Coherent Control of Dressed Matter Waves

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We demonstrate experimentally that matter waves can be coherently and adiabatically loaded and controlled in one-, two-, and three-dimensional strongly driven optical lattices. This coherent control is then used in order to reversibly induce the superfluid-Mott insulator phase transition by changing the strength of the driving. Our findings pave the way for studies of driven quantum systems and new methods for controlling matter waves.

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Even when the static properties of a quantum system are known, its response to an explicitly time-dependent variation of the system's parameters may be highly nontrivial, and qualitatively new states can appear that were absent in the original system. An atom in a radiation field, for instance, exhibits a number of fundamental phenomena such as the modification of its g factor in a radio-frequency field [1] and the dipole force acting on it in a spatially varying light field [2]. These effects can be described either in terms of Floquet quasienergy states or in the "dressed atom" picture [3]. While Floquet states represent the time-domain equivalent of the Bloch waves known from spatially periodic Hamiltonians, in the dressed atom picture the properties of the driven system arise from "dressing" the atom's electronic states with the photons of the radiation field. This concept can also be applied to macroscopic matter waves in driven periodic potentials [4], where the dressing is provided by the oscillatory motion of the lattice potential which couples to the center-of-mass motion of the matter wave through a dipole force (as opposed to other recent dressing schemes in which couplings between internal states lead to new adiabatic potentials [5,6]). In analogy to the dressed atom picture, such "dressed matter waves" can exhibit new properties and thus allow enhanced control of their quantum states. In this Letter we demonstrate coherent control of such dressed matter waves. In our experiments we adiabatically and reversibly change the quantum state of Bose-Einstein condensates (BECs) in driven optical lattices between a superfluid and a Mott insulator (MI) [7–9] by varying the amplitude of the driving. Our method offers a versatile testing ground for driven quantum systems and for new quantum control schemes for matter waves.

Cold atoms in optical lattices [8] can be described in the Bose-Hubbard (BH) model by the parameter U/J, where J is the hopping term relating to tunneling between adjacent sites and U is the on-site interaction energy [Fig. 1(a)]. If $U/J \ll 1$, tunneling dominates and the atoms are delocalized, whereas for $U/J \gg 1$ the interaction term leads to a loss of phase coherence through the formation of number-squeezed states with increased quantum phase fluctuations. At a critical value of U/J the system undergoes a quantum phase transition to a MI state. Using optical lattices one can tune U/J by changing the lattice depth [9,10], which affects both U and J through the width of the on-site wave functions [Fig. 1(a)]. Alternatively, by periodically shaking the lattice J can be suppressed [11,12]. This

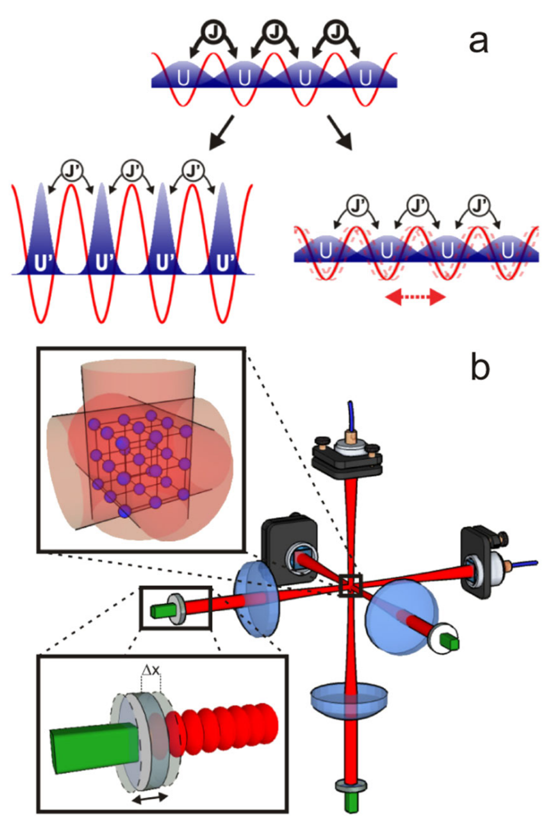

FIG. 1 (color online). (a) Principle of the coherent control of the tunneling parameter J. A condensate in a lattice is characterized by J and the on-site interaction U (above). While varying the lattice depth changes both U and J (left) to U' and J', strong driving selectively changes J to $J' = J_{\mathrm{eff}}$ (right). (b) Experimental setup for a three-dimensional driven optical lattice.

principle, related to the coherent destruction of tunneling in double-well systems [13,14], was recently experimentally demonstrated [15,16]. In the driven system a time-dependent energy term $K\cos(\omega t)\sum_j j\hat{n}_j$ is added to the Hamiltonian (with $\hat{n}_j$ the number operator on site j), and in a time-averaged description the single-particle dispersion relation

$$E(k) = -2J\cos(kd_L) \tag{1}$$

of the BH Hamiltonian (k is the quasimomentum) is still valid if J is replaced by an effective tunneling parameter

$$J_{\mathrm{eff}} = \mathcal{J}_0(K_0)J, \tag{2}$$

with d_L the lattice constant, $\mathcal{J}_0$ the zeroth-order Bessel function, and $K_0 = K/\hbar\omega$ the normalized driving strength for frequency ω and driving amplitude K (see below). It should, therefore, be possible to define an effective parameter U/J_{eff} [17] for the BH Hamiltonian.

In our experiments BECs of 6×10^4 atoms of ^{87}Rb were adiabatically loaded into the lowest energy band of an optical lattice [18]. The one, two, or three-dimensional lattices were created by focused linearly polarized laser beams ($\lambda = 842$ nm), each of which was retroreflected by a lens and a mirror [Fig. 1(b)], resulting in periodic potentials $V(x_i) = V_0\sin^2(\pi x_i/d_L)$, where V_0 is the lattice depth [19], $x_i = x,\ y,\ z$, and $d_L = \lambda/2$. The mirrors were mounted on piezoelectric actuators that allowed us to sinusoidally shake each lattice back and forth [20] with frequency $\omega/2\pi$ (up to several kHz) and amplitude Δx_i. We define a dimensionless driving strength $K_0 = K/\hbar\omega = (\pi^2/2)(\Delta x_i/d_L)(\omega/\omega_{\mathrm{rec}})$, where $\omega_{\mathrm{rec}} = \hbar\pi^2/2md_L^2 = 2\pi\times3.24$ kHz is the recoil frequency (with m the mass of the ^{87}Rb atoms).

We first show that a BEC in a driven one-dimensional (1D) lattice remains phase coherent and follows changes in K_0 (Fig. 2). After loading a BEC into the lattice ($V_0 = 18E_{\mathrm{rec}}$, where $E_{\mathrm{rec}} = \hbar\omega_{\mathrm{rec}}$ is the recoil energy), K_0 was linearly ramped from 0 to $K_0 = 2.7$ (by varying K at constant ω) in 113 ms and back to 0 in the same time. At times t [where $t = N(2\pi/\omega)$ was an integer multiple of the driving period] the lattice and the dipole trap were suddenly switched off, and the atoms were imaged on a CCD camera after 23.8 ms of free fall. The interference pattern created by atoms originating from different lattice wells (in our experiments around 40 sites were occupied) exhibited sharp peaks when the BEC was phase coherent over the entire lattice, whereas when phase coherence was lost, a featureless pattern was observed. Figure 2(a) shows that when J_{eff} is large and hence $U/J_{\mathrm{eff}} \ll 1$, phase coherence persists for several tens of milliseconds in spite of the increasingly strong driving. The appearance of a stable interference pattern proves that the BEC occupies a single Floquet state of the driven system. We also verified that while J_{eff} changes with K_0, the effective interaction U_{eff} (inferred from the widths of the on-site wave functions through the heights of the peaks in the interference pattern) remains constant [inset of Fig. 2(b)].

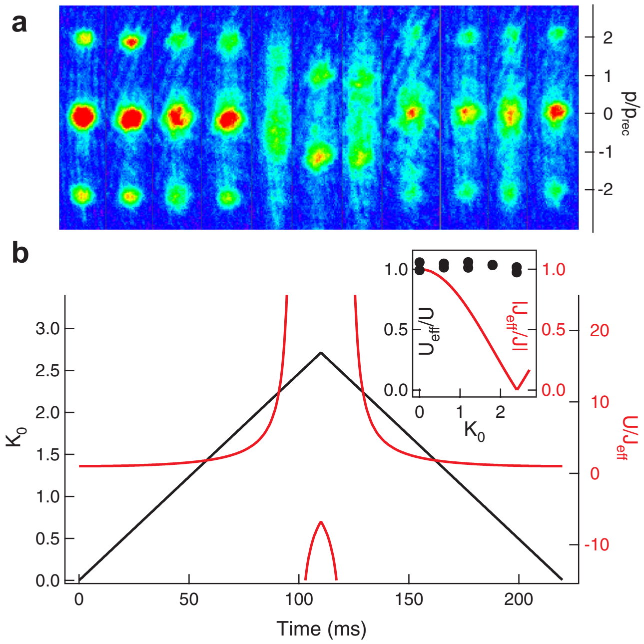

FIG. 2 (color online). Coherent control of a BEC inside a driven 1D lattice. (a) Time-of-flight interference patterns. The length scales have been converted into momentum units, with $1p_{\mathrm{rec}}$ corresponding to 138 μm. Here $\omega/2\pi = 6$ kHz and $V_0 = 18E_{\mathrm{rec}}$. (b) Time dependence of K_0 and U/J_{eff}. Inset: U_{eff}/U as a function of K_0. For comparison, the (theoretical) behavior of $|J_{\mathrm{eff}}/J|$ is also shown.

In Fig. 2, three parameter regions can be distinguished. While for $K_0 < 2.4$ [i.e., below $K_0 \approx 2.4$ for which $\mathcal{J}_0(K_0)\approx 0$] the BEC occupies a Floquet state at the center of the Brillouin zone ($k = 0$) as reflected by a dominant peak at $p = 0$ and side peaks at $\pm 2p_{\mathrm{rec}} = \pm h/d_{\mathrm{L}}$, for $K_0 > 2.4$ ($J_{\mathrm{eff}} < 0$) two peaks at $\pm p_{\mathrm{rec}}$ appear. This indicates that the Floquet state of lowest mean energy now lies at the edge of the Brillouin zone [$k = \pm p_{\mathrm{rec}}/\hbar$; see Eq. (1)] [21]. Finally, when $K_0 = 2.4$ ($J_{\mathrm{eff}} \approx 0$), $U/J_{\mathrm{eff}} \gg 1$ and phase coherence is lost. When K_0 is reduced back to 0 at the end of the cycle, the initial interference pattern is restored almost perfectly, suggesting that the response of the system to the parameter variation was approximately adiabatic.

Since adiabaticity is a key concept in physics, we studied the conditions for adiabatic following in our system more systematically. This is important as the intuitive idea of an arbitrarily slow change in one of the system's parameters allowing it to adjust its state to the instantaneous parameter values is no longer valid in driven systems [4,22], where the Floquet quasienergy spectrum exhibits avoided crossings over which a parameter scan has to be performed sufficiently fast in order for the system to perform Landau-Zener tunneling across the gap and hence adiabatically follow the energy levels of an equivalent nondriven system. The degree of adiabaticity in our experiments was measured using cycles with linear ramps from $K_0 = 0$ to $K_0 = 2.3$ [23] and back. In order to compare the results for different sets of parameters, at the end of the cycle we ramped down V_0 to $4E_{\rm rec}$ and measured the ratio of the width σ of the interference peak at $p = 0$ and σ_0 for $K_0 = 0$, which reflects any increase in energy or loss of coherence during the cycle. The main results are summarized in Fig. 3. Clearly, for fixed driving frequencies and ramp durations minimum lattice depths exist below which adiabatic ramping is impossible [Fig. 3(a)], as indicated by the sharp increase in σ/σ_0. This minimum is well defined and narrow, suggesting a transition to a chaotic regime or driving-induced interband transitions. For a given V_0 the degree of adiabaticity depends sensitively on ω [Fig. 3(b)]. Again, interband transitions may explain the breakdown for $\omega/2\pi > 6$ kHz, while, e.g., the partial breakdown between 4 kHz and 4.5 kHz cannot be explained in this way. We also investigated the dependence of σ/σ_0 on the ramp duration for constant V_0 and ω. In 1D we found an optimum ramp time of $\approx$20 ms, in agreement with the structure of the Floquet spectrum described above.

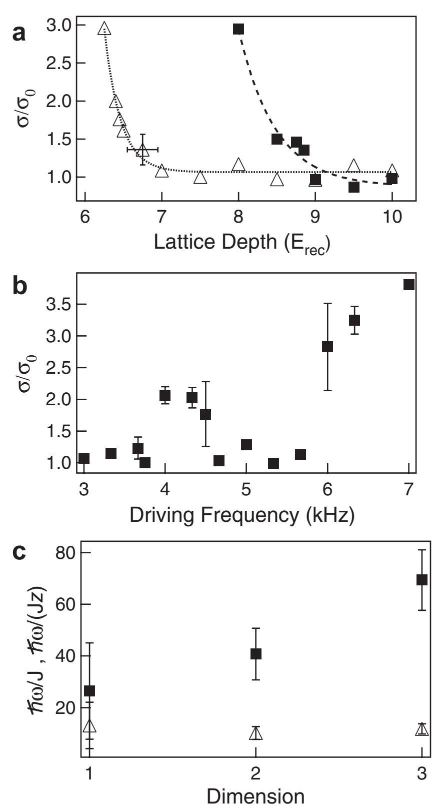

FIG. 3. Adiabaticity in driven optical lattices. (a) Dependence of σ/σ_0 (see text) on V_0 in a 1D lattice for a linear ramp $K_0 = 0$ to $K_0 = 2.3$ in 15 ms, a holding time at $K_0 = 2.3$ of 200 ms, and an identical ramp back to $K_0 = 0$. Here $\omega/2\pi = 3$ kHz (open triangles) and $\omega/2\pi = 6$ kHz (solid squares). (b) Dependence of σ/σ_0 on the driving frequency for a fixed $V_0 = 5.7E_{\rm rec}$. The ramp time was 10 ms with 2 ms holding time. (c) Minimum value of $\hbar\omega/J$ (solid squares) and $\hbar\omega/Jz$ (open triangles) for different lattice dimensions.

We performed similar tests with 2D and 3D lattices [see Fig. 3(c)] using the same ω, K_0, and driving phases for all the lattices. While adiabatic ramps were again possible for certain sets of parameters, the minimum V_0 for adiabaticity increased with dimension d and the optimum ramp time decreased to a few milliseconds. For a given ω the minimum value of the normalized ratio $\hbar\omega/Jz$ of the dimensionless parameter $\hbar\omega/J$ (where $z = 2d$ is the number of nearest neighbors) was constant at about 12 [Fig. 3(c) shows the mean values for different driving frequencies and ramp times]. While this suggests that $\hbar\omega/Jz$ indicates a value of J below which adiabatic control is possible at constant ω, we also found that changing ω and J independently for a given d does not always give the same value [as indicated by the error bars in Fig. 3(c)]. While these results are only a first step towards understanding adiabatic following of Floquet states, and more theoretical and experimental work needs to be done, they nevertheless show that there are, indeed, regions in parameter space for which adiabatic control is possible.

In order to realize the driving-induced superfluid-Mott insulator transition [17,24], we first loaded a BEC into a 3D lattice with $V_0 = 11E_{\rm rec}$ using an exponential ramp of 150 ms duration and then linearly increased K_0 from 0 to $K_0 = 1.62$ in 4 ms [25]. While in an undriven lattice with $V_0 = 11E_{\rm rec}$ the BEC is superfluid with $U/6J \approx 3.5$, for a driven lattice $U/6J_{\rm eff}$ at $K_0 = 1.62$ is ≈ 7.9, i.e., larger than the critical values 5.4 (mean field treatment [8]) and 4.9 (quantum Monte Carlo simulations [26]). Hence the system is in the Mott insulating phase [Fig. 4(a)], as indicated by a distinct loss of phase coherence [9]. When K_0 is ramped back to 0 the interference pattern reappears, proving that the transition was induced adiabatically and that the system was not excited by the driving. For a more quantitative statement we induced the MI transition in two different ways: (a) by increasing V_0 in an undriven lattice as in [9] and (b) by varying K_0 for constant V_0. Figure 4(b) clearly shows that the visibility [27] vanishes as $U/6J_{\rm eff}$ is increased and returns to its original value after ramping K_0 back to 0. The dependence of the visibility on $U/6J_{\rm eff}$ is the same for methods (a) and (b), strongly indicating that the same many-body state is reached. The independent control over $J_{\rm eff}$ also allowed us to measure the excitation spectrum of the system [10,28] by sinusoidally modulating $J_{\rm eff}$ through K_0 (rather than by modulating V_0, which also changes U). While in the superfluid regime a gapless excitation strength as a function of the modulation frequency appears, in the MI regime we find a gapped spectrum with a narrow resonance [28] [Fig. 4(c)].

Our results confirm and extend the role of cold atoms in optical lattices as versatile quantum simulators [29,30] and open new avenues for the quantum control of cold atoms, thus establishing a link to coherent control in other systems

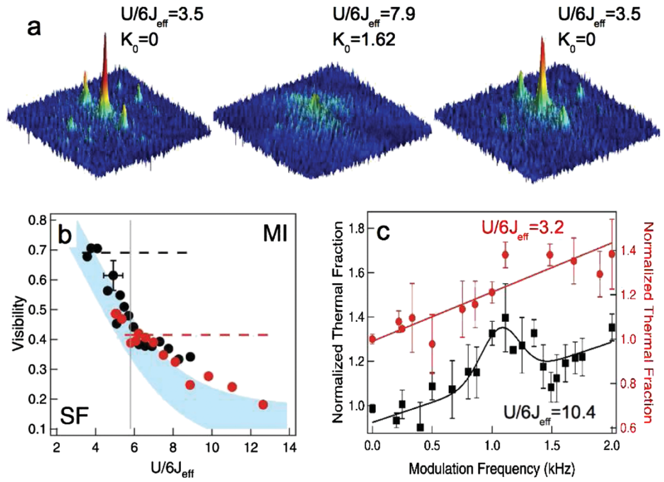

FIG. 4 (color). Driving-induced Mott insulator transition. (a) In a driven 3D lattice of constant depth ($V_0 = 11E_{\rm rec}$, $\omega/2\pi = 6$ kHz), K_0 was ramped from 0 to $K_0 = 1.62$ in 4 ms and back. (b) Visibility of the interference pattern as a function of $U/6J_{\rm eff}$ varying K_0 for $V_0 = 11E_{\rm rec}$ (black symbols) and $V_0 = 12.2E_{\rm rec}$ (red symbols) with the dotted lines indicating the respective visibilities after returning to $K_0 = 0$. The blue region is a fit (its width indicating the statistical error) to the data for varying V_0 at $K_0 = 0$. Vertical grey line: border between the superfluid (SF) and MI regions. (c) Excitation spectrum measured by modulating $J_{\rm eff}$. Plotted here is the thermal fraction (normalized to the thermal fraction without excitation) calculated from a bimodal fit to the BEC after ramping down the lattice to $V_0 = 4E_{\rm rec}$. Here $V_0 = 10.3E_{\rm rec}$, $K_0 = 0.7$ (red circles) and $V_0 = 11E_{\rm rec}$, $K_0 = 1.8$ (black squares). The modulation depths for K_0 were 0.8 and 0.3, respectively. The solid lines are fits to guide the eye, and the vertical scales have been offset for clarity.

such as molecules in laser fields [31] and Cooper pairs in Josephson qubits [32]. The principles demonstrated here can be straightforwardly extended to more than one driving frequency [33] and to more complicated lattice geometries such as superlattices [34].

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