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Dynamical Quasicondensation of Hard-Core Bosons at Finite Momenta

L. Vidmar,[1] J. P. Ronzheimer,[2,3] M. Schreiber,[2,3] S. Braun,[2,3] S. S. Hodgman,[2,3]
S. Langer,[4] F. Heidrich-Meisner,[1] I. Bloch,[2,3] and U. Schneider[2,3,5]

[1]*Department of Physics and Arnold Sommerfeld Center for Theoretical Physics, Ludwig-Maximilians-Universität München, D-80333 München, Germany*
[2]*Ludwig-Maximilians-Universität München*
[3]*Max-Planck-Institut für Quantenoptik, Garching*
[4]*Department of Physics and Astronomy, University of Pittsburgh, Pittsburgh, Pennsylvania 15213, USA*
[5]*Cavendish Laboratory, University of Cambridge, J. J. Thomson Avenue, Cambridge CB3 0HE, United Kingdom*

Long-range order in quantum many-body systems is usually associated with equilibrium situations. Here, we experimentally investigate the quasicondensation of strongly-interacting bosons at finite momenta in a far-from-equilibrium case. We prepare an inhomogeneous initial state consisting of one-dimensional Mott insulators in the center of otherwise empty one-dimensional chains in an optical lattice with a lattice constant d. After suddenly quenching the trapping potential to zero, we observe the onset of coherence in spontaneously forming quasicondensates in the lattice. Remarkably, the emerging phase order differs from the ground-state order and is characterized by peaks at finite momenta $\pm(\pi/2)(\hbar/d)$ in the momentum distribution function.

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The nonequilibrium dynamics of quantum many-body systems constitutes one of the most challenging and intriguing topics in modern physics. Generically, interacting many-body systems are expected to relax towards equilibrium and eventually thermalize [1, 2]. This standard picture, however, does not always apply.

In open or driven systems, one fascinating counterexample is the emergence of novel steady states with far-from-equilibrium long-range order, i.e., order that is absent in the equilibrium phase diagram. This includes lasers [3], where strong incoherent pumping gives rise to a coherent emission, and driven ultracold atom systems [4]. The emergence of order far from equilibrium is also studied in condensed matter systems [5] and optomechanical systems [6].

In recent years, closed quantum systems without any coupling to an environment have come into the focus of experimental and theoretical research. Experimental examples range from ultracold atoms [7–15] to quark-gluon plasmas in heavy ion collisions [16]. In closed, non-driven systems, two famous examples for the absence of thermalization [8, 9, 14] are many-body-localized [17–19] and integrable systems [20]. These peculiar systems allow for nonergodic dynamics and novel quantum phenomena.

Spontaneously emerging order is in general associated with equilibrium states at low temperatures. The canonical example is the emergence of (quasi-) long-range phase coherence when cooling an ideal Bose gas into a Bose-Einstein (quasi-) condensate [21, 22]. In this case, thermodynamics ensures that, for positive temperatures [23], the single-particle ground state becomes macroscopically occupied and thereby dictates the emerging order. Even in studies of the nonequilibrium dynamics at quantum phase transitions [24], the emergence of coherence is typically associated with gently crossing the transition from an unordered into an ordered state, and the strongest correlations and largest coherence lengths appear in the adiabatic limit [25, 26].

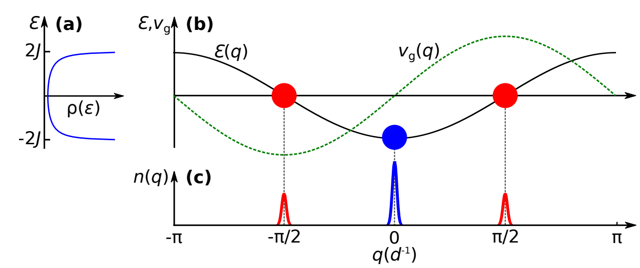

FIG. 1. *Quasicondensation of bosons.* (a) Density of states $\rho(\epsilon)$ of a homogeneous 1D lattice. (b) Dispersion $\varepsilon(q)$ (solid line) and group velocity $v_g(q)$ (dotted line) versus quasimomentum q. (c) Sketch of quasimomentum distribution $n(q)$: In equilibrium, 1D bosons quasicondense at the minimum of the band at $q = 0$, while in a sudden expansion the quasicondensation of hard-core bosons occurs at $\hbar q = \pm(\pi/2)(\hbar/d)$. This quasimomentum lies in the middle of the spectrum and is consistent with the vanishing energy per particle of this closed many-body system.

Here, in contrast, we study a condensation phenomenon of strongly interacting lattice bosons far from equilibrium. After a sudden quantum quench, we experimentally observe the spontaneous emergence of a long-lived phase order that is markedly different from the equilibrium order (cf. Fig 1). To this end, we prepare a density-one Mott insulator of strongly interacting bosons in the center of a three-dimensional (3D) optical lattice. Next, we transform the system into an array of independent one-dimensional (1D) chains, entering the regime of integrable hard-core bosons (HCBs). By suddenly quenching the confining potential along the chains to zero, we induce a sudden expansion of the cloud in a homogeneous lattice [12, 13, 15, 27] with a lattice constant d and detect the formation of a non-ground-state phase profile as a dynamical emergence of peaks at momenta $\hbar k = \pm(\pi/2)(\hbar/d)$, half way between the middle and the edge of the Brillouin zone, in time-of-flight (TOF) distributions. This finite-momentum quasicondensation was first discussed by Rigol and Muramatsu [28] (see also Refs. [29–33]), but has not been studied experimentally so far.

Ideal case. The idealized setup to study finite-momentum quasicondensates is shown in Fig. 2. We consider the Hamiltonian $H = -J\sum_j(\hat{a}^\dagger_{j+1}\hat{a}_j + \text{h.c.})$, where $\hat{a}^\dagger_j$ creates a HCB on site j of a 1D lattice. The infinitely large on-site repulsion is accounted for by the hard-core constraint $(\hat{a}^\dagger_j)^2 = 0$. The initial state is a product state $|\psi_0\rangle = \prod_{j\in L_0}\hat{a}^\dagger_j|\emptyset\rangle$, completely filling the central region of an otherwise empty and infinitely large 1D lattice. This initial state consists of $N = L_0$ localized particles with a flat quasimomentum distribution and contains no off-diagonal correlations, i.e., $\langle\hat{a}^\dagger_j\hat{a}_{j+r}\rangle = 0$ for $r \neq 0$. Surprisingly, the quasimomentum distribution $n(q) = \frac{1}{L}\sum_{j,l} e^{-iq(j-l)d}\langle\hat{a}^\dagger_j\hat{a}_l\rangle$ develops singularities at *finite* quasimomenta $\hbar q = \pm(\pi/2)(\hbar/d)$ during the expansion ($t_{\mathrm{E}} > 0$). As shown in Fig. 2(c), these singularities correspond to the emergence of *power-law* correlations

$$\langle\hat{a}^\dagger_j\hat{a}_{j+r}\rangle = \mathcal{A}(r)e^{i\Phi(r)};\ \ \mathcal{A}(r) \sim r^{-\frac{1}{2}};\ \ \Phi(r) = \pm\frac{\pi}{2}r \quad (1)$$

in each half of the expanding cloud, shown in Fig. 2(d). These power-law correlations justify the name quasicondensate [28]. Curiously, the exponent 1/2 equals the *ground-state* exponent [28, 30], even though the system is far away from equilibrium, with the energy per particle being much higher than in the ground state. In contrast to the ground state, the correlations show a running phase pattern $\Phi(r)$ with a phase difference of $\pm\pi/2$ between neighboring lattice sites, giving rise to peaks at finite quasimomenta. Coherence and quasicondensation emerge independently in the left- and right-moving halves of the cloud, corresponding to two macroscopically occupied degenerate eigenstates of the one-particle density matrix $\langle\hat{a}^\dagger_j\hat{a}_l\rangle$ that have spatial support in the left- or right-moving cloud, respectively [28]. This quasicondensation at finite quasimomenta can equivalently be seen as quasicondensation at $q = 0$ in the respective co-moving frames [32].

In one dimension, HCBs can be exactly mapped to noninteracting spinless fermions via the Jordan-Wigner transformation [34]. By virtue of this mapping, the density $n_j = \langle\hat{a}^\dagger_j\hat{a}_j\rangle$ of HCBs is identical to that of free fermions for all times, whereas the same is not true for the quasimomentum distribution [35, 36]. While the occupations of fermionic quasimomenta are constants of motion,

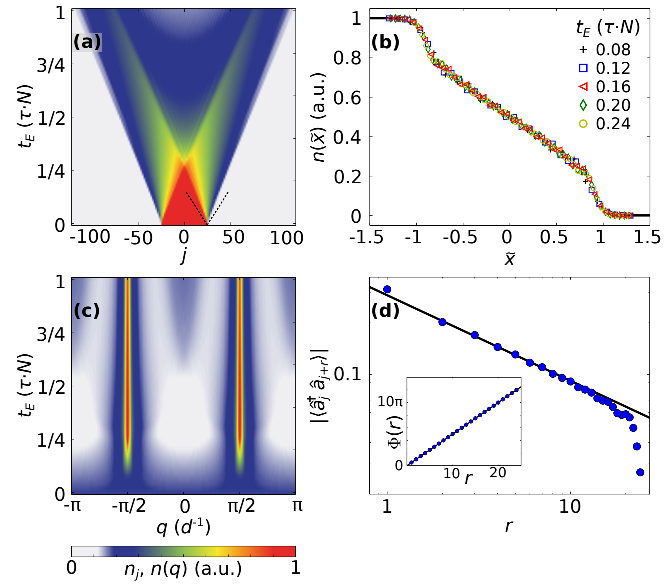

FIG. 2. *Finite-momentum quasicondensation under ideal conditions* [28]. We consider $N = 50$ initially localized HCBs. (a) Density n_j as a function of time. (b) Density as a function of the rescaled coordinate $\tilde{x} = (j-25.5)/(2t_{\mathrm{E}}/\tau)$ in the region bounded by the dashed lines in (a). For $\tilde{x} \leq 1$, the data collapse to the scaling solution [39] $n(\tilde{x}) = \arccos(\tilde{x})/\pi$. (c) Quasimomentum distribution $n(q)$ as a function of time. (d) One-particle correlations at $t_{\mathrm{E}} = 0.24N\tau$. Main panel: $|\langle\hat{a}^\dagger_j\hat{a}_{j+r}\rangle|$ at $j = 26$ (circles) and $\mathcal{A}(r) = \alpha/\sqrt{r}$ (line) with $\alpha = 0.29$ [40]. Inset: Phase pattern $\Phi(r)$.

the non-local phase factors in the Jordan-Wigner transformation give rise to the intricate momentum dynamics studied here. In the ideal case, the dynamical quasicondensates form over a time scale $t^*_{\mathrm{E}} \sim 0.3N\tau$ [28, 33], where N is the number of particles in the initial state and $\tau = \hbar/J$ denotes the tunneling time. For very long times, $n(q)$ slowly decays back to its original flat form as a consequence of the dynamical fermionization mechanism [33, 37, 38].

The dynamical quasicondensation at finite momenta is an example of a more general emergence of coherence in a sudden expansion. For instance, interacting fermions described by the Fermi-Hubbard model exhibit ground-state correlations in the transient dynamics as well [32]. Furthermore, there is a close connection to quantum magnetism, as can be seen by mapping HCBs to a spin-1/2 XX chain: The transient dynamics in each half of the expanding cloud of HCBs is equivalent to the melting of a domain-wall state [39, 41–49] of the form $|\psi_0\rangle = |\uparrow\ldots\uparrow\uparrow\downarrow\downarrow\ldots\downarrow\rangle$. For this problem, a scaling solution exists [39], which also applies to the sudden expansion at times $t < t^*_{\mathrm{E}}$. As a consequence, the densities n_j measured at different times collapse onto a single curve, as shown in Fig. 2(b). Furthermore, for spin-1/2 XX chains the emergence of power-law decaying trans-

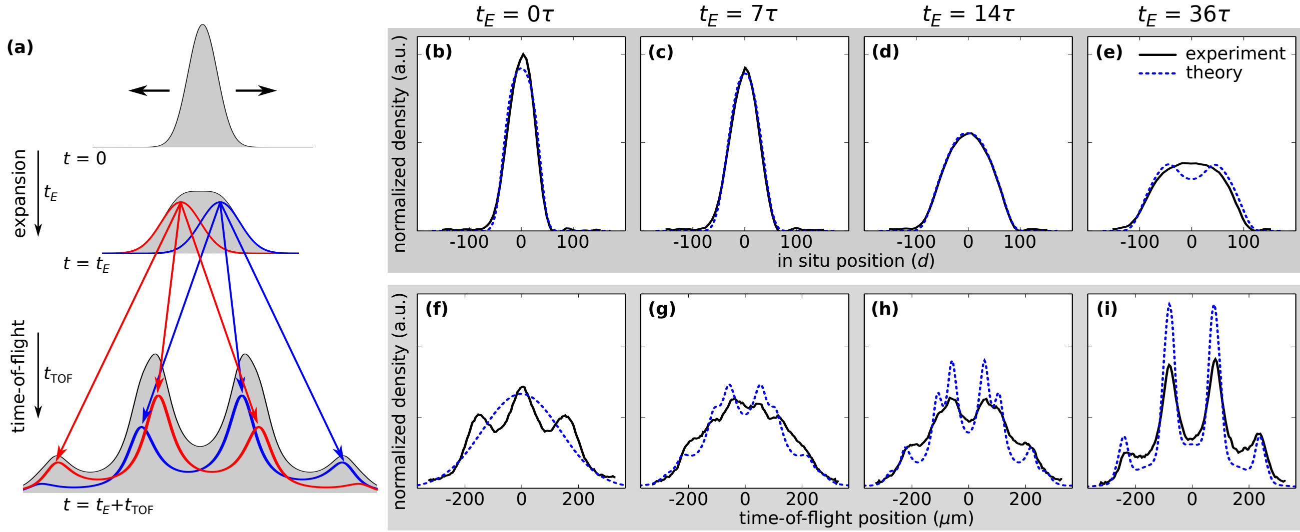

FIG. 3. *Dynamics during sudden expansion and time-of-flight sequence.* (a) Sketch of the experimental sequence: We start with a trapped gas that is released from the initial confinement at $t = 0$ (top) and then expands for a time t_{E} *in the optical lattice* (middle). At t_{E}, all potentials are removed and the distribution is measured after a finite time-of-flight time t_{TOF} (bottom). (b)-(e): *In-situ* density distributions during the sudden expansion, integrated along the y- and z-axes. The lattice constant corresponds to $d = \lambda/2 \approx 368\,\mathrm{nm}$. (f)-(i): TOF density distributions taken at $t_{\mathrm{TOF}} = 6\,\mathrm{ms}$ for the same t_{E} as in (b)-(e). All measured density distributions are averaged over 9 to 11 experimental realizations.

verse spin correlations modulated with a phase of $\pi/2$ has been derived analytically [44]. An interesting perspective onto the emergence of coherence results from noticing that both our expanding bosons and the melting domain-walls realize current-carrying nonequilibrium steady states (see [50] for a discussion).

Experimental set-up and results. The experimental set-up is identical to that employed in our previous experiment on *in-situ* density dynamics [13]. We load a Bose-Einstein condensate of approximately 10^5 ^{39}K atoms from a crossed optical dipole trap into a blue-detuned 3D optical lattice with a lattice depth of $V_0 \approx 20E_r$, where $E_r = h^2/(2m\lambda^2)$ denotes the recoil energy with atomic mass m and lattice laser wavelength $\lambda \approx 737\,\mathrm{nm}$. During the loading of the lattice, we use a magnetic Feshbach resonance to induce strong repulsive interactions between the atoms, suppressing the formation of doubly occupied sites [13]. This results in a large density-one Mott insulator in the center of the cloud. We hold the atoms in the deep lattice for a 20 ms dephasing period, during which residual correlations between lattice sites are mostly lost such that the atoms essentially become localized to individual lattice sites [51]. The expansion is initiated at $t_{\mathrm{E}} = 0$ by simultaneously lowering the lattice depth along the expansion axis in $150\,\mu\mathrm{s}$ to $V_0^x \approx 8\,E_r$ (setting $J \approx h \times 300\,\mathrm{Hz}$, $\tau \approx 0.5\,\mathrm{ms}$), and reducing the strength of the optical dipole trap to exactly compensate the anti-confinement of the blue-detuned lattice beams. This creates a flat potential along the expansion direction. Figures 3(b)-3(e) show the ballistic expansion of the *in-situ* density [13], monitored using absorption imaging. During the deep lattice period, the magnetic field is changed to tune the on-site interaction strength during the expansion to $U = 20J$. We have numerically verified that the essential features of dynamical quasi-condensation are still present for this value of U/J, with a shift of the peak position by $\approx 10\%$ towards smaller values [31, 50].

In order to measure the momentum distribution as a function of expansion time t_{E}, we employ an adapted TOF imaging technique. Directly before shutting off all lattice and trapping potentials, we rapidly increase the lattice intensity along the expansion axis for $5\,\mu\mathrm{s}$ to a depth of $33E_r$. This time is too short to affect correlations between different sites and the momentum distribution. Nonetheless, it results in a narrowing of the Wannier functions, which leads to a broadening of the Wannier envelope in the TOF density distribution and thereby facilitates the observation of higher-order peaks [50].

Figures 3(f)-3(i) contain the main result of our experiment, namely the TOF density distributions, which correspond approximately to the momentum distribution, taken at different expansion times t_{E}. In Fig. 3(f) the TOF sequence was initiated at $t_{\mathrm{E}} = 0$, i.e., directly after initiating the expansion. We observe a central peak at $k = 0$ and two higher order peaks at $\hbar k = \pm(2\pi)(\hbar/d)$, indicating a weak residual $k = 0$ coherence probably resulting from an imperfect state preparation. During the expansion, however, the momentum distribution changes fundamentally and the remnants of the initial coherence quickly vanish. Instead, new peaks at finite momenta are formed. These peaks directly signal the spontaneous formation of a different phase order. This is best seen in Fig. 3(i) at $t_{\mathrm{E}} = 36\tau$, where the finite-momenta peaks are clearly established. The observed peak positions correspond to the expected momenta close to $\hbar k = \pm(\pi/2)(\hbar/d)$ [50]. In addition, Figs. 3(g) and 3(h), taken at $t_{\mathrm{E}} = 7\tau$ and 14τ, respectively, hint at a fine structure of the emerging peaks. This structure is a consequence of the finite TOF time $t_{\mathrm{TOF}} = 6$ms, which results in the TOF distributions being a convolution of real-space and momentum-space densities. We sketch this situation in Fig. 3(a). As discussed before, the peaks in $n(k)$ close to $\hbar k = -(\pi/2)(\hbar/d)$ and $+(\pi/2)(\hbar/d)$ originate from the left- and right-moving portions of the cloud, respectively. Due to the finite t_{TOF}, the higher-order peak of the left-moving cloud with momentum $(-\pi/2 + 2\pi)(\hbar/d)$ and the main peak of the right-moving cloud with momentum $(\pi/2)(\hbar/d)$ (and vice versa) may overlap in the TOF data. A perfect overlap gives rise to single sharp peaks such as the ones present in the data shown in Fig. 3(i), while shorter expansion times, as shown in Figs. 3(g)-3(h), result in a partial overlap and additional structure (see [50] for details).

Comparison with exact time evolution. We numerically model the dynamics of 1D HCBs for realistic conditions: *(i)* The experimental set-up consists of many isolated 1D chains, which are not equivalent due to the 3D harmonic confinement. Experimentally we can only measure an ensemble average over all tubes. *(ii)* Both the finite temperature of the original 3D Bose-Einstein condensate as well as nonadiabaticities during the lattice loading result in a finite entropy, and thereby holes, in the initial state. We therefore average the results over different initial product states drawn from a thermal ensemble of a harmonically trapped 3D gas of HCBs in the atomic limit [50]. Chemical potential and temperature were calibrated to reproduce the experimental atom number and an average entropy per particle of $1.2\,k_B$ [52], thereby leaving no free parameters for the simulations. To test the consistency of the approach, we compare the average density n_j during the expansion with the *in-situ* images in Figs. 3(b)-3(e) and find a good agreement. In addition, the time evolution of the half width at half maximum of the density distribution, shown in Fig. S4 of [50], is consistent with the ballistic dynamics as previously measured in the same experimental set-up [13].

Since the momentum distribution is experimentally measured at a finite t_{TOF}, we explicitly calculate the TOF density distributions without employing the far-field approximation [50, 53] and compare the results to the experimental data in Figs. 3(f)-3(i). Remarkably, the positions and the structure of the peaks agree very well between experiment and theory, thereby supporting our two main results: The central peaks indeed correspond to a large occupation of quasimomenta close to $\hbar q = \pm(\pi/2)(\hbar/d)$, i.e., to a bunching of particles around the fastest group velocities in the middle of the single-particle spectrum. In addition, the fine structure visible for intermediate expansion times [cf. Figs. 3(g)-3(h)], which becomes more apparent in the experiment when comparing different t_{TOF} (cf. Fig. S2 in [50]), directly confirms the independent emergence of coherence in the left- and right-moving portions of the cloud. Compared to the ideal case, the presence of holes in the initial state causes a reduced visibility of the TOF density distributions [50]. Moreover, the finite initial entropy gives rise to a crossover of one-particle correlations from a power-law decay at short distances to a more rapid decay at long distances [50], similar to the effect of a finite temperature [54] in equilibrium.

We attribute the discrepancies between experimental and numerical results at short times, see Fig. 3(f), to the weak residual $k = 0$ coherence in the initial state. Additional discrepancies may arise because of a small admixture ($\lesssim 5\%$) of doublons in the initial state [13, 15] as well as small residual potentials, yet we conclude that these play a minor role [50]. Compared to the previously studied time dependence of density distributions and expansion velocities [13], the momentum distribution is more sensitive to such imperfections [50]. Performing a similar experiment with a single 1D system would allow the predicted scaling of t^*_{E} and the maximum peak height with atom number [28] to be experimentally tested.

Conclusions and outlook. We have reported experimental evidence for a far-from-equilibrium quasicondensation at finite momenta of expanding 1D HCBs in an optical lattice. The expanding particles bunch at quasimomenta close to $\pm(\pi/2)(\hbar/d)$ and the analysis of TOF distributions demonstrates the existence of two independent sources of coherence.

Whether such dynamical condensation persists in higher dimensional systems constitutes an open problem, given that the existing theoretical results are based on exact diagonalization of small systems [55] or time-dependent Gutzwiller simulations [56]. Both future experiments or advanced numerical methods (see, e.g., [57, 58]) could help clarify this question. More generally, our results raise the question of whether this type of spontaneously emerging coherence is limited to integrable systems and whether genuinely far-from-equilibrium order can also occur in generic closed many-body systems.

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