\[
\int f(\alpha) \, d\alpha = 1, \quad f \geq 0, \quad \alpha \text{ is a short notation for the pair } (q, p). \quad \text{The map } \Theta_\Delta \text{ generates the Frobenius-Perron operator acting in the space of classical density distributions}
\]

\[
\mathcal{M} f(\alpha) = \int f(\alpha') \delta(\alpha - \Theta(\alpha')) \, d\alpha'.
\]  

Since the map is not dissipative, and \( \Theta_\Delta(I) \subseteq I \), the operator \( \mathcal{M} \) preserves the probability, \( \int_I \mathcal{M} f(\alpha) \, d\alpha = \int_I f(\alpha) \, d\alpha = 1 \). The density \( f^*(\alpha) = \frac{1}{|\mathcal{M} f(\alpha)|} \) for \( \alpha \in [0, 1] \times [0, 1 - \Delta] \) and 0 elsewhere is invariant under the action of \( \mathcal{M} \) and \( \mathcal{M} f^* = f^* \). Several versions of quantum baker map on the torus are known [11, 12, 13].

We use the first form of the quantum operator proposed by Balazs and Voros [11]

\[
\mathbf{B} = F_N^* \begin{pmatrix} F_{N/2} & 0 \\ 0 & F_{N/2} \end{pmatrix},
\]

where \( F_N \) denotes the N-points discrete Fourier transformation, \( [F_N]_{kl} = \frac{1}{\sqrt{N}} e^{-2\pi i kl/N} \). Since the sloppy map \( \Theta_\Delta \) does not enjoy the symmetry of the original baker map, we will not need the symmetric quantum model introduced by Saraceno [12].

Unitary operator \( \mathbf{B} \) acts on the N dimensional Hilbert space \( \mathcal{H}_N \), where \( N \) is even.

The classical map \( \Theta_\Delta \) is irreversible, so its quantization cannot be achieved by means of unitary operators. The quantum operator \( \mathbf{A}_\Delta \) — corresponding to the classical map \( I_\Delta \) should act on the space of mixed quantum states and may be realized by a super-operator. Any super-operator \( \mathbf{A} \) which defines a completely positive map, may be written in the so-called Kraus form [6]

\[
\mathbf{A}(\rho) = \sum_{i=1}^{K} A_i \rho A_i^\dagger,
\]

where \( \rho \) is a density matrix and \( K \) is finite. If operators \( A_i \) fulfill the condition

\[
\sum_{i=1}^{K} A_i^\dagger A_i = \mathbf{1}_N,
\]

the map \( \mathbf{A}(\rho) \) is trace preserving. The classical map \( \Theta_\Delta \) transforms the bottom half of the square \( I \) into itself and shifts the top one down by \( \Delta \). The two halves of \( I \) are transformed separately. Therefore we split the phase space into bottom and top and introduce two projection operators \( D_b \) and \( D_t \), which written in the eigenbasis of position operator have the form

\[
D_b = F_N^* \begin{pmatrix} 1 & 0 \\ 0 & 0 \end{pmatrix} F_N; \quad D_t = F_N^* \begin{pmatrix} 0 & 0 \\ 0 & 1 \end{pmatrix} F_N.
\]

Notice that the super-operator \( \mathbf{A}_M(\rho) = D_b \rho D_b^\dagger + D_t \rho D_t^\dagger \) corresponds to the up/down measurement process and the Kraus operators \( A_1 = D_b \) and \( A_2 = D_t \) fulfill the condition (5). To construct a quantum shift transformation \( \mathbf{A}_\Delta \) we will use the unitary operator of translation in momentum,

\[
V |k\rangle \equiv |k + 1\rangle, \quad \mathbf{U}_N^N = \mathbf{1}_N.
\]
the the top half of $I$ by $\Delta/2$ is realized by the translation operator (7) acting on the previously measured system,

$$D'_t = V^{N \Delta/2} \cdot D_t.$$ (8)

For simplicity we assure here that exponent is integer, but the same construction works also for any real $\Delta$. Since the position of the bottom part remains unchanged, the entire quantum transformation $A_\Delta$ reads

$$A_\Delta(\rho) = D_\rho D'^\dagger_t + D'_t D^\dagger_\rho.$$ (9)

This super-operator resets to 0 the off-diagonal blocks of the $\rho$ matrix in the $p$-representation. This is related to the fact that to displac a one half of the toms we need to perform a measurement, which implies decoherence. Thus even for $\Delta = 0$ the operator $A_\Delta$ differs from ikness, but the effect of the measurement becomes negligible in the classical limit $N \to \infty$.

Using the above super-operator ($A_\Delta$) we construct the entire quantum sloppy baker map

$$B_\Delta(\rho) = A_\Delta(B_\rho B\dagger) = D_\rho B_\rho B\dagger D'^\dagger_\rho + D'_\rho B_\rho B\dagger D^\dagger_\rho.$$ (10)

Note that the Kraus operators $A_1 = D_\rho B$ and $A_2 = D'_\rho B$ fulfill condition (5) with $K = 2$.

To demonstrate that quantum system defined by (10) corresponds to the classical sloppy baker map we compare classical and quantum structures in the phase space. In order to define quantum quasi-probability distribution, we use a family of states localized at points of the square $N \times N$ lattice in the phase space constructed by means of translation operators [12]. The operator $U$ of translation in position is defined similarly to $V$,

$$U|n\rangle = |n+1\rangle, \quad U^N = 1_N,$$ (11)

where $|n\rangle$ are position eigenstates, satisfying $|n+N\rangle = |n\rangle$. As a reference state we choose arbitrarily the wave packet $|\frac{1}{2}, \frac{1}{2}\rangle$ localized in $(\frac{1}{2}, \frac{1}{2})$

$$\langle q|1/2, 1/2\rangle = (2/N)^{-1/4} e^{-\frac{(q-N/2)^2}{N-\pi n}}$$ (12)

which becomes Gaussian for $N \to \infty$. We translate it to any point $(q, p)$, where $Nq$ and $Np$ are integers ($N$ is even)

$$|q, p\rangle = V^{Nq-N/2} U^{Np-N/2} |1/2, 1/2\rangle.$$ (13)

These states allow one to define a Husimi representation in the phase space of any mixed quantum state $\rho$

$$H_\rho(q, p) = \langle q, p | \rho | q, p\rangle.$$ (14)

We analyze the evolution of an exemplary state $|\alpha\rangle$ localized at $\alpha$ and the classical transformation of the corresponding density distribution. On the left hand side of Fig. 2 we present the classical density and its image after $T = 1, 2, 5$ and $T = 30$ iterations of the Frobenius-Perron operator (2). The right hand side shows the Husimi representations (14) of the initially pure quantum state $|\alpha\rangle\langle \alpha|$ and its images after $T$ actions of the superoperator $B_\Delta$. The quantum quasi-probability distribution $H_\rho$ is localized in the same region of the phase space as the classical density distribution. Since the Husimi distribution may resolve quantum phase-space structures at the length scale of the order of $\hbar N^{-1/2}$, the classical density becomes narrower then its quantum counterpart already after first iteration.

![FIG. 3: Classical orbits of period $T = 1, 2, 3, 4$ (×) (left) are localized close to the peaks of the quantum return probability $R^T$ (right) obtained for the sloppy baker map with $\Delta = 1/4$ and $N = 96$.](image)

After 30 iterations of the classical map the density distribution is close to the invariant measure $\rho^*$. Also the quantum state $B_\Delta^{30}|\alpha\rangle\langle \alpha|$ is close to the invariant density matrix

$$\rho^* = B_\Delta(\rho^*),$$ (15)

the existence of which is guaranteed by the trace preserving condition (5). The state $\rho^*$ is localized on the rectangle $[0, 1] \times [0, 1 - \Delta]$. Moreover, it is almost isotropic on the corresponding $N(1 - \Delta)$ dimensional subspace. To show this we verified that the von Neumann entropy of the invariant state $S(\rho^*) \approx - Tr \rho^* \ln \rho^*$, is close to the maximal entropy for the $N(1 - \Delta)$ dimensional subspace of the Hilbert space $\mathcal{H}_N$.

$$S(\rho^*) \approx S_{\text{max}}^{N(1 - \Delta)/3} := \ln(N(1 - \Delta)).$$ (16)
It is instructive to look at the periodic orbits of the classical transformation $\Theta_\Delta$. They are those of the original (reversible) baker map with momentum scaled by the factor $(1 - \Delta)$,  

$$q_i^T = \frac{n}{2^T - 1}, \quad p_i^T = \frac{r(n)}{2^T - 1}(1 - \Delta) ,$$  

where $T$ denotes the length of the period, $n$ ranges from 0 to $2^T - 1$, and the symbol $r(n)$ denotes the number obtained from $n$ by reversing the order of its bits. The classical periodic orbits may be compared with the structures of the quantum return probability  

$$R^T(q, p) = \langle q, p | \mathbf{B}_\Delta | q, p \rangle / \langle q, p | q, p \rangle.$$

The function $R^T(q, p)$ measures the projection of the quantum state $|q, p\rangle$ back onto itself. As shown in Fig. 3 its maxima are indeed located in the vicinity of classical periodic orbits.

Spectral decomposition of the superoperator $\mathbf{B}_\Delta$ determines the time evolution of the system. For any trace preserving CP map (4) the operator $\mathbf{B}_\Delta$ has an eigenvalue $\lambda_1 = 1$ corresponding to the invariant state $\rho^\ast$. The spectrum is symmetric with respect to the real axis, since $\mathbf{B}_\Delta$ sends Hermitian density matrices into density matrices [14]. Not every superoperator needs to be diagonalizable, i.e. the number of eigenvectors may be smaller than the number of eigenvalues. This is the case for the superoperator of translation $\mathbf{A}_\Delta$, the spectrum of which consists of two eigenvalues 0 and 1. The multiplicity of the former is equal to $3N^2/4$ and the corresponding subspace is degenerate for any $\Delta > 0$.

Fig. 4a shows all $N^2$ eigenvalues of the linear operator $\mathbf{B}_\Delta$ for $N = 64$ and $\Delta = 1/4$. One observes a considerable spectral gap, i.e. the difference $1 - |\lambda_2|$, which determines the rate of the convergence of an initial state toward the invariant density matrix $\rho^\ast$. Moduli of the largest subleading eigenvalues influence the slope $\alpha$ of the initially linear entropy increase demonstrated in Fig. 4b. The data were obtained by averaging over a sample of 10 initially pure states drawn randomly with respect to the unique, unitarily invariant measure on the $2N - 2$ dimensional space of pure states in $\mathcal{H}_N$.

In this work we introduced an irreversible baker map and proposed a method of its quantization. On one hand, the limit $N \to \infty$ of the model may be useful to analyze the quantum-classical correspondence for chaotic, completely positive quantum maps. On the other hand, the extreme quantum regime of low $N$ may be interesting from the point of view of quantum information. Quantum baker map becomes a standard model for theoretical [15] and experimental [16] investigation of NMR quantum computing, and our generalization makes it possible to study the consequences of irreversibility in the system. The effects of the decoherence and the dynamics of entanglement in the two-qubits version of this system ($N = 4$) as well as a generalization of the model will be presented elsewhere. It is a pleasure to thank R. Alicki, A. Becker, M. Ruský, F. Minert, R. Rudnicki and D. Wójcik for fruitful discussions. This work was supported by Polish KBN grant no 2P03B-072-19.