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Recurrences in an isolated quantum many-body system

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The complexity of interacting quantum many-body systems leads to exceedingly long recurrence times of the initial quantum state for all but the smallest systems. For large systems, one cannot probe the full quantum state in all its details. Thus, experimentally, recurrences can only be determined on the level of the accessible observables. Realizing a commensurate spectrum of collective excitations in one-dimensional superfluids, we demonstrate recurrences of coherence and long range order in an interacting quantum many-body system containing thousands of particles. Our findings will enable the study of the dynamics of large quantum systems even after they have reached a transient thermal-like state.

The expectation that a non-equilibrium system evolves toward thermal equilibrium is deeply rooted in our daily experience and is one of the foundations of statistical mechanics (*1*). On the other hand, as formulated by Poincaré and Zermelo, a finite isolated physical system will recur arbitrarily close to its initial state after a sufficiently long but finite time (*2*, *3*). The reconciliation of these seemingly contradicting statements is at the heart of the emergence of irreversible processes from reversible microscopic mechanics (*4*).

The above discussion can be transferred to the quantum domain. In analogy to Boltzmann's conjecture, von Neumann formulated a quasi-ergodic theorem for the evolution of the wave function (*5*) and the equilibration of isolated quantum systems grew into an active field of research (*6*). However, a general recurrence theorem can be proven in quantum mechanics as well (*7*, *8*), and shows explicitly that the wave function returns arbitrary close to its initial state. Experimentally, recurrent behavior has been observed only in small quantum systems. Examples are the revival dynamics in the Jaynes-Cummings model (*9*) or few interacting atoms in isolated sites of an optical lattice (*10*, *11*). These studies accomplished coherent dynamics over long times and thereby revealed properties of the underlying Hamiltonian spectrum. Enabling recurrent behavior in more complex quantum many-body systems opens a window into their long-term coherent evolution by providing a strong measurable signal at times far beyond the initial relaxation. Further, it is of fundamental interest as it provides insight into the emergence of statistical ensembles from coherent unitary evolution.

For larger systems however, it becomes exponentially difficult to observe the eigenstates directly, and the complexity of the spectrum of eigenstates leads to exceedingly long recurrence times, in general prohibiting their observation. Nevertheless, for a large class of systems the essential dynamical features can be described by effective field theories with a much simpler structure (*12*, *13*). This reflects the distinction between relevant and irrelevant operators in the microscopic model and dramatically reduces the complexity of the problem from a large number of constituents to a much smaller number of populated modes. Within this level of description, quantum recurrences manifest through the return of observables dominated by these modes. Designing the whole system such that the collective excitations follow a simple commensurate spectrum makes the observation of recurrences feasible, even for many-body systems containing thousands of interacting particles.

Ultracold gases (*14*) are an ideal starting point to study these fundamental phenomena as they can be well isolated from the environment and excellent tools are available to prepare, manipulate and probe them.

As a model system we study coherence in one-dimensional (1D) superfluids. In first approximation, the corresponding many-body physics can be mapped to an effective low-energy description: a bosonic Luttinger liquid (*15*–*18*) where the collective excitations are free phonons. These phonons are directly related to the phase fluctuations observed in the interference of two 1D superfluids (*19*).

The dephasing of these collective phononic excitations leads to a loss of coherence and long range order, proceeding in a light-cone-like fashion (*20*–*23*). While the fully dephased state is described by a generalized Gibbs ensemble (*24*), the long time behavior crucially depends on the spectrum of the collective phononic modes. In an harmonic longitudinal confinement with trap frequency ω_z, these phonon frequencies are non-commensurate $\omega_j = \omega_z\sqrt{j(j+1)/2}$ (*25*), with *j* being the mode index. If the atoms are confined to a box shaped trap, the phonon frequencies become commensurate with $\omega_j \propto j$, and recurrences should be observable at short times (*26*).

We implement a box-like confinement in an atom chip double well setup (*27*) by adding hard walls to the very weak longitudinal harmonic confinement with the help of a blue detuned optical dipole potential (*28*). To observe recurrent dynamics we prepare a thermal equilibrium state in a coupled double well potential (*29*). Typical samples have a linear density of about 70 atoms per μm and, depending on the box size, between 2300 and 4800 atoms in each well, resulting in an interaction energy per atom close to 1 kHz. The tunneling coupling J is tuned to a regime where the phases $\theta_{1,2}$ of the two superfluids lock, creating a system with a strongly correlated relative phase field $\varphi = \theta_1 - \theta_2 \left(\cos(\varphi) \approx 1\right)$. To initiate the non-equilibrium dynamics, the coupling is rapidly ramped to zero leaving the two gases to evolve independently (Fig. 1).

We observe the subsequent dynamics by matter-wave interferometry (*27*) which gives direct access to the spatially resolved relative phase $\varphi(z)$ between the two superfluids. To study the coherence and long range order we evaluate the two-point correlation function (*23*)

$$C\left(\bar{z} = |z - z'|\right) = \cos\left(\varphi(z) - \varphi(z')\right), \quad (1)$$

with the expectation value taken over many experimental realizations. Further, we average over all points z, z' within the central part of the clouds that fulfill $\bar{z} = |z - z'|$.

A typical temporal evolution of the phase correlations in a box trap of length $L = 49\,\mu m$ is shown in Fig. 2A. Before the quench at $t = 0$ the relative phase between the superfluids is locked and correlations are close to unity over the whole sample. Immediately after decoupling, this long range order decays, reflecting the dephasing of the collective excitations. At the first minimum in Fig. 2B the initial state is completely dephased and the system is indistinguishable from a thermal state (see inset). For a system with incommensurately spaced modes this dephased state persists for a long time, showing the emergence of statistical properties from the unitary quantum evolution (*22*). In our system however, mode frequencies are designed to be commensurate and two partial recurrences of phase coherence are clearly visible in the subsequent evolution.

To understand the recurrence time we need to look at the dispersion relation of excitations. For a perfect hard-wall box confinement, $\omega_j = \frac{c\pi}{L} j$, where c is the speed of sound and *j* is the mode index (*28*). These modes are not only commensurate, but also equally spaced, facilitating a recurrence at the earliest possible time $t = \frac{2\pi}{\Delta\omega} = \frac{2L}{c}$ with $\Delta\omega$ being the energy spacing between the modes. At this time the lowest lying mode has finished a full rotation whereas the higher energy modes have all performed an integer number of turns bringing all excitations back to their initial configuration. Half way to this full recurrence the system rephases to the mirrored initial state. As we initially start from a nearly flat relative phase profile and our observable C (Eq. 1) is insensitive to the transformation $\varphi(z) \to \varphi(-z)$ this point is equivalent to the full recurrence. Therefore, the expected recurrence time for the correlations is given by $t_{rec} = \frac{L}{c}$ (blue and red bars in Fig. 2B), which agrees well with the observed peaks in coherence.

Although two-point correlations return close to the initial state, the relative phase field does not recur. The time evolution of the coherence factor $\cos(\varphi)$ shows no recurrence (Fig. 2C). It relaxes during the initial dephasing dynamics and stays close to zero from then on. This is equivalent to observing that the ensemble averaged interference picture shows no revival of high contrast fringes (Fig. 2C, inset). At the time of the recurrences, the interference fringes return close to their straight initial state, displaying long range coherence throughout the system; however, for each distinct realization they are shifted by a random global phase. The reason behind this global phase is a small random atom number imbalance between the two wells. This imbalance originates from the thermal fluctuations of the initial state, imperfections in the experiment and quantum noise relevant for lower temperatures. It leads to an inevitable population of the $k = 0$ mode (*30*) resulting in a global phase accumulation that is different for each realization and therefore to vanishing interference contrast in the ensemble average. In contrast, the phase correlations (Eq. 1) are insensitive to a global offset of the phase $\varphi(z)$, and the recurrences of excitations on top of the field can be observed. This illustrates that recurrent behavior can be hidden below a global phase diffusion, which necessitates the measurement of, at least, two-point correlations to reveal the underlying coherent dynamics.

To confirm the scaling of the recurrence time with the size of the system, we vary the length of the box potential by changing the position of the dipole trap walls. As the system size is increased the recurrence is shifted to later times (Fig. 3A). For this comparison the time axis was rescaled by the theoretical prediction for the speed of sound to make measurements with slightly different atomic densities comparable. Extracting the exact times of the first and second recurrence by fitting the peaks in the data (*28*) reveals the linear scaling with L (Fig. 3B).

Although the first two recurrences are clearly visible, for all these measurements the height of the recurrences is rapidly damped (see Fig. 2, A and B) and observing a third recurrence becomes infeasible in most cases. To probe the decay of the recurrences in more detail, we studied the evolution for different initial temperatures in the $L = 49\,\mu m$ trap. Increasing the initial temperature, the height of the observed recurrence with respect to the correlations of the initial state decreases rapidly (Fig. 4). The temperatures for this analysis are extracted from the full distribution functions of interference contrasts (*31*) for the completely dephased state in between the recurrences (*28*).

To understand the origin of this damping with temperature we considered different theoretical descriptions of our system. We first investigated the low-energy effective description by solving the Luttinger liquid Hamiltonian (dashed lines in Fig. 4) . This model describes the free propagation of phononic excitations on top of a stationary background density. For a homogeneous background it would give perfect recurrences of phase coherence. For the comparison to the experimental data we consider an inhomogeneous background density that reflects our box-like potential. In addition, we took into account the typical spread in total particle number and particle number imbalance between the wells, measured independently for the respective samples (*28*). The shot-to-shot fluctuations of the speed of sound induced by this spread together with the inhomogeneous background density constitute the only source of damping within the model. Although the recurrence times are well described as seen in Fig. 3B, the damping in this description is too weak to explain the experimental observations even for the lowest temperatures.

As a second model (solid lines in Fig. 4) we numerically simulated the dynamics using the Gross-Pitaevskii equation for finite temperature initial states (*28*). This description goes beyond the free phononic excitations of the Luttinger liquid and takes interactions between these quasi-particles into account. It agrees well with our experimental findings (Fig. 4). This shows that physics beyond the low-energy description becomes relevant and indicates that phonon-phonon interactions are responsible for the observed damping. As the temperature is increased these processes get more important as higher energy modes are populated.

This illustrates that the observation of recurrences enables insight into the non-equilibrium dynamics of quantum many-body systems (*32*). Their shape and magnitude are sensitive probes into coherent dynamics far beyond the time scales of the initial relaxation. This makes them a versatile tool for the verification of coherence in quantum simulators and to test the regimes of validity for approximate models. Moreover, the long coherent evolution probed through recurrences magnifies the effects of small perturbations whose influence on the short term dynamics are negligible. In our model case, the decay of recurrences clearly shows that physics beyond the Luttinger liquid sets in rather fast. A combined study looking at recurrences and mode occupations will shed light onto the fundamental quantum processes in the relaxation of 1D systems.

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SUPPLEMENTARY MATERIALS

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Materials and Methods
Figs. S1 to S7
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Data File S1

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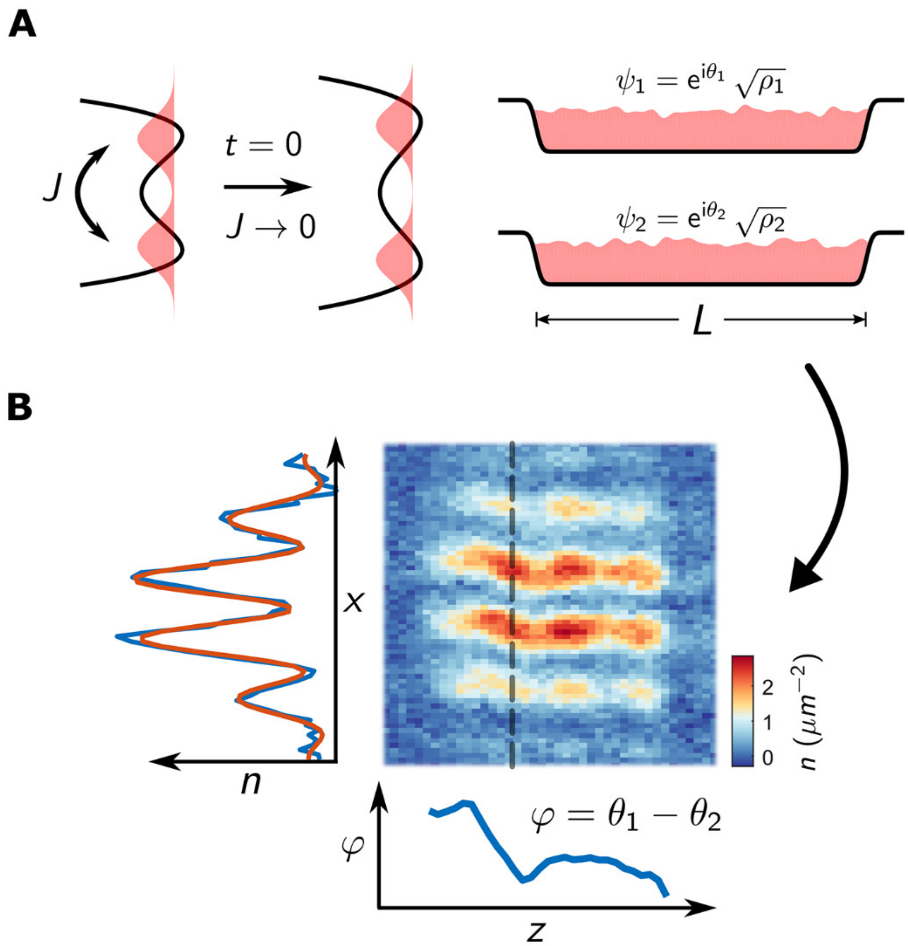

Fig. 1. Schematics of the experiment and the measurement process. (**A**) Two coupled 1D superfluids $\psi_{1(2)}$ in a double well potential are taken out of equilibrium by a sudden quench of the tunneling coupling J to zero. Along the longitudinal direction z the system is confined to a box shaped potential of variable length L. (**B**) The evolution of the decoupled system is measured through matter-wave interference in time-of-flight. A typical interference picture showing the atomic density n is given. A sinusoidal function is fitted to the local interference fringes for each position z. On the left we show an example of such a fit (red) for one slice (blue) indicated by the dashed line in the image; the x coordinate being the spatial direction of the double well separation. This gives access to the spatially resolved relative phase $\varphi(z) = \theta_1(z) - \theta_2(z)$ between the two superfluids (bottom).

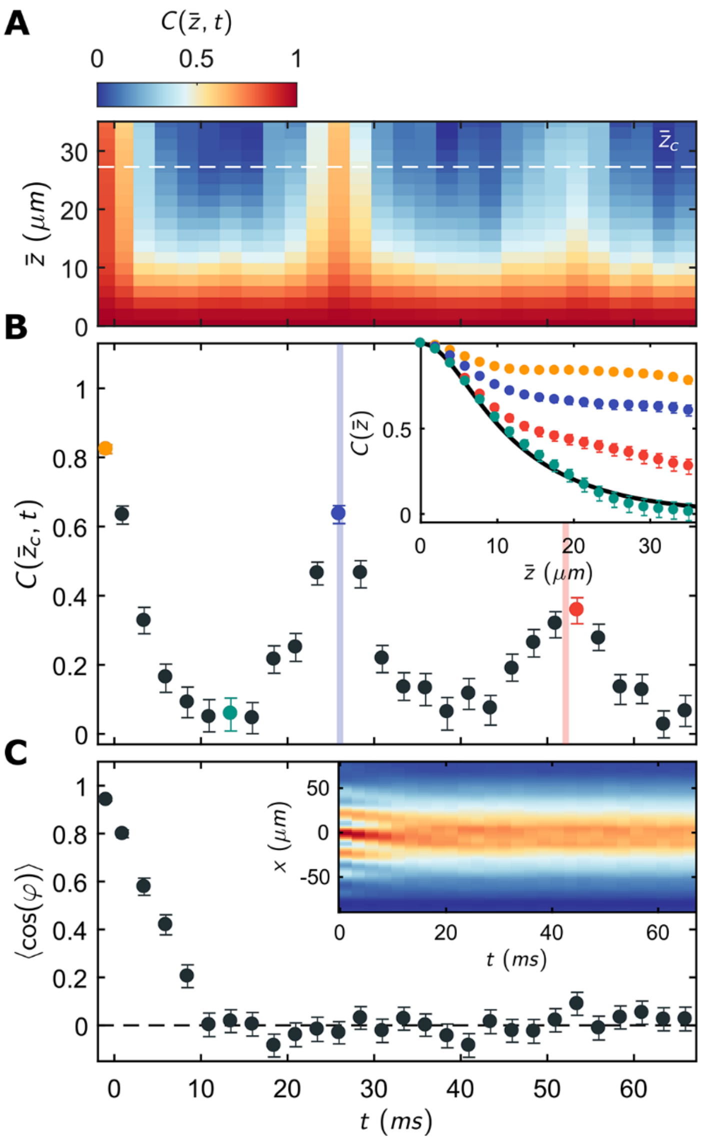

Fig. 2. Dynamics after decoupling. (**A**) Temporal evolution of the phase correlation function $C(\bar{z},t)$ after decoupling in a $49\mu m$ long box potential. (**B**) A cut at $\bar{z}_c = 27.3\mu m$ with the error bars giving the 68% confidence interval obtained from a bootstrap. The first and second recurrence are clearly visible and occur at the expected times $t_{rec} = L/c$ and $2t_{rec}$, as indicated by the vertical blue and red bars respectively. The distance $\bar{z}_c$ is chosen such that it is considerably smaller than the size of the system but long enough that the correlations after the initial dephasing are low; this is to facilitate a high recurrence visibility. The inset shows the phase correlation function at the first (blue) and second (red) recurrence and compares them to the correlations in between the recurrences (green) and the correlations in the initial state (orange). The corresponding points in (B) are colored accordingly. A thermal fit to the correlations in between the recurrences is given by the solid black line. (**C**) Evolution of the coherence factor $\cos(\varphi)$ in the center of the trap at $z = 0$, showing no sign of a recurrence (see main text for details). Owing to the near homogeneity of the system the behavior is similar over the whole sample. Therefore, in contrast to the phase correlation function the averaged interference picture shows no recurrent behavior as shown in the inset (integrated along the *z*-axis).

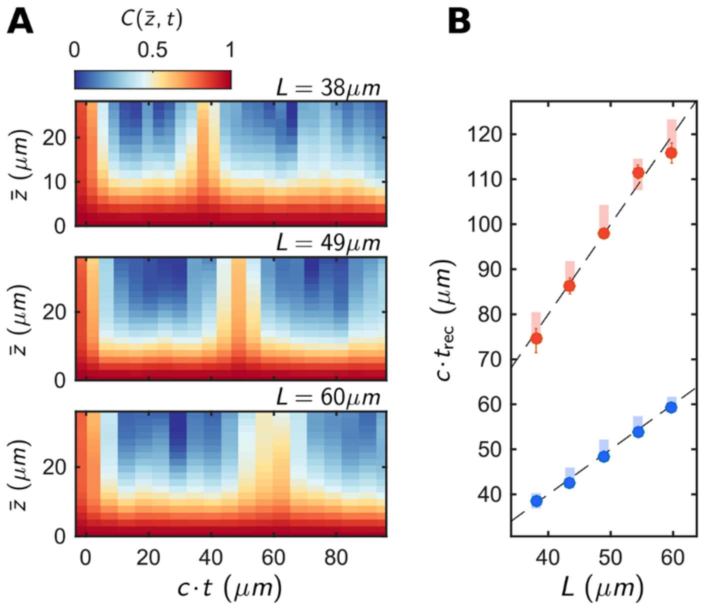

Fig. 3. Comparison of the recurrences for different box lengths. (**A**) Phase correlations for three different box lengths $L = 38, 49,$ and $60\,\mu m$ (top to bottom). The time axis is rescaled with the theoretical prediction for the speed of sound *c*. (**B**) Recurrence time over the box length extracted from the phase correlations at $\bar{z}_c = 27.3\,\mu m$ (*28*). For each box length the time of the first (blue) and second (red) recurrence is plotted and compared to the ideal linear scaling (dashed lines). The error bars give the 68% confidence interval obtained from a bootstrap whereas the shaded bars indicate the predictions of the Luttinger liquid model for our particular trap (*28*). The vertical extension of the bars corresponds to the uncertainty of the decoupling time whereas the horizontal extension is chosen arbitrarily.

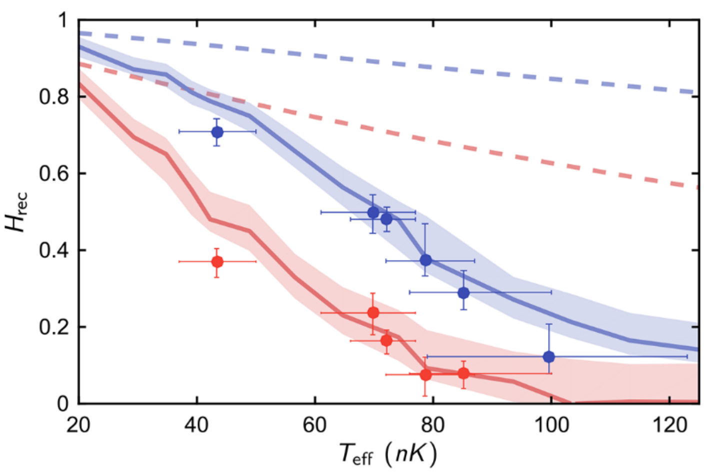

Fig. 4. Temperature dependence of the recurrence height. The data points represent measurements of the height H_{rec} of the first (blue) and second (red) recurrence in a box of length $L = 49\,\mu m$ for different effective temperatures T_{eff} of the relative degrees of freedom. The height is extracted from fitting the peaks of the correlation function at $\bar{z}_c = 27.3\,\mu m$ and the error bars give the 68% confidence interval obtained from a bootstrap. A fit of the full distribution function of measured interference contrasts at $t = t_{rec}/2$ gives the effective temperature. The solid lines are results of GPE simulations analyzed in the same way as the experiment. The shaded area indicates the uncertainty caused by the limited experimental statistics (1σ deviation). The dashed lines are the predictions of the Luttinger liquid model for the experimental parameters (*28*). The experimental shot-to-shot fluctuations of the atom number as well as the inhomogeneous confinement are incorporated in both theoretical calculations and constitute the only source of damping in the Luttinger liquid model.

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