

Evgeny N. Bulgakov,<sup>1,2</sup> Ingrid Rotter,<sup>1</sup> and Almas F. Sadreev<sup>1,2,3</sup>

<sup>1</sup>*Max-Planck-Institut für Physik komplexer Systeme, D-01187 Dresden, Germany*

<sup>2</sup>*Kirensky Institute of Physics, 660036 Krasnoyarsk, Russia*

<sup>3</sup>*Department of Physics and Measurement,  
Technology Linköping University, S-581 83 Linköping, Sweden*

(Dated: July 4, 2006)

We consider the effective Hamiltonian of an open quantum system, its biorthogonal eigenfunctions  $\phi_\lambda$  and define the value  $r_\lambda = (\phi_\lambda|\phi_\lambda)/\langle\phi_\lambda|\phi_\lambda\rangle$  that characterizes the phase rigidity of the eigenfunctions  $\phi_\lambda$ . In the scenario with avoided level crossings,  $r_\lambda$  varies between 1 and 0 due to the mutual influence of neighboring resonances. The variation of  $r_\lambda$  may be considered as an internal property of an *open* quantum system. In the literature, the phase rigidity  $\rho$  of the scattering wave function  $\Psi_C^E$  is considered. Since  $\Psi_C^E$  can be represented in the interior of the system by the  $\phi_\lambda$ , the phase rigidity  $\rho$  of the  $\Psi_C^E$  is related to the  $r_\lambda$  and therefore also to the mutual influence of neighboring resonances. As a consequence, the reduction of the phase rigidity  $\rho$  to values smaller than 1 should be considered, at least partly, as an internal property of an open quantum system in the overlapping regime. The relation to measurable values such as the transmission through a quantum dot, follows from the fact that the transmission is, in any case, resonant with respect to the effective Hamiltonian. We illustrate the relation between phase rigidity  $\rho$  and transmission numerically for small open cavities.

PACS numbers: 03.65.Yz, 73.23.-b, 73.63.Kv, 05.60.Gg

## I. INTRODUCTION

The advances in the nanotechnology make it possible to produce small quantum dots with desired controllable properties. The analogy between such a system and an atom has proved to be quite close<sup>1</sup>. Since however no natural restrictions for choosing the control parameters exist, the quantum dots may show new properties that we cannot obtain from studies on atoms. For example, Fano resonances have been observed experimentally in quantum dots<sup>2</sup>, and the Fano parameters may be complex<sup>3</sup>. Theoretically, the Fano parameter expresses the interference between the resonant part of the transmission and a smooth (direct) nonresonant part. It is real when the resonance part is caused by the existence of an isolated resonance state<sup>4</sup> as it is the case usually in atoms. When however the resonant part itself results from the interference of, e.g., two neighboring resonance states, the Fano parameter becomes complex<sup>5</sup>. The same may appear in a cavity due to the absorption at the walls<sup>6</sup>.

Further, it has been stated<sup>7</sup> that the line shape of a Fano resonance may be affected by some dephasing<sup>8</sup> that may be caused by intrinsic sources (e.g., from electron-electron interactions) as well as by extrinsic sources (e.g., radiation, magnetic impurities)<sup>9</sup>. For a quantitative study, the phase rigidity<sup>10,11</sup>

$$\rho = \frac{\int dr \Psi(r)^2}{\int dr |\Psi(r)|^2} = e^{2i\theta} \frac{\int dr (|\operatorname{Re}\tilde{\Psi}(r)|^2 - |\operatorname{Im}\tilde{\Psi}(r)|^2)}{\int dr (|\operatorname{Re}\tilde{\Psi}(r)|^2 + |\operatorname{Im}\tilde{\Psi}(r)|^2)} \quad (1)$$

has been introduced that characterizes the degree to which the wave function  $\Psi$  is really complex (the phase  $\theta$  arises from a transformation so that  $\operatorname{Re}\tilde{\Psi}$  and  $\operatorname{Im}\tilde{\Psi}$  are orthogonal). An experimental and theoretical study<sup>6</sup> showed that the two different mechanisms, dephasing and

dissipation, are equivalent in terms of their effect onto the evolution of Fano resonance line shapes.

Another interesting observation in the experimental results on quantum dots<sup>2,7</sup> is that, as a function of the gate voltage that controls the transparencies of the point contacts, the widths of the observed resonances behave non-monotonic. The conductance peaks start as narrow Breit-Wigner resonances when the quantum dot is pinched off, then widen as the contacts are opened into resonances exhibiting the Kondo effect<sup>7</sup>. As the contacts are opened further, the resonances become more narrow and have the Fano form with some background conductance. A similar result is obtained in a study on a tunable microwave scattering device<sup>12</sup>. The explanation given in Ref.<sup>7</sup> is that diffraction at the contacts to the quantum dots is strongest at intermediate point contact transparencies, leading to large sticking probabilities.

The non-monotonic increase of resonance widths as a function of the degree of opening the system is however a typical feature of open quantum systems<sup>13</sup>. It was found first in theoretical nuclear reaction studies by using the formalism of the effective non-hermitian Hamiltonian  $H_{\text{eff}}$ <sup>14</sup>, then in atoms<sup>15</sup> and in microwave cavities<sup>16</sup> where it is proven experimentally<sup>6,12,17</sup>. Also the transmission through a double quantum dot is studied as a function of the coupling strength between dot and attached leads<sup>18</sup>. In the theoretical studies, the non-monotonic behavior of the resonance widths is caused by the width bifurcation that may appear at the avoided crossings of resonance states in the complex energy plane. Further studies of open quantum systems showed that the real and imaginary parts of the eigenfunctions of  $H_{\text{eff}}$  evolve more or less independently from one another in the avoided level crossing scenario<sup>19</sup> and that long-range correlations oc-

cur between the different states of the system<sup>20</sup>.

The effective Hamiltonian used in the random matrix interpolation between the standard ensembles with real and complex matrix elements, is  $H(\alpha) = H_0 + \alpha H_1$  where  $H_0$  and  $H_1$  are real and complex random hermitian matrices, respectively, and  $\alpha$  is a crossover parameter<sup>21</sup>. For such an ensemble, the eigenvector elements acquire correlations between the elements of the same eigenvector<sup>10,22</sup> and between different eigenvectors<sup>23</sup>. For individual systems, such a crossover may be observed already in a billiard with only two attached waveguides<sup>24</sup>. The wave functions in the cross-over regime show long-range correlations<sup>11</sup> like the eigenvectors of the Hamiltonian  $H(\alpha)$ . These long-range correlations of the wave functions in the real-to-complex crossover have recently been measured in an open microwave billiard<sup>25</sup>.

Thus, the results obtained experimentally and theoretically for the profile of the transmission peaks, for the intensity fluctuations and also for the long-range correlations of the wave functions point to the fact that the system through which the transmission occurs may be essentially different from the original closed quantum system without attached leads. The differences arise from the interaction of neighboring resonance states via the continuum of scattering states. In quantum dots, these differences can be traced experimentally by varying the gate voltage that controls the transparencies of the point contacts.

In the following, we will study the phases of the eigenfunctions of the effective Hamiltonian in the avoided level crossing scenario in detail. We will show that they are not fixed but vary due to the mutual influence of neighboring resonance states. In this spirit, a neighboring resonance state causes some "perturbation" for the considered state which is similar to that caused by an impurity. There is however an important difference between these two cases. In contrast to the perturbation by an impurity, the mutual interaction between neighboring states can not be avoided. It is an internal property of an open quantum system in the regime of overlapping resonances.

In Sect. II, we repeat some general properties of the interplay between the internal and external interaction in an open quantum system. While the internal interaction is of standard type, the external interaction occurs additionally via the common continuum of scattering wave functions. The amplitude for the transmission through a quantum dot will also be given. It is resonant in relation to the effective Hamiltonian  $H_{\text{eff}}$  of the open quantum system. Then we define the phase rigidity  $r_\lambda$  of the eigenfunctions  $\phi_\lambda$  of  $H_{\text{eff}}$ . Further, we consider the scattering wave functions  $\Psi_C^E$  being solutions of  $(H - E)\Psi_C^E = 0$  in the total function space. In the interior of the system, the  $\Psi_C^E$  can be represented by the  $\phi_\lambda$  and, as a consequence, the phase rigidity  $\rho$  of the scattering wave functions  $\Psi_C^E$  is related to that of the eigenfunctions  $\phi_\lambda$ . In Sect. III, we provide numerical results for the relation between transmission and phase rigidity  $\rho$  of the scattering wave functions for three special cases. The results

are summarized and some conclusions are drawn in the last section.

## II. PHASE RIGIDITY OF THE EIGENFUNCTIONS AND OF THE SCATTERING WAVE FUNCTION IN AN OPEN QUANTUM SYSTEM

The relation between a closed and the corresponding open quantum system is as follows. A closed quantum system is described by a hermitian Hamilton operator  $H_B$  the eigenvalues and eigenfunctions of which contain the internal interaction  $u$  of the discrete states. When embedded into the common continuum of scattering states, the discrete eigenstates of the closed system turn over into resonance states with a finite lifetime. The effective Hamiltonian  $H_{\text{eff}}$  of the open quantum system contains  $H_B$  as well as an additional term<sup>13</sup> that describes the coupling of the resonance states to the common environment,

$$H_{\text{eff}} = H_B + \sum_C V_{BC}(E^+ - H_C)^{-1}V_{CB}. \quad (2)$$

Here  $V_{BC}$ ,  $V_{CB}$  stand for the coupling matrix elements between the *eigenstates* of  $H_B$  and the environment that may consist of different continua  $C$ , e.g. the scattering waves propagating in the left ( $C = L$ ) and right ( $C = R$ ) leads attached to a quantum dot<sup>26</sup>. They are described by the Hamiltonian  $H_C$ .  $H_{\text{eff}}$  is non-hermitian, its eigenvalues  $z_\lambda$  and eigenfunctions  $\phi_\lambda$  contain the external interaction  $v$  of the resonance states via the continuum ( $v$  is used here instead of the concrete values  $V_{BC}$  and  $V_{CB}$ ). While  $u$  causes level repulsion in energy,  $v$  is responsible for the bifurcation of the widths of the resonance states.

As long as  $u/v \gg 1$ , the spectroscopic properties of the open system are similar to those of the corresponding closed system. The lifetimes (widths) of the resonance states are, as a rule, comparable in value, i.e. the states exist at the same time and can influence one another. Usually, the states avoid crossing in the complex energy plane in a similar manner as it is well known from the avoided crossings of the discrete states of a closed system. When however  $u/v \ll 1$ , the resonance states do no longer exist at the same time due to their different lifetimes owing to widths bifurcation<sup>13,19</sup>. They do not cross therefore in the complex energy plane. Most interesting is the situation  $u \approx v$ . The interplay between  $u$  and  $v$  may cause unexpected and even counterintuitive results such as the non-monotonic dependence of the resonance widths on the degree of opening the system<sup>13,15-18</sup>.

Since the effective Hamiltonian  $H_{\text{eff}}$  depends explicitly on the energy  $E$ , so do its eigenvalues  $z_\lambda$ . The energy dependence is weak, as a rule, in an energy interval that is determined by the width of the resonance state. The solutions of the fixed-point equations  $E_\lambda = \text{Re}(z_\lambda)|_{E=E_\lambda}$  and of  $\Gamma_\lambda = -2\text{Im}(z_\lambda)|_{E=E_\lambda}$  are numbers that coincide approximately with the poles of

the  $S$  matrix. The width  $\Gamma_\lambda$  determines the time scale characteristic of the resonance state  $\lambda$ . The amplitude for the transmission through a quantum dot is<sup>26</sup>

$$t = -2\pi i \sum_{\lambda} \langle \xi_L^E | V | \phi_{\lambda} \rangle \frac{\langle \phi_{\lambda} | V | \xi_R^E \rangle}{E - z_{\lambda}} \quad (3)$$

where the scattering wave functions in the leads are denoted by  $\xi_C^E$ . According to (3), the transmission is *resonant* in relation to  $H_{\text{eff}}$ . In (3), the eigenvalues  $z_{\lambda}$  with their full energy dependence are involved.

The eigenfunctions  $\phi_{\lambda}$  of  $H_{\text{eff}}$  are complex and biorthogonal,

$$\langle \phi_{\lambda} | \phi_{\lambda'} \rangle \equiv \langle \phi_{\lambda}^* | \phi_{\lambda'} \rangle = \delta_{\lambda, \lambda'} \quad (4)$$

$$\langle \phi_{\lambda} | \phi_{\lambda} \rangle \equiv A_{\lambda} \geq 1 \quad ; \quad |\langle \phi_{\lambda} | \phi_{\lambda'} \rangle| \equiv B_{\lambda}^{\lambda'} \geq 0. \quad (5)$$

The value

$$\begin{aligned} r_{\lambda} &= \frac{\int dr \phi_{\lambda}^2}{\int dr |\phi_{\lambda}|^2} = \frac{1}{A_{\lambda}} \\ &= \frac{\int dr (|\text{Re} \phi_{\lambda}(r)|^2 - |\text{Im} \phi_{\lambda}(r)|^2)}{\int dr (|\text{Re} \phi_{\lambda}(r)|^2 + |\text{Im} \phi_{\lambda}(r)|^2)} \end{aligned} \quad (6)$$

is a measure for the biorthogonality of the eigenfunctions  $\phi_{\lambda}$  of  $H_{\text{eff}}$  and for their phase rigidity,  $0 \leq r_{\lambda} \leq 1$ . It is large for an isolated resonance state at the energy  $E = E_{\lambda}$ , where the transmission probability has a peak. In the regime of overlapping resonance states however, where avoided level crossings appear<sup>19</sup>,  $r_{\lambda}$  is usually small. Approaching a branch point in the complex energy plane where two eigenvalues  $z_{\lambda}$  and  $z_{\lambda'}$  coalesce, we have<sup>15,19,27</sup>  $A_{\lambda(\lambda')} \rightarrow \infty$  and  $r_{\lambda(\lambda')} \rightarrow 0$ , and further

$$|\phi_{\lambda} \rangle \rightarrow \pm i |\phi_{\lambda' \neq \lambda} \rangle. \quad (7)$$

Here, the widths bifurcate: with further increasing  $v$ , one of the states aligns with a scattering state (channel wave function  $\xi_{C=R}^E$  or  $\xi_{C=L}^E$  in the one-channel transmission) and becomes short-lived while the other one becomes long-lived. For large  $v$ , we have therefore only one long-lived resonance state and, as for non-overlapping resonance states,  $r_{\lambda} \rightarrow 1$  at the energy  $E = E_{\lambda}$  of the long-lived state.

The scattering wave function  $\Psi_C^E$  is solution of the Schrödinger equation  $(H - E)\Psi_C^E = 0$  in the total function space that consists of the discrete states of the quantum dot and of the scattering states  $\xi_C^E$  in the attached leads.  $H$  is hermitian. The solution reads<sup>13</sup>

$$\Psi_C^E(r) = \xi_C^E(r) + \sum_{\lambda} \Omega_{\lambda}^C(r) \frac{\langle \phi_{\lambda} | V | \xi_C^E \rangle}{E - z_{\lambda}} \quad (8)$$

where  $\Omega_{\lambda}^C(r) = [1 + (E^+ - H_C)^{-1}V] \phi_{\lambda}(r)$  is the wave function of the resonance state  $\lambda$ . According to (8), the eigenfunctions  $\phi_{\lambda}(r)$  of the effective Hamiltonian  $H_{\text{eff}}$  give the main contribution to the scattering wave function  $\Psi_C^E(r)$  in the interior of the cavity. Here,

$$\Psi_C^E(r) \rightarrow \sum_{\lambda} \phi_{\lambda}(r) \frac{\langle \phi_{\lambda} | V | \xi_C^E \rangle}{E - z_{\lambda}}. \quad (9)$$

According to the relation (9), the phase rigidity  $\rho$  of the scattering wave function, Eq. (1), is generally determined to the values of the phase rigidity  $r_{\lambda}$  of the individual resonance states  $\lambda$ . While  $r_{\lambda}$  characterizes the phase rigidity of the special resonance state  $\lambda$ ,  $\rho$  contains the  $r_{\lambda}$  from the different states that contribute to the scattering wave function  $\Psi_C^E$  in the considered energy region. This difference between  $\rho$  and  $r_{\lambda}$  is illustrated best by the following example. Approaching a branch point in the complex energy plane, the contributions to (9) from the two states  $\lambda$  and  $\lambda'$  with the relation (7) of their wave functions to each other, cancel each other. The phase rigidity  $\rho$  of the scattering wave function is determined therefore by the contributions of other states that, as a rule, are relatively far from one another. That means, in spite of  $r_{\lambda} = r_{\lambda'} = 0$ ,  $|\rho|$  might be large at the branch point.

We mention that the scattering wave function (9) has a similar structure as the transmission amplitude (3). Both expressions consist of a sum over overlapping resonance states with the weight factor  $\langle \phi_{\lambda} | V | \xi_C^E \rangle / (E - z_{\lambda})$ . We expect therefore some correlation between both values. This relation, being trivial for isolated resonances, is of special interest in the regime of overlapping resonances. The transmission may be large not only at the positions  $E = E_{\lambda}$  of the resonance states  $\lambda$  as in the case of isolated resonances. It may be enhanced in a larger energy region when  $|\rho| < 1$ , i.e. when some states  $\lambda$  are (partly) aligned with the scattering states  $\xi_C^E$  so that  $r_{\lambda} < 1$  for them. In such a case, the matrix elements  $\langle \xi_C^E | V | \phi_{\lambda} \rangle$  may be large not only at  $E = E_{\lambda}$  but in a larger energy region. For illustration, we will provide in the following section some numerical results for  $|t|$  and  $|\rho|$  for two model quantum billiards with a small number of states the transmission through which is studied earlier, and for a realistic chaotic quantum billiard. The correlation between  $|t|$  and  $1 - |\rho|$  in the overlapping regime can clearly be seen in all cases.

### III. NUMERICAL EXAMPLES

#### A. 1d-system

For illustration, let us first consider the relation between the degree of resonance overlapping and the phase rigidity in the transmission through a 1d-system, see Fig. 1. The widths of the states in the middle  $E = 0$  of the spectrum are larger than those near to the thresholds  $E = \pm 2$  as can be seen immediately from the profile of the transmission peaks. Therefore, also the degree of resonance overlapping increases towards the middle of the spectrum. The results show that the transmission increases with increasing resonance overlapping and, correspondingly, the phase rigidity decreases. That means, the transmission increases when the resonance states align with the scattering wave function which is expressed by the decreasing of  $\rho$ .

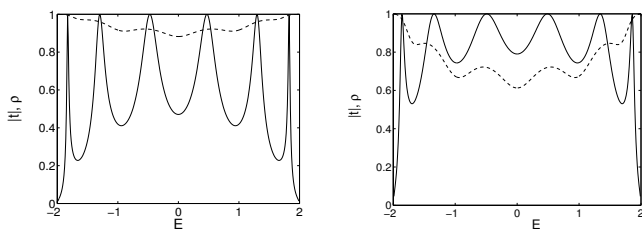


FIG. 1: The transmission  $|t|$  (full lines) and phase rigidity  $|\rho|$  (dashed lines) for a chain of 6 sites as a function of energy with  $v=0.5$  (left) and  $v=0.7$  (right). For the model see Ref. <sup>26</sup>.

## B. Double quantum dot

Let us now consider a double quantum dot consisting of two identical single dots that are connected by an internal wire. Suppose both dots and the wire have one state each. At  $v = v_c$ , the eigenvalues  $z_\lambda$  and  $z_{\lambda'}$  of the effective Hamiltonian coalesce and a branch point in the complex energy plane appears<sup>18,19</sup>. When  $v < v_c$ , the states repel each other in energy while widths bifurcation starts beyond the branch point where  $v > v_c$ . In Fig. 2, we show the transmission  $|t|$  and the landscape of  $|\rho|$  over energy  $E$  and coupling strength  $v$  (for fixed  $u$ ). The smallest value of  $|\rho|$  ( $\rho = 0$ ) is reached when the transmission is maximal with a plateau  $|t| = 1$  (for the plateau compare Fig. 4 in<sup>18</sup>). At  $v = v_c$ ,  $|\rho|$  is relatively large due to the fact that the contributions of the two states  $\lambda$  and  $\lambda'$  cancel each other. Beyond  $v_c$ , the phase rigidity increases further up to its maximal value 1.

## C. Sinai billiard

In Fig. 3, we show  $|t|$  and  $|\rho|$  calculated for a realistic chaotic system with many states and many avoided level crossings. The high correlation between the two values  $|t|$  and  $1 - |\rho|$  can be seen also under these conditions. In some energy intervals,  $|t|$  is near to 1. Here  $|\rho|$  is small, in agreement with the results shown in Fig. 2.

## IV. CONCLUSIONS

From the results obtained, we conclude that  $|\rho| < 1$  characterizes the scenario of overlapping resonances in which the resonance states interact with one another and avoided (and true) crossings appear, generally. In this scenario, the system can not be described by a hermitian Hamilton operator that provides rigid phases of its eigenfunctions ( $r_\lambda = 1$ ). Rather,  $H_{\text{eff}}$  is non-hermitian, its eigenfunctions are biorthogonal and  $r_\lambda$  varies with energy and coupling strength. Since the scattering wave function and the eigenfunctions of the effective Hamiltonian are related to one another according to (9), also  $\rho$  varies between 1 and 0 in the regime of overlapping resonances. That means, the reduction of the phase rigidity  $\rho$

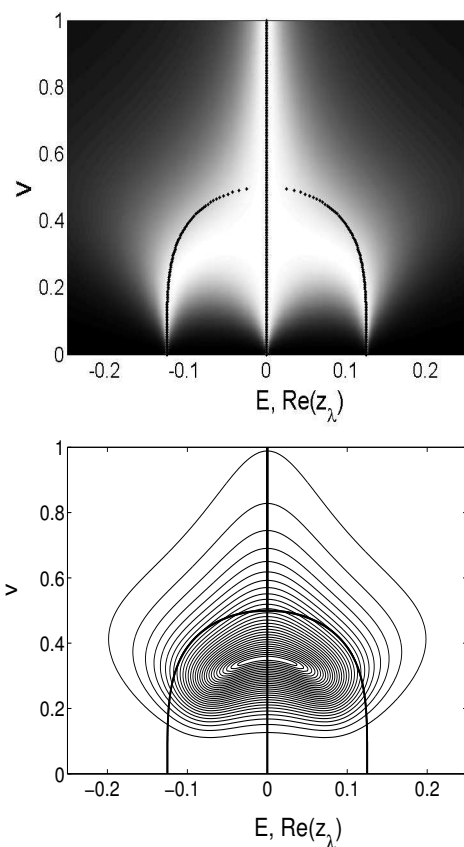


FIG. 2: The transmission  $|t|$  (top) and the landscape of the phase rigidity  $|\rho|$  (bottom, thin lines) for a double quantum dot over energy  $E$  and coupling strength  $v$ . The distance between the contour lines is  $\Delta|\rho| = 1/30$ . The minimal value  $\rho = 0$  is surrounded by a high density of contour lines. The highest shown contour line corresponds to  $|\rho| = 1 - 1/30$ . The  $\text{Re}(z_\lambda)$  of the three eigenstates (thick lines in both panels of the figure) are calculated at  $E = 0$ . The branch point is at  $v_c = 1/2$ ,  $E_c = 0$ .  $u = \sqrt{2}/16$ . Around  $v = 0.345$ , the phase rigidity is minimal and the transmission maximal with a plateau  $|t| = 1$  (for the plateau compare Fig. 4 in Ref. <sup>18</sup>).

of the scattering wave function is caused, at least partly, by the mutual interaction of neighboring resonances. It is an internal property of an open quantum system in the overlapping regime and can not be avoided.

The reduction of the phase rigidity  $\rho$  is, in the one-channel transmission, an expression for the (partial) alignment of two of the eigenfunctions  $\phi_{\lambda,\lambda'}$  of the effective Hamiltonian  $H_{\text{eff}}$  with the two scattering wave functions  $\xi_{C=R,L}^E$  in the overlapping regime. Due to this alignment (described by  $r_{\lambda,\lambda'} < 1$ ), the transmission through the system is enlarged in a certain energy region around the energies  $E_{\lambda,\lambda'}$  of the two resonance states. As can be seen from Fig. 2 and the corresponding discussion in<sup>18</sup>, the profile of the transmission is, when  $\rho \approx 0$ , completely different from that through two isolated resonance states. It is plateau-like. This result explains the correlation between  $|t|$  and  $1 - |\rho|$  which can be seen in all our numerical

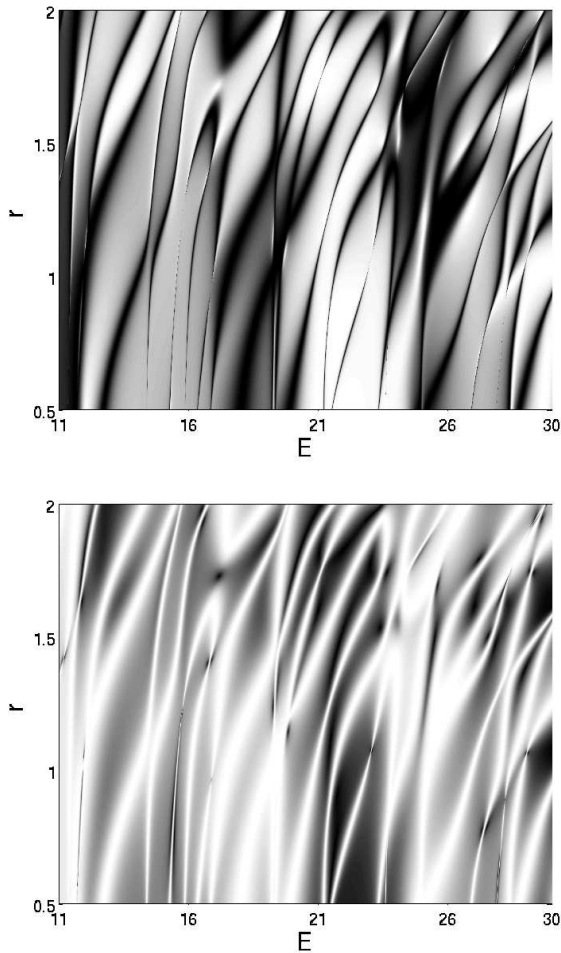


FIG. 3: The transmission  $|t|$  (top) and phase rigidity  $|\rho|$  (bottom) for a Sinai billiard over energy  $E$  and radius  $r$  of the circular disk (size of the billiard:  $x = 4$ ,  $y = 5$  in units of the width of the leads). The calculations are performed in the tight-binding lattice model<sup>28</sup>.

studies (Figs. 1, 2 and 3). Both values are, according to Eqs. (3) and (9), characterized by the contributions from a sum of overlapping resonance states.

Since the variation of  $r_\lambda$  and  $|\rho|$  with energy follows from the interaction of the resonance states via the common continuum of scattering wave functions, it is related, generally, to the profile of Fano resonances. The relation is however unique only with respect to  $r_\lambda$ . For example, the mutual influence of neighboring resonances is, according to the results presented in<sup>5,30</sup>, maximum when two eigenvalues of the effective Hamiltonian coincide. The line profile of two completely overlapping resonance states is, in the one-channel case and up to the background term, given by<sup>5,31</sup>

$$S = 1 - 2i \frac{\Gamma_d}{E - E_d + \frac{i}{2}\Gamma_d} - \frac{\Gamma_d^2}{(E - E_d + \frac{i}{2}\Gamma_d)^2} \quad (10)$$

where  $E_1 = E_2 \equiv E_d$  and  $\Gamma_1 = \Gamma_2 \equiv \Gamma_d$ . According to this equation, the line profile of the two resonances differs strongly from that of isolated resonances<sup>5</sup> (determined by the rigid value  $r_\lambda = 1$ ). In correspondence to this result, we have  $r_\lambda = 0$  in the one-channel case when two resonance states completely overlap. However  $|\rho|$  might be large in this case as for weakly overlapping resonance states, as discussed above. This result underlines once more the difference between the value  $r_\lambda$  characteristic of a special resonance state  $\lambda$  of the system, and  $\rho$  characteristic of the behavior of a sum of resonance states in the overlapping regime.

Thus, the avoided crossings of resonance states represent one of the sources for the dephasing<sup>8</sup> observed in experimental data. The variation of  $r_\lambda$  and  $|\rho|$  between 0 and 1 is characteristic of an open quantum system with overlapping resonance states, as stated above. Therefore, dephasing and dissipation should be equivalent in terms of their effect onto the Fano profile. Such a result is observed experimentally, indeed<sup>6</sup>.

As to averaged values, we mention the following results obtained earlier. The distributions of  $|\rho|$  with energy and ensemble averaging calculated for large chaotic cavities, do not depend on the concrete shape of the cavity<sup>29</sup> since the averaging smears the different contributions to  $|\rho|$ . As a result, the distribution of  $|\rho|$  is characterized, in such a case, only by the number of channels<sup>11</sup>.

We underline that the interaction of neighboring resonance states via the common continuum which is considered here as source for the dephasing<sup>8</sup>, is fundamentally different from, e.g., the electron-electron interaction or the interaction due to some radiation. It is rather of nonlinear geometrical origin, related to branch points in the complex energy plane, as can be seen in the following manner. On the one hand, the branch points determine the physical properties of an open quantum system in the overlapping regime as discussed in<sup>18,19,27</sup>. They are, generally, related to the avoided crossings of resonance states which demonstrate the mutual influence of resonance states. On the other hand, as shown in the present paper, the mutual interaction of neighboring resonance states is accompanied by phase changes of the eigenfunctions  $\phi_\lambda$  of  $H_{\text{eff}}$  as well as of the scattering wave functions  $\Psi_C^E$  inside the cavity. In any case, the interaction of neighboring resonance states via the common continuum and the resulting dephasing can not be neglected in an open quantum system in the regime of overlapping resonances.

**Acknowledgments:** ENB and AFS thank the MPI PKS in Dresden for hospitality. This work has been supported by RFBR grant 05-02-97713 "Enisey".

\* ben@tnp.krasn.ru; rotter@mpipks-dresden.mpg.de; almsa@ifm.liu.se; almas@tnp.krasn.ru

- <sup>1</sup> A. Yacoby, M. Heiblum, D. Mahalu, and H. Shtrikman, *Phys. Rev. Lett.* **74**, 4047 (1995); E. Buks, R. Schuster, M. Heiblum, D. Mahalu, V. Umansky, and H. Shtrikman, *Phys. Rev. Lett.* **77**, 4664 (1996); R. Schuster, E. Buks, M. Heiblum, D. Mahalu, V. Umansky, and H. Shtrikman, *Nature (London)* **385**, 417 (1997).
- <sup>2</sup> J. Göres, D. Goldhaber-Gordon, S. Heemeyer, M.A. Kastner, H. Shtrikman, D. Mahalu, and U. Meirav, *Phys. Rev. B* **62**, 2188 (2000).
- <sup>3</sup> K. Kobayashi, H. Aikawa, S. Katsumoto, and Y. Iye, *Phys. Rev. Lett.* **88**, 256806 (2002)
- <sup>4</sup> U. Fano, *Phys. Rev.* **124**, 1866 (1961).
- <sup>5</sup> A.I. Magunov, I. Rotter and S.I. Strakhova, *Phys. Rev. B* **68**, 245305 (2003).
- <sup>6</sup> S. Rotter, U. Kuhl, F. Libisch, J. Burgdörfer, and H.J. Stöckmann, *Physica E* **29**, 325 (2005).
- <sup>7</sup> A.A. Clerk, X. Waintal, and P.W. Brouwer, *Phys. Rev. Lett.* **20**, 4636 (2001).
- <sup>8</sup> The term 'dephasing' is used here according to the definition given in Refs.<sup>7,9</sup>
- <sup>9</sup> P.W. Brouwer and C.W.J. Beenakker, *Phys. Rev. B* **55**, 4695 (1997).
- <sup>10</sup> S.A. van Langen, P.W. Brouwer, and C.W.J. Beenakker, *Phys. Rev. E* **55**, 1 (1997).
- <sup>11</sup> P.W. Brouwer, *Phys. Rev. E* **68**, 046205 (2003).
- <sup>12</sup> S. Rotter, F. Libisch, J. Burgdörfer, U. Kuhl, and H.J. Stöckmann, *Phys. Rev. E* **69**, 046208 (2004).
- <sup>13</sup> I. Rotter, *Rep. Prog. Phys.* **54**, 635 (1991); J. Okołowicz, M. Płoszajczak, and I. Rotter, *Phys. Rep.* **374**, 271 (2003).
- <sup>14</sup> P. Kleinwächter and I. Rotter, *Phys. Rev. C* **32**, 1742 (1985).
- <sup>15</sup> A.I. Magunov, I. Rotter and S.I. Strakhova, *J. Phys. B* **32**, 1669 (1999) and **34**, 29 (2001).
- <sup>16</sup> E. Persson, K. Pichugin, I. Rotter, and P. Šeba, *Phys. Rev. E* **58**, 8001 (1998); P. Šeba, I. Rotter, M. Müller, E. Persson, and K. Pichugin, *Phys. Rev. E* **61**, 66 (2000); I. Rotter, E. Persson, K. Pichugin, and P. Šeba, *Phys. Rev. E* **62**, 450 (2000).
- <sup>17</sup> E. Persson, I. Rotter, H.J. Stöckmann, and M. Barth, *Phys. Rev. Lett.* **85**, 2478 (2000); H.J. Stöckmann, E. Persson, Y.H. Kim, M. Barth, U. Kuhl, and I. Rotter, *Phys. Rev. E* **65**, 066211 (2002).
- <sup>18</sup> I. Rotter and A.F. Sadreev, *Phys. Rev. E* **69**, 066201 (2004).
- <sup>19</sup> I. Rotter and A.F. Sadreev, *Phys. Rev. E* **71**, 036227 (2005).
- <sup>20</sup> C. Jung, M. Müller, and I. Rotter, *Phys. Rev. E* **60**, 114 (1999).
- <sup>21</sup> A. Pandey and M.L. Mehta, *Commun. Math. Phys.* **87**, 449 (1983).
- <sup>22</sup> V.I. Fal'ko and K.B. Efetov, *Phys. Rev. B* **50**, R11267 (1994); *Phys. Rev. Lett.* **77**, 912 (1996).
- <sup>23</sup> S. Adam, P.W. Brouwer, J.P. Sethna, and X. Waintal, *Phys. Rev. B* **66**, 165310 (2002).
- <sup>24</sup> A.F. Sadreev and K.F. Berggren, *Phys. Rev. E* **70**, 026201 (2004).
- <sup>25</sup> Y.H. Kim, U. Kuhl, H.J. Stöckmann, and P.W. Brouwer, *Phys. Rev. Lett.* **94**, 036804 (2005).
- <sup>26</sup> A.F. Sadreev and I. Rotter, *J. Phys. A* **36**, 11413 (2003).
- <sup>27</sup> I. Rotter, *Phys. Rev. E* **64**, 036213 (2001).
- <sup>28</sup> S. Datta, *Electronic transport in mesoscopic systems*, Cambridge University Press, 1995.
- <sup>29</sup> E.N. Bulgakov and I. Rotter, *Phys. Rev. E* **73**, 066222 (2006).
- <sup>30</sup> M. Müller, F.M. Dittes, W. Iskra and I. Rotter, *Phys. Rev. E* **52**, 5961 (1995).
- <sup>31</sup> I. Rotter, *Phys. Rev. E* **68**, 016211 (2003).