Sensitivity of Nucleus-Nucleus Cross Sections and Atomic-Electron Effects in Dissipative Heavy-Ion Collisions

Sergey Yu. Kun
Department of Theoretical Physics, RSPhysSE, IAS, The Australian National University, Canberra ACT 0200, Australia and MPI für Kernphysik, Heidelberg D-69029, Germany

(Received 20 November 1998; revised manuscript received 30 June 1999)

We argue that spontaneous microchannel S-matrix correlations and slow spin decoherence result in a sensitivity of the cross sections for strongly dissipative heavy-ion collisions to an arbitrarily small perturbation. Such a sensitivity implies that atomic electrons should influence energetic ($\simeq$100 MeV) heavy-ion reactions. The atomic-electron effects are predicted to be $\simeq$100% of the magnitude of the non-self-averaged oscillating component of the nucleus-nucleus cross sections.

PACS numbers: 25.70.Lm, 05.45.–a, 24.60.Ky, 24.60.Lz

In the modern quantum theory of highly excited strongly interacting systems it is commonly assumed (see, e.g., Ref. [1]) that the decoherence time, i.e., the time it takes to lose all the initial phase correlations, is the shortest time scale of the problem. Application of this idea to the theory of quantum chaotic scattering and complex quantum collisions proceeding through the formation and decay of a highly excited intermediate system implies the absence of a correlation between the reaction amplitudes pertaining to different total spin and exit microchannel quantum numbers [1–4].

The idea of a rapid decoherence has been successfully developed and applied in many instances, in particular in the context of the random matrix theory [1]. Yet one also does not *a priori* expect its overall applicability. In particular, strong spin and microchannel correlations (MC) have been revealed [5–7] in strongly dissipative heavy-ion collisions (DHIC). These MC manifest themselves in the non-self-averaging of excitation function oscillations in these processes.

In this Letter we argue that the spontaneous MC and slow spin decoherence [8,9], which can explain [10] the non-self-averaging of excitation function oscillations, result in the cross sections for DHIC being sensitive to an infinitesimally small perturbation. In particular, we predict that atomic electrons should influence energetic ($\simeq$100 MeV) heavy-ion reactions. These atomic-electron effects in DHIC are estimated to be $\simeq$100% of the magnitude of the non-self-averaged oscillating component of the nuclear cross sections and therefore should be easily measured.

There is convincing evidence that the effects of complexity and stochasticity in nuclear systems are shared by other many-body systems [1]. Therefore the spontaneous MC and the extreme sensitivity of the cross sections should be expected for other complex quantum collisions, e.g., atomic, molecular, electron-ion, and atomic cluster collisions.

It follows from our consideration that the precondition for a sensitivity of the cross sections of DHIC is a finite spin decoherence width, i.e., the presence of quantum chaos and spin decoherence [8–10]. Physically this decoherence width determines a new energy scale for quantum many-body systems. This new scale is analogous to the Thouless energy in disordered systems. Beyond this scale the random-matrix theory of quantum many-body systems ceases to apply [1]. Our analysis gives a value of the decoherence width to be about 2–3 orders of magnitude less than the spreading width [1,11]. This implies that the time it takes to reach complete mixing within the available phase space or Hilbert space is 2–3 orders of magnitude longer than that evaluated from the characteristic values of spreading widths [1,11].

We treat the DHIC in terms of the formation and decay of a highly excited intermediate nuclear system (INS). The normalized S-matrix elements are taken to be of the form $\overline{S}^J_{\overline{a}\,\overline{b}}(E) = (\Gamma D/2\pi)^{1/2}\sum_\mu \overline{\gamma}^{J\overline{a}}_\mu \overline{\gamma}^{J\overline{b}}_\mu/(E - E^J_\mu + i\Gamma/2)$ [9,10]. Here, E is the total energy, J is the total spin, Γ and D are the total decay width and the average level spacing, respectively, and E^J_μ are the resonance energies of the highly excited ($D/\Gamma \ll 1$) INS. The $\overline{a}(\overline{b})$ indices specify intrinsic microstates of the reaction partners in the entrance (exit) channels, and $\overline{\gamma}^{J\overline{a}(\overline{b})}_\mu = \sum_j B^J_{\mu_j}\overline{\xi}^{J\overline{a}(\overline{b})}_j$, where $\overline{\gamma}^{J\overline{a}(\overline{b})}_\mu$ are the normalized partial width amplitudes for transitions between the entrance (exit) channel and the resonance states, ϕ^J_μ, due to the interaction V_n. The $\overline{\xi}^{J\overline{a}(\overline{b})}_j$ are the corresponding coupling amplitudes between the entrance (exit) channel and the Slater determinants X^J_j. Both $\overline{\gamma}^{J\overline{a}(\overline{b})}_\mu$ and $\overline{\xi}^{J\overline{a}(\overline{b})}_j$ are normalized random Gaussian variables with a mean value of zero. The quantities $B^J_{\mu_j}$ are the orthogonal matrices arising from the expansion $\phi^J_\mu = \sum_j B^J_{\mu_j}X^J_j$ of the eigenstates ϕ^J_μ of the INS over the X^J_j [9].

Let us consider the regime of slow spin decoherence, $g\beta \leq \Gamma$ [7,10], where β is the spin decoherence width and g is the number of partial waves contributing to the reaction. Following [10] we use the decomposition $\overline{S}^J_{\overline{a}\,\overline{b}}(E) = \Delta S^{J,I}_{\overline{a}\,\overline{b}}(E) + \delta S^{J,I}_{\overline{a}\,\overline{b}}(E)$, where I is the average spin and $\delta S^{J,I}_{\overline{a}\,\overline{b}}(E)$ and $\Delta S^{J,I}_{\overline{a}\,\overline{b}}(E)$ are given by Eqs. (12) and (13) in Ref. [10]. The correlation between $\overline{S}^{J}_{\overline{a}\,\overline{b}}(E)$ and $\overline{S}^{J}_{\overline{a}'\overline{b}'}(E)$ with $\overline{a} \neq \overline{a}'$ and/or $\overline{b} \neq \overline{b}'$ originates from the correlation between $\delta S^{J,I}_{\overline{a}\,\overline{b}}(E)$ and $\delta S^{J,I}_{\overline{a}'\overline{b}'}(E)$ [10]. From Eqs. (9), (14), (15), and (16) in Ref. [10] we obtain $\overline{\delta S^{J,I}_{\overline{a}\,\overline{b}}(E) \delta S^{J,I}_{\overline{a}'\overline{b}'}(E)^*} = \overline{|\delta S^{J,I}_{\overline{a}\,\overline{b}}(E)|^2} = \overline{|\delta S^{J,I}_{\overline{a}'\overline{b}'}(E)|^2} = 2|J - I|\beta/(\Gamma + 2|J - I|\beta)$, i.e., $\delta S^{J,I}_{\overline{a}\,\overline{b}}(E) = \delta S^{J,I}_{\overline{a}'\overline{b}'}(E) \equiv \delta S^{J,I}(E)$ for all $\overline{a} \neq \overline{a}'$ and $\overline{b} \neq \overline{b}'$. This MC originates from the infinitesimally small ($|\lambda| \to 0$) entrance-exit correlation $\overline{\overline{\xi}^{J\overline{a}}_i \overline{\xi}^{J'\overline{b}}_j} = \lambda K^{JJ'}_{ij}(\overline{a}, \overline{b})/N^{1/2}$ [8,9], where $K^{JJ'}_{ij}(\overline{a}, \overline{b})$ are symmetric matrices whose elements are of the order of ± 1, and $N \to \infty$ is the dimension of the Hilbert space. It should be noted that the $K^{JJ'}_{ij}(\overline{a}, \overline{b})$ matrices with different $(\overline{a}, \overline{b})$ are different and independent. Yet the $\delta S^{J,I}_{\overline{a}\,\overline{b}}(E)$ with different $(\overline{a}, \overline{b})$ are identical. Clearly this is possible only if the limit $|\lambda| \to 0$ is taken. Indeed then, and only then, the $K^{JJ'}_{ij}(\overline{a}, \overline{b})$ matrices with different $(\overline{a}, \overline{b})$ become indistinguishable so that we can consider these matrices to be $(\overline{a}, \overline{b})$ independent: $K^{JJ'}_{ij}(\overline{a}, \overline{b}) = K^{JJ'}_{ij}(\overline{a}', \overline{b}') = K^{JJ'}_{ij}$. This is a precondition for obtaining the $(\overline{a}, \overline{b})$ independence of the correlation coefficient $\rho^{JJ'}_{\mu\nu} = \overline{\gamma^{J\overline{a}}_\mu \gamma^{J'\overline{b}}_\nu}$. On the other hand, switching off the correlation between $\overline{\xi}^{J\overline{a}}_i$ and $\overline{\xi}^{J'\overline{b}}_j$ does not result in the vanishing of $\rho^{JJ'}_{\mu\nu}$ if the limit $|\lambda| \to 0$ is taken simultaneously with the limit $N \to \infty$ [8,9].

We observe that the correlation between $\delta S^{J,I}_{\overline{a}\,\overline{b}}(E)$ and $\delta S^{J,I}_{\overline{a}'\overline{b}'}(E)$ occurs spontaneously. Therefore the detailed energy dependence of the $\delta S^{J,I}(E)$ is set up at random, reflecting the extreme sensitivity of the cross sections to an arbitrarily small perturbation.

The sensitivity of the cross sections for DHIC implies that, e.g., atomic electrons should influence the nucleus-nucleus collision. In particular, we predict that the cross sections should be different for the same kinetic energy but for different charges of the ions in the incident beam. For example, Refs. [6,7] report measurements of the cross sections for $^{19}\mathrm{F}^{+8} + {}^{89}\mathrm{Y}$ strongly dissipative collisions. Let us consider the same collision but with the charge of the ^{19}F ions being +7. We neglect the total spin of the system of inner electrons moving in the nuclear Coulomb field without perturbing the INS. We also do not distinguish the positions of the c.m. of the INS and that of the whole system. In zero approximation the outer electron is moving in the nuclear Coulomb field screened by the inner electrons and this outer electron does not perturb the INS. We take the orbital momentum of the electron to be zero, $l = 0$, and also neglect the electron spin and the spins of the colliding nuclei. In zero approximation, the wave function of the system is $f^{J=J_n,M=0}_{J_n\mu,nl=0} = \varphi^{m_e=0}_{nl=0}\phi^{Jm_n=0}_\mu$. Here, $\varphi^{m_e=0}_{nl=0} \equiv \varphi_{nl=0}$ (n, l, m_e are the principal, orbital, and magnetic quantum numbers) is the electron Coulomb wave function, and $\phi^{Jm_n=0}_\mu \equiv \phi^J_\mu$ is the wave function of INS with the spin $J_n = J$ and its z projection $m_n = M = 0$, where the z axis is chosen along the direction of incident beam. A rotationally invariant perturbation introduced due to the electron-nuclear interaction has the form $V_{\mathrm{pert}} = (Z_{\mathrm{eff}}/Z)(e^2/R)(\sum_{p=1}^{Z} R/|\mathbf{R} - \mathbf{r}_p| - Z) \simeq V^{(1)} + V^{(2)}$, where $V^{(1)} = (Z_{\mathrm{eff}}e^2/RZ)\sum_p(\mathbf{R}\mathbf{r}_p)/R^2, V^{(2)} = (Z_{\mathrm{eff}}e^2/2RZ)\sum_p[3(\mathbf{R}\mathbf{r}_p)^2/R^4 - r_p^2/R^2]$. Here, $\mathbf{R}$ and $\mathbf{r}_p$ are the radius vectors of the outer electron and the protons, $Z_{\mathrm{eff}} = Z - N_e + 1$, Z is the charge of INS, and N_e is the total number of the electrons. The strength of the perturbation is determined by the matrix elements $\langle f^{J=J_n,M=0}_{J_n\mu,nl=0}|V^{(1)} + V^{(2)}|f^{J'M'}_{J'_n\mu',n'l'}\rangle = \delta_{JJ'}\delta_{0M'}(K^{(1)} + K^{(2)})$, where $f^{J'M'}_{J'_n\mu',n'l'}$ are given in terms of the electron wave function and the nuclear wave function coupled to total spin $J' = J$ with z projection $M' = 0$. The evaluation of the $K^{(1,2)}$ involves (i) the straightforward calculation of the electron matrix elements, and (ii) the calculation of the nuclear many-body matrix elements for the single-particle operators. These nuclear matrix elements can be calculated following the method [12]. We expand $\phi^{Jm_n}_\mu$ over the nuclear shell-model states. We further expand these shell-model states into products of single-particle proton states, and of the shell-model states of the rest of $A - 1$ nucleons. We calculate the single-particle proton matrix elements using an infinite well approximation. Applying the random-matrix theory [12] and neglecting the nucleon spin-orbital interaction we obtain

$$(K^{(1)})^2 \sim \delta_{J,|J'_n \pm 1|}\delta_{l',1}(Z_{\mathrm{eff}}/n)^6(e^2/a)^2(R_n/a)^2 D\Gamma_{\mathrm{spr}}/2\pi \times [(E^J_\mu - E^{J'}_{\mu'})^2 + \Gamma^2_{\mathrm{spr}}/4].$$

Here a is the Bohr radius, R_n is the radius of INS, and $\Gamma_{\mathrm{spr}} \sim 5-10$ MeV [11]. The above estimate is obtained for $n - n' = 1$, when the electron matrix elements have maximal values. For the intrinsic excitations $\sim 15-20$ MeV of the INS $^{19}\mathrm{F} + {}^{89}\mathrm{Y}$ we have $D \sim 10^{-10}$ MeV. Estimating $n \sim (3N_e)^{1/3}$ we obtain $K^{(1)}/D \sim \pm 10^{-4}$. We also calculate $K^{(2)}$, for which $V^{(2)}$ couples the nuclear states with $|J_n - J'_n| = 0, 2$. Neglecting the tunneling of the outer electron through the Coulomb barrier created by the inner electrons we obtain $K^{(2)} \sim (Z_{\mathrm{eff}}R_n/a)K^{(1)} \sim \pm 10^{-7}D$. Denoting $f^{J=J_n,M=0}_{J_n\mu,nl=0} = f^J_\mu$ (which corresponds to the $^{19}\mathrm{F}^{+8} + {}^{89}\mathrm{Y}$ DHIC) we have $\tilde{f}^J_\mu = f^J_\mu + \delta f^{J(1)}_\mu + \delta f^{J(2)}_\mu$, where $\tilde{f}^J_\mu$ is the wave function of the system obtained by taking into account the electron-nucleus interaction. Within our approximation, this is the wave function of the intermediate system formed in the $^{19}\mathrm{F}^{+7} + {}^{89}\mathrm{Y}$ DHIC. Using the first order of the perturbation theory we obtain $\langle \delta f^{J(1,2)}_\mu | \delta f^{J(1,2)}_\mu \rangle \sim (K^{(1,2)}/D)^2 \sim 10^{-8}, 10^{-14}$, respectively. In the absence of the electron-nucleus interaction, $\gamma^{J\overline{a}(\overline{b})}_\mu = \langle \chi^J_{\overline{a}(\overline{b})}\varphi_{nl=0}|V_n|f^J_\mu\rangle = \overline{V}_n\langle \chi^J_{\overline{a}(\overline{b})}|(V_n/\overline{V}_n)|\phi^J_\mu\rangle$, where $\overline{V}_n \sim 1$ MeV is a characteristic strength of the residual nuclear interaction and $\chi^J_{\overline{a}(\overline{b})}$ is the entrance (exit) channel wave function. In the presence of the electron-nucleus interaction, $\tilde{\gamma}^{J\overline{a}(\overline{b})}_\mu = \langle \chi^J_{\overline{a}(\overline{b})}\varphi_{nl=0}|V_n + V^{(1)} + V^{(2)}|\tilde{f}^J_\mu\rangle = \gamma^{J\overline{a}(\overline{b})}_\mu + \delta\gamma^{J\overline{a}(\overline{b})}_\mu + \Delta\gamma^{J\overline{a}(\overline{b})}_\mu$, where

$$\delta\gamma^{J\overline{a}(\overline{b})}_\mu = \langle \chi^J_{\overline{a}(\overline{b})}\varphi_{nl=0}|V_n|(\delta f^{J(1)}_\mu + \delta f^{J(2)}_\mu)\rangle = \langle \chi^J_{\overline{a}(\overline{b})}|V_n|\delta\phi^J_\mu\rangle,$$

$$\delta\phi^J_\mu = \langle \varphi_{nl=0}\,|\,\delta f^{J(2)}_\mu\rangle,$$

$$\Delta\gamma^{J\overline{a}(\overline{b})}_\mu = \langle \chi^J_{\overline{a}(\overline{b})}\varphi_{nl=0}|V^{(1)} + V^{(2)}|\tilde{f}^J_\mu\rangle \sim \langle \chi^J_{\overline{a}(\overline{b})}\varphi_{nl=0}|V^{(2)}|\varphi_{nl=0}\phi^J_\mu\rangle \sim (Z_{\mathrm{eff}})^4 \times (e^2/N_e a)(R_n/a)^2 \left\langle \chi^J_{\overline{a}(\overline{b})} \,\middle|\, \sum_p r_p^2/ZR_n^2 \,\middle|\, \phi^J_\mu \right\rangle.$$

Since $\langle \delta\phi^J_\mu\,|\,\delta\phi^J_\mu\rangle \sim (K^{(2)}/D)^2 \sim 10^{-14}$, then $\delta\gamma^{J\overline{a}(\overline{b})}_\mu \sim 10^{-7}\gamma^{J\overline{a}(\overline{b})}_\mu$, while $\Delta\gamma^{J\overline{a}(\overline{b})}_\mu < 10^{-10}\gamma^{J\overline{a}(\overline{b})}_\mu$, where we have taken into account that $(\sum_p r_p^2/ZR_n^2)$ is a smoother function of the nucleon coordinates than the residual interaction $(V_n/\overline{V}_n)$. Altogether we have $\tilde{\gamma}^{J\overline{a}(\overline{b})}_\mu = \langle \chi^J_{\overline{a}(\overline{b})}|V_n|\tilde{\phi}^J_\mu\rangle$, where $\tilde{\phi}^J_\mu = \phi^J_\mu + \delta\phi^J_\mu$ with $\langle \phi^J_\mu\,|\,\tilde{\phi}^I_{\mu'}\rangle = \delta_{JI}[\delta_{\mu\mu'} + (1-\delta_{\mu\mu'})\mathcal{O}(K^{(2)}/D)]$ and $K^{(2)}/D \sim \pm 10^{-7}$. Therefore one does not expect a detectable difference between the cross sections for $^{19}F^{+7} + {}^{89}Y$ and $^{19}F^{+8} + {}^{89}Y$ DHIC. Yet, because of the extreme sensitivity of the nucleus-nucleus cross sections, this difference is predicted to be $\simeq$100% of the magnitude of the non-self-averaged oscillating component of the cross section. In order to support this prediction we calculate the correlation coefficient, $\kappa = \overline{\delta S^{J,I}(E)\delta\tilde{S}^{J,I}(E)^*}/\overline{|\delta S^{J,I}(E)|^2}$, where $\delta S^{J,I}_{\overline{a}\,\overline{b}}(E) = (\Gamma D/2\pi)^{1/2}\sum_\mu R^{J,I}_\mu/(E - E^J_\mu + i\Gamma/2)$ corresponds to the $^{19}F^{+8} + {}^{89}Y$ DHIC and $\delta\tilde{S}^{J,I}_{\overline{a}\,\overline{b}}(E)$ to the $^{19}F^{+7} + {}^{89}Y$ DHIC. We rewrite R^{JI}_μ [see Eq. (8.3) in Ref. [8]] in the equivalent form

$$R^{JI}_\mu = (1/N^{1/2})\sum_k (u_k/N^{1/2}) \sum_{\nu\neq\nu'} q^{JI}_{\mu\nu}(k)q^{JI}_{\mu\nu'}(k)M^{I,k}_{\nu\nu'}, \tag{1}$$

where

$$q^{JI}_{\mu\nu}(k) = \langle \overline{\phi}^{J,k}_\mu\,|\,\overline{\phi}^{I,k}_\nu\rangle,$$

$$\overline{\phi}^{J(I),k}_{\mu(\nu)} = N^{1/2}\sum_i C^{J(I)}_{\mu(\nu)i}U_{ik}Y_i,$$

$$C^{J(I)}_{\mu(\nu)i} = \langle \phi^{J(I)}_{\mu(\nu)}\,|\,Y_i\rangle = \sum_j B^{J(I)}_{\mu(\nu)j}[T^{-1}]^{J(I)}_{j;i},$$

$$Y_i = \sum_{Jj} T^J_{i;j}X^J_j,$$

$$M^{I,k}_{\nu\nu'} = (N^{1/2}\langle \phi^I_\nu\,|\,Z_k\rangle)(N^{1/2}\langle \phi^I_{\nu'}\,|\,Z_k\rangle),$$

$$Z_k = \sum_i U_{ik}Y_i.$$

In Eq. (1), we have omitted the diagonal sum $(\sum_{\nu=\nu'}\ldots)$ [see Eq. (8.3) in Ref. [8]] since its contribution to $\delta S^{J,I}(E)$ is negligible provided that $\Gamma/D \ll (\beta/D)^2$. The orthogonal T matrix diagonalizes the symmetric $K^{JJ'}_{ij}$ matrix. This T matrix generates a new set of random variables $\eta^{\overline{a}(\overline{b})}_j = \sum_{Ji} T^J_{j;i}\overline{\xi}^{J\overline{a}(\overline{b})}_i$. An orthogonal U matrix diagonalizes the symmetric A matrix: $(UAU^T)_{ij} = u_i\delta_{ij}$ with $A_{ij} = (\eta^{\overline{a}}_i\eta^{\overline{b}}_j + \eta^{\overline{a}}_j\eta^{\overline{b}}_i)/2^{1/2}$. The expression for $\delta\tilde{S}^{J,I}_{\overline{a}\,\overline{b}}(E)$ is similar to that for $\delta S^{J,I}_{\overline{a}\,\overline{b}}(E)$ but with $\tilde{\phi}^{J(I)}_{\mu(\nu)}$ instead of $\phi^{J(I)}_{\mu(\nu)}$. Accordingly, we have to change $B \to \tilde{B}$, $C \to \tilde{C}$, $\overline{\phi}^{J(I),k}_{\mu(\nu)} \to \tilde{\phi}^{J(I),k}_{\mu(\nu)}$. However, the K, T, A, U matrices, $\overline{\xi}$'s and η's, and the X^J_i, Y_i, Z_k basis sets remain unchanged.

We can show that there is no correlation between the signs of (i) u_k and $q^{JI}_{\mu\nu}(k)$, (ii) u_k and $M^{I,k}_{\nu\neq\nu'}$, (iii) $q^{JI}_{\mu\nu}(k)$ and $q^{JI}_{\mu\nu'}(k)$ with $\nu \neq \nu'$ [8,9], (iv) $q^{JI}_{\mu\nu}(k)$ and $M^{I,k}_{\nu\neq\nu'}$, and (v) $q^{JI}_{\mu\nu}(k)$ and $\tilde{M}^{I,k}_{\nu\neq\nu'}$. This yields

$$\overline{R^{JI}_\mu\tilde{R}^{JI}_{\tilde{\mu}}} = 2\sum_{\nu\neq\nu'}\sum_{\tilde{\nu}\neq\tilde{\nu}'} \overline{q^{JI}_{\mu\nu}(k)\tilde{q}^{JI}_{\tilde{\mu}\tilde{\nu}}(k)}^{\,k}\;\overline{q^{JI}_{\mu\nu'}(k)\tilde{q}^{JI}_{\tilde{\mu}\tilde{\nu}'}(k)}^{\,k} \times \overline{M^{I,k}_{\nu\nu'}\tilde{M}^{I,k}_{\tilde{\nu}\tilde{\nu}'}}^{\,k},$$

where $\overline{(\cdots)}^k = (1/N)\sum_k(\cdots)$ denotes the k ensemble averaging [8,9]. Considering $(N^{1/2}\langle \phi^I_\nu\,|\,Z_k\rangle) \sim \pm 1$ and $(N^{1/2}\langle \tilde{\phi}^I_\nu\,|\,Z_k\rangle) \sim \pm 1$ to be Gaussian random variables and taking into account that Z_k is a complete basis set we obtain

$$\overline{M^{I,k}_{\nu\nu'}\tilde{M}^{I,k}_{\tilde{\nu}\tilde{\nu}'}}^{\,k} = \langle \phi^I_\nu\,|\,\tilde{\phi}^I_{\tilde{\nu}}\rangle\langle \phi^I_{\nu'}\,|\,\tilde{\phi}^I_{\tilde{\nu}'}\rangle + \langle \phi^I_\nu\,|\,\tilde{\phi}^I_{\tilde{\nu}'}\rangle\langle \phi^I_{\nu'}\,|\,\tilde{\phi}^I_{\tilde{\nu}}\rangle \simeq \delta_{\nu\tilde{\nu}}\delta_{\nu'\tilde{\nu}'} + \delta_{\nu\tilde{\nu}'}\delta_{\nu'\tilde{\nu}}.$$

In order to evaluate $\overline{q^{JI}_{\mu\nu}(k)\tilde{q}^{JI}_{\tilde{\mu}\tilde{\nu}}(k)}^{\,k}$ we consider the relation [8]: $(1/N^{1/2})\sum_k q^{JI}_{\mu\nu}(k) = N^{1/2}\langle \phi^J_\mu\,|\,\phi^I_\nu\rangle = N^{1/2}\delta_{JI}\delta_{\mu\nu}$. In the limit $N\to\infty$, the $q^{JI}_{\mu\nu}(k)$ with fixed $J \neq I$, μ, ν and running k index constitute an ensemble of uncorrelated random variables [8]. This means that the nonvanishing of individual $q^{JI}_{\mu\nu}(k)$ results from the nonvanishing of the uncertainty $N^{1/2}\delta_{JI}\delta_{\mu\nu}$ in the limit $N\to\infty$. Let us ask what is the sign of $N^{1/2}\delta_{JI}\delta_{\mu\nu}$? Since it is the same as the sign of $\delta_{JI}\delta_{\mu\nu} = 0$, then there is no preference in favor of either "+" or "−." Therefore, in the limit $N\to\infty$, the signs of $N^{1/2}\langle \phi^J_\mu\,|\,\phi^I_\nu\rangle$ are set up at random. This also implies that the signs of $N^{1/2}\langle \phi^J_\mu\,|\,\phi^I_\nu\rangle$ should be sensitive to an extremely small perturbation, e.g., such as occurs due to the interaction of the outer electron with INS. Accordingly, we can write $\mathrm{sgn}(N^{1/2}\langle \phi^J_\mu\,|\,\phi^I_\nu\rangle) = t^{JI}_{\mu\nu}\,\mathrm{sgn}(N^{1/2}\langle \tilde{\phi}^J_\mu\,|\,\tilde{\phi}^J_\nu\rangle)$, where $|t^{JI}_{\mu\nu}| = 1$ but $t^{JI}_{\mu\nu}$ has random signs with respect to μ, ν indices, e.g., $t^{JI}_{\mu\nu} = (-1)^{\mu+\nu}$. Since $(1/N^{1/2})\sum_k \tilde{q}^{JI}_{\mu\nu}(k) = N^{1/2}\langle \tilde{\phi}^J_\mu\,|\,\tilde{\phi}^I_\nu\rangle = N^{1/2}\delta_{JI}\delta_{\mu\nu}$, we have $\tilde{q}^{JI}_{\tilde{\mu}\tilde{\nu}}(k) = t^{JI}_{\tilde{\mu}\tilde{\nu}}\,\mathrm{sgn}[q^{JI}_{\tilde{\mu}\tilde{\nu}}(k)]\,|\tilde{q}^{JI}_{\tilde{\mu}\tilde{\nu}}(k)|$. Taking into account that $\overline{q^{JI}_{\mu\nu}(k)q^{JI}_{\tilde{\mu}\tilde{\nu}}(k)}^{\,k} = \delta_{\mu\tilde{\mu}}\delta_{\nu\tilde{\nu}}\overline{[q^{JI}_{\mu\nu}(k)]^2}^{\,k}$ [8] and

$\overline{[q^{JI}_{\mu\nu}(k)]^2}^k = \overline{[\tilde{q}^{JI}_{\mu\nu}(k)]^2}^k$ we obtain $\overline{q^{JI}_{\mu\nu}(k)\tilde{q}^{JI}_{\tilde{\mu}\tilde{\nu}}(k)}^k = \delta_{\mu\tilde{\mu}}\delta_{\nu\tilde{\nu}}t^{JI}_{\mu\nu}\overline{[q^{JI}_{\mu\nu}(k)]^2}^k$ if $|q^{JI}_{\tilde{\mu}\tilde{\nu}}(k)| = |\tilde{q}^{JI}_{\tilde{\mu}\tilde{\nu}}(k)|$ and $\overline{q^{JI}_{\mu\nu}(k)\tilde{q}^{JI}_{\tilde{\mu}\tilde{\nu}}(k)}^k = (2/\pi)\delta_{\mu\tilde{\mu}}\delta_{\nu\tilde{\nu}}t^{JI}_{\mu\nu}\overline{[q^{JI}_{\mu\nu}(k)]^2}^k$ if $q^{JI}_{\tilde{\mu}\tilde{\nu}}(k)$ and $\tilde{q}^{JI}_{\tilde{\mu}\tilde{\nu}}(k)$ are uncorrelated Gaussian variables. This yields

$$\begin{aligned}\overline{R^{JI}_{\mu}\tilde{R}^{JI}_{\tilde{\mu}}} &\sim \delta_{\mu\tilde{\mu}}\sum_{\nu\neq\nu'} t^{JI}_{\mu\nu}t^{JI}_{\mu\nu'}\overline{[q^{JI}_{\mu\nu}(k)]^2}^k\,\overline{[q^{JI}_{\mu\nu'}(k)]^2}^k \\ &\sim \delta_{\mu\tilde{\mu}}\sum_{\nu}\{\overline{[q^{JI}_{\mu\nu}(k)]^2}^k\}^2 \sim \delta_{\mu\tilde{\mu}}D/\beta|J-I|\,.\end{aligned}$$

Now the calculation of the correlation coefficient between $\delta S^{J,I}(E)$ and $\delta\tilde{S}^{J,I}(E)$ is straightforward and we obtain $\kappa \sim D(\Gamma + 2|J-I|\beta)/\beta^2(J-I)^2 \sim D\Gamma/\beta^2 \sim 10^{-6}$ for $D = 10^{-10}$ MeV, $\Gamma = 0.1$ MeV, and $\beta = 3.5$ keV [7,10]. The absence of the correlation between $\delta S^{J,I}(E)$ and $\delta\tilde{S}^{J,I}(E)$ means that the cross section oscillations for $^{19}F^{+7} + {}^{89}Y$ and $^{19}F^{+8} + {}^{89}Y$ DHIC are uncorrelated. Therefore the atomic-electron effects are predicted to be $\simeq$100% of the magnitude of the non-self-averaged oscillating component of the nuclear cross section. For $^{19}F + {}^{89}Y$ DHIC the relative magnitude of this oscillating component is $\simeq\pm15\%$ [6,7] so that the atomic-electron effects should be detectable as reliably as the effect of the non-self-averaging of excitation function oscillations for DHIC [5–7]. At the same time the decaying cross section energy autocorrelation functions [10] are predicted to be insensitive and indistinguishable for the $^{19}F^{+7} + {}^{89}Y$ and $^{19}F^{+8} + {}^{89}Y$ DHIC. Both the sensitivity of the cross sections and the decaying energy autocorrelation functions are manifestations of quantum chaos [13,14] in strongly dissipative nucleus-nucleus collisions.

The precondition for a sensitivity of the cross sections of DHIC is a finite β, i.e., the presence of quantum chaos and spin decoherence [8–10]. Indeed, in the limit $\beta/\Gamma \to 0$, $\overline{|\delta S^{J,I}(E)|^2} \propto \beta/\Gamma \to 0$ and $\overline{S}^{J}_{\overline{a}\overline{b}}(E) \to \Delta S^{J,I}_{\overline{a}\overline{b}}(E)$ resulting in both insensitivity and self-averaging of the cross sections. The cross sections are also insensitive and self-averaging in the compound nucleus limit $\Gamma/\beta \to 0$, since in this limit the MC vanishes [10], leading to the randomness conditions in the statistical model of compound nucleus scattering [15,16], theory of quantum chaotic scattering, and random-matrix theory [1–4,15–19].

In conclusion, we have shown that spontaneous microchannel S-matrix correlations and slow spin decoherence lead to a sensitivity of the cross sections for DHIC to an arbitrarily small perturbation in the limit when this perturbation vanishes. Such a sensitivity implies that atomic electrons should influence the cross sections for DHIC. These atomic-electron effects are predicted to be $\simeq$100% of the magnitude of the non-self-averaged oscillating component of the nucleus-nucleus cross section. New experiments are being proposed to test the predicted possibility of controlling energetic ($\simeq$100 MeV) nucleus-nucleus reactions by changing the atomic-electron environment.

I am grateful to Sasha Melnitchenko for support which enabled me to work on this paper. I thank Hans Weidenmüller for useful discussions and suggestions. The work under the revised version of this Letter has been supported by the Humboldt Fellowship and MPI Fellowship, Heidelberg.

[1] T. Guhr, A. Müller-Groeling, and H. A. Weidenmüller, Phys. Rep. **299**, 189 (1998), and references therein.

[2] R. Blümel and U. Smilansky, Phys. Rev. Lett. **60**, 477 (1988).

[3] R. Blümel and U. Smilansky, Phys. Rev. Lett. **64**, 241 (1990).

[4] U. Smilansky, in *Chaos and Quantum Physics*, Proceedings of the Les Houches Summer School, Session LII, edited by M. J. Giannoni, A. Voros, and J. Zinn-Justin (North-Holland, Amsterdam, 1991), p. 372, and references therein.

[5] A. De Rosa *et al.*, Phys. Lett. **160B**, 239 (1985); G. Pappalardo, Nucl. Phys. **A488**, 395c (1988); A. De Rosa *et al.*, Phys. Rev. C **37**, 1042 (1988); **40**, 627 (1989); **44**, 747 (1991); Wang Qi *et al.*, Chin. Phys. Lett. **10**, 656 (1993); Wang Qi *et al.*, Chin. J. Nucl. Phys. **15**, 113 (1993); F. Rizzo *et al.*, Z. Phys. A **349**, 169 (1994); Wang Qi, *et al.*, High Energy Phys. Nucl. Phys. **18**, 25 (1994); Lu Jun *et al.*, Chin. Phys. Lett. **12**, 661 (1995); Wang Qi *et al.*, High Energy Phys. Nucl. Phys. **20**, 289 (1996); Lu Jun *et al.*, Chin. J. Nucl. Phys. **18**, 91 (1996); Wang Qi *et al.*, Phys. Lett. B **388**, 462 (1996); I. Berceanu *et al.*, Phys. Rev. C **57**, 2359 (1998); Wang Qi, Nucl. Phys. Rev. **15**, 74 (1998); Tian Wendong *et al.*, High Energy Phys. Nucl. Phys. **23**, 334 (1999); I. Berceanu *et al.*, Pramana J. Phys. **53**, 419 (1999); I. Berceanu *et al.* (unpublished); Wang Qi *et al.* (unpublished).

[6] T. Suomijärvi *et al.*, Phys. Rev. C **36**, 181 (1987).

[7] S. Yu. Kun *et al.*, Z. Phys. A **359**, 263 (1997).

[8] S. Yu. Kun, Z. Phys. A **357**, 255 (1997).

[9] S. Yu. Kun, Z. Phys. A **357**, 271 (1997).

[10] S. Yu. Kun and A. V. Vagov, Z. Phys. A **359**, 137 (1997).

[11] O. Bohigas and H. A. Weidenmüller, Annu. Rev. Nucl. Part. Sci. **38**, 421 (1988).

[12] B. R. Barrett, S. Shlomo, and H. A. Weidenmüller, Phys. Rev. C **17**, 544 (1978).

[13] R. Blümel, Phys. Rev. Lett. **73**, 428 (1994).

[14] R. Blümel, Nucl. Phys. **A572**, 9 (1994).

[15] T. Ericson, Ann. Phys. (N.Y.) **23**, 390 (1963).

[16] J. J. M. Verbaarschot, H. A. Weidenmüller, and M. R. Zirnbauer, Phys. Rep. **129**, 367 (1985), and references therein.

[17] V. V. Flambaum, A. A. Gribakina, and G. F. Gribakin, Phys. Rev. A **54**, 2066 (1996).

[18] Y. V. Fyodorov and H.-J. Sommers, J. Math. Phys. (N.Y.) **38**, 1918 (1997), and references therein.

[19] Y. V. Fyodorov and Y. Alhassid, Phys. Rev. A **58**, R3375 (1998).