Coherent control of dressed matter waves

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By moving the pivot of a pendulum rapidly up and down one can create a stable position with the pendulum’s bob above the pivot rather than below it [1]. This surprising and counterintuitive phenomenon is a widespread feature of driven systems and carries over into the quantum world. Even when the static properties of a quantum system are known, its response to an explicitly time-dependent variation of its parameters may be highly non-trivial, and qualitatively new states can appear that were absent in the original system. In quantum mechanics the archetype for this kind of behaviour is an atom in a radiation field, which exhibits a number of fundamental phenomena such as the modification of its $g$-factor in a radio-frequency field [2] and the dipole force acting on an atom moving in a spatially varying light field [3]. These effects can be successfully described in the so-called dressed atom picture [4]. Here we show that the concept of dressing can also be applied to macroscopic matter waves [5], and that the quantum states of ”dressed matter waves” can be coherently controlled. In our experiments we use Bose-Einstein condensates in driven optical lattices and demonstrate that the many-body state of this system can be adiabatically and reversibly changed between a superfluid and a Mott insulating state [6, 7, 8] by varying the amplitude of the driving. Our setup represents a versatile testing ground for driven quantum systems, and our results indicate the direction towards new quantum control schemes for matter waves.

An atom in a radiation field can be described in the dressed atom picture [4] (or in equivalent approaches using, e.g., Floquet quasienergy states) in which the modified 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 [5], where the ”dressing” is provided by the oscillatory motion of the lattice potential. In analogy to the dressed atom picture, such ”dressed matter waves” can exhibit new properties absent in the original system and thus allow enhanced control of its quantum states. Here we demonstrate that matter waves can be adiabatically transferred into a well-defined Floquet quasienergy state of a driven periodic potential while preserving their quantum coherence.

Cold atoms in optical lattices [7] can be described in the Bose-Hubbard 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 (see Fig. 1a). When $U/J$ is small, tunneling dominates and the atoms are delocalized over the lattice, whereas a large value means that the inter-

FIG. 1: a, Principle of the coherent control of the tunneling parameter $J$. A condensate in a lattice is characterized by the tunneling parameter $J$ and the on-site interaction $U$ (above). While changing the lattice depth results in a variation of both $U$ and $J$ (left) to $U'$ and $J'$, strong driving selectively changes $J$ to $J' = J_{\text{eff}}$ (right). b, Experimental setup for a three-dimensional driven optical lattice. The periodic potentials in the three spatial directions are produced by retro-reflecting a focused laser beam off a mirror mounted on a piezo-electric actuator.
struction of tunneling in double-well systems [12, 13], was recently experimentally demonstrated [14, 15]. In the driven system \( J \equiv J_0(K_0)J \), with \( J_0 \) the zeroth-order Bessel function and \( K_0 \) the driving strength (defined below). This suggests that it should be possible to use \( J_{\text{eff}} \) in the many-body Hamiltonian describing a BEC in a lattice and hence to define an effective parameter \( U/J_{\text{eff}} \) [16]. In the following, we show that this assumption is borne out by experiment.

In our experiments we created Bose-Einstein condensates of \( 6 \times 10^4 \) atoms of \(^{87}\text{Rb} \), which were then adiabatically loaded into the lowest energy band of an optical lattice [17]. The 1, 2 or 3-dimensional lattices were realized by focusing linearly polarized laser beams (\( \lambda = 842 \text{ nm} \)) onto the BEC. Each lattice beam was retro-reflected by a combination of a lens and a mirror (see Fig. 1b), resulting in periodic potentials \( V(x_i) = V_0 \sin^2(\pi x_i/d_L) \) along the three spatial directions, where \( d_L = \lambda/2 \) is the lattice constant, \( V_0 \) the lattice depth and \( x_i = x, y, z \). The mirrors were mounted on piezo-electric actuators that allowed us to sinusoidally shake each optical lattice back and forth [18] with frequency \( \omega \) (up to several kHz) and amplitude \( \Delta x_i \). We define a dimensionless driving strength \( K_0 = K/\hbar \omega = (\pi^2/2)(\Delta x_i/d_L)(\omega/\omega_{\text{rec}}) \), where \( \omega_{\text{rec}} = \hbar \pi^2/2m d_L^2 = 2\pi \times 3.24 \text{ kHz} \) is the recoil frequency (with \( m \) the mass of the \(^{87}\text{Rb} \) atoms).

In order to show that a BEC in a driven lattice (a) maintains its phase coherence and (b) adiabatically follows changes in the \( K_0 \), we performed preliminary experiments in one-dimensional lattices (see Fig. 2). After loading a BEC into the lattice \( (V_0 = 18 E_{\text{rec}}, \text{where } E_{\text{rec}} = \hbar \omega_{\text{rec}} \text{ is the recoil energy, } K_0 \text{ was linearly increased from } 0 \text{ to } K_0 = 2.7 \text{ in } 113 \text{ ms and back to } 0 \text{ in the same time. At times } t \text{ (where } t = N \times 2\pi/\omega \text{ 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 \text{ ms of free fall. The interference pattern created by atoms originating from different lattice wells (in our experiments around } 40 \text{ sites were occupied) consisted of well-defined peaks when the condensate was phase coherent over the entire lattice, whereas when phase coherence was lost a broader, featureless pattern was observed. Fig. 2a shows that when } J_{\text{eff}} \text{ is large and hence } U/J_{\text{eff}} \ll 1, \text{ the phase coherence persists for several tens of milliseconds in spite of the strong driving. The appearance of a stable, well-defined interference pattern proves that the BEC occupies a single Floquet state of the driven system and adiabatically follows that state as } K_0 \text{ is varied. We also verified that while } J_{\text{eff}} \text{ changes with } K_0, \text{ the effective interaction parameter } U_{\text{eff}} \text{ (inferred from the the relative height of the side-peaks in the interference pattern) remains constant (see the inset of Fig. 2b).}

While for \( K_0 < 2.4 \) (for \( K_0 \approx 2.4, J_0(K_0) \approx 0 \)) the condensate occupies a Floquet state with quasimomentum \( q = 0 \) at the center of the Brillouin zone as reflected by an interference pattern with a dominant peak at zero momentum and sidepeaks at \( \pm 2 \rho_{\text{rec}} = \pm 2 \times \hbar/\lambda, \) for \( K_0 > 2.4 \) (where \( J_{\text{eff}} \) is negative) two peaks at \( \pm \rho_{\text{rec}} \) appear. This indicates that the Floquet state of lowest mean energy now corresponds to \( q = \pm \rho_{\text{rec}} \) at the edge of the Brillouin zone. Finally, when \( K_0 = 2.4 \) and hence \( J_{\text{eff}} \approx 0, \) \( U/J_{\text{eff}} \gg 1 \) and phase coherence is lost due to increased quantum phase fluctuations (22). 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 adiabatic.

Since adiabaticity is a key concept in physics, we studied the conditions for adiabaticity 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 at all times is no longer valid in driven systems [5, 19]. The degree of adiabaticity in our experiments was measured by performing cycles with triangular ramps for various lattice depths, driving frequencies, ramp durations and dimensionalities of the lattice. In order to compare the results for different sets of parameters, at the end of the cycle we ramped down \( V_0 \) to 4 \( E_{\text{rec}} \), and measured the ratio of the width \( \sigma \) of the interference peak at \( p = 0 \) and its width \( \sigma_0 \) for \( K_0 = 0 \), which reflects any increase in energy and/or loss of coherence during the cycle. The main results are summarized in Fig. 3. Clearly, for fixed driving frequencies and ramp durations there exist
minimum lattice depths below which no adiabatic ramping is possible (Fig. 3a), as indicated by the sharp increase in $\sigma/\sigma_0$ below those values. This minimum is well-defined and narrow and suggests a transition to a chaotic regime or interband transitions induced by the driving. We also found that for a given lattice depth the degree of adiabaticity depends sensitively on the driving frequency (Fig. 3b). Again, interband transitions may be responsible for the breakdown of adiabaticity at frequencies above 6 kHz, while other features such as the partial breakdown between 4 kHz and 4.5 kHz cannot be explained in this way. We also investigated the dependence of the degree of adiabaticity on the ramp duration keeping $V_0$ and $\omega$ constant. We found that there exists an optimum ramp time of around 20 ms (depending slightly on $V_0$ and $\omega$).

Furthermore, we performed adiabaticity tests with two- and three-dimensional lattices (see Fig. 3c) for which the BEC was loaded into a 2D or 3D lattice as described above for the 1D case, and the driving strength of the lattices was then ramped up and down (using the same frequencies, phases and driving strengths for all the lattices). Again, adiabatic ramps were possible for certain sets of parameters. In particular, the minimum depth for adiabaticity increased with increasing dimension $d$, whereas the optimum ramp time decreased to a few milliseconds. Defining a dimensionless parameter $h\omega/J$, we found that for a given $\omega$ the minimum value of the normalized ratio $h\omega/J_z$ (where $z = 2d$ is the number of nearest neighbours) is constant at about 12 (Fig. 3c shows the mean values for different driving frequencies and ramp times). While this suggests that the ratio $h\omega/J_z$ is useful for describing the borderline in $J$ below which adiabatic control is possible at constant $\omega$, we also found that changing $\omega$ and $J$ independently for a given $d$ does not always give the same value (as indicated by the error bars in Fig. 3c). As already seen in Fig. 3b, the conditions for adiabatic following depend on $J$ and $\omega$ in a more complicated way. 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 large regions in parameter space for which adiabatic control is possible.

We now turn to the driving-induced superfluid-Mott insulator transition effected through an adiabatic variation of $K_0$. We first loaded a BEC into a 3D lattice of depth $V_0 = 11 E_{\text{rec}}$ and then linearly ramped $K_0$ from 0 to $K_0 = 1.6$. While in an undriven lattice at $11 E_{\text{rec}}$ the BEC is in the superfluid regime with $U/6J \approx 3.5$, the effective Bose-Hubbard parameter $U/6J_{\text{eff}}$ for the driven lattice at $K_0 = 1.6$ is around 7.9, which is larger than the critical value of 5.4 and hence the system is in the Mott insulating phase (see Fig. 4a). In this region, we observe a distinct loss of phase coherence in the interference pattern. When $K_0$ is ramped back to 0, the interference pattern re-appears, proving that the transition was induced adiabatically and that the BEC was not excited by the driving. In order to have a more quantitative indication of the phase transition, we have induced the Mott insulator transition in two different ways: (a) by increasing $V_0$ in an undriven lattice as in and (b) by varying $K_0$ for a fixed lattice depth. In Fig. 4b, one clearly sees that the visibility of the interference pattern vanishes as $U/6J_{\text{eff}}$ is increased (and returns to its original value after ramping $K_0$ back to 0, as indicated by the horizontal dashed lines). The dependence of the visibility on $U/6J_{\text{eff}}$ is the same for methods (a) and (b), strongly indicating that in both cases the same many-body state is reached. The independent control over $J_{\text{eff}}$ also allowed us to measure the excitation spectrum of the system by sinusoidally modulating $J_{\text{eff}}$ (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 Mott insulator regime
we find a gapped spectrum (Fig. 4c).

Our results confirm and extend the role of cold atoms in optical lattices as versatile quantum simulators [23, 24] and open new avenues for the quantum control of cold atoms, thus establishing a link to coherent control in other systems such as molecules in laser fields [25] and Cooper pairs in Josephson qubits [26]. The principles demonstrated here can be straightforwardly extended to more than one driving frequency [27] and to more complicated lattice geometries such as superlattices [28].

Methods

Driven optical lattices

The driven or "spatially shaken" lattices (see Fig. 1a) were realized by mounting the retro-reflecting mirror for each

lattice on a piezo-electric actuator (Queensgate Instruments, model MTP15). These actuators were powered by three phase-locked Stanford function generators producing a sinusoidal signal, the amplitude of which could be controlled between 0 and 10 volts. The response of the mirror-actuator couples had been previously checked in an interferometric setup for the range of driving frequencies used in the experiments (between 3 kHz and 7 kHz). Furthermore, the actuators could be calibrated in situ using two different methods:

1. By observing the interference pattern of a condensate released from a driven one-dimensional lattice after a few milliseconds. We repeated this experiment for increasing values of the driving amplitude until the interference pattern was completely dephased. This amplitude then corresponded to the point where $J_{\text{eff}} = 0$ and hence $K_0 = 2.4$. Previously we had checked that the spatial amplitudes of the oscillations of the actuator-mirror couples were linear in the driving voltage (as measured at the connections of the actuators), so having calibrated the voltage $V$ for which $K_0(V) = 2.4$ we could extrapolate to the other values.

2. By observing the free expansion of a condensate in a driven lattice [14]. The condensate was allowed to freely expand in the lattice direction by switching off one of the dipole traps, and the width of the condensate was observed in situ after a fixed expansion time. In this way, the Bessel-function renormalization of the tunneling parameter $J_{\text{eff}} = J_0(K_0)V$ could be directly measured.

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Competing financial interests

The authors declare that they have no competing financial interests.
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