Local random wave model for semiclassical fractal structure of chaotic resonance states

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Abstract. The semiclassical structure of resonance states of classically chaotic scattering systems with partial escape is investigated. We introduce a local randomization on phase space for the baker map with escape, which separatesthe smallest multifractal scale from the scale of the Planck cell. This allows for deriving a semiclassical description of resonance states based on a local random wave model and conditional invariance. We numerically demonstrate that the resulting classical measures perfectly describe resonance states of all decay rates γ for the randomized baker map. Comparison to the baker map without randomization shows very good agreement of the multifractal structures. Quantitative differences indicate that a semiclassical description for systems without randomization must take into account that the multifractal structures persist down to the Planck scale.

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1. Introduction

Understanding the correspondence of quantum and classical dynamics is an essential task in complex systems [1]. Most prominently, the semiclassical structure of quantum eigenstates in closed systems is related to classically invariant phase-space structures [2, 3]. For classically ergodic dynamics almost all quantum eigenstates converge weakly towards the uniform measure on phase space as proven by the quantum ergodicity theorem [4, 5, 6, 7, 8]. This uniform limit is also established for quantum maps [9, 10, 11]. The amplitude distribution of eigenstates is described by random wave models, first introduced for quantum billiards [2], in which a complex Gaussian distribution of wave amplitudes leads to a universal exponential distribution of the intensities [11, 12, 13, 14, 15, 16, 17, 18].

In contrast, for systems with loss of particles or intensity in some interaction region [19] the semiclassical limit is not fully understood. Such scattering systems, e.g., the three-disk system [20, 21, 22] or optical microcavities [23], are described by resonance poles and the corresponding resonance states ψ which have a decay rate γ. The distribution of decay rates is well described by a fractal Weyl law in systems
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with full escape [24, 25, 26, 27, 28, 29, 30, 31, 32], and has been studied in systems with partial escape [33, 34, 35, 36]. Here we focus on the semiclassical structure of resonance states in chaotic systems with escape. They are described by multifractal measures on phase space [37, 38, 39, 40, 41, 42, 43, 44, 45, 46, 47, 48, 49] with exponentially distributed intensity fluctuations [50]. These measures are conditionally invariant under the classical dynamics with escape [51, 19] and strongly depend on the corresponding quantum decay rate \( \gamma \) [40, 47, 48]. Such questions are also of interest in mixed systems and in non-Hermitian PT-symmetric systems [52, 53, 54, 55].

For fully chaotic systems there exist approximative descriptions of resonance states with arbitrary decay rates \( \gamma \) for full [47] and partial escape [48], however, the precise semiclassical limit is not known.

The main difficulty in describing chaotic resonance states, according to our understanding, is the fact that their multifractal structures persist all the way down to the scale of a Planck cell. Since the multifractal structure has its origin in the classical dynamics, there is no separation of the smallest classical scale and the scale of a Planck cell, not even when going to the semiclassical limit.

In this paper we introduce a chaotic model system with partial escape for which the smallest multifractal scale and the scale of the Planck cell are separated. This allows for deriving a semiclassical description of resonance states. Specifically, we define the randomized baker map with escape, where in addition to the quantum baker map and escape from phase space, local randomization operations on phase space with an adjustable scale are applied. Resonance states show multifractal structures on scales larger than the scale of randomization and are uniformly fluctuating on smaller scales. We introduce classical measures, which (i) are conditionally invariant with some desired decay rate \( \gamma \) on scales larger than the scale of randomization and (ii) obey an extremum principle derived from a local random wave model in each randomization region. We numerically demonstrate that for all decay rates \( \gamma \) these classical measures perfectly describe resonance states in the semiclassical limit. Comparison to the baker map without randomization shows very good agreement of the multifractal structures. quantitative differences indicate that a semiclassical description for systems without randomization must take into account that the multifractal structures persist down to the Planck scale.

The paper is organized as follows. In Sec. 2 we review the classical baker map and its symbolic dynamics. Section 3 introduces the randomized quantum baker map with escape and discusses the structure of its resonance states. In Sec. 4 we determine classical measures from a local random wave model with the constraint of conditional invariance. We conjecture that these classical measures describe the semiclassical structure of resonance states. Numerical support for the conjecture is presented in Sec. 5. A comparison to the quantum baker map without randomization is made in Sec. 6 showing qualitative similarity, but quantitative difference. Section 7 gives an outlook and in Appendix C a random matrix model with escape is discussed.

This paper touches many topics that Fritz Haake has worked on [1], as it studies the semiclassical limit of a quantum model combining classical deterministic chaotic dynamics with random matrices. Specifically, the local randomization procedure is related to the blurring of phase-space resolution studied by Fritz [56, 57, 58]. Furthermore the present paper deals with open quantum systems, another field to which Fritz contributed substantially in terms of scattering systems [59, 60], as well as in terms of coupling to a thermal environment [61, 62]. It would have been a pleasure to discuss the results with Fritz and to carefully listen to his insights.
2. Classical baker map with escape

In this section we introduce the baker map, one of the paradigmatic models for chaos, combined with escape from phase space [63, 19]. Its symbolic dynamics will be used to define a phase-space partition with rectangles that will be relevant for the randomized quantum baker map with escape introduced in the next section.

2.1. Baker map with escape

The baker map with \( n \) stripes is defined on the two-torus \( T = [0,1) \times [0,1) \) as follows [64]. Consider the vertical rectangles \( A_k = \left[ k/n, (k+1)/n \right) \times [0,1) \) for \( k \in \{0, \ldots, n-1\} \). The \( n \)-baker map compresses each of these rectangles along the \( p \)-direction by the factor \( 1/n \) and stretches it along the horizontal \( q \)-direction by the factor \( n \), after which they are stacked on top of each other,

\[
B_n : T \to T
A_k \ni (q,p) \mapsto (qn - k, (p + k)/n). \tag{1}
\]

Note that \( B_n \) is volume preserving on \( T \).

For the baker map with escape we consider a loss of intensity which is constant in each of the rectangles \( A_k \). In particular, let \( r \in \mathbb{R}_+^n \) and define the reflectivity function \( r : T \to \mathbb{R}_+ \) as \( A_k \ni (q,p) \mapsto r_k \), see Fig. 1(a). Extending the phase space with the intensity space \( \mathbb{R}_+ \), the map with escape is defined by

\[
B_{n,r} : T \times \mathbb{R}_+ \to T \times \mathbb{R}_+
((q,p), I) \mapsto (B_n(q,p), I \cdot r(q,p)), \tag{2}
\]

i.e., the intensity in \( A_k \) is changed by the factor \( r_k \). Often \( r_k \leq 1 \) is considered, leading to classical decay under the forward map [19]. On the other hand, considering \( r_k \geq 1 \) for some (or all) \( A_k \) implies growth of intensities [65]. The simple structure of the baker map with escape allows for determining the natural decay rates under forward and backward dynamics analytically. The natural measure is uniform along the unstable \( q \)-direction such that the natural decay rate is given by \( \gamma_{\text{nat}} = -\ln\left( \frac{1}{n} \sum_{k=0}^{n-1} r_k \right) \) [66]. The natural measure of the backward dynamics is uniform along the stable \( p \)-direction, such that the natural growth rate of the inverse map is \( \gamma_{\text{inv}} = \ln\left( \frac{1}{n} \sum_{k=0}^{n-1} \frac{1}{r_k} \right) \) [65, 48].

Another important classical decay rate is the average decay of a typical ergodic orbit.

![Figure 1](image-url)

**Figure 1.** (a) Illustration of reflectivities \( r_k \) in rectangles \( A_k \) for the 3-baker map with escape. (b) Phase-space partition by rectangles \( [\alpha_{m_p,m_q}] \) with \( m_p = 1 \) and \( m_q = 2 \) shown by dashed lines. The rectangle \([1 \cdot 02]\) and its iterate \([10 \cdot 2]\) are highlighted, where the latter intersects with 3 rectangles of the partition.
[34], the so-called typical decay rate $\gamma_{\text{typ}} = -\frac{1}{n} \sum_{k=0}^{n-1} \ln r_k$. One has the ordering $\gamma_{\text{nat}} \leq \gamma_{\text{typ}} \leq \gamma_{\text{inv}}$, with equality if all $r_k$ are identical.

2.2. Symbolic dynamics and phase-space partition

The baker map has a complete symbolic dynamics in terms of the Markov partition $\{A_k\}_{k=0}^{n-1}$. Therefore it is equivalent to a shift on symbolic sequences, which is recalled in the following. Let $S$ be the right shift operator on the symbolic space $Z_n = \{0,1,\ldots,n-1\}^2$, which consists of bi-infinite sequences $\alpha = \{\alpha_i\}_{i=-\infty}^{\infty} = \ldots\alpha_{-2}\alpha_{-1}\cdot\alpha_0\alpha_1\alpha_2\ldots$. The shift operator $S : Z_n \to Z_n$ acts on such sequences as

$$S(\alpha) = \ldots\alpha_{-2}\alpha_{-1}\alpha_0\cdot\alpha_1\alpha_2\ldots,$$

i.e., it shifts the center of this sequence by one such that $(S(\alpha))_i = \alpha_{i+1}$. There is a mapping from the space of symbolic sequences to the torus, $J_n : Z_n \to \mathbb{T}$,

$$J_n(\alpha) = (q,p) := \left(\sum_{i=0}^{\infty} \frac{\alpha_i}{n^{i+1}}, \sum_{i=0}^{\infty} \frac{\alpha_{-(i+1)}}{n^{i+1}}\right),$$

which is bijective on a suitable subset of $Z_n$ [40]. Thus, $J_n$ conjugates the symbolic shift with the baker map, $B_n = J_n \circ S \circ J_n^{-1}$, and they have equivalent dynamics.

The symbolic space allows to define a partition of the torus in terms of rectangles on phase space by finite sequences $\alpha_{m_p,m_q} = \{\alpha_i\}_{i=-m_p}^{m_q-1}$ of length $m_p + m_q$, see illustration in Fig. 1(b). For such a finite sequence a subset of $Z_n$ is defined as

$$[\alpha_{m_p,m_q}] := \{\beta \in Z_n : \beta_i = \alpha_i \ \forall -m_p \leq i \leq m_q - 1\},$$

which contains all infinite sequences with the central part being equal to $\alpha_{m_p,m_q}$. By construction $J_n([\alpha_{m_p,m_q}])$ is a rectangle on $\mathbb{T}$ and a subset of $A_{\alpha_0}$. There are $n^{m_p+m_q}$ such non-overlapping rectangles of size $\Delta p = 1/n^{m_p}$ and $\Delta q = 1/n^{m_q}$, see Fig. 1(b). In the following we will use the term rectangle for both, the sets of symbolic sequences $[\alpha_{m_p,m_q}]$ and their image $J_n([\alpha_{m_p,m_q}])$ on phase space.

3. Randomized quantum baker map with escape

In this section we first give the standard quantization of the baker map with escape. Then we introduce a randomization operator that acts locally on all phase-space rectangles, described by finite symbolic sequences, as introduced in Sec. 2.2. The resonance states of this randomized baker map are multifractal down to the scale of randomization and uniformly fluctuating on smaller scales. Thus, the proposed randomization separates the smallest classical scale from the scale of a Planck cell. This allows for a well-defined semiclassical limit, where the size of a Planck cell becomes arbitrarily small compared to all classical scales.

3.1. Quantum baker map with escape

The baker map can be quantized with geometric quantization [67, 68]. For this the Hilbert space $\mathbb{C}^N$ with $N = n^N$ is considered and the time-evolution operator on $\mathbb{C}^N$ is defined in position basis as

$$B_n := F^{-1}_n \text{diag}(F_{n/N}, \ldots, F_{n/N}).$$
Here, the matrices $F_M$ denote the discrete Fourier transform, which for arbitrary $M$ are given by

$$(F_M)_{ml} := M^{-\frac{1}{2}} e^{-\frac{2\pi i}{M} (m+\frac{1}{2}) (l+\frac{1}{2})}.$$  

Note that the exponents in the definition of $F$ ensure for the closed baker map that symmetries of $B_n$ are respected by the eigenfunctions of $B_n$ [68].

The quantization of the baker map with escape is composed of the closed quantum map $B_n$ and a reflectivity operator $R_r$. For constant escape from the rectangles $A_k$, as in the classical map, one has

$$R_r := \text{diag}(\sqrt{r_0} \mathbf{1}_{N/n}, \ldots, \sqrt{r_{n-1}} \mathbf{1}_{N/n}).$$  

The quantization of the baker map with escape, Eq. (2), is then given by

$$B_{n,r} := B_n R_r,$$

which has been extensively studied for full and partial escape [39, 40, 41, 42, 31, 45].

3.2. Randomization

We introduce a procedure to randomize the quantum dynamics of the $n$-baker map. As randomization regions we consider the rectangles $[\alpha_{\ell_p, \ell_q}]$ from a finite symbolic sequence of length $\ell_p$ in backward and $\ell_q$ in forward direction, see Eq. (5). The randomization within each of these $n^{\ell_p+\ell_q}$ rectangles is quantum mechanically achieved with random matrices $\mathcal{M}_{\text{cue}}$ of dimension $N_{\text{cue}}$, where we choose the circular unitary ensemble (CUE) which has no symmetries. Then the dimension $N$ of the full Hilbert space is given by $N = n^{\ell_q+\ell_p} N_{\text{cue}}$. The randomization operator in a basis corresponding to these rectangles is given by the block-diagonal operator

$$\hat{U}_{\text{cue}, \ell_p, \ell_q} = \text{diag}(\mathcal{M}_{\text{cue}}, \ldots, \mathcal{M}_{\text{cue}}),$$

consisting of $n^{\ell_q+\ell_p}$ CUE-matrices with dimension $N_{\text{cue}}$. In order to represent this operator in position basis we define a block matrix of $n^{\ell_q}$ Fourier transformations,

$$U_{\ell_q} = \text{diag}(F_{N/n^{\ell_q}}, \ldots, F_{N/n^{\ell_q}}),$$

acting on vertical phase-space regions that consist each of $n^{\ell_p}$ randomization regions. The randomization operator in position basis is then given by

$$U_{\ell_q}^{-1} \hat{U}_{\text{cue}, \ell_p, \ell_q} U_{\ell_q}.$$

Altogether, we define the randomized baker map with escape as the composition of randomization, escape and baker map,

$$B_{n,r,\ell_p,\ell_q} := B_n R_r U_{\ell_q}^{-1} \hat{U}_{\text{cue}, \ell_p, \ell_q}.$$  

3.3. Resonance states

Right resonance states of the randomized baker map with escape are defined by

$$B_{n,r,\ell_p,\ell_q} \psi = e^{i\theta-\gamma/2} \psi,$$

where $\gamma$ is the decay rate, i.e., the norm of $\psi$ decays by a factor of $e^{-\gamma}$ when $B_{n,r,\ell_p,\ell_q}$ is applied. The phase $\theta$ and decay rate $\gamma$ are related to the energy and width of resonance poles in autonomous scattering [69, 70]. The corresponding left resonance states are not discussed separately, as one can show that they are right resonance
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Figure 2. Phase-space distributions of resonance states and classical measures for randomization regions on increasingly finer levels (a) $(\ell_p, \ell_q) = (0, 1)$, (b) $(\ell_p, \ell_q) = (1, 2)$, and (c) $(\ell_p, \ell_q) = (2, 3)$ for the randomized 3-baker map with $r = (0.2, 0.01, 1)$. Shown is (left) Husimi function $\mathcal{H}_\psi$ for the resonance state $\psi$ with decay rate $\gamma$ closest to $\gamma_{\text{typ}}$ for $N = 25515$, (middle) projected weights $w_\psi(\{\alpha_{\ell_p+1, \ell_q}\})$ on the level of the randomization subregions, and (right) corresponding classical measure $\mu_{\ell_p, \ell_q}^{\psi}$ derived in Sec. 4.

states for the related operator, where $\mathcal{R}_r$ is replaced with $\mathcal{R}_r^{-1}$ in Eq. (13), i.e., with inverted reflectivity constants, $1/r_k$.

As we are interested in the phase-space distribution of resonance states $\psi$, we consider the Husimi function $\mathcal{H}_\psi$ [71] which is a smooth probability density. It is defined using the overlap of the state $\psi$ with a coherent state $c(q, p)$ centered at some phase-space point $(q, p)$,

$$\mathcal{H}_\psi(q, p) = N \| \langle c(q, p) | \psi \rangle \|^2. \quad (15)$$

In the following numerical illustrations the width of the coherent states is chosen to be symmetric in phase space.

We illustrate in Fig. 2(a) in the left panel a resonance state for the randomized
3-baker map with \( n = 3 \) stripes, reflectivities \( r = (0.2, 0.01, 1) \) and \( N = 35 \cdot 3^6 = 25515 \) with decay rate closest to \( \gamma_{\text{typ}} \). Here the randomization regions are defined by \( (\ell_p, \ell_q) = (0, 1) \), i.e., they correspond to the rectangles \( A_k \). One observes that the resonance state is not uniformly fluctuating in each of the \( n = 3 \) randomization regions. Instead, we find a finer substructure, where each randomization region has \( n = 3 \) subregions in \( p \)-direction, leading to a \( 3 \times 3 \) grid. This subregion structure is discussed in more detail in Appendix A.

Within each subregion the resonance state visually fluctuates as in closed chaotic quantum maps and we expect an exponential distribution around the average, as studied in Ref. [50] for quantum maps with escape and without randomization. The subregion structure remains the same for increasing \( N \) and the Planck cell becomes arbitrarily small compared to it, giving a well-defined semiclassical limit. This is in contrast to a usual quantum map with escape, where even for arbitrarily large \( N \) the phase-space structure of a resonance state is multifractal down to the scale of a Planck cell.

In Figs. 2(b) and (c) the randomization regions are defined on finer scales by \( (\ell_p, \ell_q) = (1, 2) \) and \( (\ell_p, \ell_q) = (2, 3) \), respectively. The resonance states (left panel) are again uniformly fluctuating on \( n \) subregions of each randomization region, leading to a \( n^\ell_p + 1 \times n^\ell_q \)-grid structure. Again, for increasing \( N \) the Planck cell size becomes arbitrarily small compared to the phase-space structures of resonance states.

For increasingly finer randomization levels we observe in Figs. 2(a)–(c) that the phase-space structure on smaller scales resembles (qualitatively) the phase-space structure on larger scales. For example the region \( (q, p) \in [1/3, 2/3] \times [2/3, 1] \) in Fig. 2(b) looks similar to the full phase space in Fig. 2(a). In the limit of \( \ell_p, \ell_q \to \infty \) with sufficiently large \( N \) this leads to multifractal structures.

Note as an aside, that if for a fixed Hilbert space dimension \( N \), the number \( n^\ell_p + \ell_q \) of randomization regions is increased, the dimension \( N_{\text{cue}} = N/n^\ell_p + \ell_q \) of the random matrices becomes smaller as well as the number of Planck cells \( N_{\text{sub}} = N_{\text{cue}}/n = N/n^{\ell_p + 1 + \ell_q} \) within each randomization subregion. In Figs. 2(a)–(c) this leads to \( N_{\text{sub}} = 2835 \), \( N_{\text{sub}} = 315 \), and \( N_{\text{sub}} = 35 \), respectively. Numerically, this limits the number of levels that can be explored, while still probing the semiclassical limit.

### 3.4. Weight of resonance state on rectangular phase-space regions

In order to compare the phase-space structure of a resonance state with some classical measure, we choose rectangular phase-space regions onto which we project the resonance state. A suitable projector is defined, by first projecting the position representation to the interval corresponding to the extent of the rectangle in \( q \)-direction. The second step is a transformation to the momentum representation and projection on the interval corresponding to the extent of the rectangle in \( p \)-direction.

For convenience we use the rectangles \( [\alpha_{mp,mq}] \) partitioning the phase space based on the finite symbolic sequences of the baker map. We define a projector \( P_{[mp,mq]} \) for all \( n^m_{p} + n^m_{q} \) possible finite sequences \( \alpha_{mp,mq} \) with \( \alpha_i \in \{0, \ldots, n - 1\} \),

\[
P_{[\alpha_{mp,mq}]} = \prod_{n_{\text{rect}}} \prod_{n_{\text{rect}}} {I_{\text{rect}}^{N_{\text{rect}}}} \prod_{n_{\text{rect}}} {I_{\text{rect}}^{N_{\text{rect}}}} \prod_{n_{\text{rect}}} {I_{\text{rect}}^{N_{\text{rect}}}} \prod_{n_{\text{rect}}} {I_{\text{rect}}^{N_{\text{rect}}}} \prod_{n_{\text{rect}}} {I_{\text{rect}}^{N_{\text{rect}}}} \prod_{n_{\text{rect}}} {I_{\text{rect}}^{N_{\text{rect}}}}
\]

where \( N_{\text{rect}} \) is the dimension of the Hilbert space of each rectangle and \( N = n^m_{p} + n^m_{q} \), \( N_{\text{rect}} \) is the total Hilbert space dimension. Here the rectangles \( [\alpha_{mp,mq}] \) are ordered according to the index \( I(\alpha_{mp,mq}) = \sum_{i=0}^{n^m_{p}} \alpha_i - m_p n^i + \)
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\[ \sum_{i=0}^{m-1} \alpha_{m-1-i} n^i \]

which counts the rectangles first in the leftmost column from bottom to top, then in the second column, and so on. The projected weight of a state \( \psi \) on a rectangle \( [\alpha_{m_p, m_q}] \) is then given by

\[ w_\psi([\alpha_{m_p, m_q}]) = \| P_{[\alpha_{m_p, m_q}]} \psi \|^2. \]  

(17)

Alternatively, one may determine the weights by integrating the Husimi function over the corresponding rectangle on \( \mathbb{T} \),

\[ w_\psi([\alpha_{m_p, m_q}]) = \int J_n([\alpha_{m_p, m_q}]) H_\psi(q,p) \, d\mu_L. \]

(18)

Note that the weights of both approaches agree in the semiclassical limit \( N \rightarrow \infty \).

Convergence is slower when using the Husimi function, as the weight in rectangles with small weights is strongly influenced by neighboring rectangles with large weights due to the finite resolution of the Husimi function.

In Figs. 2(a-c) in the middle panel the weight on all subregion rectangles \( [\alpha_{\ell_p+1, \ell_q}] \) of a resonance state is illustrated. This can be compared to classical measures in the right panel, derived in Sec. 4, showing excellent agreement on a qualitative level. A quantitative comparison will be presented in Sec. 5.

We emphasize that while Fig. 2 shows results for just one exemplary decay rate \( \gamma_{\text{typ}} \), the weight in each subregion rectangle depends on the decay rate \( \gamma \). It is the goal of this paper to derive classical measures for arbitrary decay rates. In the next section we consider randomization regions with arbitrary but fixed \( \ell_p \) and \( \ell_q \) and make a conjecture for the semiclassical limit \( N \rightarrow \infty \).

4. Semiclassical structure of resonance states

In this section we present a conjecture for the semiclassical structure of resonance states of the randomized baker map. To this end we introduce classical measures, which (i) are conditionally invariant with some desired decay rate \( \gamma \) on scales larger than the scale of randomization and (ii) obey an extremum principle derived from a local random wave model in each randomization region.

4.1. Conditional invariance on rectangles

Semiclassically, resonance states of quantum maps with escape converge to conditionally invariant measures of the same decay rate \([39, 40, 47]\). A measure \( \mu \) on \( \mathbb{T} \) is conditionally invariant \([51]\) with respect to the classical baker map \( B_n \) with escape \( r \), if there exists a decay rate \( \gamma \in \mathbb{R} \) with

\[ \int_{B_n^{-1}(A)} r \, d\mu = e^{-\gamma} \mu(A) \quad \forall A \subset \mathbb{T}. \]  

(18)

For constant reflectivity factors \( r_k \) from the rectangles \( A_k \), as considered in this paper, this condition simplifies for subsets of \( A_k \) to

\[ r_k \mu(C) = e^{-\gamma} \mu(B_n(C)) \quad \forall C \subset A_k. \]  

(19)

In general, Eq. (18) holds for all measurable sets \( A \subset \mathbb{T} \) up to arbitrary small scales, which is related to the multifractal structure of conditionally invariant measures. However, for the randomized quantum baker map, as introduced in the last section, we observe no structure of resonance states on scales finer than the subregions of the randomization regions \( [\alpha_{\ell_p, \ell_q}] \), see Figs. 2(a-c) left panel. For this reason, conditional invariance cannot be satisfied on arbitrarily fine scales. This motivates to introduce the following scale dependent definition. We call a measure \( \mu \) on \( \mathbb{T} \) conditionally
invariant down to the scale \((\ell_p, \ell_q)\) with respect to the baker map with escape \(B_{n,r}\), if Eq. (19) holds for all sets \(C = J_n([(\alpha_{\ell_p, \ell_q}]])\).

We assume for the randomized quantum baker map, that the resonance states converge in the semiclassical limit to measures which satisfy conditional invariance down to the scale \((\ell_p, \ell_q)\) of the randomization regions and are constant on the subregions of scale \((\ell_p + 1, \ell_q)\). Such measures are shown in Figs. 2(a-c) right panel and are derived in this section.

From conditional invariance down to the scale of randomization we derive constraints for \(\mu\). Recall that the randomization regions are defined as the rectangles \([\alpha_{\ell_p, \ell_q}] = [\alpha_{-\ell_p} \ldots \alpha_{-1} \cdot \alpha_0 \alpha_1 \ldots \alpha_{\ell_q-1}]\) with \(\alpha_i \in \{0, \ldots, n-1\}\) for all \(-\ell_p \leq i \leq \ell_q - 1\). Their iterates are given by \([S(\alpha_{\ell_p, \ell_q}]) = [\alpha_{-\ell_p} \ldots \alpha_{-1} \cdot \alpha_0 \alpha_1 \ldots \alpha_{\ell_q-1}]\) - these iterates correspond to rectangles, which are stretched by a factor \(n\) in \(q\)-direction and compressed by \(1/n\) in \(p\)-direction. Each of these iterates overlaps with exactly \(n\) of the randomization regions (and vice versa), see illustration in Fig. 1(b). Hence, the original and the iterated randomization rectangles can be partitioned as follows,

\[
\bigcup_{j=0}^{n-1}[j\alpha_{\ell_p, \ell_q}] = \bigcup_{j=0}^{n-1}[S(\alpha_{\ell_p, \ell_q})] \quad (20)
\]

where for each \(j \in \{0, \ldots, n-1\}\) the smaller rectangles \([j\alpha_{\ell_p, \ell_q}] := [\alpha_{-\ell_p} \ldots \alpha_{-1} \cdot \alpha_0 \alpha_1 \ldots \alpha_{\ell_q-1}]\) and \([S(\alpha_{\ell_p, \ell_q})] := [\alpha_{-\ell_p} \ldots \alpha_{-1} \cdot \alpha_0 \alpha_1 \ldots \alpha_{\ell_q-1}]\) are the subregions of the randomization regions on which the measure \(\mu\) is constant.

Conditional invariance down to the scale of the randomization regions \([\alpha_{\ell_p, \ell_q}]\) implies, by setting \(C = J_n(\alpha_{\ell_p, \ell_q})\) in Eq. (19) with \(C \subset A_{\ell_p}\), that

\[
r_{\ell_p, \ell_q} \mu((\alpha_{\ell_p, \ell_q})) = e^{-\gamma} \mu([S(\alpha_{\ell_p, \ell_q})])
\]

for all multi-indices \(\alpha_{\ell_p, \ell_q}\). Inserting the partitions from Eq. (20) we get

\[
r_{\ell_p, \ell_q} \mu\left(\bigcup_{j=0}^{n-1}[j\alpha_{\ell_p, \ell_q}]\right) = e^{-\gamma} \mu\left(\bigcup_{j=0}^{n-1}[S(\alpha_{\ell_p, \ell_q})] \right)
\]

This gives the constraints of conditional invariance on the weights of all subregion rectangles with scales \((\ell_p + 1, \ell_q)\),

\[
r_{\ell_p, \ell_q} \sum_{j=0}^{n-1} \mu([j\alpha_{\ell_p, \ell_q}]) = e^{-\gamma} \sum_{j=0}^{n-1} \mu([S(\alpha_{\ell_p, \ell_q})]), \quad (23)
\]

holding for all possible finite symbolic sequences of the form \(\alpha_{\ell_p, \ell_q} = \alpha_{-\ell_p} \ldots \alpha_{-1} \cdot \alpha_0 \alpha_1 \ldots \alpha_{\ell_q-1}\) with \(\alpha_i \in \{0, \ldots, n-1\}\). Altogether, these are \(n^{\ell_p+\ell_q}\) constraints on the subregion weights.

Normalization of the measure \(\mu\)

\[
\mu(T) = \mu\left(\bigcup_{j=0}^{n-1} \bigcup_{\alpha_{-\ell_p}=0}^{n-1} \bigcup_{\alpha_0=0}^{n-1} \bigcup_{\alpha_{\ell_q-1}=0}^{n-1}[j\alpha_{\ell_p, \ell_q}]\right) = 1 \quad (24)
\]

adds one more constraint

\[
\sum_{j=0}^{n-1} \sum_{\alpha_{-\ell_p}=0}^{n-1} \sum_{\alpha_0=0}^{n-1} \sum_{\alpha_{\ell_q-1}=0}^{n-1} \mu([j\alpha_{\ell_p, \ell_q}]) = 1. \quad (25)
\]

Thus, conditional invariance down to scale \((\ell_p, \ell_q)\), Eq. (23), and normalization, Eq. (25), give \(n^{\ell_p+\ell_q}+1\) constraints on the weights of all \(n^{\ell_p+\ell_q}\) subregion rectangles. This leads to the question, how quantum mechanics distributes the weight over the subregion rectangles under the restrictions of Eqs. (23) and (25).
4.2. Local random wave model

It is well known that eigenfunctions of closed chaotic systems are described by Berry’s random wave model [2] for billiards [12, 13, 14, 15] and maps [16, 17, 11, 18]. There are also extensions of this model including constraints by two-point correlations [72, 73, 74] and systems with a mixed phase space [75, 76]. Based on a local variant of the random wave model, we derive in the following the semiclassical structure of resonance states of the randomized quantum baker map and its dependence on the decay rate $\gamma$.

We introduce a local random wave model for each randomization region as follows. Let $[\alpha_{\epsilon_p},\epsilon_q]$ be one of these randomization regions and let $N_{\text{cue}}$ be the number of sites of a random vector, corresponding to the number of Planck cells in the randomization region. The a priori unknown weight $\mu([\alpha_{\epsilon_p},\epsilon_q])$ of the randomization region is assumed to be distributed on these sites according to independent exponentially distributed site weight factors $f_i \geq 0$ for each $i \in \{0,\ldots,N_{\text{cue}} - 1\}$,

$$P(f_i) \, df_i = N_{\text{cue}} \, e^{-N_{\text{cue}} f_i} \, df_i,$$

such that the weight at site $i$ is given by $f_i \mu([\alpha_{\epsilon_p},\epsilon_q])$. Such an exponential distribution of the factors $f_i$ follows from assuming Gaussian distributed complex amplitudes of random vectors describing a given randomization region. The mean value of the exponential distribution is $1/N_{\text{cue}}$ such that on average $\sum_{i=0}^{N_{\text{cue}} - 1} f_i = 1$ holds. If the randomization region is the entire phase space $T$, which has measure $\mu(T) = 1$, this is equivalent to the usual random wave model for chaotic maps, see e.g., Ref. [11].

For smaller randomization regions $[\alpha_{\epsilon_p},\epsilon_q]$ it describes the local fluctuations within each randomization region, but has no restrictions on the weights $\mu([\alpha_{\epsilon_p},\epsilon_q])$ of each randomization region.

Let $\{f_0,\ldots,f_{N_{\text{cue}} - 1}\}$ be a set of site weight factors for one of the randomization regions, which exactly fulfills the normalization $\sum_{i=0}^{N_{\text{cue}} - 1} f_i = 1$. All such sets have the same probability, as the product of the individual exponential distributions is a constant, $\prod_i e^{-N_{\text{cue}} f_i} = e^{-N_{\text{cue}} \sum_i f_i} = e^{-N_{\text{cue}}}$.

Each randomization region consists of $n$ subregions $[j\alpha_{\epsilon_p},\epsilon_q]$ with $j \in \{0,\ldots,n-1\}$, where each subregion has $N_{\text{sub}}$ sites if $N_{\text{cue}} = nN_{\text{sub}}$ is chosen appropriately. We define a subregion weight factor $F_j$ by

$$F_j = \frac{\mu([j\alpha_{\epsilon_p},\epsilon_q])}{\mu([\alpha_{\epsilon_p},\epsilon_q])}.$$  

The $F_j$ are thus given by the sum over all corresponding site weight factors $f_i$. For a suitable ordering of the $f_i$ this can be expressed as

$$F_j = \sum_{i=0}^{N_{\text{sub}} - 1} f_j N_{\text{sub}} + i \quad \forall j = 0,\ldots,n - 1.$$

Note that normalization of the subregion weight factors, $\sum_{j=0}^{n-1} F_j = \sum_{i=0}^{N_{\text{cue}} - 1} f_i = 1$, follows from the above normalization of the site weight factors $f_i$.

The probability $P\{F_j\}$ for a set $\{F_0,\ldots,F_{n-1}\}$ of subregion weight factors in a randomization region is given by integrating over the probability of all sets of site weight factors, $\{f_0,\ldots,f_{N_{\text{cue}} - 1}\}$, fulfilling Eqs. (28) for all $j$. Recall that the probability of the sets $\{f_0,\ldots,f_{N_{\text{cue}} - 1}\}$ is uniformly distributed, as shown above. Therefore the probability $P\{F_j\}$ is proportional to the volume occupied by the sets $\{f_0,\ldots,f_{N_{\text{cue}} - 1}\}$, fulfilling Eqs. (28) for all $j$. Specifically, Eq. (28) for a given $j$ defines a $(N_{\text{sub}}-1)$-dimensional standard simplex within the $N_{\text{sub}}$-dimensional space.
In this product we insert the definition of the subregion weight factors, Eq. (27), and thus has a volume proportional to \( F_j^{N_{\text{sub}}-1} \). For example, for \( N_{\text{sub}} = 3 \) it is a triangle in the 3D space \((f_{3j}, f_{3j+1}, f_{3j+2})\) with corners \((F_j, 0, 0), (0, F_j, 0), (0, 0, F_j)\) and 2D volume proportional to \( F_j^2 \). Thus, the total volume is proportional to an n-fold product of such simplices and we obtain for the probability of a set of subregion weight factors,

\[
P(\{F_j\}) \propto \prod_{j=0}^{n-1} F_j^{N_{\text{sub}}-1}.
\]  

(29)

We now extend this argument to all \( n^{\ell_p+\ell_q} \) randomization regions \([\alpha_{\ell_p, \ell_q}]\), each with \( n \) subregion weight factors \( F_j\alpha_{\ell_p, \ell_q} \) from the local random wave model. Assuming independence, the probability for a set of all subregion weight factors is proportional to

\[
P(\{F_j\alpha_{\ell_p, \ell_q}\}) \propto \prod_{\alpha_{\ell_p, \ell_q}} \prod_{j=0}^{n-1} F_j^{N_{\text{sub}}-1} = \left( \prod_{\alpha_{\ell_p, \ell_q}} \prod_{j=0}^{n-1} F_j \alpha_{\ell_p, \ell_q} \right)^{N_{\text{sub}}-1}.
\]  

(30)

In the semiclassical limit, \( N_{\text{sub}} \to \infty \), this distribution approaches a delta function at the maximum of the product of the subregion weight factors, \( \prod_{\alpha_{\ell_p, \ell_q}} \prod_{j=0}^{n-1} F_j \alpha_{\ell_p, \ell_q} \). In this product we insert the definition of the subregion weight factors, Eq. (27),

\[
F_j \alpha_{\ell_p, \ell_q} = \frac{\mu(j \alpha_{\ell_p, \ell_q})}{\sum_{l=0}^{n-1} \mu(l \alpha_{\ell_p, \ell_q})},
\]  

(31)

using the subregion weights \( \mu(j \alpha_{\ell_p, \ell_q}) \) and the weight of the corresponding randomization region \( \mu([\alpha_{\ell_p, \ell_q}]) = \sum_{l=0}^{n-1} \mu(l \alpha_{\ell_p, \ell_q}) \). This allows for writing the product of subregion weight factors as a product of relative weights,

\[
\Pi(\mu) = \prod_{\alpha_{\ell_p, \ell_q}} \prod_{j=0}^{n-1} \frac{\mu(j \alpha_{\ell_p, \ell_q})}{\sum_{l=0}^{n-1} \mu(l \alpha_{\ell_p, \ell_q})}.
\]  

(32)

Thus the local random wave model for each randomization region leads to the following extremum principle: The subregion weights \( \mu(j \alpha_{\ell_p, \ell_q}) \) maximize the product \( \Pi(\mu) \) of relative weights, Eq. (32), under the constraints of conditional invariance, Eq. (23), and normalization, Eq. (25). They can be determined, e.g., by using the method of Lagrange multipliers leading to a set of coupled nonlinear equations. These subregion weights \( \mu(j \alpha_{\ell_p, \ell_q}) \) define a classical measure \( \mu^{\ell_p, \ell_q} \) on \( T \), that is conditionally invariant down to the scale of the randomization regions, determined by \( \ell_p \) and \( \ell_q \), and that depends on the decay rate \( \gamma \).

In Figs. 2(a–c) in the right panel examples of such classical measures \( \mu^{\ell_p, \ell_q} \) for three increasing levels of randomization \( \ell_p \in \{0, 1, 2\} \) with \( \ell_q = \ell_p + 1 \) are shown. We find quite good agreement with the structure of the resonance states (left panel) and the projected weights within each subregion rectangle (middle panel). Slight deviations are observable in (c), where the subregion dimension \( N_{\text{sub}} = 35 \) is quite small and the projected weights are still influenced by the fluctuations of an individual resonance state.
4.3. Semiclassical conjecture

The following conjecture is formulated in analogy to quantum ergodicity in closed chaotic maps [11, 77]. For finite $\ell_p$ and $\ell_q$ we conjecture that the classical measures $\mu_{\ell_p,\ell_q}$ are the semiclassical limit measures of resonance states $\psi$ of the randomized quantum baker map with escape $B_{n,r,\ell_p,\ell_q}$, which decay with the rate $\gamma$. In particular, let $B_{n,r,\ell_p,\ell_q}^N$ be the randomized quantum baker with dimension $N$. If $\psi_N$ is a sequence of eigenstates, $B_{n,r,\ell_p,\ell_q}^N \psi_N = \lambda_N \psi_N$, with decay rates $\gamma_N = -\ln |\lambda_N|^2 N \rightarrow \infty$, then for all observables $a$ on $T$ with quantization $\text{op}_N(a)$ we expect

$$\langle \psi_N | \text{op}_N(a) | \psi_N \rangle \rightarrow \int_T a \, d\mu_{\ell_p,\ell_q}. \quad (33)$$

This means, that the quantum expectation values converge to the classical ones given by $\mu_{\ell_p,\ell_q}$. Note that similar to quantum ergodicity there could be exceptional states or sequences, such that convergence in Eq. (33) is expected just for a suitable subsequence of density one. A quantitative analysis of the quantum-classical agreement is given in Sec. 5.

4.4. Properties of the classical measure $\mu_{\ell_p,\ell_q}$

We list some relevant properties and remarks on the classical measures $\mu_{\ell_p,\ell_q}$:

(i) Numerics for increasing $\ell_p,\ell_q$ suggests that in the limit $\ell_p,\ell_q \rightarrow \infty$ the classical measures $\mu_{\ell_p,\ell_q}$ converge weakly to some limiting measure $\mu^*$, see Appendix B.

(ii) The product $\prod(\mu)$ of relative weights, Eq. (32), is maximal if all relative weights are equally given by $1/n$. Assuming conditional invariance one can show that this implies that the measures are also uniform within the sets $[\alpha_0,\ell_q]$, i.e., along the $p$-direction. The subregion weights follow as $\mu([\alpha_{\ell_p+1,\ell_q}]) = \prod_{i=0}^{\ell_q-1} r_{\alpha_i}^{-1}/N$ with some normalization $N$. This measure is conditionally invariant with decay rate $\gamma_{\text{inv}}$, the natural growth rate of the inverse map, see Sec. 2.1.

(iii) For the natural decay rate $\gamma_{\text{nat}}$ one needs to rewrite the product $\prod(\mu)$ using conditional invariance, Eq. (23), with decay rate $\gamma_{\text{nat}}$, leading after reordering to

$$\Pi(\mu) = \prod_{\alpha_{\ell_p,\ell_q}} \prod_{j=0}^{n-1} r_{\alpha_{\ell_p-1,\ell_q}} e^{\gamma_{\text{nat}}} \frac{\mu([\alpha_{\ell_p,\ell_q}])}{\sum_{l=0}^{n-1} \mu([\alpha_{\ell_p,\ell_q}])}. \quad (34)$$

This product is maximal if all relative weights, now with respect to iterated randomization regions, are equally given by $1/n$. One can show that this implies that all subregion weights are given by $\mu([\alpha_{\ell_p+1,\ell_q}]) = \prod_{i=0}^{\ell_q-1} r_{\alpha_i}/N$ with some normalization $N$, i.e., the measure is uniform in $q$-direction. This measure is indeed conditionally invariant with decay rate $\gamma_{\text{nat}}$.

(iv) The classical measure $\mu_{\ell_p,\ell_q}$ for the special case $\gamma_{\text{nat}}$ ($\gamma_{\text{inv}}$) is consistent with the general observation that at the (inverse) natural decay rate resonance states are uniform along the unstable (stable) direction [19, 48].
(v) We mention as an aside, that if one applies this extremum principle to closed chaotic systems, it leads to the uniform distribution, which is consistent with quantum ergodicity [4, 5, 6].

(vi) It is possible to apply the local random wave model combined with conditional invariance to random matrices with partial escape, consistently retrieving the results of Ref. [50, App. D]. The details are given in Appendix C.

5. Numerical verification

In this section we present quantitative numerical support for the semiclassical conjecture given in the last section, going beyond the qualitative comparison of Fig. 2.

5.1. Weights on classically large regions

In order to test the semiclassical conjecture, Eq. (33), we choose as observables the projection operators on classically large regions compared to a Planck cell. Specifically, we use squares given by symbolic sequences $[\alpha_{\ell,\ell}]$ with $\ell = 1$ and $\ell = 2$. For different systems and Hilbert space dimensions we consider all resonance states $\psi_i$ and show how the projected weights $w_{\psi_i}([\alpha_{\ell,\ell}])$, Eq. (17), depend on their quantum decay rate $\gamma_i$. This is compared to the classical measures $\mu_{\ell_p,\ell_q}([\alpha_{\ell,\ell}])$ on these squares and their

![Figure 3](image-url)

**Figure 3.** Comparison of the 9 projected weights $w_{\psi_i}([\alpha_{1,1}])$ of each resonance state $\psi_i$ vs. its decay rate $\gamma_i$ (dots) and classical measures $\mu_{\ell_p,\ell_q}([\alpha_{1,1}])$ (lines) for 3-baker map with $r = (0.2, 0.01, 1)$ and randomization level $(\ell_p, \ell_q) = (0, 1)$ for Hilbert space dimension (a) $N = 2835$ and (b) $N = 25515$. The same data is shown on linear (left) and logarithmic scale (middle) as well as a magnification (right). Decay rates $\gamma_{\text{nat}}$, $\gamma_{\text{typ}}$, and $\gamma_{\text{inv}}$ are indicated (dotted lines).
dependence on $\gamma$. The Hilbert space dimension $N$ is in all cases chosen to be a multiple of the number $n^{\ell_p+1+\ell_q}$ of subregions.

In Fig. 3 we consider the 3-baker map with $r = (0.2, 0.01, 1)$ and randomization level $(\ell_p, \ell_q) = (0, 1)$ for increasing Hilbert space dimension. For $\ell = 1$ there are $n^{2\ell} = 9$ quantum and classical weights which are compared. We observe that the quantum weights (dots) scatter symmetrically around the classical measures (lines). The width of the distribution decreases with increasing Hilbert space dimension for each of the considered squares (compare (a) and (b)), supporting the semiclassical conjecture. This also holds for small weights, as can be seen on the logarithmic scale (middle panel).

At the natural decay rates $\gamma_{\text{nat}}$ and $\gamma_{\text{inv}}$ the measures are constant along one direction, see Sec. 4.4, and several weights are identical in agreement with the properties of the resonance states.

Figure 4 shows the comparison for smaller randomization regions $(\ell_p, \ell_q) = (1, 2)$ and weights on smaller phase-space regions $[\alpha_{\ell, \ell}]$ with $\ell = 2$. There are $n^{2\ell} = 81$ such regions out of which 9 are presented in the figure, corresponding to the phase-space region $(q, p) \in [1/3, 2/3] \times [2/3, 1]$, see Fig. 2(a). Again, quantum and classical weights show excellent agreement. The width of each distribution is wider than in Fig. 3(b), as the observables are smaller phase-space regions. For larger Hilbert space dimension

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure4.png}
\caption{Same as Fig. 3(b), but for randomization level $(\ell_p, \ell_q) = (1, 2)$ and for 9 selected squares from $[\alpha_{2,2}]$ with symbolic sequence $\alpha = i2 \cdot 1j$ and $i, j \in \{0, 1, 2\}$.}
\end{figure}

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure5.png}
\caption{Comparison of $n^{2\ell}$ projected weights $w_{\psi_i}(\alpha_{\ell, \ell})$ of each resonance state $\psi_i$ (dots) and classical measures $\mu_{\ell_p}^{\ell_q}([\alpha_{\ell, \ell}])$ (lines) for 2-baker map with $r = (0.1, 1)$ and randomization level $(\ell_p, \ell_q) = (\ell - 1, \ell)$ for (a) $\ell = 1$ and (b) $\ell = 2$ for Hilbert space dimension $N = 25600$. Decay rates $\gamma_{\text{nat}}, \gamma_{\text{typ}},$ and $\gamma_{\text{inv}}$ are indicated (dotted lines).}
\end{figure}
$N$ it is expected to decrease. We emphasize that the dependence on $\gamma$ qualitatively resembles the results for $(\ell_p, \ell_q) = (0, 1)$ shown in Fig. 3, but now on a smaller scale in phase space. This indicates again the multifractal phase-space structure, which emerges for smaller randomization regions, as observed in Fig. 2.

Additionally, we show the quantum-classical comparison for the 2-baker map with $r = (0.1, 1)$ for two randomization levels in Fig. 5. Again, we observe excellent agreement for all decay rates.

5.2. Jensen-Shannon divergence

In order to further analyze the semiclassical convergence of resonance states we consider the Jensen-Shannon divergence between quantum and classical measures,

$$d_{JS}(\mu_{qm}, \mu_{cl}) = H\left(\frac{\mu_{qm} + \mu_{cl}}{2}\right) - \frac{H(\mu_{qm}) + H(\mu_{cl})}{2},$$

(35)

where $H$ is the Shannon entropy

$$H(\mu) = -\sum_{\alpha, \ell} \mu([\alpha, \ell]) \ln \mu([\alpha, \ell]),$$

(36)

evaluated on rectangles $[\alpha, \ell]$ on level $\ell$. The quantum measure $\mu_{qm}$ is determined by the projected weights $w_{\psi}(\alpha_{\ell, \ell})$ of a resonance state $\psi$ and the classical measure $\mu_{cl}$ is given by $\mu_{\ell_p \ell_q}$ for $\gamma = \gamma_i$. The square root of $d_{JS}$ is a distance on the set of measures. Weak convergence of the quantum measure to the classical measure implies

Figure 6. Jensen-Shannon divergence $d_{JS}$ between quantum and classical measures vs. $\gamma$ for baker map with $r = (0.2, 0.01, 1)$ and randomization on level (a) $(\ell_p, \ell_q) = (0, 1)$ for increasing $N \in \{315, 2835, 25515\}$, and (b) $(\ell_p, \ell_q) = (1, 2)$ for $N \in \{324, 2835, 25515\}$ evaluated on level $\ell = \ell_q$. 
that for finite $\ell$ the Jensen-Shannon entropy must converge to zero in the semiclassical limit $N \to \infty$.

In Fig. 6 we illustrate $d_{JS}$ versus $\gamma$ for increasing $N$ for two different randomization levels. For all decay rates $\gamma$ the Jensen-Shannon divergence decreases with increasing $N$. This supports the semiclassical weak convergence conjectured in Sec. 4.3.

We remark that we tested the semiclassical conjecture qualitatively and quantitatively, for many different randomized baker maps, i.e., different $n$, different $r$, and also different combinations of $\ell_p$ and $\ell_q$. In all cases we found quantum-classical agreement in support of the semiclassical conjecture.

6. Quantum baker map with escape

In this section we compare the structure of resonance states of the usual quantum baker map with escape, i.e., without randomization, to the randomized baker map with escape. In Fig. 7(a) in the left panel the Husimi function of a resonance state of the usual 3-baker map with escape $B_{n,r}$, Eq. (9), is shown. This is qualitatively similar to the Husimi functions of the randomized baker map in Fig. 2(c) for small randomization regions. Here, we compare it to the classical measures $\mu_{n,\ell_p,\ell_q}^{\ell_p,\ell_q}$ for even smaller randomization, $(\ell_p, \ell_q) = (4, 5)$, see middle panel in Fig. 7(a). These measures

![Figure 7](attachment:figure7.png)

**Figure 7.** 3-baker map without randomization with $r = (0.2, 0.01, 1)$ and $N = 25191$. (a) Husimi function of resonance state with decay rate closest to $\gamma_{typ}$ (left), corresponding classical measure $\mu_{4,5}$ (middle), and product measure $\mu_{prod}$ based on Ref. [48] (right). The same colormap is used for all densities. (b) Comparison of the 9 projected weights $w_\psi(\alpha_{1,1})$ of each resonance state $\psi$ (dots) to classical measures $\mu_{4,5}$ (solid lines) and $\mu_{prod}$ (dashed lines). The same data is shown on linear (left) and logarithmic scale (middle), as well as a magnification (right).
approximate the limit $\ell_p, \ell_q \to \infty$, see Appendix B. We observe qualitative agreement between the multifractal structures of both measures.

Quantitative differences can be seen in Fig. 7(b), where quantum and classical weights on the rectangles $[\alpha_\ell, \ell]$ with $\ell = 1$ are shown for all decay rates $\gamma$. For example, in the left panel the largest quantum weights do not scatter symmetrically around the corresponding classical weight. Similar deviations can also be seen in the right panel. Therefore the semiclassical limit measures $\mu_{\ell_p, \ell_q}^{\gamma}$ of the randomized baker map just approximately describe resonance states of the baker map without randomization. Numerics suggest that this does not change, neither for larger $N$ in the baker map with escape, nor for classical measures with larger $\ell_p, \ell_q$, see Appendix B.

Additionally, Fig. 7 shows a comparison to the product measures, proposed as approximate descriptions of resonance states in Ref. [48]. In Fig. 7(a) in the right panel the product measure visually shows differences on small scales. In Fig. 7(b), where the product measure is shown as a dashed line, it deviates substantially, e.g., in the middle panel for the smallest weight at intermediate $\gamma$.

There are two main conclusions from the comparison of the baker map with and without randomization. First of all, in both models resonance states have qualitatively similar multifractal structures on scales larger than the scale of randomization. Secondly, however, we observe that these models have slightly different structures of their resonance states, even when the scale of randomization is chosen very small. Thus it remains open, if and how the local random wave model and the classical measures $\mu_{\ell_p, \ell_q}$ can be generalized to systems without randomization.

7. Conclusions and outlook

In conclusion, we derive a semiclassical description for resonance states of randomized baker maps with escape based on conditional invariance and a local random wave model. We present numerical support which indicates that the conjectured measures $\mu_{\ell_p, \ell_q}^{\gamma}$ are indeed the semiclassical limit measures.

This model with randomization acts as a bridge towards a full understanding of resonance states in chaotic systems with escape. The following generalizations are of interest: (i) The limit from partial to full escape can be investigated for the randomized baker map and compared with the existing approximative description for systems with full escape [47]. (ii) Generalized baker maps with stripes of arbitrary width allow for studying non-uniformly hyperbolic systems. The randomization regions could either be given by finite symbolic sequences of this map leading to different sizes (and straightforward modifications of the local random wave model) or be given by regions of identical sizes as in this work. (iii) Randomization of other chaotic maps and for smooth reflectivity functions, as they occur in optical microcavities, should be studied. Classical measures, based on the local random wave model and conditional invariance, might again perfectly describe their resonance states. Furthermore, these classical measures might give rise to the best available approximations for the resonance states of non-randomized maps, similar to our observation for the baker map in Sec. 6.

Finally, we observe that it makes a quantitative difference, whether the multifractal structure of resonance states ends at the scale of a Planck cell or at the scale of randomization, even when the resonance states are compared on much larger classical scales. This insight should be helpful for future work on semiclassical theories for resonance states.
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Appendix A. Substructure of randomization regions

In this section we discuss the observation that resonance states of the randomized baker map have substructure within the randomization regions, see Fig. 2 in Sec. 3.3.

First we illustrate in Fig. A1 how a resonance state \( \psi \) evolves under sequential application of the operators for randomization \( U_{\ell_p,\ell_q}^{\text{cue}} \), escape \( R_r \), and baker map \( B_n \), that appear in the randomized baker map with esape \( B_{n,r,\ell_p,\ell_q} \), Eq. (13). The same resonance state as in Fig. 2(a, left panel) is used, where each of the 3 randomization regions is divided along the \( p \)-direction into 3 subregions. After application of the randomization operator, \( U_{\ell_p,\ell_q}^{\text{cue}} \psi \), each of the 3 randomization regions has again 3 subregions, but now along the \( q \)-direction. Application of the escape operator \( R_r \) reduces the weight in each rectangle \( A_k \) by the factor \( r_k \). In particular, for \( r = (0.2, 0.01, 1) \) the weight in the middle rectangle is strongly decreased. After application of the quantum baker map \( B_n \) the original 3×3-structure is recovered due to stretching in \( q \)-direction and compressing in \( p \)-direction. The norm, however, is reduced by a global factor \( e^{\gamma} \) compared to the norm of \( \psi \).

Naively, one would expect to have no substructure in a randomization region after application of the randomization operator \( U_{\ell_p,\ell_q}^{\text{cue}} \). This is indeed the case, if it is applied to an arbitrary state. In Fig. A1, however, it is applied to an eigenstate of the randomized baker map, Eq. (13), making a substructure in the randomization region before and after the application of \( U_{\ell_p,\ell_q}^{\text{cue}} \) necessary.

Note that cyclic permutations of the three operators in \( B_{n,r,\ell_p,\ell_q} \) would lead to resonance states with phase-space structures corresponding to Fig. A1(b) or (c), while the spectrum would not change. The operators \( U_{\ell_p,\ell_q}^{\text{cue}} \) and \( R_r \) commute for \( \ell_q \geq 1 \), such that their exchange neither changes spectrum nor resonance states.

![Figure A1. Husimi function of a resonance state \( \psi \), as in Fig. 2(a, left panel) and sequential application of the operators for randomization \( U_{\ell_p,\ell_q}^{\text{cue}} \), escape \( R_r \), and baker map \( B_n \).](image-url)
Appendix B. Convergence of classical weights for $\ell_p, \ell_q \to \infty$

In the following we briefly present numerical evidence on the weak convergence of the classical measures $\mu_{\ell_p, \ell_q}^{r, s}(A_\ell, \ell)$ for increasing $\ell_p$ and $\ell_q$. Fixing a set $A \subset \mathbb{T}$ we are interested if the limit $\lim_{\ell_p, \ell_q \to \infty} \mu_{\ell_p, \ell_q}^{s}(A)$ exists. For simplicity $A$ is chosen as one of the rectangles $[\alpha_\ell, \ell]$ at $\ell = 1$ in the following.

In Fig. B1 we illustrate the convergence of these classical weights versus $\gamma$ for increasing level of randomization. We consider the cases (a) $\ell_q = \ell_p + 1$ and (b) $\ell_q = \ell_p$ for increasing $\ell_p$. We observe that these curves converge already for the small $\ell_p$ presented here. Furthermore, the classical measures for the largest $\ell_p$ agree very well between (a) and (b). This suggests, that there is a unique limit measure, $\mu_\gamma^*$, towards which $\mu_{\ell_p, \ell_q}^{r, s}$ converges weakly, regardless how the limit $\ell_p, \ell_q \to \infty$ is taken. A proof of this observation is an open task.

Appendix C. Random matrix model with partial escape

In this section we apply the local random wave model combined with conditional invariance to resonance states of a random matrix model with partial escape. This leads to a straightforward derivation of the analytic expressions for their structure and its dependence on $\gamma$, given in Ref. [50, App. D].

Such models are defined as the composition of a random matrix and an escape operator [69, 78, 79, 80, 50],

$$\mathcal{U}_r := \mathcal{U}^{\text{res}}_{\ell_p=0, \ell_q=0} \mathcal{R}_r,$$

i.e., randomization takes place on the entire phase space, $\ell_p = 0$, $\ell_q = 0$. Here we consider a reflectivity operator $\mathcal{R}_r$ with constant escape from the rectangles $A_k$, see Eq. (8). Resonance states are uniformly fluctuating in subregions given by the rectangles $A_k$, see Fig. C1(a), even though randomization takes place on the whole phase space.

We describe the structure of these resonance states by a classical measure $\mu$ that is constant on each rectangle $A_k$ with weight

$$\mu_k := \mu(A_k).$$
According to the local random wave model, Eq. (32), we maximize the product of relative weights,

$$\Pi(\mu) = \prod_{k=0}^{n-1} \frac{\mu_k}{\sum_{l=0}^{n-1} \mu_l} = \prod_{k=0}^{n-1} \mu_k, \quad (C.3)$$

where normalization $$\sum_{k=0}^{n-1} \mu_k = 1$$ is used (and actually a global random wave model). The constraint of conditional invariance follows from Eq. (18) applied to the entire phase space $$T$$,

$$\sum_{k=0}^{n-1} r_k \mu_k = e^{-\gamma} \sum_{k=0}^{n-1} \mu_k \quad \Leftrightarrow \quad \sum_{k=0}^{n-1} (e^{\gamma r_k} - 1) \mu_k = 0. \quad (C.4)$$

Using the method of Lagrange multipliers we define a cost function maximizing $$\ln \Pi(\mu)$$

$$f(\{\mu_k\}) = \sum_{k=0}^{n-1} \ln \mu_k - \lambda \left( \sum_{k=0}^{n-1} \mu_k - 1 \right) - \nu \left( \sum_{k=0}^{n-1} (e^{\gamma r_k} - 1) \mu_k \right), \quad (C.5)$$

with Lagrange multipliers $$\lambda$$ and $$\nu$$. Setting the partial derivatives ($$\partial/\partial \mu_k$$) $$f(\{\mu_k\})$$ to zero we get a set of $$n$$ equations,

$$\frac{1}{\mu_k} - \lambda - \nu (e^{\gamma r_k} - 1) = 0. \quad (C.6)$$

Multiplying each equation by $$\mu_k$$ and summing over all $$k$$ gives $$\lambda = n$$. The weights $$\mu_k$$ are thus given by,

$$\mu_k = \frac{1}{n + \nu (e^{\gamma r_k} - 1)}. \quad (C.7)$$

This agrees with the expression for the density in Ref. [50, Eq. (D10)], while Eq. (C.4) agrees with Ref. [50, Eq. (D9)] and can be used for determining the remaining Lagrange multiplier $$\nu$$ corresponding to the variable $$\xi^*$$ in Ref. [50] (up to a factor $$n$$).

In Fig. C1 we illustrate the structure of resonance states of a random matrix with escape from $$n = 5$$ rectangles. In (a) the Husimi function (left) and the corresponding classical measure $$\mu$$ (right) are shown for $$\gamma_{\text{typ}}$$. In (b) agreement of quantum and classical weights for all $$\gamma$$ and all 5 rectangles $$A_k$$ is demonstrated.
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