherent superposition of Rydberg single-electron states [2]. The possibility to produce spatially localized electron wave packets has been generating considerable interest among atomic, molecular, and optical researchers [3] because it allows one to study the Bohr correspondence principle limit of the Rydberg states and the connection between the orbits of a classically chaotic system and the motion of single-particle quantum wave packets.

In this paper we address the quantum-classical correspondence problem for a highly excited many-body intermediate system (IS) with strongly overlapping resonances created in complex quantum collisions (CQC). We show that the precondition for producing stable many-body wave packets is slow spin decoherence. In fact, one requires the decoherence width [4,5] to be much smaller than anticipated in the random-matrix theory (RMT) of many-body systems [6–9] and the theory of quantum chaotic scattering [10–12]. Therefore the many-body coherent states identified in this paper are intimately related to the significant deviations from RMT. As time proceeds the decoherence gradually destroys the highly excited coherent condensate (many-body wave packet) resulting in the regime described by RMT. In the limiting case of fast decoherence our approach fully recovers the RMT of many-body systems. The spin coherence between the highly excited states of the IS occurs spontaneously [4,5] resembling the spontaneous breaking of rotational symmetry and describing the collapse of the initially delocalized incoherent superposition of the extended resonance eigenstates into the spatially localized many-body wave packets. The following slow decoherence, resulting in slow restoration of rotational symmetry and the spreading of the rotating wave packets, is closely related to localization within the infinite number of orthogonal subspaces of the Hilbert space [4,5].

As an example, we demonstrate that the significant deviations from RMT in $^{12}C+^{24}Mg$ elastic and inelastic CQC [13,14] can be interpreted in terms of coherently rotating wave packets. For these CQC, high intrinsic excitations and, therefore, a strong overlap of resonances of the IS manifest themselves in the fact that maxima in the cross sections for elastic and different inelastic channels occur at different energies. This is reflected unambigously, in a model-independent way, by the insignificant, 0.17 ± 0.15 [14], cross-channel correlation coefficients in these CQC. In contrast, the excitation of isolated molecular resonances of the IS manifests itself in maxima in the cross sections for elastic and different inelastic channels occurring at the same energies and resulting in large, $\simeq 1$, cross-channel correlation coefficients [15,16]. The above model-independent arguments are consistent with the theoretical evaluation of the average level spacing D of the IS. The average total energy of the IS, $E\simeq 35\,\text{MeV}$, consists of the deformation energy, which is mainly given by the Coulomb energy ($\simeq$13 MeV) of the two touched ions, the rotational energy ($\simeq$10 MeV) and the intrinsic excitation energy E^*. We have $E^*\geqslant 10\,\text{MeV}$ corresponding to $D\leqslant 10^{-5}\,\text{MeV}$ while, as we shall see, the total width of the resonance levels $\Gamma\simeq 0.4\,\text{MeV}$. This means that, even though the intrinsic energy constitutes only $\simeq$30% of the total energy, at any fixed energy of the incident beam the observable cross sections are formed by means of interference between $\geqslant 10^4$ partial amplitudes corresponding to the decay of $\geqslant 10^4$ strongly overlapping resonance levels of the IS. It should be stressed that even for an unrealistically large, e.g., 10, average number of complete revolutions of the IS before its decay, the measured cross sections would still be formed by means of interference between $\geqslant 10^3$ partial amplitudes corresponding to the decay of $\geqslant 10^3$ strongly overlapping resonances of the IS.

There is convincing evidence that the effects of complexity and stochasticity in nuclear systems are shared by other many-body systems [9]. Therefore it should be of interest to experimentally search for a quantum-classical correspondence in, e.g., atomic, molecular and atomic cluster collisions. It should be noted that the manifestation of quantum coherence for hot C_{60} atomic clusters has been unambiguously demonstrated in recent experiments [17]. This supports the conceptual similarity in the quantum mechanical treatment of atomic clusters and heavy-ion collisions [18].

It should be stressed that the identification of distinct deviations from RMT in this paper does not imply its general usefulness in studying highly excited many-body systems. However such identifications do allow one to specify a new time scale, inverse decoherence width [4,5], for many-body systems. Physically this time scale is analogous to the inverse Thouless energy for single-particle motion in disordered systems. but it is completely different from the inverse spreading width in many-body systems [9]. For times shorter than this scale, the RMT of quantum many-body systems ceases to apply.

For a fixed total spin J, the S matrix can be written in the form [9,12] $S^J_{ab}(E)=\langle S^J_{ab}(E)\rangle+\delta S^J_{ab}(E)$ with $\langle \delta S^J_{ab}(E)\rangle=0$. The angle brackets denote an average over an energy interval $\Delta E\simeq\hbar/t_{\rm dir}$ with $t_{\rm dir}\simeq R/v$, where R is the characteristic linear size of the interaction region and v is the relative velocity of the collision partners. Consequently, $\langle S^J_{ab}(E)\rangle$ corresponds to potential scattering or direct reactions. These fast processes do not proceed through the formation of a long-lived IS. For the ^{12}C+^{24}Mg CQC analyzed in this paper we have $\Delta E\simeq 5$ MeV. The fluctuations, i.e., the rapid variations with energy E, originate from $\delta S^J_{ab}(E)$ corresponding to the formation of a relatively long-lived IS. The formation of a deformed IS is supported by the two-center shell model and by density-functional calculations for heavy-ion [15,16] and atomic cluster [19] CQC. The formation of a deformed IS in CQC is also supported by time-dependent Hartree-Fock calculations [20]. The above decomposition of the S matrix is meaningful if $\Delta E\gg\Gamma$, where Γ is a characteristic scale for the energy variations of $\delta S^J_{ab}(E)$. As we shall see, this condition is consistent with the results of our interpretation of the ^{12}C+^{24}Mg CQC, yielding $\Gamma=0.4$ MeV.

We use the unitary S-matrix representation [21]. The fluctuating S-matrix elements $\delta S^J_{ab}(E)=S^{J,{\rm res}}_{ab}(E)-\langle S^{J,{\rm res}}_{ab}(E)\rangle$, originate from the resonance part of the full S matrix $S^{J,{\rm res}}_{ab}(E)=-2i\exp(i\phi^J_{ab})t^{J,{\rm res}}_{ab}(E)$, where $t^{J,{\rm res}}_{ab}(E)=\Sigma_{\mu\mu'}\gamma^{Ja}_\mu[\mathcal{D}(J)^{-1}]_{\mu\mu'}\gamma^{Jb}_{\mu'}$ and $[\mathcal{D}(J)]_{\mu\mu'}=\delta_{\mu\mu'}(E-E^J_\mu)+i\Sigma_c\gamma^{Jc}_\mu\gamma^{Jc}_{\mu'}$. Here, $\phi^J_{ab}=\delta^J_a+\delta^J_b$ and $\delta^J_{a(b)}$ are the energy smooth potential phase shifts in the entrance (exit) channels, E^J_μ are the resonance energies and $\gamma^{Ja(b)}_\mu$ are the partial width amplitudes. The channel label, $b=(\overline{b},\tilde{b})$, carries the indices of the intrinsic state, $\overline{b}$, of the collision partners and $\tilde{b}=(l_b,j_b)$, where l_b is the orbital momentum and j_b is the channel spin. The cross section energy autocorrelation function (EAF) is given in terms of the S or t matrix EAF: $c^{JJ'}(\varepsilon)=\langle \delta\overline{t^J_{ab}}(E)\,\delta\overline{t^{J'}_{a'b'}}(E+\varepsilon)^*\rangle$, where $\overline{a}=\overline{a'}$, $\overline{b}=\overline{b'}$, $\delta\overline{t^J_{ab}}(E)=\delta t^J_{ab}(E)/[\langle|\delta t^J_{ab}(E)|^2\rangle]^{1/2}$ are the normalized δt-matrix elements $\delta t^J_{ab}(E)=t^{J,{\rm res}}_{ab}(E)-\langle t^{J,{\rm res}}_{ab}(E)\rangle$, and the angle brackets denote energy averaging over an interval $\mathcal{I}\gg\Gamma$. Considering the $\gamma^{Ja(b)}_\mu$ as Gaussian random variables [7] we calculate $c^{JJ'}(\varepsilon)$ by averaging over an ensemble of γ's for a fixed $E=\overline{E}$ rather than performing energy averaging. The equivalence of energy averaging and ensemble averaging has been justified in Refs. [22,23].

The procedure of ensemble averaging is considerably simplified under the diagonal approximation for $[\mathcal{D}(J)]_{\mu\mu'}$, $\Sigma_c\gamma^{Jc}_\mu\gamma^{Jc}_{\mu'}=(\Gamma/2)\delta_{\mu\mu'}[1\pm\mathcal{O}(1/N^{1/2})]\pm\Gamma(1-\delta_{\mu\mu'})\mathcal{O}(1/N^{1/2})\simeq\delta_{\mu\mu'}\Gamma/2$, where $N\gg1$ is the number of open channels and $\Gamma=2\Sigma_c(\gamma^{Jc}_\mu)^2$. This approximation is suggested by the noncorrelation condition [7], $\overline{\gamma^{Jc}_\mu\gamma^{Jc}_{\mu'}}=\delta_{\mu\mu'}\overline{(\gamma^{Jc}_\mu)^2}$, where the overbars stand for ensemble averaging. We obtain $\delta\overline{t^J_{ab}}(E)=(\Gamma D/2\pi)^{1/2}\Sigma_\mu\overline{\gamma}^{J\overline{a}}_\mu\overline{\gamma}^{Jb}_\mu/(E-E^J_\mu+i\Gamma/2)$, where we have neglected the J dependence of Γ and D. Here, the normalized $\overline{\gamma}^{J\overline{a}(\overline{b})}_\mu=\gamma^{Ja(b)}_\mu/[\overline{(\gamma^{Ja(b)}_\mu)^2}]^{1/2}$ are taken to be $\tilde{a}(\tilde{b})$ independent. The smooth $(\tilde{b})$ dependence of $\overline{\gamma}^{Jb}_\mu$ is justified due to the continuum correlation [4,5], implying that $(\tilde{a},\tilde{b})$ dependence of $\delta S^J_{ab}(E)$ is taken into account by the potential phase shifts ϕ^J_{ab}. We use the correlation relation [5]

$$\overline{\overline{\gamma}^{J\overline{a}}_\mu\overline{\gamma}^{Jb}_\mu\overline{\gamma}^{J'\overline{a}}_\nu\overline{\gamma}^{J'b}_\nu}-\overline{\overline{\gamma}^{J\overline{a}}_\mu\overline{\gamma}^{Jb}_\mu}\;\overline{\overline{\gamma}^{J'\overline{a}}_\nu\overline{\gamma}^{J'b}_\nu}=(1/\pi)D\beta|J-J'|/[(E^J_\mu-E^{J'}_\nu-\hbar\omega(J-J'))^2+\beta^2(J-J')^2], \tag{1}$$

where ω is the angular velocity of the coherent rotation and $\beta\gg D$ is the spin decoherence width. We obtain

$$c^{JJ'}(\varepsilon)=\Gamma/[\Gamma+\beta|J-J'|+i\hbar\omega(J-J')-i\varepsilon]. \tag{2}$$

The intensity of the correlation (1) is maximal for $E^J_\mu-E^J_{\rm rot}=E^{J'}_\nu-E^{J'}_{\rm rot}$, i.e., for the same intrinsic excitation energies of the IS having different spin values. Here, $E^J_{\rm rot}=\hbar^2J^2/2\mathcal{J}$, $\mathcal{J}$ is the moment of inertia of the IS, $E^J_{\rm rot}-E^{J'}_{\rm rot}\simeq(J-J')\hbar\omega$, $\omega=\hbar(J+J')/2\mathcal{J}\simeq\hbar I/\mathcal{J}$, $|J-I|$, $|J'-I|\ll I$ and I is the average spin of the IS.

Since the diagonal approximation for $[\mathcal{D}(J)]_{\mu\mu'}$ violates unitarity of the S matrix the above derivation is questionable. The calculation of $c^{JJ'}(\varepsilon)$ for the $t^{J({\rm res})}_{ab}(E)$ corresponding to a unitary S matrix, i.e., beyond the diagonal approximation for $[\mathcal{D}(J)]_{\mu\mu'}$, proceeds through the following main steps. (i) Following Ref. [7] we expand $[\mathcal{D}(J)^{-1}]_{\mu\mu'}$ into a series over the nondiagonal part of $[\mathcal{D}(J)]_{\mu\mu'}$. (ii) We perform ensemble averaging and resummations within each individual $\delta t^J_{ab}(E+\varepsilon)$ and $\delta t^{J'}_{a'b'}(E)^*$, in complete analogy with Ref. [7]. (iii) Using Eq. (1) we perform ensemble averaging of the spin off-diagonal contributions, i.e., of the cross-contracted pairs of $\overline{\gamma}$'s [7] carrying different spin values. The resultant expression for $c^{JJ'}(\varepsilon)$ is also given by Eq. (2) but with $\Gamma\simeq ND/2\pi$, where $\hbar/\Gamma$ is the average lifetime of the IS.

We consider first the collision of spinless particles in the entrance and exit channels. The cross section has the form $\sigma_{\overline{ab}}(E,\theta)=\sigma^d_{\overline{ab}}(\theta)+\sigma^{\rm osc}_{\overline{ab}}(E,\theta)+2\,{\rm Re}[f^d_{\overline{ab}}(\theta)^*f^{\rm osc}_{\overline{ab}}(E,\theta)]$, where $\sigma^d_{\overline{ab}}(\theta)=|f^d_{\overline{ab}}(\theta)|^2$ and $\sigma^{\rm osc}_{\overline{ab}}(E,\theta)=|f^{\rm osc}_{\overline{ab}}(E,\theta)|^2$. The energy independent (smooth) $f^d_{\overline{ab}}(\theta)$ and $\sigma^d_{\overline{ab}}(\theta)$ are determined by the energy averaged S-matrix elements. The energy dependence of the cross section originates from the amplitude $f^{\rm osc}_{\overline{ab}}(E,\theta)=\Sigma_J(2J+1)W(J)^{1/2}\exp(i\Phi J)\delta\overline{t^J_{\overline{ab}}}(E)P_J(\theta)$ with $\langle f^{\rm osc}_{\overline{ab}}(E,\theta)\rangle=0$.

Here $\Phi \equiv \Phi_{\overline{ab}} = d\phi^{J}_{\overline{ab}}/dJ$ is the deflection angle due to the J dependence of the potential phase shifts, $P_J(\theta)$ are the Legendre polynomials and the partial reaction probability, $W(J) \equiv W_{\overline{ab}}(|J-I|/d) = \langle |\delta t^{J}_{\overline{ab}}(E)|^2 \rangle$, is taken in the J window form with I being the average spin and $d \ll I$ the J-window width. We calculate the cross section energy autocorrelation function (EAF) [6] and obtain

$$C(\varepsilon,\theta) = [|\rho(\varepsilon,\theta)|^2 + 2\sigma^{d}_{\overline{ab}} \mathrm{Re}\,\rho(\varepsilon,\theta)]/\langle \sigma_{\overline{ab}}(E,\theta)\rangle^2,$$

where

$$\rho(\varepsilon,\theta) = \langle \delta f^{\mathrm{osc}}_{\overline{ab}}(E+\varepsilon,\theta)\,\delta f^{\mathrm{osc}}_{\overline{ab}}(E,\theta)^{*}\rangle,$$

$$\langle \sigma_{\overline{ab}}(E,\theta)\rangle = \sigma^{d}_{\overline{ab}}(\theta) + \langle \sigma^{\mathrm{osc}}_{\overline{ab}}(E,\theta)\rangle.$$

Consider the Fourier component of $C(\varepsilon,\theta)$. We obtain

$$\int_{-\infty}^{\infty} d\varepsilon\, \exp(-i\varepsilon t/\hbar)\mathrm{Re}\,\rho(\varepsilon,\theta)$$

$$\propto \int_{-\infty}^{\infty} d\varepsilon\, \exp(-i\varepsilon t/\hbar)\rho(\varepsilon,\theta) \propto P(t,\theta)$$

[24], where $P(t,\theta)$ is the time power spectrum (TPS) of the collision. Since $\delta f^{\mathrm{osc}}_{\overline{ab}}(E,\theta)$ are Gaussian stochastic processes we also have $\int_{-\infty}^{\infty} d\varepsilon \exp(-i\varepsilon t/\hbar)|\rho(\varepsilon,\theta)|^2 \propto \int_0^{\infty} d\tau P(\tau,\theta)P(t+\tau,\theta)$. This demonstrates that $C(\varepsilon,\theta)$ does indeed contain information about the TPS of CQC. In calculating $P(t,\theta)$ we use Eq. (2) and the asymptotic form of the Legendre polynomials. For $\theta d \geqslant 1$, we obtain $P(t,\theta) \propto P^{(+)}(t,\theta) + P^{(-)}(t,\theta)$ with

$$P^{(\pm)}(t,\theta) \simeq \exp(-\Gamma t/\hbar)[1-\exp(-2\beta t/\hbar - 2\Delta)]/[1-\exp[i(\omega t - \Phi \mp \theta) - \beta t/\hbar - \Delta]|^2 \tag{3}$$

and $\Delta = 1/d$. For $t < \hbar/\beta$, $P(t,\theta)$ shows maxima at $t_m \simeq (2\pi m + \Phi \pm \theta)/\omega \geqslant 0$ with the widths $\delta t_m \simeq (\beta t_m/\hbar + \Delta)/\omega$ and $m = 0,1,\ldots,$ demonstrating macroscopiclike rotation of the IS due to the quantum spin coherence. The same physical picture, the lighthouse effect, was obtained in Ref. [25]. However, unlike Ref. [25] which addressed the regime of isolated resonances, in this paper the macroscopic rotation (3) has been obtained in spite of the high intrinsic excitations and strong overlap of resonance levels of the IS providing $\beta \leqslant \Gamma$ and $d \geqslant 2-3$. As time proceeds, the decoherence results in the spreading of the rotating wave packets and, for $t \gg \hbar/\beta$, $P(t,\theta) \to \exp(-\Gamma t/\hbar)$. Clearly, for $\Gamma/\beta \ll 1$ or $d \to 0$, $P(t,\theta) \to \exp(-\Gamma t/\hbar)$ and $C(\varepsilon,\theta) \propto 1/(1+\varepsilon^2/\Gamma^2)$ recovering the RMT result [6,9,12]. Therefore, formation of the coherently rotating many-body wave packets in CQC should manifest itself in distinct deviations of the EAF from the Lorentzian angle-independent shape predicted by RMT.

It can be shown that $P(t,\theta)$ has the form (3) also for the collision products having nonzero intrinsic spins. The only difference is that Φ will depend on the intrinsic spins of the collision partners.

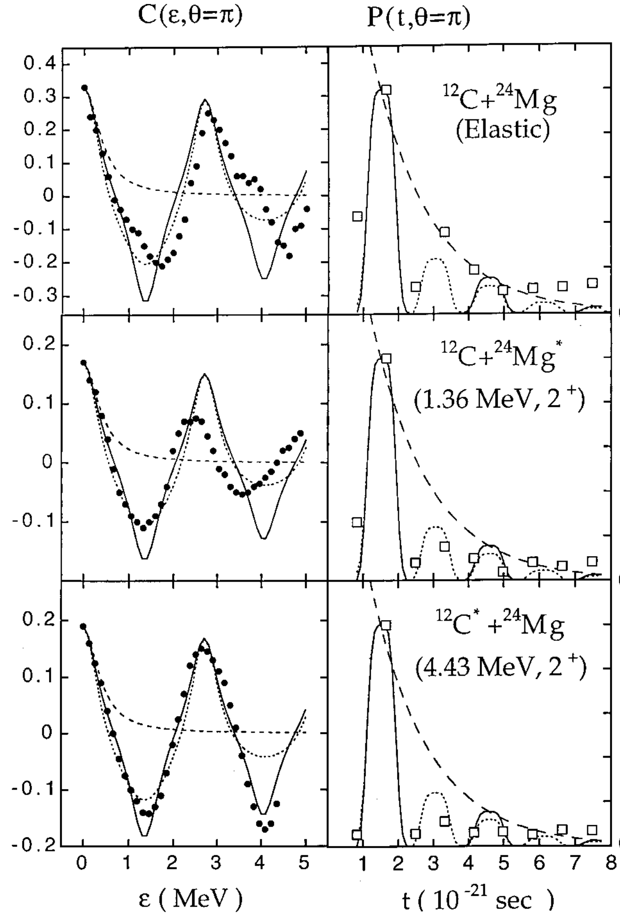

FIG. 1. Experimental (dots and squares) [14] and theoretical energy autocorrelation functions (left) and time power spectra (right) for $^{12}C+^{24}Mg$ elastic and inelastic collisions. The time power spectra are in arbitrary units. The solid lines are theoretical fits with $\beta_{>} = \beta_{<} = \beta$ and the dotted lines are theoretical fits with $\beta_{>} \gg \beta_{<}$ (see text). The dashed lines are Lorentzians (left) and exponents (right) obtained in the RMT limit $\Gamma \ll \beta$ with $\Gamma = 0.4$ MeV.

Although our method is expected to apply to atomic, molecular, and atomic cluster collisions currently the only available data with sufficient accuracy are heavy-ion data. We analyze, as an example, $^{12}C+^{24}Mg$ collisions at $\theta = \pi$ [13]. This choice is motivated by previous studies of the same system [24] indicating slow decoherence, $\beta/\Gamma \ll 1$, in these CQC for higher energy regions. This also enables us to demonstrate the consistency of the present interpretation and the energy independence of the decoherence rate over the large, 12–43 MeV c.m. energy range. In Fig. 1 we present $C(\varepsilon,\pi)$ obtained [14] from the data [13] by employing the Pappalardo trend reduction method with the energy averaging interval $\Delta E_{\mathrm{c.m.}} = 4.26$ MeV. The experimental EAFs correspond to the c.m. energy range 12.27–22.8 MeV for the elastic channel, and to 11.33–26.4 MeV and 16.53–27.47 MeV for the $^{12}C+^{24}Mg^{*}$ and $^{12}C^{*}+^{24}Mg$ inelastic channels, respectively. The EAFs oscillate with a period of $\simeq 2.6-2.8$ MeV demonstrating distinct deviations from the Lorentzian shape predicted by the RMT. Since $C(\varepsilon=0,\pi) = 1 - y_d^2$ [6], where $y_d = \sigma^{d}_{\overline{ab}}(\pi)/\langle \sigma_{\overline{ab}}(E,\pi)\rangle$, then, from

Fig. 1, $y_d > 0.8$. Therefore, $P(t,\pi) \propto \int_0^\infty d\varepsilon \cos(\varepsilon t/\hbar) C(\varepsilon,\pi)$ for all the three channels. The main common feature of the normalized experimental TPSs in Fig. 1 is the strong maxima at $t \simeq 1.6 \times 10^{-21}$ s. Since $\Phi \leqslant \theta_{gr} \simeq 30°$ [13] it can be neglected for backwards angles. The fits in Fig. 1 are obtained using Eq. (3) with $\Phi = 0$, $d = 3$, $\Gamma = 0.4$ MeV, $\beta = 10$ keV, $\hbar\omega = 1.35$ MeV for all the three channels.

In Fig. 1 we present the fits of the EAFs, $C(\varepsilon,\pi) \propto \mathrm{Re}\,\rho(\varepsilon,\pi) \propto \mathrm{Re} \int_0^\infty dt \exp(i\varepsilon t/\hbar) P(t,\pi)$, obtained with the set of parameters derived from the fit of the TPSs. For $\theta > |\Phi|$, $\theta d > 1$, $\hbar\omega > \beta$, an analytical estimate of $\rho(\varepsilon, \theta)$, which can be used for the evaluation of the angle dependence of the $C(\varepsilon,\theta)$, yields $\rho(\varepsilon,\theta) \propto \rho^{(+)}(\varepsilon,\theta) + \rho^{(-)}(\varepsilon,\theta) + B$. Here $\rho^{(\pm)}(\varepsilon,\theta) = \pi \exp(-a^{(\pm)} b)/\sinh(\pi b)(\hbar\omega - i\beta)$, $a^{(\pm)} = (\Phi \pm \theta \mp \pi + i\Delta)$, $b = (\Gamma - i|\varepsilon|)/(\hbar\omega - i\beta)$, and $B = 2/(i\varepsilon - \Gamma)$. In order to account for the noticeable second peak at $t \simeq 3.2 \times 10^{-21}$ s in the TPS for the elastic channel we take the decoherence width $\beta_>$ between the states with different parities (odd $|J - J'|$ values) to be greater than that $\beta_<$ between the states with the same parities (even $|J - J'|$ values) [24]. The corresponding fits with $\beta_> = 0.4$ MeV, $\beta_< = 10$ keV and with the same remaining parameters are presented in Fig. 1. One can see that the fits of the data for all the three channels do not provide a conclusive indication for the faster restoration of inversion symmetry as comparing to the rate of the restoration of rotational symmetry of the IS.

For $E_{\mathrm{c.m.}} \simeq 15 - 20$ MeV, $I \simeq 15$ [13]. Calculating ω for the moment of inertia corresponding to two touched spheres with the radii of ^{12}C and ^{24}Mg we obtain $\hbar\omega \simeq 1.5$ MeV. This supports our interpretation yielding $\hbar\omega = 1.35$ MeV which suggests $\simeq$(2:1) deformation of the IS formed in the collision.

The present identification of an extremely small decoherence width, $\beta = 10$ keV, between the partial width amplitudes of a highly excited IS is in sharp contrast with the conventional theories of highly excited many-body systems. Within these theories [9], the lower limit of β should be given by $\beta \geqslant \Gamma_{\mathrm{spr}}$, where Γ_{spr} is the spreading width. The inverse spreading width, $\hbar/\Gamma_{\mathrm{spr}}$, has a meaning of the energy relaxation time, i.e., the time it takes the two-body interaction to equally distribute energy between the particles of the system. For highly excited nuclear systems, $\Gamma_{\mathrm{spr}} \geqslant 5$ MeV [9], so that the conventional estimate would be $\beta \geqslant 5$ MeV. This would mean that the decoherence time $\hbar/\beta$ is much smaller than the average lifetime $\hbar/\Gamma$, resulting in the RMT limit $\Gamma \ll \beta$ in sharp contrast with the data shown in Fig. 1.

It should be noted that, for $\theta \gg |\Phi|$, the position of the first peak in the TPS (3) is determined by the scattering angle: $t_{\mathrm{peak}} \simeq \theta/\omega = 4.9\theta \times 10^{-22}$ s. This offers the possibility for an unambiguous check of the present interpretation by studying experimentally the θ dependence of $C(\varepsilon,\theta)$ and their Fourier components $P(t,\theta)$. In view of the sharp deviation of our interpretation from the conventional theory of highly excited many-body systems, such an experimental check is highly desirable.

The significant deviations from the RMT manifesting themselves in the oscillating EAFs have been also found for heavier systems [26]. This suggests that the quantum-classical correspondence in CQC may reflect general features of highly excited many-body systems rather than being of a system specific origin.

In conclusion we have proposed a novel method to study quantum-classical correspondence in highly excited many-body systems. The formation of classical-like wave packets and slow decoherence are shown to be intimately related to the significant deviations from RMT. The quantum-classical correspondence has been illustrated by analyzing ^{12}C+^{24}Mg elastic and inelastic collisions. New experiments are being proposed for an unambiguous test of the quantum-classical correspondence in CQC.

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