Transport of a Bose Gas in 1D Disordered Lattices at the Fluid-Insulator Transition

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We investigate the momentum-dependent transport of 1D quasicondensates in quasiperiodic optical lattices. We observe a sharp crossover from a weakly dissipative regime to a strongly unstable one at a disorder-dependent critical momentum. In the limit of nondisordered lattices the observations suggest a contribution of quantum phase slips to the dissipation. We identify a set of critical disorder and interaction strengths for which such critical momentum vanishes, separating a fluid regime from an insulating one. We relate our observation to the predicted zero-temperature superfluid-Bose glass transition.

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The transport in low-dimensional superfluids and superconductors is strongly affected by the presence of disorder, isolated defects, or even a periodic lattice. The superflow tends to become unstable for increasing velocities and decay via phase-slip nucleation, a mechanism that is particularly strong in 1D because of the large quantum and thermal fluctuations [1]. This mechanism is relevant for different systems, such as superfluid He in porous materials [2–4], superconducting nanowires [5–9], or ultracold atoms [10–15].

In particular, disorder has been identified as the main source of dissipation in superconductors and superfluid He. By employing strongly disordered nanowires close to the superconductor-insulator transition, some degree of control of the quantum phase-slip nucleation rate was demonstrated [9], and models of dissipation due to disorder have been developed [7,16]. A good control of the disorder is now available in ultracold atom systems. Experiments are starting to address the open questions about the superfluid-Bose glass transition [17–21] and have studied the effect of a controlled disorder on the transport of 3D Bose-Einstein condensates [22,23]. A study of the momentum- and disorder-dependent transport in the strongly fluctuating 1D environment is, however, still missing.

In this work we experimentally address this problem with 1D ultracold atomic bosons in quasiperiodic lattices, which allow us to simulate a controllable disorder and tunable interaction. We start our investigation from the limit of nondisordered lattices, where suitable theoretical models for phase slips are available. By exciting a motion with variable momentum p in systems with relatively large density, we observe a rather sharp transition from a weakly dissipative regime at low p to a strongly unstable one at large p, in contrast to what was observed in low-density systems [13,14]. Measurements of the momentum- and interaction-dependent dissipation suggest a relevant role of quantum phase slips. We then find that a weak disorder tends to increase the dissipation and reduce the critical momentum p_c for the instability. We observe that for a given interaction strength there is a critical disorder strength above which p_c vanishes, which indicates the crossover into an insulating regime. From a set of different measurements we find a crossover line in the interaction-disorder plane that is compatible with theoretical estimates for the superfluid-Bose glass transition at $T = 0$ [24–27].

In the experiment we employ an ensemble of 1D quasicondensates of ^{39}K atoms with tunable repulsive interaction [28], moving in a harmonic trap in the presence of a quasiperiodic lattice [29]. The system is realized by splitting a 3D Bose-Einstein condensate into a few hundreds of 1D quasicondensates with a deep 2D lattice in the horizontal plane. Each subsystem contains on average 50 atoms and has longitudinal (transverse) trapping frequency $\omega_z = 2\pi \times 150$ Hz ($\omega_\perp = 2\pi \times 50$ kHz). Along the longitudinal direction, a quasiperiodic lattice is created by superimposing two laser standing waves with incommensurate wavelengths ($\lambda_1 = 1064$ nm, $\lambda_2 = 859$ nm). The first lattice is stronger and sets the tunneling energy $J = h \times 150$ Hz, while the weaker secondary lattice sets the amplitude Δ of the diagonal disorder [30]. For $\Delta > 2J$ all single-particle eigenstates of the first lattice band are exponentially localized as in a truly disordered system [31,32]. The Bose-Hubbard interaction energy U can be varied in the range $(0.3–10)J$ by adjusting the atomic scattering length at a Feshbach resonance [33]. The mean atom number per site n, which scales approximately as $U^{-1/3}$, varies in the range of 2 to 4. From the width of the momentum distribution of the weakly interacting quasicondensates, we estimate an upper limit for the equilibrium temperature of $k_B T \simeq 6J$ [34].

To study the transport, the trap center along the vertical direction is suddenly displaced by a small amount $z_0 = 3.9(2)$ μm by switching off a magnetic-field gradient. In the absence of any dissipation, the atoms would oscillate with a frequency $\omega^* = \omega_z\sqrt{m/m^*} \simeq 2\pi \times 90$ Hz, where $m^* \simeq 2.8m$ is the atomic effective mass in the lattice. After a variable waiting time, all potentials are suddenly switched off and the momentum distribution $\rho(p)$ is recorded after a free expansion.

We started our investigation with nondisordered lattices, i.e., $\Delta = 0$, where theoretical models are available. A typical observation of the evolution of $\rho(p)$ is shown in Fig. 1 and compared to the solution of the semiclassical equations of motion [34]. At short times, the displacement of the peak momentum p_0 can be approximated with a damped oscillation $p_0(t) = m^*\omega^{*2}z_0/\omega' \sin(\omega' t)e^{-\gamma^* t}$, where $\omega' = \sqrt{\omega^{*2} - \gamma^{*2}}$ and $\gamma^* = \gamma m/m^*$, with a damping rate $\gamma = 2\pi \times (20\text{–}300)$ Hz. At longer times, as p_0 increases towards the center of the Brillouin zone ($p = h/2\lambda_1$), we observe a sudden increase of γ. This causes a stopping of the increase of p_0, followed by a decay towards zero which can again be fit with a constant damping rate of the order of 1 kHz. A corresponding change of behavior is shown by the width of $\rho(p)$ (see the inset of Fig. 1), which stays constant until p_0 increases, indicating a negligible heating of the system, while it rapidly increases at the instability point.

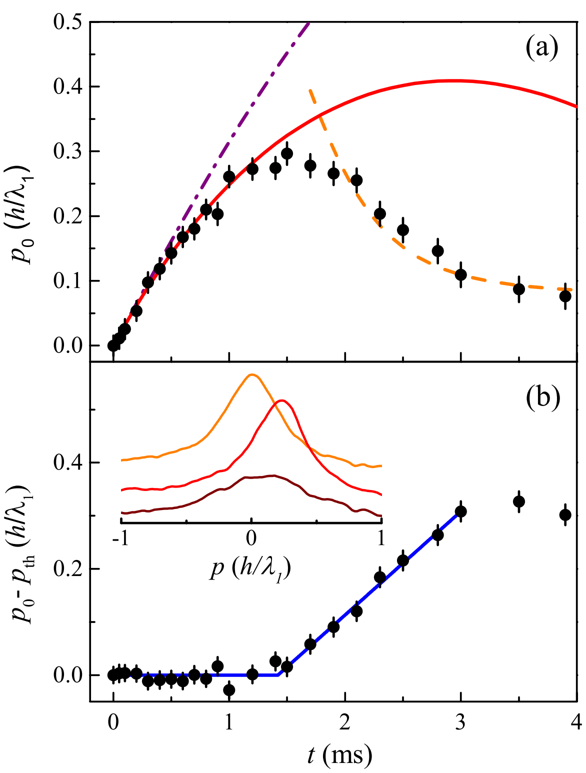

FIG. 1 (color online). Transport in nondisordered lattices. (a) Time evolution of the peak momentum for $U = 1.26J$ and $n = 3.6$. The experimental data (dots) are fitted at short times with a damped oscillation with $\gamma/2\pi = 135(10)$ Hz (continuous line) and at later times with $\gamma'/2\pi = 600(50)$ Hz (dashed line). The dash-dotted line is the expected oscillation in the absence of damping. (b) The difference between the fit to the initial damped motion and the experimental data (dots) is fitted (continuous line) to estimate the critical momentum. The inset shows $\rho(p)$ at three different times: $t = 0$, $t = 0.8$ ms, $t = 3.5$ ms, from top to bottom. The error bars represent the squared sum of statistical and systematic uncertainties.

This observation is in qualitative agreement with theoretical models [12,35–38] predicting two different regimes of quantum and thermal phase slips, in two different temperature regimes separated by a crossover temperature $k_B T_0 = c\sqrt{nJU}$, where c is a velocity-dependent factor smaller than unity [12,34,38]. For $T < T_0$ quantum phase slips dominate, with an exponential scaling of the nucleation rate with the interaction energy, density, and momentum as $\Gamma_Q \propto \exp[-7.1\sqrt{nJ/U}(\pi/2 - p\lambda_1/2\hbar)^{5/2}]$. For $T > T_0$, thermal activation of phase slips dominates, with a rate $\Gamma_T \propto \exp[-4nJ/3k_B T(\pi/2 - p\lambda_1/2\hbar)^3]$ [35]. In the framework of these models, the weak dependence of γ on p observed in previous experiments with low-density ($n \simeq 1$) 1D bosons in lattices [13,14] was justified by the small prefactor in the exponential scaling with p. Similarly, the smaller initial γ observed in our experiment can be attributed mainly to the fraction of the system with lower density. Our γ are indeed comparable to those of a previous experiment [13]. We actually observe an asymmetry in $\rho(p)$ that supports the idea of an inhomogeneous damping (inset of Fig. 1). The sudden instability is instead presumably due to the higher-n fraction, for which the theoretical expressions above predict a fast exponential increase of γ with p. Numerical calculations indicate that in a trap such instability should not be affected by the low-n component [12].

We estimate a critical momentum p_c separating the initial regime of weaker dissipation from the strongly unstable regime, by linearly fitting the difference between the experiment and the fit of the initial oscillation, as shown in Fig. 1(b). The measured p_c features a clear decrease when increasing U at constant J, while γ increases, as shown in Fig. 2. Eventually, p_c approaches zero as U approaches the predicted critical value for the

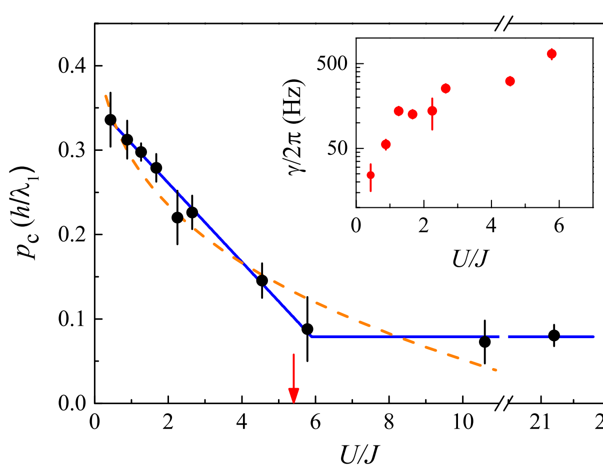

FIG. 2 (color online). Critical momentum for nondisordered lattices (dots) versus the interaction energy. The continuous line is a linear fit, the arrow marks the critical U/J for the superfluid-Mott insulator transition for $n = 2$, and the dashed line is the estimated p_c from the quantum phase-slips model. Inset: Initial damping rate γ.

Mott insulator ($U_c/J = 2 \times 2.674$ for the calculated mean occupation $n = 2$ [39]). Actually, even deep into the insulating regime we observe a small but finite p_c of the order of the inverse size of the system, as already observed [14]. By a piecewise fit of the data, we obtain a critical interaction that is comparable to theory: $U_c/J = 5.9(2)(4)$, where the uncertainties are statistical and systematic, respectively. These observations lead to the conclusion that also in 1D the onset of the Mott regime can be detected from a vanishing of p_c, as in 3D systems [14]. In 1D the transport is, however, clearly dissipative also for $p < p_c$, as expected.

The decrease of p_c and the corresponding increase of γ with U suggest a quantum activation of phase slip, since only Γ_Q has a direct dependence on U in the exponential. Since phase-slips models for γ in our large p and inhomogeneous n are not available, we tentatively compare our data to the complete expressions for Γ_Q [37] and Γ_T [35] regimes. In the spirit of Ref. [35], we estimate p_c by imposing that the nucleation rate gets larger than the experimental damping rate ($\approx 2\pi \times 1$ kHz). An unknown prefactor in the calculations is adjusted to match a single experimental datum at $U/J = 4.5$. The quantum phase-slip model predictions are in relatively good agreement with the experiment, as shown in Fig. 2. A similar analysis with the thermal model predicts instead an essentially constant p_c at constant T (see [34] for more discussion). We note that this result is not fully supported by the estimated T somewhat above T_0 although it is of the same order of magnitude; a careful verification of the role of quantum and thermal phase slips is left to future studies.

Let us now turn to the transport in the presence of disorder. We have, in particular, studied the weakly interacting regime, $U/J < 3$, where p_c for the nondisordered lattice can be very precisely measured. The experiment is performed as before, except for a finite Δ that is introduced together with the main lattice. Figure 3 shows how a small Δ results in a moderate increase of γ, but also in an anticipated instability. Both changes can be related to the idea that transport in disorder is dominated by the weakest hopping links, resulting in a smaller effective $J(\Delta)$ that in turn produces an increase of the phase-slip nucleation rates above, due to their exponential dependence on J [34]. A related phase-slip model developed for disordered superconductors indicates, indeed, a nucleation rate scaling exponentially with Δ [7], but it was derived in a different range of parameters and cannot be applied directly to our system. An important observation shown in Fig. 4 is that, for a fixed U, p_c features a clear decreasing trend for increasing Δ. Above a critical disorder strength Δ_c of the order of the total interaction energy per atom nU, p_c stops decreasing and stays constant at a small value close to that observed in the Mott-insulator regime. This is actually the regime where a weakly interacting Bose glass is predicted to appear, since the disorder can overcome the delocalization effect of the interaction [24,25]. The data in Fig. 4

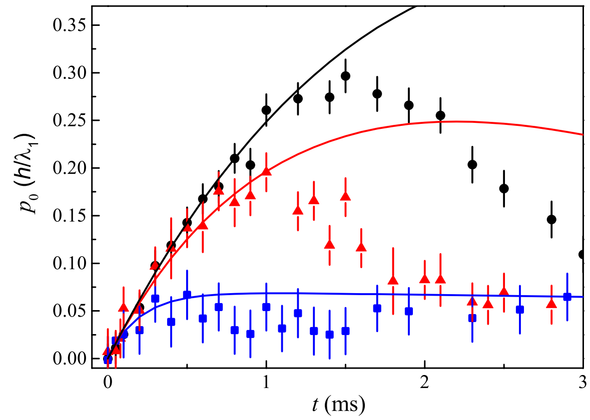

FIG. 3 (color online). Transport in disordered lattices. Time evolution of the peak momentum for $U = 1.26J$ for $\Delta/J = 0$ (dots), $\Delta/J = 3.6$ (triangles), and $\Delta/J = 10$ (squares). The lines are fits of the semiclassical motion to the initial oscillation. The fitted damping rates are $\gamma/2\pi = 130(10)$ Hz, $\gamma/2\pi = 250(30)$ Hz, and $\gamma/2\pi = 1.1(6)$ kHz, respectively.

show also that the decrease of p_c is accompanied by an increase of the rms momentum width δp at equilibrium (i.e., at $t = 0$), which is essentially the inverse of the correlation length ξ, towards a saturation value. δp starts to increase well before p_c has reached its minimum, indicating that the vanishing of p_c signals the onset of a strongly insulating phase, with a correlation length $\xi \simeq d$. Note that the observed p-dependent dynamics suggests that a simpler method with a fixed observation time, used in strongly interacting disordered systems [20], might underestimate the critical disorder strength for the insulating regime.

Motivated by the possibility of discriminating the fluid regime from the insulating one, we have studied how Δ_c evolves with U. For each U, we estimated Δ_c with a piecewise fit of the decreasing $p_c(\Delta)$, as shown in Fig. 4.

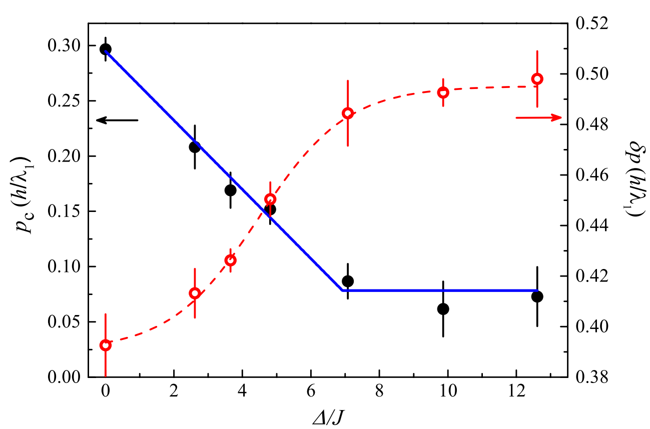

FIG. 4 (color online). Critical momentum p_c (full circles) and initial rms momentum width δp (open circles) for a fixed interaction energy ($U/J = 1.26$) and increasing disorder strength. A linear fit (continuous line) is used to estimate Δ_c, while the dashed line is a sigmoidal fit of δp.

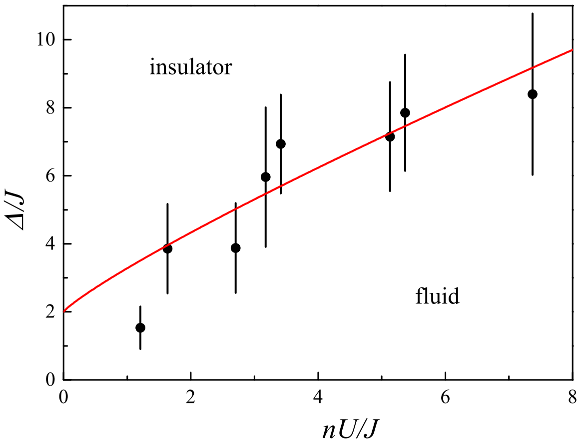

FIG. 5 (color online). Critical disorder to enter the insulating phase versus interaction energy. The experimental data from the critical momentum (dots) are fitted with the model described in the text (line). The uncertainty is dominated by a 20% error on the calibration of Δ [34].

The summary of these measurements in Fig. 5 shows a clear increase of Δ_c with U, indicating that the critical momentum of more strongly interacting systems is less affected by the disorder. The increase of Δ_c is actually fully justified, since the critical disorder strength to enter the Bose glass phase from the superfluid in the regime of weak interactions is expected to scale as $\Delta_c/J = A(E_{\rm int}/J)^{\alpha}$, where $E_{\rm int} \simeq nU$ is the total interaction energy per atom, while A and α are coefficients of the order of unity [24–26]. In the absence of an analytical model for the superfluid-Bose glass transition in a quasiperiodic lattice, we fit the experimental data with $(\Delta_c - 2)/J = A(nU/J)^{\alpha}$ to account for the critical $\Delta/J \simeq 2$ for localization in the noninteracting system. This choice is supported by the results of the density matrix renormalization group study in [27]. The fit gives an exponent $\alpha = 0.86(22)$ and a coefficient $A = 1.3(4)$. In the fit we excluded the data point for $\Delta/J < 2$, which should be described by a different mechanism of competition between the miniband structure of the quasiperiodic lattice and the interaction energy [22].

The exponent is compatible with the mean-field theory prediction $\alpha = 1$ for correlated Gaussian disorder in the so-called Thomas-Fermi regime, where $E_{\rm int}$ is larger than the typical disorder correlation energy E_c [25]. For the quasiperiodic lattice we estimate indeed an upper bound $E_c/J \simeq 0.7$ [34]. The observation is, however, not incompatible with the prediction $\alpha = \alpha(U) < 1$ found in disorder models that include corrections beyond the mean field [26]. We obtain a comparable exponent, although with a different prefactor A, from a similar analysis of the crossover in δp, also in agreement with previous experiments for very small U [19]. It is interesting to note that many current models for the superfluid-Bose glass transition at $T = 0$ are essentially based on the evolution of the same phase-slip nucleation rate that seems to be responsible for the dynamics observed in the present work [40–42]. A careful assessment of finite-size and finite-T effects is, however, required to establish the relation between the observed critical line and the theoretical fluid-insulator transition.

In conclusion, we have studied the momentum-dependent transport of 1D disordered bosons. We have employed the vanishing of the critical momentum for the observed instability to locate the fluid-insulator transition driven by disorder, across the interaction-disorder plane. The present study was for weak interactions and fixed equilibrium temperature. Future work should explore the role of temperature, also in connection with models for the many-body localization [43], and try to establish a link with the Luttinger-liquid theory for the superfluid-Bose glass transition for generic U and Δ [17,18,40–42]. In this context, the extension of the techniques used here to smaller momenta might allow us to probe the predicted universal scalings in lattices [37,38] and in disorder [7].

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[1] N. Giordano, Phys. Rev. Lett. **61**, 2137 (1988).
[2] R. Toda, M. Hieda, T. Matsushita, N. Wada, J. Taniguchi, H. Ikegami, S. Inagaki, and Y. Fukushima, Phys. Rev. Lett. **99**, 255301 (2007).
[3] J. Taniguchi, R. Fujii, and M. Suzuki, Phys. Rev. B **84**, 134511 (2011).
[4] T. Eggel, M. A. Cazalilla, and M. Oshikawa, Phys. Rev. Lett. **107**, 275302 (2011).
[5] A. Bezryadin, C. N. Lau, and M. Thinkham, Nature (London) **404**, 971 (2000).
[6] K. Yu. Arutyunova, D. S. Golubevc, and A. D. Zaikin, Phys. Rep. **464**, 1 (2008).
[7] S. Khlebnikov and L. P. Pryadko, Phys. Rev. Lett. **95**, 107007 (2005).
[8] A. T. Bollinger, R. C. Dinsmore III, A. Rogachev, and A. Bezryadin, Phys. Rev. Lett. **101**, 227003 (2008).
[9] O. V. Astafiev, L. B. Ioffe, S. Kafanov, Yu. A. Pashkin, K. Yu. Arutyunov, D. Shahar, O. Cohen, and J. S. Tsai, Nature (London) **484**, 355 (2012).
[10] A. Smerzi, A. Trombettoni, P. G. Kevrekidis, and A. R. Bishop, Phys. Rev. Lett. **89**, 170402 (2002).
[11] L. Fallani, L. De Sarlo, J. E. Lye, M. Modugno, R. Saers, C. Fort, and M. Inguscio Phys. Rev. Lett. **93**, 140406 (2004).
[12] E. Altman, A. Polkovnikov, E. Demler, B. I. Halperin, and M. D. Lukin, Phys. Rev. Lett. **95**, 020402 (2005).
[13] C. D. Fertig, K. M. OHara, J. H. Huckans, S. L. Rolston, W. D. Phillips, and J. V. Porto, Phys. Rev. Lett. **94**, 120403 (2005).
[14] J. Mun, P. Medley, G. K. Campbell, L. G. Marcassa, D. E. Pritchard, and W. Ketterle, Phys. Rev. Lett. **99**, 150604 (2007).

[15] D. McKay, M. White, M. Pasienski, and B. DeMarco, Nature (London) **453**, 76 (2008).

[16] M. Albert, T. Paul, N. Pavloff, and P. Leboeuf, Phys. Rev. Lett. **100**, 250405 (2008).

[17] T. Giamarchi and H. J. Schulz, Phys. Rev. B **37**, 325 (1988).

[18] M. P. A. Fisher, P. B. Weichman, G. Grinstein, and D. S. Fisher, Phys. Rev. B **40**, 546 (1989).

[19] B. Deissler, M. Zaccanti, G. Roati, C. D'Errico, M. Fattori, M. Modugno, G. Modugno, and M. Inguscio, Nat. Phys. **6**, 354 (2010).

[20] M. Pasienski, D. McKay, M. White, and B. DeMarco, Nat. Phys. **6**, 677 (2010).

[21] B. Gadway, D. Pertot, J. Reeves, M. Vogt, and D. Schneble, Phys. Rev. Lett. **107**, 145306 (2011).

[22] J. E. Lye, L. Fallani, C. Fort, V. Guarrera, M. Modugno, D. S. Wiersma, and M. Inguscio Phys. Rev. A **75**, 061603 (R) (2007).

[23] D. Dries, S. E. Pollack, J. M. Hitchcock, and R. G. Hulet, Phys. Rev. A **82**, 033603 (2010).

[24] P. Lugan, D. Clement, P. Bouyer, A. Aspect, M. Lewenstein, and L. Sanchez-Palencia, Phys. Rev. Lett. **98**, 170403 (2007).

[25] L. Fontanesi, M. Wouters, and V. Savona, Phys. Rev. Lett. **103**, 030403 (2009).

[26] R. Vosk and E. Altman, Phys. Rev. B **85**, 024531 (2012).

[27] G. Roux, T. Barthel, I. P. McCulloch, C. Kollath, U. Schollwöck, and T. Giamarchi, Phys. Rev. A **78**, 023628 (2008).

[28] G. Roati, M. Zaccanti, C. D'Errico, J. Catani, M. Modugno, A. Simoni, M. Inguscio, and G. Modugno, Phys. Rev. Lett. **99**, 010403 (2007).

[29] L. Fallani, J. E. Lye, V. Guarrera, C. Fort, and M. Inguscio, Phys. Rev. Lett. **98**, 130404 (2007).

[30] M. Modugno, New J. Phys. **11**, 033023 (2009).

[31] G. Roati, C. D'Errico, L. Fallani, M. Fattori, C. Fort, M. Zaccanti, G. Modugno, M. Modugno, and M. Inguscio, Nature (London) **453**, 895 (2008).

[32] S. Aubry and G. Andre, Ann. Isr. Phys. Soc. **3**, 133 (1980).

[33] C. D'Errico, M. Zaccanti, M. Fattori, G. Roati, M. Inguscio, G. Modugno, and A. Simoni, New J. Phys. **9**, 223 (2007).

[34] See Supplemental Material at http://link.aps.org/supplemental/10.1103/PhysRevLett.111.115301 for detailed information about the experiment and the analysis.

[35] A. Polkovnikov, E. Altman, E. Demler, B. Halperin, and M. D. Lukin, Phys. Rev. A **71**, 063613 (2005).

[36] J. Schachenmayer, G. Pupillo, and A. J. Daley, New J. Phys. **12**, 025014 (2010).

[37] I. Danshita and A. Polkovnikov, Phys. Rev. A **85**, 023638 (2012).

[38] I. Danshita, Phys. Rev. Lett. **111**, 025303 (2013).

[39] I. Danshita and A. Polkovnikov, Phys. Rev. A **84**, 063637 (2011).

[40] E. Altman, Y. Kafri, A. Polkovnikov, and G. Refael, Phys. Rev. B **81**, 174528 (2010).

[41] Z. Ristivojevic, A. Petkovic, P. Le Doussal, and T. Giamarchi, Phys. Rev. Lett. **109**, 026402 (2012).

[42] L. Pollet, N. V. Prokof'ev, and B. V. Svistunov, Phys. Rev. B **87**, 144203 (2013).

[43] I. L. Aleiner, B. L. Altshuler, and G. V. Shlyapnikov, Nat. Phys. **6**, 900 (2010).