I. Introduction and motivation for the experiment

For highly excited interacting many-body systems the independent particle picture has very little validity when the mean spacing between the many-body energy levels is much smaller than the spacing of the single-particle levels [1,2]. For high excitations, the interaction leads to a quick decay of the single-particle as well as of collective modes [3] which are not eigenstates of the total Hamiltonian of the system. This decay results in a formation of highly complicated many-body configurations. Each of these many-body states is characterized by uniform occupation of all accessible parts of phase space and sharing of energy between many particles of the system. The characteristic time for the formation of such ergodic, independent of the initial conditions, many-body states is given by the inverse spreading width, $\tau_{\text{erg}} = h/\Gamma_{\text{spr}}$ [3].

Consider a highly excited many-body system, whose spectrum obeys Wigner-Dyson statistics, for the time interval $t \gg \tau_{\text{erg}}$. Does ergodicity of all individual many-body eigenstates necessarily imply that their superposition is incoherent random superposition? Can superposition of spatially extended ergodic modes of a highly excited many-body system produce localized or non-equilibrium non-ergodic patterns? In the absence of a theory for phase randomization in an isolated (disconnected from a heat bath) systems one may conventionally rely on the hypothesis emerging from the foundations and modern developments of the random matrix theory (RMT) of highly excited many-body systems. This hypothesis implies that the energy relaxation, i.e. the mere formation of ergodic individual many-body configurations, is a sufficient condition for a phase randomization between these ergodic eigenstates [3]. If true, this conjecture should validate universal applicability of the RMT for the energy interval $\Delta E \leq \Gamma_{\text{spr}}$ and for the time interval $t \geq \tau_{\text{erg}}$, accordingly.

Consider the decay of a highly excited many-body system with strongly overlapping resonances, $\Gamma \gg D$, where $h/\Gamma \gg \tau_{\text{erg}}$ is the average life-time and $D$ is the mean level spacing of the system. This regime, $D \ll \Gamma \ll \Gamma_{\text{spr}}$, is known as a regime of Ericson fluctuations [4,5] for the decay of equilibrated nuclear, atomic and molecular systems and in coherent electron transport through nanostructures [3]. Suppose that RMT universally applies for $t \gg \tau_{\text{erg}}$. This implies absence of correlations between transition amplitudes (partial width amplitudes), known as Bethe’s random signs hypothesis, for the decay of different ergodic states to either the same or different quantum micro-channels [3]. Consider, e.g., a strongly dissipative heavy-ion collision (DHIC) characterized by a high intrinsic excitation energy ($\geq 15$ MeV) of the double (deformed) intermediate system. Since for nuclear systems $\Gamma_{\text{spr}} \simeq 5$ MeV and, for DHIC, $\Gamma \simeq 100$ keV [6] we deal with the decay of a superposition of ergodic strongly overlapping ($\Gamma \gg D$) many-body configurations. Then the RMT hypothesis, $\tau_{\text{deph}} \leq \tau_{\text{erg}}$ with $\tau_{\text{deph}}$ being the phase randomization (dephasing) time between ergodic states, implies that the cross sections for the DHIC, summed over a very large number of partial cross sections, corresponding to different micro-states of the reaction fragments, should show a smooth energy dependence with the characteristic energy variation $\Gamma_{\text{spr}} \simeq 5$ MeV. In contrast, experimental studies [7-9] present overwhelming evidence for the persistence of rapid ($\simeq 100$ keV) energy oscillations in the cross sections for DHIC. This manifests itself in the correlations between different transition amplitudes indicating that the phase randomization between individual ergodic configurations should be a much slower process than energy relaxation ($\tau_{\text{deph}} \gg \tau_{\text{erg}}$) in sharp contrast with the RMT hypothesis.

In attempting to interpret the non-self-averaging of excitation function oscillations in DHIC one faces a non-straightforward task of realization of Wigner’s dream [10], namely to modify RMT by taking into account level-level and channel-channel correlations between the transition amplitudes. Such a possible modification has been presented in Refs. [11,12] in terms of spontaneous coherence and slow phase randomization in highly excited many-body systems. While RMT develops “a new statistical mechanics” (Dyson) of, by purpose, fully equilibrated finite systems, the work [14,15] discusses critical and non-equilibrium phenomena in finite highly excited many-body systems.

It has been found [11-13] that a precondition for micro-channel correlations (MC) in complex quantum collisions is $\tau_{\text{deph}} \gg \tau_{\text{erg}}$. Physically, $\tau_{\text{deph}}$ sets up a new time scale for quantum many-body systems. For times $t < \tau_{\text{deph}}$, the RMT ceased to apply even though $t \approx h/\Gamma \gg \tau_{\text{erg}}$. Since the physical picture [11-13] for the MC is in a sharp contrast with the RMT and the theory of quantum chaotic scattering [3] it is highly desirable to have an additional independent test of the approach [11-13]. Such a possibility does indeed exist. It has been argued [14] that the spontaneous origin of MC should result in the cross sections for DHIC being sensitive to an infinitesimally small perturbation.

There is convincing evidence that the effects of complexity and stochasticity in nuclear systems are
shared by other microscopic and mesoscopic many-body systems [3]. Therefore the spontaneous MC and the extreme sensitivity should be expected for other complex quantum collisions, e.g., atomic, molecular, and atomic cluster collisions.

The discussion of Ref. [14] has not taken into account different distributions of electro-magnetic fields, defects etc. within different independently prepared target foils. How might the consideration [14] apply in the presence of such differently distributed “target-environmental” perturbations within different targets?

Consider a simple case of spinless reaction partners in the entrance and exit channels. A generalization for the case of the reaction partners having intrinsic spins is straightforward. Then the measured cross section, per a single target nucleus and for a fixed single exit micro-channel \( \bar{b} \) (microscopic states of the reaction products), is given by

\[
\sigma_b(E, \theta) = (1/\mathcal{N}) \sum_{j=1}^{\mathcal{N}} \sigma_b^{(j)}(E, \theta),
\]

where

\[
\sigma_b^{(j)}(E, \theta) = |f_b^{(j)}(E, \theta)|^2.
\]

Here index \((j)\) labels individual target nuclei participating in the collision, whose number is \(\mathcal{N} \gg 1\), and \(f_b^{(j)}(E, \theta)\) is the amplitude of a collision involving \((j)\) target nucleus. The difference between \(f_b^{(j)}(E, \theta)\) with different \((j)\) originates from a nonuniform distribution of “target-environmental” perturbations. This introduces different local perturbations, \(V_j \neq V_i\), in the purely nuclear Hamiltonian \(H\) of highly excited nuclear molecules created in the collision of the incident ion with different \((j \neq i)\) target nuclei. We evaluate the strength of the “target-environmental” perturbations to be of the order of the atomic electron effects [14] in DHIC. We employ the perturbation theory [14] and use the decomposition \(f_b^{(j)}(E, \theta) = f_b(E, \theta) + \delta f_b^{(j)}(E, \theta)\), where \(f_b(E, \theta)\) is the collision amplitude in the absence of the “target-environmental” perturbations. We also drop the incoherent sum \((1/\mathcal{N}) \sum_{j=1}^{\mathcal{N}} |\delta f_b^{(j)}(E, \theta)|^2\). This sum is about fourteen orders of magnitude smaller than \(\sigma_b(E, \theta)\). We obtain

\[
\sigma_b(E, \theta) = |f_b(E, \theta)|^2 - (1/\mathcal{N}) \sum_{j=1}^{\mathcal{N}} |\delta f_b^{(j)}(E, \theta)|^2 \rightarrow |F_b(E, \theta)|^2,
\]

where \(|(1/\mathcal{N}) \sum_{j=1}^{\mathcal{N}} |\delta f_b^{(j)}(E, \theta)|^2 \leq 10^{-14}|F_b(E, \theta)|^2\), and \(F_b(E, \theta)\) is the collision amplitude corresponding to the Hamiltonian \((H+v)\) with \(v = (1/\mathcal{N}) \sum_{j=1}^{\mathcal{N}} V_j\).

It is reasonable to assume that a distribution of the local “target-environmental” perturbations \(V_j\) is random throughout the target. This means that \(\delta f_b^{(j)}(E, \theta)\) with different \((j)\) have random phases. In this case we have

\[
|F_b(E, \theta) - f_b(E, \theta)| \sim (1/\mathcal{N})^{1/2}|\delta f_b^{(j)}(E, \theta)| \sim (1/\mathcal{N})^{1/2}10^{-7}|f_b(E, \theta)|,
\]

where we used the estimate \(|\delta f_b^{(j)}(E, \theta)| \sim 10^{-7}|f_b(E, \theta)|\) from Ref. [14].

Suppose we perform two independent measurements with two different targets. The “target-environmental” perturbations, \(V_j\) in the first target and \(V_i\) in the second one, are different. The cross sections are given by the different amplitudes, \(F_b(E, \theta)\) and \(\tilde{F}_b(E, \theta)\), corresponding to different Hamiltonians, \((H+v)\) and \((H+v')\), accordingly. Then we have

\[
|F_b(E, \theta) - \tilde{F}_b(E, \theta)| \sim (1/\mathcal{N})^{1/2}10^{-7}|F_b(E, \theta)|.
\]

Therefore one does not expect a detectable difference for the cross sections measured with two different targets. Indeed, such a detectable difference does not occur if one considers \(\sigma_b(E, \theta)\) for a single fixed \(\bar{b}\) independently from the cross sections for the decay to other \(\bar{b}' \neq \bar{b}\) micro-channels. However, as suggested in Ref. [14], the situation may change drastically for the cross sections summed over very large number of exit micro-channels. This is the case for DHIC where the collision products have high excitation energies and the measured cross section, \(\sigma(E, \theta) = \sum_{\bar{b}} \sigma_b(E, \theta)\), is the sum over very large number of micro-channels, \(N_b \gg 1\).

The above consideration suggests that the spontaneous MC might lead to up to 100% difference between the non-self-averaging oscillating components of the cross sections for DHIC for two measurements with different targets. The key element in the interpretation of the spontaneous coherence, non-self-averaging and extreme sensitivity in complex quantum collisions is introduction of the infinitesimally small off-diagonal MC between different model transition amplitudes which couple model single-particle states (Slater determinants) of the quasi-bound IS and the continuum states [11-14]. It has been argued that the limit of the vanishing of this infinitesimally small correlation properly supplemented by the limit of the infinite dimensionality of the Hilbert space does not destroy correlation between different physical transition amplitudes which couple the
many-body configurations of the IS and the continuum states. As a result, the highly-excited thermalized ($\hbar/\Gamma \gg \tau_{\text{erg}}$) matter displays coexistence of two distinct phases. The decay of the disordered phase is associated with the $\Delta S_j^J$-matrix [11], where $J$ is the total spin of the IS, and, thereby, with the amplitude $\Delta F_j(E, \theta)$ which is a linear combination of $\Delta S_j^J$ with different $J$. Since $\Delta F_j(E, \theta)$ with different $\hbar \neq \hbar'$ do not correlate, this disordered phase does not contribute to the MC producing the stable reproducible self-averaging, i.e. energy smooth, background in cross sections. The non-self-averaging, i.e. micro-channel correlations, and sensitivity originate from decay of the ordered phase corresponding to the micro-channel independent $\delta S^J$-matrix [13,14] and, thereby, the micro-channel independent $\delta F(E, \theta)$. It is this micro-channel independent $\delta F(E, \theta)$ which is so sensitive and, therefore, non-reproducible due to the spontaneous origin of the MC so that $|\delta F(E, \theta) - \delta \tilde{F}(E, \theta)| \sim |\delta F(E, \theta)| \sim |F_0(E, \theta)|$, where $\delta F(E, \theta)$ and $\delta \tilde{F}(E, \theta)$ correspond to different targets with different distributions of “target-environmental” perturbations.

It follows from the above consideration and Ref. [14] that the non-self-averaging energy oscillating component of the cross section is determined by the “target-environmental” perturbations averaged over the whole target and not by the micro-channel differences unless these differences produce different perturbations of the Hamiltonian of hot double intermediate system. Indeed the detailed energy dependence of the non-self-averaging energy oscillating component of the cross section is determined by the amplitude $\delta F(E, \theta)$ which does not depend on the micro-channel indices. The case when the micro-channel differences in the entrance channel do produce different perturbations of the Hamiltonian of the intermediate system is considered in Ref. [14].

Pictorially, the sensitivity of the $\delta S^J$-matrix and $\delta F(E, \theta)$ resembles the sensitivity of the direction of the spontaneous magnetization vector, below the Curie point, to the direction of an infinitesimally small external magnetic field.

II. Experimental method

In order to test the sensitivity two independent measurements of excitation functions for the strongly dissipative collision for the same reaction system of $^{19}\text{F}^{+}\text{Nb}$ have been carried out at the China Institute of Atomic Energy (CIAE), Beijing. In these measurements, the $^{19}\text{F}^{+}$ beam was provided by the HI-13 tandem accelerator. The beam incident energies were varied from 102 to 108 MeV in steps of 250 keV. For both measurements the same accelerator parameters and the same electronic and acquisition systems were selected. The same two sets of gas-solid $(\Delta E - E)$ telescopes, with a charge resolution $Z/\Delta Z \geq 30$ and an energy resolution $\Delta E \leq 0.4$ MeV, were set at 38° and 53°. The $\Delta E$ detector is an ionization chamber filled with P10 gas at a pressure of 103 mb, the residual energy $E$ is deposited in a Si position sensitive detector with a thickness of 1000 $\mu$m, a size of 8×47 mm and a marked position resolution of 0.5 mm. The solid angles of the two telescopes are 1.80 msr and 2.62 msr, respectively. Count rates in the experiment were less than 10 counts per sec. so that a pile-up problem did not occur.

In Fig. 1 we present a typical $(\Delta E - E)$ scatter-plot obtained at $E_{\text{lab}}(^{19}\text{F})=100.25$ MeV. It is seen that the projectile-like fragments from the $^{19}\text{F}^{+}\text{Nb}$ reaction can be separated. For the F fragments direct and quasi-elastic processes constitute the major contribution into the cross section. For the Ne fragments there was no a sufficient statistics. Therefore we restrict our analysis to the cross sections of the N and O products of the $^{19}\text{F}^{+}\text{Nb}$ DHIC.

To avoid a possible effect of the carbon build up in the target, we analyse events with $E_{\text{lab}}(\text{N}) \geq 50$ MeV and $E_{\text{lab}}(\text{O}) \geq 55$ MeV for $\theta_{\text{lab}} = 38^\circ$, and with $E_{\text{lab}}(\text{N}) \geq 40$ MeV and $E_{\text{lab}}(\text{O}) \geq 40$ MeV for $\theta_{\text{lab}} = 53^\circ$ (see Fig. 1). In Fig. 2 we present a $(\Delta E - E)$ scatter-plot for the fragments from the $^{19}\text{F}^{+}^{12}\text{C}$ reaction at $E_{\text{lab}}(^{19}\text{F}) = 100.25$ MeV. Our measurement shows that the cross sections are negligible for the N and O outgoing energies $\geq 45$ MeV for $\theta_{\text{lab}} = 38^\circ$, and $\geq 40$ MeV for $\theta_{\text{lab}} = 53^\circ$. We also note that, for $E_{\text{lab}}(^{19}\text{F}) = 108$ MeV and $\theta_{\text{lab}} = 53^\circ$, the production of the N and O fragments with the outgoing energy $\geq 39$ MeV in the $^{19}\text{F}^{+}^{12}\text{C}$ reaction is kinematically forbidden. Since the energies of the N and O yields in our measurements $\geq 50$ MeV for $\theta_{\text{lab}} = 38^\circ$ and $\geq 40$ MeV for $\theta_{\text{lab}} = 53^\circ$ we conclude that the carbon build up does not produce uncontrolled errors and does not affect our data for the cross sections of the N and O products of the $^{19}\text{F}^{+}\text{Nb}$ DHIC.

In Fig. 3 we present, as an example, a typical energy spectrum of the dissipative N yield for $E_{\text{lab}}(^{19}\text{F})=103.25$ MeV and $\theta_{\text{lab}} = 38^\circ$ produced in the $^{19}\text{F}^{+}\text{Nb}$ DHIC for one run in the second experiment (see Fig. 5). Therefore the counting rate for the correspondent data in Fig. 6 is about as twice as higher than that in Fig. 3.
In Fig. 4 we present angular distributions for the N outgoing fragments measured at \( E_{\text{lab}}(^{19}\text{F})=100 \) MeV and 105 MeV. The angular distributions are strongly forward peaked due to the major contribution of direct fast reaction processes at the forward, near grazing, angles.

In the two measurements we used different, independently prepared, self-supporting \(^{93}\text{Nb}\) target foils with the thickness \( \simeq 70 \, \mu\text{g/cm}^2 \). Both the target foils were produced by the sputtering method. The thickness of each of the two foils was determined by the spectrophotometry. It was found that the difference in thickness of the two foils \( \leq 5 \, \mu\text{g/cm}^2 \). This difference results in different stopping energy losses in the two different targets. However, this itself should not affect reproducibility of the cross sections since this difference in stopping energy losses \( \sim 15 \, \text{keV} \) is smaller than the energy spread \( \sim 50 \, \text{keV} \) in the beam and additional energy spread \( \sim 150 \, \text{keV} \) in the target.

Absolute cross sections were not determined, though great care was taken to ensure no spurious sources of oscillations were introduced into the relative cross sections. The stability of the beam direction was controlled as follows: (i) TV monitor screen was used before each energy step to check and correct the position of the beam spot on the target. (ii) Two silicon detectors were placed at \( \theta_{\text{lab}} = \pm 12^\circ \). (iii) The beam charge was collected using a Faraday cup placed at \( \theta = 0^\circ \) and was compared with the counting rates of the silicon detectors. The data were normalized both with respect to the count rates of each of the silicon detectors and the integrated beam current. All the three normalizations produced the relative cross sections, for each individual experiment, which agree within the statistical errors, \( 1/N^{1/2} \), where \( N \) is a count rate. We have taken 5 repeat points (one repetition for 5 different energies measured) for the first experiment (target) and 21 repeat points (one repetition for 21 different energies) for the second experiment (target). Before to repeat each point the TV monitor screen was used to check and correct a position of the beam spot. All the repeated points demonstrated the reproducibility, within the statistical errors, for both individual experiments (targets). This reproducibility is demonstrated in Fig. 5 for the two runs in the second experiment. Such a reproducibility for the two runs for the same targets indicates that no damages of the targets, which could bring about uncontrolled spurious effects, occurred in our experiments. All the above procedures indicate that the systematic uncertainties do not seem to be present and the data errors can be evaluated as statistical only.

### III. Experimental results

The cross sections \( \sigma(E) \) for the products N and O in the \(^{19}\text{F}+^{93}\text{Nb}\) DHIC are presented in Fig.6, where the error bars are statistical only. Although Fig. 6 presents energy integrated yields over the wide, \( \sim 25 \) MeV, ranges of the dissipative spectra (i.e. these yields are summed over huge number of different final micro-channels of the highly excited collision products) the characteristic non-self-averaging oscillating structures of the excitation functions in DHIC can be visually identified.

Taking into account that an energy resolution of our \( (\Delta E - E) \) telescopes \( \leq 0.4 \) MeV, from Fig. 3 we find that possible cross section energy variations due to the lower energy cut-off are \( <1\% \) for the dissipative N yield at \( \theta_{\text{lab}} = 38^\circ \). From the energy spectra of the dissipative O yield we found that the lower energy cut-off also produces negligible, \( <1\% \), cross section energy variations for the O reaction products at \( \theta_{\text{lab}} = 38^\circ \).

From Fig. 6 we notice that, for some incident energies, the cross sections measured for two different target foils are different. A statistical significance of this non-reproducibility is discussed in Section IV.

For \( E_{\text{lab}}(^{19}\text{F})=105 \) MeV, the total excitation energy of the double intermediate system is \( E=87 \) MeV. It consists of the deformation energy \( E_{\text{def}} \), the rotational energy \( E_{\text{rot}} \) and the intrinsic excitation energy \( E^* \). The deformation energy is mainly given by the Coulomb energy of the two touched ions which yields \( E_{\text{def}} \simeq 43 \) MeV. We calculate the average rotational energy for a rigid body moment of inertia of the two touched ions with the average angular momentum \( J = (J_{\text{cr}} + J_{\text{gr}})/2 \), where \( J_{\text{cr}} \) and \( J_{\text{gr}} \) are the critical and the grazing angular momenta, respectively. In our case, \( J_{\text{cr}} = 40 \) and \( J_{\text{gr}} = 53 \) in h units. We have \( E_{\text{rot}}=27 \) MeV and \( \hbar \omega = 1.2 \) MeV, where \( \omega \) is the average angular velocity of the double intermediate system. We have \( E^*=17 \) MeV which corresponds to the average level spacing \( D \sim 10^{-11} \) MeV and the total width for evaporation from the excited double intermediate system \( \Gamma \leq 0.1 \) keV (see Fig. 7 in Ref. [4]). Accordingly, the average time it takes for the hot intermediate system to evaporate one nucleon is about \( 6 \times 10^{-18} \) sec. This corresponds to \( \sim 2000 \) complete revolutions of the intermediate system with \( \hbar \omega = 1.2 \) MeV. This is about three orders of magnitude larger than a typical average number (\( \sim 1 - 3 \)) of complete revolutions of the hot double intermediate system before its disintegration into two fragments [7,8,9]. This indicates that the production of the projectile like ejectiles is a primary binary process which is not affected by the nucleon
IV. Tests of a statistical significance of the data

One possibility to find out if the oscillations in the individual excitation functions measured for each of the two different targets (Fig. 6) are true oscillations is to calculate the experimental normalized variances of the oscillations, $C(\varepsilon = 0)$. Here $C(\varepsilon) = \langle \Delta \sigma(E + \varepsilon) \Delta \sigma(E) \rangle$ is a cross section energy autocorrelation function, $\Delta \sigma(E) = (\sigma(E) / \sigma(E) > -1)$ is a relative oscillating yield, and $\langle \sigma(E) \rangle$ is an energy averaged smooth cross section which was obtained from the best second order polynomial fit of the original data. For the two independent measurements of the N oscillating yields (Fig. 6) at $\theta = 53^\circ$ we obtain $C(\varepsilon = 0) = 0.015 \pm 0.0035$ for the first target and $C(\varepsilon = 0) = 0.017 \pm 0.004$ for the second one, where the uncertainties are due to the finite data range only [5]. For the O oscillating yields at $\theta = 53^\circ$ we have $C(\varepsilon = 0) = 0.012 \pm 0.003$ for the first target and $C(\varepsilon = 0) = 0.016 \pm 0.0037$ for the second one. This is to be compared with the quantities $1/N$, which represent $C(\varepsilon = 0)$ corresponding only to statistical uncertainties due to the finite average counting rate $N$. For the N yield we have $1/N = 0.004$ and for the O yield $1/N = 0.0035$. Therefore, for $\theta = 53^\circ$, the experimental values of $C(\varepsilon = 0)$ are larger by a factor of 3 than $1/N$ expected based on finite statistics. Similarly, for the two independent measurements of the N oscillating yields (Fig. 6) at $\theta = 38^\circ$ we obtain $C(\varepsilon = 0) = 0.0024 \pm 0.0006$ for the first target and $C(\varepsilon = 0) = 0.0028 \pm 0.0007$ for the second one. For the O oscillating yields at $\theta = 38^\circ$ we have $C(\varepsilon = 0) = 0.0024 \pm 0.0006$ for the first target and $C(\varepsilon = 0) = 0.0022 \pm 0.00055$ for the second one. These values are larger by a factor of 3 than corresponding average inverse counting rates $(1/N = 0.0008$ for the N products and $1/N = 0.0007$ for the O products) at $\theta = 38^\circ$. The above analysis indicates that the oscillations shown in Fig. 6 are true oscillations and do not result from insufficient statistics.

Another indication for the statistical significance of the oscillations in Fig. 6 can be revealed from the analysis of probability distributions of the properly scaled cross section relative deviations, $(\sigma_i/ < \sigma_i > -1)/(1/N_i)^{1/2}$, from the energy smooth background $< \sigma(E) >$. Here, $\sigma_i = \sigma(E_i)$, $< \sigma_i > = < \sigma(E_i) >$ is an energy averaged smooth cross section obtained from the best second order polynomial fit of the data, and $N_i$ is the counting rate for the $E_i$ energy step. Suppose that the cross section energy oscillations in Fig. 6 are not true oscillations but originate from the finite count rate. If this would be the case then the probability distribution of $(\sigma_i/ < \sigma_i > -1)/(1/N_i)^{1/2}$ should be a Gaussian distribution with zero expectations and unit standard deviation (variance). Gaussian distributions and the actual probability distributions of absolute values of the measured cross section deviations from the energy smooth background are presented in Figs. 7, 8 and 9. One observes that the experimental probability distributions are systematically wider than Gaussian distribution with unit standard deviation. Also 21% of all the deviations exceed three standard deviations (Fig. 9).

As a first step in evaluation of the statistical significance of the non-reproducibility (Fig. 6) we calculate correlation coefficients, $\varrho$, between the corresponding oscillating yields produced with the two different target foils:

$$\varrho = (1/n) \sum_{i=1}^{n} (\sigma_i^{(1)} / < \sigma_i^{(1)} > -1)(\sigma_i^{(2)} / < \sigma_i^{(2)} > -1)/[C^{(1)}(\varepsilon = 0)C^{(2)}(\varepsilon = 0)]^{1/2},$$

where

$$C^{(1,2)}(\varepsilon = 0) = (1/n) \sum_{i=1}^{n} (\sigma_i^{(1,2)} / < \sigma_i^{(1,2)} > -1)^2,$$

$\sigma_i^{(1,2)} = \sigma^{(1,2)}(E_i)$, $< \sigma_i^{(1,2)} > = < \sigma^{(1,2)}(E_i) >$, and $n$ is a number of energy steps. Indices (1, 2) correspond to the first and second measurement (target), accordingly. We find $\varrho = 0.24 \pm 0.06$ for the N products at $\theta_{lab} = 38^\circ$, $\varrho = 0.23 \pm 0.06$ for the O products at $\theta_{lab} = 38^\circ$, $\varrho = 0.09 \pm 0.022$ for the N products at $\theta_{lab} = 53^\circ$ and $\varrho = 0.06 \pm 0.015$ for the O products at $\theta_{lab} = 53^\circ$, where the uncertainties are due to the finite data range only [5]. This indicates that the non-self-averaging non-reproducible components of the cross sections for the two measurements oscillate around each other in nearly statistically independent uncorrelated way.

Consider a probability distribution of

$$[\sigma_1(E) - \sigma_2(E)]/[(\delta \sigma_1^2 + \delta \sigma_2^2 + 2\rho \delta \sigma_1 \delta \sigma_2)]^{1/2},$$

where $\rho$ is the correlation coefficient between the cross section energy oscillations in Figs. 6 and 7 at $\theta = 53^\circ$. The observed deviations are too large by a factor of 3 than the statistical uncertainties inside the finite average counting rate only [5].
where

\[ \delta\sigma_{1,2}^2 = \langle 1/n \sum_{i=1}^{n} (\sigma_{i}^{(1,2)} - <\sigma_{i}^{(1,2)}>)^2, \]

\[ \rho = \langle 1/n \sum_{i=1}^{n} (\sigma_{i}^{(1)} - <\sigma_{i}^{(1)}>)(\sigma_{i}^{(2)} - <\sigma_{i}^{(2)}>) / \delta\sigma_{1}\delta\sigma_{2}, \]

and indices (1, 2) stand for the first and second measurement (target), respectively.

Suppose that the non-reproducibility of the cross section energy oscillations in Fig. 6 is not a true effect but originate from the finite count rates. If this would be the case then the probability distribution of the quantity (4) with

\[ \delta\sigma_{1,2}^2 = \langle 1/n \sum_{i=1}^{n} <\sigma_{i}^{(1,2)}>^2 / N_{i}^{(1,2)}, \]

should be a Gaussian distribution with zero expectation and unit standard deviation (variance). In Eq. (7) \( N_{i}^{(1,2)} \) are the counting rates for the \( E_{i} \) energy step in the first and second measurements, accordingly. Gaussian distributions with unit standard deviation and the actual experimental probability distributions of absolute values of the quantities (4) with \( \delta\sigma_{1,2}^2 \) given by Eq. (7) are presented in Figs. 10 and 11. One observes that the experimental probability distributions are systematically wider than Gaussian distribution with unit standard deviation. A level of the non-reproducibility exceeds three standard deviations for 18% of all the cross section differences measured (Fig. 11). This indicates that the non-reproducibility of the cross section energy oscillations (Fig. 6) measured with different target foils of nominally the same thickness is of a statistical significance.

On the contrary, the two runs for the same target foil (Fig. 5) produce the reproducible cross section energy oscillations (see Figs. 10 and 11).

Our data indicate a strong correlation between the N and O dissipative yields for each of the individual measurement (target). For example, for \( \theta_{lab} = 38^\circ \), a correlation coefficient between the N and O cross sections is \( \rho = 0.6 \pm 0.15 \) for the first target and \( \rho = 0.61 \pm 0.15 \) for the second one, where the uncertainties are due to the finite data range only [5].

Such a strong (\( \approx 0.5 \sim 0.9 \)) correlation between strongly dissipative yields with different charges or different masses is a characteristic feature of most of the systems measured [7-9]. This suggests that our data, for each individual measurement (target), are of a similar character as those reported in Refs. [7-9]. A strong correlation between dissipative yields with different charges and different masses is consistent with the interpretation of the MC [9,14].

Measurements of the excitation function oscillations, for the single target foil, with distinction of isotopes for the yields from the \(^{19}\text{F}^{+}\text{V}\) strongly dissipative collisions were reported in ref. [7], Wang Qi et al. (1996). These data demonstrate a strong, \( \approx 0.8 \sim 0.9 \), correlation between different isotopes. Assuming that the present data, for each of the two individual measurements (targets), are of a similar character as those reported in ref. [7], Wang Qi et al. (1996) one would expect that a level of non-reproducibility for different target foils for isotopes of N and O would be similar to that observed without distinction of isotops (Fig. 6).

It is clear that the random matrix theory and statistical theory of Ericson fluctuations are of no help for the interpretation of the experimental results reported in this paper. Indeed, the theory of Ericson fluctuations is conceptually based on the statistical model which disregards outright micro-channel correlations [3,4]. Therefore, the necessary conditions for applicability of the random matrix theory, statistical model and theory of Ericson fluctuations to the interpretation of the data reported here must be (i) absence of oscillations in the cross sections, i.e. energy smooth excitation functions for each of the individual measurement, and (ii) reproducibility of these energy smooth cross sections in the measurements with different target foils. Both of these conditions are not met for the data sets reported in this paper.

A quantitative interpretation of the energy oscillations in the individual data sets in terms of the spontaneous self-organisation, non-equilibrium micro-channel correlation phase transitions and anomalously slow phase randomization in highly excited strongly interacting finite quantum many-body systems [15] will be presented in a separate communication.

V. Conclusion
In conclusion, the two independent measurements with different target foils of nominally the same thickness indicate statistically significant non-reproducibility of the cross sections for the $^{19}\text{F} + ^{93}\text{Nb}$ DHIC. The non-reproducibility is consistent with the recent theoretical arguments on spontaneous coherence, slow phase randomization and anomalous sensitivity in finite highly excited quantum systems. If this non-reproducibility is confirmed in future experiments it will signal that a realization of Wigner’s dream [10], a theory for the transition amplitude correlations, will require conceptual revision of modern understanding of microscopic and mesoscopic quantum many-body systems.

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Figure Captions

Fig. 1. $\Delta E - E$ scatter-plots obtained in the $^{19}\text{F} + ^{93}\text{Nb}$ dissipative heavy-ion collision at $\theta_{lab} = 38^\circ$ (left panel) and $\theta_{lab} = 53^\circ$ (right panel) at $E_{lab} = 100.25$ MeV. The Fig. also shows energy gates used in the analysis.

Fig. 2. $\Delta E - E$ scatter-plots obtained in the $^{19}\text{F} + ^{12}\text{C}$ reaction at $\theta_{lab} = 38^\circ$ (left panel) and $\theta_{lab} = 53^\circ$ (right panel) at $E_{lab} = 100.25$ MeV.

Fig. 3. Energy spectrum of $Z=7$ dissipative fragments produced in the $^{19}\text{F} + ^{93}\text{Nb}$ dissipative heavy-ion collision at $\theta_{lab} = 38^\circ$ and $E_{lab} = 103.25$ MeV for one run in the second experiment (see Fig. 5).

Fig. 4. Angular distributions of N dissipative yield of the $^{19}\text{F} + ^{93}\text{Nb}$ dissipative heavy-ion collision at $E_{lab} = 100$ MeV and 105 MeV. The solid and dashed lines are for the eye guide.

Fig. 5. Excitation functions for the N and O yields of the $^{19}\text{F} + ^{93}\text{Nb}$ strongly dissipative heavy-ion collisions obtained in the two runs (triangles and crossed circles) for the same single target foil in the second experiment. The error bars are statistical only.

Fig. 6. Excitation functions for the N and O yields of the $^{19}\text{F} + ^{93}\text{Nb}$ strongly dissipative heavy-ion collisions obtained in the two independent experiments. Full dots correspond to the first experiment and open squares to the second one. The error bars are statistical only.

Fig. 7. Probability distributions of absolute values of the cross section relative deviations from the energy smooth background obtained in the first measurement (dashed histograms). Solid histograms are Gaussian distributions with unit standard deviation expected based on the finite count rates only (see text).

Fig. 8. The same as in Fig. 7 but for the second measurement with different target (see text).

Fig. 9. The same as in Figs. 7 and 8 but for the sum of all 8 sets of the individual probability distributions from Figs. 7 and 8.

Fig. 10. Probability distributions of absolute values of the properly scaled differences between the cross sections obtained in the two measurements with different target foils (dashed histograms). Dotted histograms are probability distributions of absolute values of the properly scaled differences between the cross sections obtained in the two runs with the same target foil for the second measurement. Gaussian distributions (solid histograms) with unit standard deviation expected based on the finite count rates only (see text).

Fig. 11. The same as in Fig. 10 but for the sum of all 4 sets of the individual probability distributions from Fig. 10.