

Quantum-gas microscopy and Talbot interferometry of the Bose-glass phase

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Disordered potentials fundamentally affect transport and coherence in quantum systems, giving rise to a Bose-glass phase in interacting bosonic systems—an insulating yet compressible phase lacking long-range coherence. Directly measuring a reduced coherence length of the Bose glass has been an outstanding challenge. We address this by employing Talbot interferometry combined with single-atom-resolved detection in a quantum-gas microscope. Using ultracold bosonic atoms in a two-dimensional lattice with site-resolved, reproducible disorder, we identify the Bose-glass phase through *in situ* density distributions and particle-number fluctuations, quantified via the Edwards-Anderson parameter, and through the visibility of interference patterns after time of flight. By driving the system across the Bose-glass phase, we further observe signatures of nonergodic dynamics. Our studies provide a starting point to further explore disordered systems in and out of equilibrium, and are relevant for understanding the dynamics and stability of disordered and glasslike quantum states in solid-state systems.

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I. INTRODUCTION

Disordered potentials govern electron transport and thermalization in real materials [1–3] and have also been studied in many-body quantum systems with ultracold atoms [4–7], in both experiment and theory. For noninteracting quantum systems, destructive interference from scattering in the disordered medium leads to Anderson localization [8], a phenomenon that has been observed, for example, in photonic lattices [9,10] and in cold-atom systems [11–15]. In general, the addition of interactions gives rise to a complex set of phenomena including, at low energies, the emergence of the Bose glass [16–19] and spin glasses [20] in bosonic and fermionic systems. The Bose glass has been examined using various theoretical approaches. Mean-field theories work less well near phase transitions because of statistical fluctuations and this also affects disordered systems [21]. While numerical methods can capture transitions for smaller system sizes in the presence of disorder, they face challenges with finite-size scaling [22], making the Bose glass an interesting system to study experimentally using a quantum simulator.

The Bose glass is both insulating and compressible [23,24], it has a gapless energy spectrum, and, unlike a superfluid, it has no long-range phase coherence, but instead it develops “superfluid puddles” that are not connected to each other.

It was initially believed to exist solely as a ground-state phase [16,24]; however, recent studies suggest that it can also occur at finite temperature [19,25,26]. In earlier studies, excitations in a quantum gas following a sudden change (quench) of optical speckle patterns have been observed [27], consistent with the existence of the Bose-glass phase, alongside a reduction of coherence [28]. Another important question is how disorder affects ergodicity and thermalization. In quantum many-body systems, establishing thermal states locally requires the propagation of correlations, which is restricted to exponentially long timescales when the wave functions are localized, either in an insulating or a disordered system. Such phenomena were observed with fermionic and bosonic atoms [29–31]. Signatures of nonergodic behavior have also recently been observed in the Bose glass in a two-dimensional optical quasicrystal [32], by probing the adiabaticity of the Bose-glass-to-superfluid transition.

In this study, we probe bosonic atoms in a two-dimensional square optical lattice onto which we project a known, controllable, and repeatable disorder light potential at the scale of individual sites [30]. Single-atom-resolved detection gives access to the *in situ* distribution and particle fluctuations, allowing us to obtain a local measurement of the Edwards-Anderson parameter [33–36]. We study a regime between the superfluid and Mott-insulating phases in the presence of disorder that shows increased values of the Edwards-Anderson parameter and a reduced visibility in the interference pattern after a time-of-flight expansion. These are complemented by a quantitative measurement of the coherence length using Talbot interferometry. We find that the coherence length reduces with increasing disorder strength, consistent with the formation of “superfluid puddles.” Finally, we investigate adiabaticity during transitions between the different regions in

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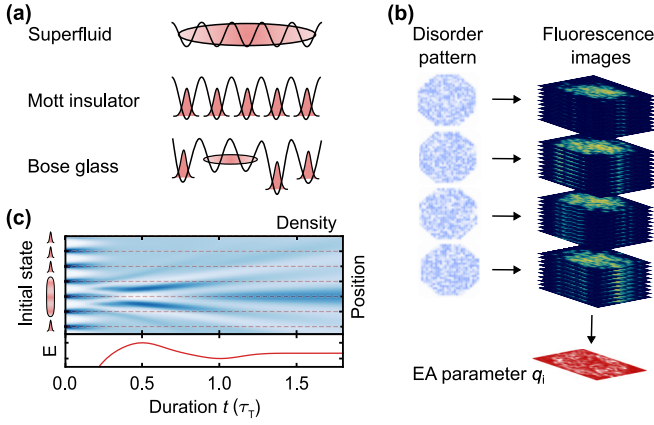


FIG. 1. Characterizing phases and coherence in the disordered Bose-Hubbard model. (a) Illustrations of the phases of the disordered Bose-Hubbard model highlighting the degree of coherence between lattice sites. (b) Schematic to calculate the Edwards-Anderson (EA) parameter. We record 10 fluorescence images for a number of disorder patterns at the same disorder strength and average the occupation on each lattice site to obtain disorder averages; see Eq. (2). (c) We probe the coherence length of the system along one axis using an interferometry technique based on the Talbot effect. Top panel: Time evolution of the sketched one-dimensional initial state after switching off the horizontal lattice. Initially phase-coherent regions show (shifted) revivals after (half-) integer Talbot times, τ_T . Horizontal dashed lines indicate the positions of the lattice sites. Bottom panel: Total energy after recapturing the atoms in the original lattice after a duration t . The decay time of the resulting oscillations can be related to the coherence length of the initial state (see the Appendix).

the phase diagram, observing nonadiabatic behavior when crossing into the disordered phase, consistent with properties of the Bose-glass phase.

II. BOSE-HUBBARD MODEL WITH DISORDER

The disordered Bose-Hubbard model is described by the Hamiltonian

$$\hat{H} = -J \sum_{\langle i,j \rangle} \hat{a}_i^\dagger \hat{a}_j + \sum_i \frac{U}{2} \hat{n}_i (\hat{n}_i - 1) + \sum_i (\epsilon_i - \mu) \hat{n}_i + \sum_i \Delta_i \hat{n}_i, \quad (1)$$

where J is the tunneling strength between neighboring lattice sites, $\hat{a}_i^\dagger$ and $\hat{a}_i$ are the bosonic creation and annihilation operators on neighboring sites i and j , respectively, U is the on-site interaction energy, $\hat{n}_i = \hat{a}_i^\dagger \hat{a}_i$ is the number operator, ϵ_i is the local energy shift from the harmonic confinement, and μ is the chemical potential. Δ_i is a local potential that creates random energy offsets on each lattice site, with a uniform box distribution $\Delta_i \in [0, \Delta]$, where Δ is the disorder strength. The model can give rise to a superfluid, a Mott-insulating and a Bose-glass phase [see illustration in Fig. 1(a)]. In the strongly interacting regime, $U \gg J$, a Mott-insulating phase exists for weak disorder $\Delta < U$. When the disorder is larger than the energy gap, $\Delta > E_g$, the state becomes gapless [16]. In the weakly interacting regime, the disorder results in localization of atoms and a shorter-range phase coherence compared to a

superfluid. In a two-dimensional infinite system, it is expected that the Bose-glass phase exists between the Mott insulator and superfluid for any nonzero value of Δ [24].

III. EXPERIMENTAL DETAILS

In our experiments [37], a cloud of ultracold ^{87}Rb atoms is held in a single antinode of a vertical optical lattice (wavelength $\lambda = 1064$ nm and depth $V_z = 25 E_r$, where $E_r = \hbar^2/2m\lambda^2$ is the recoil energy and m is the mass of a ^{87}Rb atom). After evaporative cooling, a superfluid of ~ 200 atoms is prepared in shallow horizontal lattices ($V_x = V_y \leq 10 E_r$). We then vary U/J by changing V_x and V_y along the horizontal axes. We use fluorescence imaging to obtain the *in situ* site occupations of the cloud, which has a diameter of ~ 16 sites. Typical temperatures are $k_B T \sim 0.1 U$, measured using radial fits of these occupations in a Mott-insulating state [38].

To create disordered systems, we project a repulsive light potential generated by a blue-detuned laser incident on a digital micromirror device (DMD). At each lattice site, this potential [pictured in Fig. 1(b)] has a random amplitude, $\Delta_i \in [0, \Delta]$, that follows a uniform box distribution (see the Appendix). Our method of creating disorder patterns is in contrast to most other ultracold-atom experiments on disordered systems, which typically use speckle patterns, quasirandom, or quasicrystalline potentials [12,39,40]. Using the DMD, we can generate not only reproducible disorder patterns, but also maintain phase alignment with the horizontal lattices, which would otherwise drift over time [37].

IV. TIME-OF-FLIGHT VISIBILITY

In the weakly interacting regime, the presence of disorder should result in a reduced coherence, which we initially measure via the visibility after time of flight [41]. This allows us to map out the parameter space for different disorder strengths and lattice depths, within a reasonable experimental runtime. We later complement this with a quantitative measurement of the coherence length via Talbot interferometry (see Sec. VI) for selected parameters.

In a time-of-flight measurement, the long-range phase coherence of the superfluid results in peaks in the interference pattern along the horizontal lattice directions. We quantify the visibility of the interference peaks, $\mathcal{V} = (n_{\max} - n_{\min})/(n_{\max} + n_{\min})$, by comparing the number of atoms in the peaks [$n_{\max}$, orange boxes in Fig. 2(a)] to those in regions diagonal from and equidistant to the central peak ($n_{\min}$, black boxes) [43]. For $\Delta = 0$, the visibility drops for increasing U/J [see bottom row of Fig. 2(b)], indicating a reduction of global phase coherence as a Mott-insulating state starts to form in the center of the cloud. Interestingly, when we increase Δ , we also observe a decrease in $\mathcal{V}$ at U/J below the superfluid-to-Mott-insulator transition, which occurs at $(U/J)_c = 16.7$ [22,44]. In the absence of thermal atoms, this reduction in $\mathcal{V}$ is consistent with the formation of the Bose glass. Our observation of a reduced $\mathcal{V}$ with increasing Δ cannot be attributed to trivial thermal excitations, as we demonstrate in Sec. VII. There is no clear transition point in our system due to the underlying harmonic confinement, resulting in a spatially varying chemical potential, and finite-size effects. We also compare

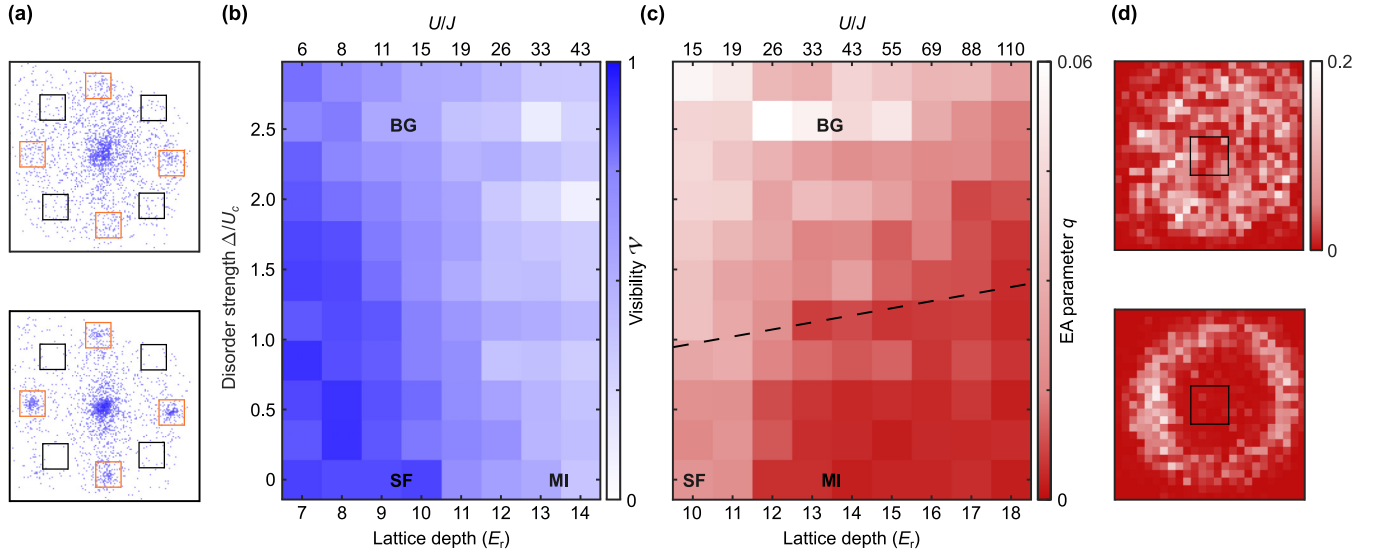


FIG. 2. Probing the disordered Bose-Hubbard model. (a) Average occupations of 10 time-of-flight images at $U/J = 15$ and $\Delta/U_c = 2.8$ (top), $\Delta/U_c = 0$ (bottom), where U_c is the interaction energy at $(U/J)_c$, plotted on a logarithmic color scale for clarity. Orange and black boxes highlight the regions used to determine $n_{\max}$ and $n_{\min}$, respectively, used to calculate the visibility, $\mathcal{V}$. (b) Visibility as a function of lattice depth and disorder strength. Our analysis identifies superfluid (SF), Mott-insulating (MI), and Bose-glass (BG) phases. (c) Edwards-Anderson (EA) parameter given by Eq. (2) averaged over the central 5×5 lattice sites [black box in bottom panel of (d)]. The dashed line indicates $\Delta/U = 1$, above which the disorder strength is larger than the energy gap in the ideal Mott insulator [24]. (d) Site-resolved EA parameter, q_i , at $U/J = 55$, using four different disorder realizations with 10 *in situ* images each, for $\Delta/U_c = 2.8$ (top) and $\Delta/U_c = 0$ (bottom). Error bars of $\mathcal{V}$ and q are shown in Figs. S2 and S3 (see the Supplemental Material [42]).

the visibility measurements to results from quantum Monte Carlo simulations (see Supplemental Material [42], which also includes Refs. [45–50]), which show the same general trend.

V. EDWARDS-ANDERSON PARAMETER

The visibility is not a reliable observable for distinguishing the Mott insulator from the Bose glass, as both lack long-range phase coherence [rightmost columns of Fig. 2(b)]. Therefore, we employ the Edwards-Anderson parameter q_i , which uses the atom distribution of the *in situ* images [35,36],

$$q_i = \overline{\langle \hat{n}_i \rangle^2} - \overline{\langle \hat{n}_i \rangle}^2. \quad (2)$$

Here, the brackets denote the thermal or ensemble average for a single disorder pattern, while the overline represents an average over different disorder patterns. Specifically, our data sets are divided into four subsets, one for each disorder pattern [Fig. 1(b)]. For each lattice site i , we determine q_i , which is zero when the occupations resulting from each disorder pattern are equal, and positive when the occupations are different from each other. For consistency, the $\Delta = 0$ data were also divided into four subsets, and the small but finite values of q_i reflect thermal fluctuations and atom-number variations between these subsets rather than genuine disorder effects.

For $\Delta = 0$, our measurements show $q_i \approx 0$ in the center of the system [bottom panel of Fig. 2(d)], limited by thermal excitations at finite temperature, indicative of the Mott insulator. We observe a circular region of nonzero values of q_i around the edge of the cloud. This is due to the entropy redistribution to the edge of the cloud resulting from the spatially varying chemical potential and fluctuations in the total atom

number between repetitions of the experiment. As we increase Δ , the Mott-insulating region decreases until q_i starts to rise in the center, indicating that no Mott-insulating component remains. To capture the transition from Mott insulator to Bose glass, we average q_i over the central 5×5 lattice sites, q , for increasing Δ [Fig. 2(c)]. When $\Delta > U$, the Bose glass is expected to form across the whole system [24]. This predicted behavior is closely matched by our experimental data, where q increases at disorder strengths $\Delta/U > 1$ [black dashed line in Fig. 2(c)]. As we analyze only the center of the system, where the effect of the harmonic confinement is negligible, we see a clearer transition here compared to the global measurement of the visibility [Fig. 2(b)]. We observe $q \approx 0.06$ for our maximum Δ , although the maximum possible value is $q = 0.25$. Theoretical studies [35] have shown that at a finite temperature close to ours ($k_B T \sim 0.1 U$), the Edwards-Anderson parameter can still distinguish well between glassy phases and thermal fluctuations. However, its magnitude is significantly reduced from $q = 0.25$ [35,51], consistent with our experimental observations.

The Edwards-Anderson parameter is unsuitable to distinguish between the superfluid and Bose glass [see leftmost columns of Fig. 2(c)] for lattice depths lower than $12 E_r$ ($U/J = 25.5$) where a superfluid component is present. This disordered superfluid phase has an inhomogeneous density, resulting in a nonzero q_i , making it indistinguishable from the Bose glass. Therefore, our measurements of the visibility and of the Edwards-Anderson parameter complement each other. Together, they reveal a region in the phase diagram with properties consistent with the Bose glass. In the measurements that follow, we study the coherence of this region in more detail.

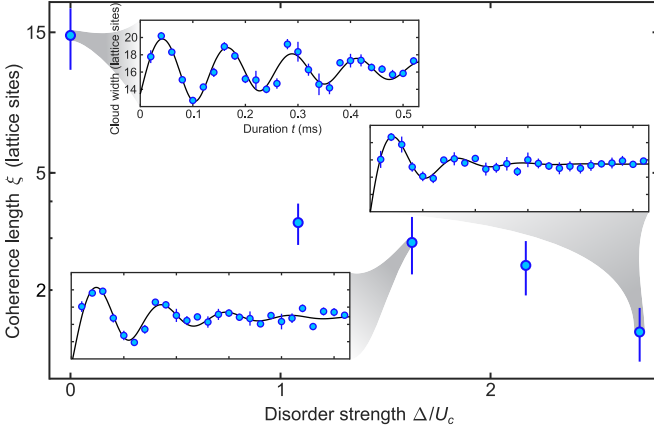


FIG. 3. Talbot interferometry. Coherence length ξ in units of the lattice spacing a_ℓ for increasing disorder strength at constant lattice depth ($V_{x,y} = 9 E_r$, $U/J = 11$). The coherence length is extracted from the fitted decay time of the Talbot interferometry measurements (insets) by comparing to a numerical calculation (see the Appendix). Error bars represent the 68% confidence intervals of the fitted decay times. Each data point of the curves in the insets results from the average of five images, each containing ~ 200 atoms. The error bars are the standard error. Black lines are fits with a damped sine, the decay time of which is used to extract ξ (see the Appendix).

VI. TALBOT INTERFEROMETRY

A distinct feature of the Bose glass is the absence of long-range phase coherence, while maintaining short-range coherence over a small number of sites. To measure the coherence length, we employ an interferometry technique based on the Talbot effect, following an earlier experimental demonstration [52]. It involves briefly switching off the lattice potential in one horizontal axis for a duration t , before rapidly turning it back on to project the matter wave onto the lattice. At integer multiples of the Talbot time, $\tau_T = 2ma_\ell^2/h = 123 \mu\text{s}$, where a_ℓ is the lattice spacing and h is Planck's constant, an initially coherent matter wave interferes constructively at the positions of the lattice sites. At noninteger multiples of τ_T , the positions of the interfering atoms are no longer aligned with the lattice sites, causing a sudden increase in potential energy when the lattice is turned back on. After a thermalization time, any potential energy manifests itself as an increase in the cloud temperature, resulting in periodic minima of the cloud width at integer multiples of τ_T and periodic maxima at half-integer multiples [see sketch in Fig. 1(c)]. We employ this method to quantitatively probe how the coherence length ξ of an initially superfluid system is affected by the addition of disorder. To measure the cloud width, we calculate the average distance of each atom from the center, after a thermalization time of 250 ms at $V_{x,y} = 6 E_r$. The disorder potential is switched off simultaneously with the lattice and remains off during thermalization and imaging. The resulting curve shows a damped oscillation as a function of t (insets in Fig. 3), where the presence of n half periods of oscillation indicates phase coherence over n lattice sites [52]. For $\Delta = 0$, the oscillation shows a weak decay with all peaks pronounced, indicating long-range phase coherence. When adding disorder, we observe that the oscillations decay more rapidly. At the

maximum experimentally accessible disorder strength, only the first two peaks are distinguishable from zero. To quantify the coherence length, we fit the data with an exponentially damped oscillation (solid lines in the insets of Fig. 3). We compare the fitted decay times to a numerical calculation, which use a density matrix, $\rho(i, j) \sim \exp(-|i - j|/\xi)$, with exponentially decaying coherence; see the Appendix. While for $\Delta = 0$ we observe a coherence length $\xi = 15(3)a_\ell$, the coherence length decreases to $\xi = 1.4(3)a_\ell$ for our maximum available disorder strength; see Fig. 3. This observation is consistent with the characteristics of a Bose glass, which is expected to exhibit coherence locally in small regions, but not across the entire system [16,53]. Our analysis shows the advantage of the Talbot interferometry measurements to give a quantitative measurement of the coherence length, in contrast to the time-of-flight visibility (Fig. 2). In the Mott-insulating phase, the oscillations from the Talbot interferometry measurements are significantly more damped (see the Supplemental Material [42]) compared to those shown in Fig. 3, indicating a coherence length that is smaller than a lattice site.

VII. ERGODICITY, ADIABATICITY, AND COHERENCE

In a finite-size system, a phase transition can be crossed without significant heating if both phases are ergodic. Signatures of nonergodicity in the Bose-glass-to-superfluid transition were recently observed using a quasicrystalline potential [32]. Here, we explore in a square lattice how the Mott-insulator-to-superfluid transition is affected by the presence of disorder. Using the data shown in Figs. 2(b) and 2(c), we draw an approximate phase diagram for our system [Fig. 4(a)] to serve as a guide for the subsequent measurements.

Initially, we perform Talbot interferometry for a set of points in the superfluid phase ($U/J = 5.5$), for three different disorder strengths. We find that the addition of the disorder potential does not result in significant heating by holding a disordered superfluid for 400 ms [points “b” and “c” in Fig. 4(a)] and observe long-range coherence comparable to the case without disorder (point “a”). Importantly, this relates back to our earlier measurements, where the reduced visibility at large disorder strengths [Fig. 2(b)], as well as the quantitative observation of the reduction of coherence length in the Talbot measurements (Fig. 3), can then be interpreted as signatures of the emerging Bose glass.

To probe adiabaticity, we first verify that transitioning to the Mott insulator and back to the superfluid at $\Delta = 0$ [trajectory “a₁” in Fig. 4(a)] restores a comparable coherence length, shown in Fig. 4(b). This was done by increasing V_x and V_y to reach $U/J = 25.5$ in 200 ms, and vice versa. Next, we investigate how entering the region where we identified the Bose glass affects the system's ability to restore phase coherence. Starting with a superfluid with weak disorder and increasing V_x and V_y to reach $U/J = 10.8$ in 200 ms, and vice versa [trajectory “b₁” in Fig. 4(a)], we observe a significantly reduced coherence length. We conduct two additional measurements, starting from a superfluid with stronger initial disorder to reach $U/J = 14.6$ and $U/J = 25.5$ [trajectories “c₁” and “c₂” in Fig. 4(a)]. These measurements show a further reduction in

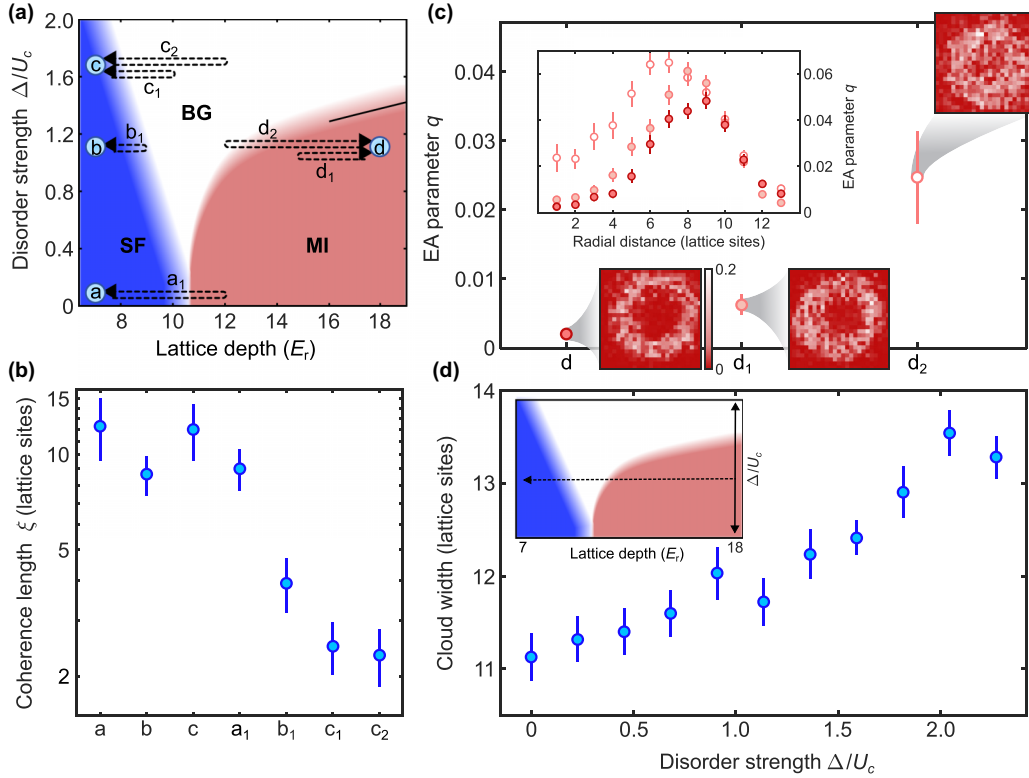


FIG. 4. Ergodicity and adiabaticity. (a) Approximate phase diagram of the disordered Bose-Hubbard model interpreted from the datasets in Fig. 2. The blue shaded area highlights the region where the entire cloud is in the superfluid phase and the red shaded region shows where the center of the cloud is Mott insulating. The blue circles show the points and the arrows show the trajectories used for the data sets shown in (b) and (c). (b) Coherence length ξ , measured using Talbot interferometry for the points and trajectories shown in (a) (for details, see the text). Each data point results from fits of four repetitions with the same parameters. (c) EA parameter q for three points and trajectories as shown in (a). Each measurement is an average over four disorder patterns with 10 images each. (d) Cloud width after one Talbot time, τ_T , after a ramp from the Mott-insulator regime to superfluid regime for increasing disorder strength (see inset where the dashed arrow indicates the trajectory for a data point at a particular disorder strength).

coherence length, shown in Fig. 4(b), consistent with a larger proportion of the cloud entering the Bose glass compared to “ b_1 ”, indicating the inability of the Bose-glass component to adiabatically transition back into the superfluid.

The observed loss of coherence cannot be attributed to intensity noise on the lattice lasers as we were able to restore the initial coherence when transitioning from the Mott insulator to the superfluid (trajectory “ a_1 ”). We have also verified that the coherence does not change when extending the duration over which we change V_x and V_y for the presented trajectories from 200 to 500 ms, which is the longest accessible timescale in our experiment.

Next, we explore the transition from the identified Bose-glass region to the Mott insulator. We first measure the Edwards-Anderson parameter in a Mott insulator in the presence of weak disorder [$U/J = 110$, $\Delta/U_c = 1.1$; point “ d ” in Fig. 4(a)], and find $q \approx 0$ in the center of the cloud, as shown in Fig. 4(c). When we decrease the lattice depth to reach $U/J = 55$ and return to $U/J = 110$ (trajectory “ d_1 ”), we observe an increase of q in the outer region of the cloud, which we attribute to localization occurring here first due to the lower chemical potential caused by the harmonic confinement. We then perform a similar measurement, changing instead to $U/J = 25.5$ and back (trajectory “ d_2 ”), to cross fully into the Bose-glass regime. When attempting to transi-

tion back to the original state, we observe an increase of q over the entire cloud [see Fig. 4(c)], again indicating nonadiabaticity.

Lastly, we investigate the effect of crossing the Bose glass at increasing disorder strengths, starting in the Mott insulator and finishing in the superfluid; see inset in Fig. 4(d). Specifically, we prepare the system at $U/J = 110$ and transition to $U/J = 5.5$ within 200 ms, before using Talbot interferometry to measure the cloud width after τ_T , at the first minimum of the oscillation. We observe an increase in cloud width for increasing disorder strengths, shown in Fig. 4(d), indicating a loss of coherence. We attribute this to excitations in the system due to the expected nonergodicity [32] when crossing an increasingly large Bose-glass region in the phase diagram.

VIII. CONCLUSION

In summary, we identified the Bose-glass phase experimentally in both weakly and strongly interacting regimes in a two-dimensional square lattice with controllable and reproducible disorder potentials. Using Talbot interferometry across the superfluid-to-Bose-glass transition, we observed a change from long- to short-range coherence. Additionally, we found that transitioning in and out of the Bose glass is nonadiabatic on the accessible experimental timescales. A similar

behavior was also observed in a quasicrystalline lattice at the superfluid-to-Bose-glass transition [32], which opens opportunities to compare and contrast the behavior of these systems and random disorder on Bravais lattices. It may be possible to cross transitions adiabatically on longer timescales, by reducing trap losses and heating, via ultralow vacuum pressure [54], and optical lattices with reduced intensity noise. This would enable the exploration of ergodic and nonergodic phases in a single model, for both weak and strong interactions, through scaling of the adiabatic timescales for traversing transitions or restoring system properties.

Our low-energy studies, combined with Talbot interferometry, provide a starting point to further explore disordered systems in and out of equilibrium, and particularly to probe the timescales to establish phase coherence in these systems. This is of relevance to understanding disordered and glasslike quantum states which have been observed in, e.g., cavity polaritons [55], quantum magnets [56], and overdoped cuprates [57]. Future studies with ultracold atoms could focus on the influence of disorder in different lattice geometries [58,59] and lower dimensions [60,61], or further explore the stability of many-body localization [30] and the onset of quantum avalanches [62]. Furthermore, controlled disorder is important in studying frustration and the stability of the spin-glass phase in the Fermi-Hubbard model [63–66]. By extending the system sizes, it is possible to study the similarities and differences between disordered and quasicrystalline systems further, e.g., by realization of the two-dimensional (2D) extension of the Aubry-Andre model [67,68], which could feature mobility edges and critical states and give further insights into the nature of many-body localization and the Bose glass [32].

Note added. Recently, we became aware of a new study of the disordered Bose-Hubbard model using a superconducting processor [69].

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L.K. and C.P. conducted the experiments and analyzed the data, with support from T.H., E.H., A.L.R., and S.K. Also, C.W.D., P.K., and A.J.D. developed the quantum Monte Carlo simulations, while T.H. led the simulations of the Talbot interferometry. All authors contributed to interpreting the results and writing the manuscript.

There are no competing interests to declare.

DATA AVAILABILITY

The data that support the findings of this article are openly available [70]; embargo periods may apply.

APPENDIX: EXPERIMENTAL DETAILS AND DATA ANALYSIS

1. System preparation and detection

We use the same quantum-gas microscope setup as detailed in our previous work [37]. Initially, we prepare a Bose-Einstein condensate of ~ 200 atoms in a single antinode of an optical lattice (wavelength $\lambda = 1064$ nm) in the vertical direction, with a depth of $V_z = 25 E_r$. To prepare a disordered system in the weakly interacting regime, we turn on the horizontal lattices, V_x and V_y , to the desired depth, $V_f \leq 10 E_r$, using an s-shaped ramp within 200 ms, and then turn on the disorder potential within 200 ms again using an s-shaped ramp. To prepare a disordered system in the strongly interacting regime, we first turn on the horizontal lattices to reach $V_{x,y} = 10 E_r$, before adding the disorder while simultaneously increasing V_x and V_y to their final values, $V_f > 10 E_r$, in a further 200 ms. We detect the atoms by collecting their fluorescence light with a high-NA microscope objective and use the Lucy-Richardson deconvolution method [71] to infer the site occupations.

2. Generation of disorder patterns

We create the repulsive disorder potentials using a laser ($\lambda = 666$ nm) incident on a digital micromirror device (DMD), similar to our previous work [37]. The light from the DMD is overlapped with the optical path used for fluorescence imaging of the atoms and projected through the microscope onto the atoms in the 2D plane. The magnification is such that approximately 18×18 pixels on the DMD correspond to one lattice site. We create a potential at each lattice site [pictured in Fig. 1(b)], with random amplitude $\Delta_i \in [0, \Delta]$, that follows a uniform box distribution. The disorder pattern covers a region with a diameter of 25 lattice sites.

A motorized flip mirror can direct the light pattern from the DMD onto a monitor camera. We employ an iterative intensity feedback algorithm to shape the Gaussian beam incident on the DMD into the target disorder pattern. The light pattern is blurred by the microscope objective when it is projected onto the atoms. To compensate for this, we convolve the monitor camera images with the point-spread function of the objective to effectively optimize the potential as seen by the atoms. Once the DMD pattern has been calculated, it is saved and displayed continuously. The disorder strength is controlled solely by changing the laser power. As detailed in our previous work [37], we measure the phase drift of the optical lattice between experimental realizations by analyzing the position of single atoms in the fluorescence image. This information is then used to shift the position of the DMD pattern for the next measurement, ensuring that the disorder potential remains aligned with the optical lattice. The estimated phase drift from one experimental realization to the next is $< 0.05 a_\ell$.

3. Talbot data analysis and simulation

To analyze the Talbot measurements, we obtain the cloud energy via the cloud width after a thermalization time of 250 ms at $V_{x,y} = 6 E_r$. We measure the cloud width by averaging the distances of each atom from the center of the cloud

after thermalization in two dimensions. We then use the fit function

$$R(t) \sim \sqrt{E_1 \cos(\omega t + \phi) \exp\left(-\frac{t}{\tau}\right) + E_\infty}, \quad (\text{A1})$$

where R is the cloud width, E_1 is the oscillation amplitude, $\omega = 2\pi/\tau_T$ is the Talbot frequency, ϕ is a phase shift, τ is the decay time related to the coherence length, and E_∞ is the long-time limit. This allows us to relate the measured Talbot decay time to the coherence in the system.

We compare the results of the Talbot measurements to single-particle simulations based on exact diagonalization in one spatial dimension and thereby relate the measured decay time τ to the coherence length ξ of the quantum state.

We start the Talbot simulations with a single-particle reduced density matrix $\rho_{x,x'}(t)$,

$$\begin{aligned} \rho_{x,x'}(t=0) &= \sum_{i,j} w_i(x) w_j^*(x') \sqrt{n_i n_j} \\ &\times \exp\left(-|i-j| \frac{a_\ell}{\xi}\right) |x\rangle \langle x'|, \end{aligned} \quad (\text{A2})$$

where i, j are the indices of the lattice sites, $w_i(x)$ is the Wannier function of the lowest band localized at site i , and n_i is the mean density of site i . This results in an exponential decay of the single-particle correlation function,

$$g^{(1)} = \langle a_i^\dagger a_j \rangle \sim \exp(-|i-j| a_\ell / \xi), \quad (\text{A3})$$

between lattice sites i and j . Note that Eq. (A2) describes perfect coherence within one lattice site, even for very small ξ , and that also the situation where the coherence length ξ is much larger than the system size L is covered.

We then time evolve $\rho_{x,x'}$ under the free single-particle propagator $U_t = e^{-iH_0 t/\hbar}$ for a duration t to obtain $\rho_{x,x'}(t) = U_t \rho_{x,x'} U_t^\dagger$ (Fig. 5), and measure the energy E of the final state in the lattice: $E(t) = \text{Tr}[\rho(t) H_\ell]$, where $H_\ell = p^2/(2m) + V \cos^2(\pi x/a_\ell)$.

Our cloud is roughly circular with radius $8a_\ell$ resulting in a (weighted) mean chain length of $13a_\ell$. Starting with $L = 13$ equally occupied sites (ground-band Wannier states, $V = 10 E_\tau$), we find Talbot oscillations (see Fig. 5) similar in shape to the experimental ones (Fig. 3). From the numerical data, we extract the decay time τ with the model $E(t) = E_1 \cos(\omega t + \phi) \exp(-t/\tau) + E_\infty$, where E_1 is the oscillation amplitude, ω the Talbot frequency, ϕ a phase shift, and E_∞ the long-time limit. Plotting τ versus ξ , we find the expected linear relationship $\tau/\tau_T = \xi/(2a_\ell)$ plus finite-size corrections for large coherence lengths. The relation between the Talbot decay time τ and initial coherence length ξ is well described by

$$\frac{\tau}{\tau_T} = \left(\frac{2a_\ell}{\xi} + 0.22 \right)^{-1}. \quad (\text{A4})$$

Experimentally, the thermalization in two dimensions after the one-dimensional Talbot evolution averages over the different

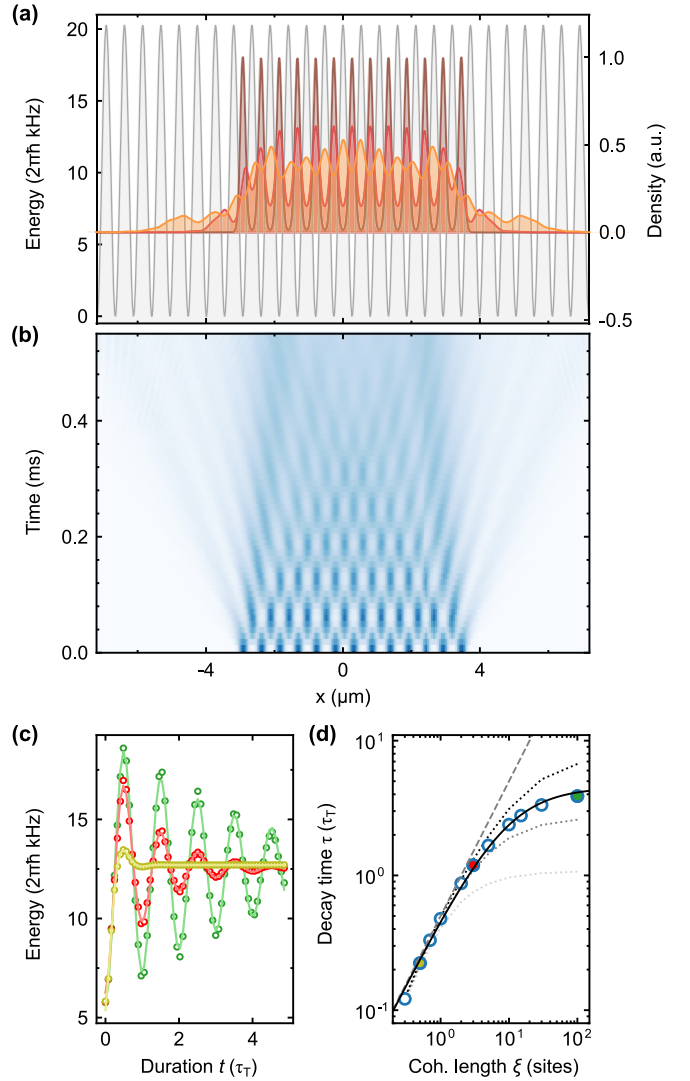


FIG. 5. Numerical simulation of Talbot interference. (a) Density distribution after free time evolution for $t = [0, 1, 2.5] \tau_T$ starting from a homogeneous lowest-band distribution over $L = 13$ sites with coherence length $\xi = 3 a_\ell$. (b) The full time evolution of this density distribution shows (shifted) revivals after (half-) integer Talbot times with decreasing contrast. (c) Total energy after recapturing the atoms with the lattice after duration t . The case from (a) and (b) (red) with $\xi = 3 a_\ell$ is compared to initial conditions with $\xi = 0.5 a_\ell$ (yellow) and $100 a_\ell$ (green). The solid lines are fits with exponentially damped harmonic oscillations. (d) Relation between coherence length ξ and damping of the Talbot oscillations τ from numerics for $L = 13$ [blue points, solid fillings indicate curves in (c)]. For low ξ , we find the expected $\tau = \xi/2 \times \tau_T / a_\ell$ (dashed), while at large ξ , the Talbot signal is dominated by finite-size effects [solid line, Eq.(A4)]. Dashed lines show results for $L = 20, 10, 5$ (decreasing darkness).

chain lengths in our system. We numerically checked that explicitly adding the energy of the one-dimensional chains of different length results in the same behavior as Eq. (A4). We have compared results with and without the trapping potential of the vertical lattice and found no significant effect over the timescales probed in the experiment.

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