

Interaction dependent heating and atom loss in a periodically driven optical lattice

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Periodic driving of optical lattices has enabled the creation of novel band structures not realizable in static lattice systems, such as topological bands for neutral particles. However, especially driven systems of interacting bosonic particles often suffer from strong heating. We have systematically studied heating in an interacting Bose-Einstein condensate in a driven one-dimensional optical lattice. We find interaction-dependent heating rates that depend both on the scattering length and the driving strength and identify the underlying resonant intra- and interband scattering processes. By comparing experimental data and theory, we find that for driving frequencies well above the trap depth, the heating rate is dramatically reduced by the fact that resonantly scattered atoms leave the trap before dissipating their energy into the system. This mechanism of Floquet evaporative cooling offers a powerful strategy to minimize heating in Floquet engineered quantum gases.

Introduction.— Floquet engineering, the coherent control of quantum systems by means of time-periodic driving, enables the realization of novel band structures and many-body phases beyond what is possible in static systems [1–12]. It has become an important tool for studies of quantum gases [13], where it e.g. enables the breaking of time-reversal symmetry and thereby the realization of topological bands even for charge-neutral particles [5, 10–12]. In the form of infrared laser pulses, time-periodic driving can give rise to novel effects in traditional condensed matter systems, such as graphene-like systems [14–16] or high-temperature superconductors [17–19]. It also lies at the heart of the recently realized discrete time crystals [20–25].

Despite those recent accomplishments, successfully combining periodic driving with interactions remains a major experimental challenge: In a driven system energy is not conserved, as the system can absorb or emit energy from or into the drive. Therefore, for any fully ergodic driven system, there can only be one steady state, namely the fully mixed density matrix corresponding to an infinite temperature state [26–39]. While this scenario could be avoided by using non-ergodic systems, such as e.g. many-body localized states [40], their use cannot solve the problem in general, as many interesting phases, such as fractional quantum Hall states, are typically ergodic. Therefore, one has to find setups and parameter regimes that allow experimental studies of novel, driven phases on intermediate timescales before the unavoidable heating dominates.

In this work, we experimentally study loss rates of condensed atoms in a driven optical lattice as a function of both driving and interaction strength and can thereby distinguish single-particle from interaction effects. Single-particle heating occurs via discrete single- or multiphoton interband resonances [41] that can easily be avoided. This is in contrast to two-particle processes, which in one- or two-dimensional lattices are always resonant, as collisions can convert arbitrary energies into

transverse excitations [32, 34, 35]. This is in stark contrast to three-dimensional lattices where these processes can be suppressed [28, 32, 34]. We focus on the two experimentally most relevant driving regimes: For low shaking frequencies ω much smaller than the resonance frequency to the first excited band but above the bandwidth of the lowest band, the tunneling matrix element of the lowest band is effectively renormalized by a Bessel function (corresponding to dynamic localization [42, 43]). At the same time, multiphoton resonances are weak as they require many photons. This regime is typically employed for engineering artificial gauge fields [3, 5, 12]. The second regime lies between the two lowest single-photon single-particle resonances. Here, the dispersion relation can acquire two separate minima that can be exploited to study the formation of symmetry-broken domains [4, 6]. We find that for large driving frequencies heating is strongly reduced by the fact that scattered particles with energy $\sim \hbar\omega$ typically leave the trap before dissipating the absorbed energy into the system.

Experimental setup.— We load an almost pure Bose-Einstein condensate (BEC) of about 4×10^5 ^{39}K atoms into the lowest band of a one-dimensional lattice with lattice constant $a = 425$ nm, which is created by interfering two blue-detuned laser beams with a wavelength of $\lambda = 736.8$ nm at an angle of $\theta = 120^\circ$, see Fig. 1(a). Then, we shake the lattice position by periodically modulating the frequency of one of the two laser beams. The atoms feel a periodic inertial force in the frame co-moving with the lattice, which is given by $F_x(t) = -(K/a) \cos(\omega t)$, where we have introduced the driving amplitude K . In order to avoid strong, non-adiabatic excitations to higher bands during the switch-on of the modulation, we continuously ramp the driving amplitude in 10 ms to its desired value. After a variable shaking duration, we determine the heating and losses induced by the drive by measuring the remaining atom number in the BEC. To this end we abruptly stop the drive after an integer number of shaking cycles, immediately followed by

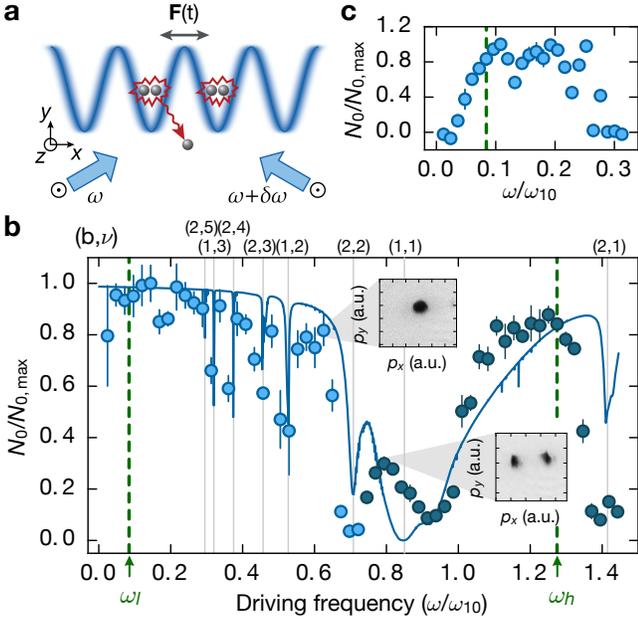


FIG. 1. **Schematic of the experiment and frequency scan.** (a) Two lattice beams with linear out-of-plane polarization intersect at an angle of 120° to form a one dimensional lattice of “pancakes”. By periodically modulating the frequency of one of the lattice beams we can shake the lattice, i.e., modulate its position. (b) Normalized atom number after modulating for 50 ms (for $\omega/\omega_{10} > 0.7$) or 100 ms (for $\omega/\omega_{10} < 0.7$) with variable frequency at a driving strength of $\alpha \simeq 0.9$. Error bars indicate the standard error of the mean from four measurements per data point. The solid blue line shows the theoretically expected single-particle excitations to higher bands. Thin lines mark the resonance positions of single-, and multiphoton transitions to higher bands labeled by (b, ν) . In the frequency region from roughly $0.7\omega_{10}$ to $1.1\omega_{10}$ we observe a splitting of the BEC due to two degenerate minima in the lowest dressed band, which is included in the theory curve. The insets show raw quasimomentum images of the BEC. (c) Zoom into the regime of small shaking frequencies with $\alpha = 2.2$ and 200 ms shaking duration.

bandmapping in the static lattice and 15 ms of time-of-flight (TOF). This TOF is long enough to dilute any thermal background such that we can reliably determine the remaining number of condensed atoms. For all measurements we use a lattice depth of $V_0 = 11.0(3) E_r$, where $E_r = \hbar^2 / (8Ma^2) \approx \hbar \times 7.1$ kHz is the effective recoil energy of this lattice with M being the mass of ^{39}K . In this static lattice, the first interband excitation at zero momentum appears at a frequency of $\omega_{10} = 2\pi \times 41.6(5)$ kHz.

Frequency scan.— Single-particle transitions only occur at specific resonances where the shaking frequency ω fulfills a multiphoton resonance condition $\nu\hbar\omega \simeq \Delta_{b0}(q)$, with ν being an integer and $\Delta_{b0}(q)$ denoting the separation of the lowest band to the b^{th} excited band at a given quasimomentum q . To ensure that we avoid these resonances, we measure the remaining BEC atom number after shaking with variable frequency at a dimensionless driving strength $\alpha \equiv K/(\hbar\omega)$ and scattering

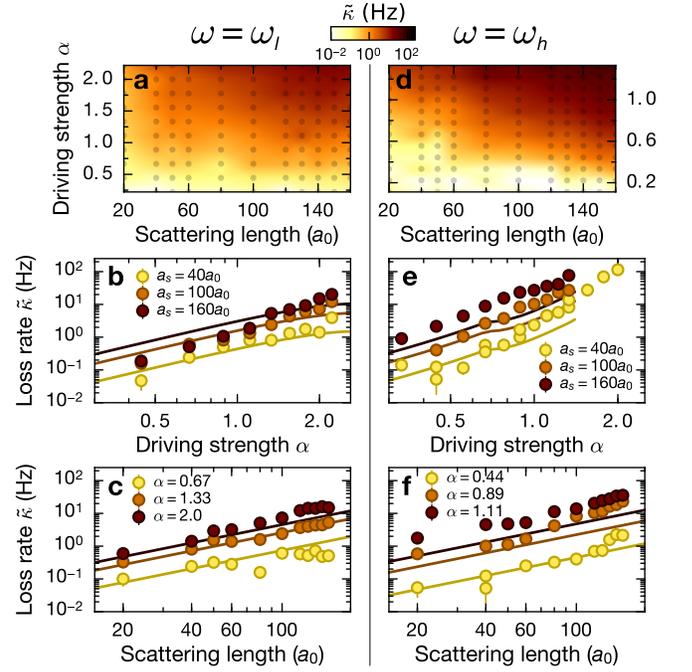


FIG. 2. **Loss rates in the presence of periodic driving.** (a,d) Effective loss rates for different driving amplitudes and scattering lengths. Each dot corresponds to a single lifetime measurement. The shaking frequency is (a-c) $\omega = \omega_l$ and (d-f) $\omega = \omega_h$. (b,e) Crosscuts at fixed scattering lengths. The solid lines correspond to the theoretically predicted scattering rates and error bars indicate fit uncertainties. (c,f) Corresponding crosscuts at fixed driving strengths. Theory lines in (b,c) assume $f\beta\hbar\omega = 10$ to account for the thermalization of scattered atoms (see text).

length $a_s = 60a_0$, with a_0 being Bohr’s radius, using a Feshbach resonance at 400 G [44], see Fig. 1. The solid blue line shows the result of a numerical single-particle simulation assuming a Gaussian width of the BEC in momentum space of $\Delta q = 0.2\pi/a$ (for method see [41, 45]). While the resonances at large frequencies are clearly visible, multiphoton resonances at small driving frequencies are highly suppressed. For our subsequent lifetime measurements we choose $\omega_l = 0.084\omega_{10}$ and $\omega_h = 1.27\omega_{10}$ (green dashed lines) as low and high frequency, far away from all single-particle resonances.

Experimental loss rates.— The total loss rate of condensate atoms in our system is given by summing over background losses in the static system, characterized by a lifetime τ , and heating and losses induced by lattice shaking. We assume that all losses happen on a sufficiently slow timescale such that the system heats up, but stays in global thermal equilibrium, and describe the condensed part using the Thomas-Fermi approximation. We have verified independently that the cloud size indeed shrinks according to the decreasing number of condensed atoms N_0 [45]. Within this approximation, the driving induced loss rate of condensed atoms due to two-particle collisions takes the form $-\kappa N_0^{7/5}$ [45]. Including the back-

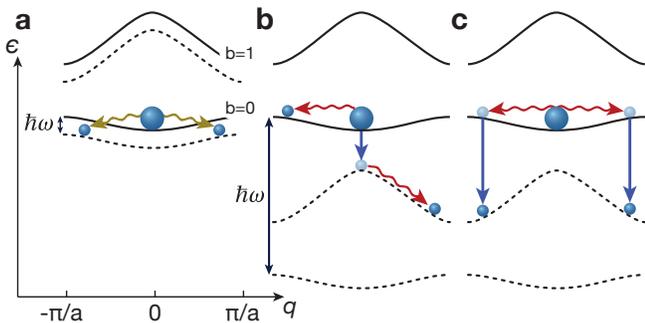


FIG. 3. Examples of two-particle scattering channels. The lowest two bands of a schematic lattice dispersion are sketched by solid lines, Floquet modes shifted by $-\hbar\omega$ ($m = -1$) are depicted by dashed lines. The condensate is represented by a large sphere, scattered particles by small spheres. The pair of yellow wiggly arrows in (a) denotes a two-photon scattering process, where the atoms absorb two photons, while the red wiggly lines in (b,c) denote a zero-photon (ordinary) collision between two atoms and blue arrows describe single-photon interband transitions. (a) When the driving frequency is much smaller than the band gap, the dominant loss process are two-photon intraband collisions. (b,c) For driving frequencies larger than the band gap, the leading (subleading) excitation channels combine one (two) single-photon interband transitions with zero-photon collisions. Note that only one photon number m is associated with the whole system and not one per particle as suggested, for simplicity, in the diagrams.

ground losses $-N_0/\tau$ of the static system we obtain

$$N_0(t) = N_0(0) e^{-\frac{t}{\tau}} \left[1 + N_0(0)^{\frac{2}{5}} \kappa \tau \left(1 - e^{-\frac{2t}{5\tau}} \right) \right]^{-\frac{5}{2}}. \quad (1)$$

We measure τ independently for each scattering length in the static lattice. In order to form a more intuitive quantity, we define a scaled loss rate $\tilde{\kappa} = \kappa N_0(0)^{2/5}$ such that the initial driving-induced losses scale as $\propto \tilde{\kappa} N_0(0)$. As shown in Fig. 2, both stronger interactions and larger driving strengths lead to dramatically higher loss rates for the BEC.

Theoretical description.— In order to identify and estimate the dominant heating channels associated with two-particle scattering, we start by describing a homogeneous system in the Floquet space of time-periodic states, where an integer Fourier index m describes the change in “photon” number relative to a large classical background, i.e., $-m$ counts the number of absorbed photons. In this dressed-atom-like picture, the dynamics is generated by the quasienergy operator Q . Within the subspace of a given relative photon number m , it acts like $Q_{m,m} = H^{(0)} + m\hbar\omega$, whereas the coupling between subspaces m' and m corresponds to an $(m - m')$ photon process which is captured by $Q_{m',m} = H^{(m'-m)}$. Here $H^{(\nu)} = \frac{1}{T} \int_0^T dt e^{i\nu\omega t} H(t)$ denotes the ν^{th} Fourier component of the time-dependent Hamiltonian $H(t)$. The time-averaged Hamiltonian $H^{(0)}$ describes a dispersion relation $\varepsilon_b(k_x) + E_{\perp}(k_y, k_z)$, with effective band struc-

ture $\varepsilon_b(k_x)$ and transverse kinetic energy $E_{\perp} = \hbar^2(k_y^2 + k_z^2)/(2M)$, as well as interactions.

In Fig. 3 we sketch the lowest two bands $b = 0, 1$ for the relative photon numbers $m = 0, -1$ (solid and dashed lines, respectively), given by $\varepsilon_b(k_x) + m\hbar\omega$. The diagrams depict three examples of relevant scattering channels, where two particles (small spheres) are excited out of the condensate into the states $|b, \mathbf{k}\rangle$ and $|b', -\mathbf{k}\rangle$ absorbing $\nu = m - m'$ photons, as indicated by the number of particles transferred to the dashed bands ($\nu = 1$ in b and $\nu = 2$ in a and c). The resonance condition for two-particle excitations can be written as

$$\varepsilon_b(k_x) + \varepsilon_{b'}(-k_x) - \nu\hbar\omega - 2\varepsilon_0(0) = -2E_{\perp} \leq 0. \quad (2)$$

Here, we have separated the transverse kinetic energy $2E_{\perp} = E_{\perp}(k_y, k_z) + E_{\perp}(-k_y, -k_z)$ created in the scattering process on one side of the equation. As there is no lattice potential along these directions, the transverse kinetic energy can take arbitrary non-negative values. As a consequence, in Fig. 3 all states with total energy below the original BEC are accessible.

For the smaller driving frequency ω_l , the tunneling in the lowest band is modified by a Bessel-function $J_{\text{eff}} = \mathcal{J}_0(\alpha)J_0$ [45, 46], where $\mathcal{J}_{\nu}(\alpha)$ is a Bessel function and J_0 the tunneling matrix element of band $b = 0$. In this regime, scattering particles to an excited band would require absorbing a large number of photons. Therefore, the dominant heating channel is *intraband* scattering with small ν . Processes with odd ν are forbidden by symmetry for a condensate with zero momentum, and only acquire finite values due to the momentum spread of the condensate [45]. For a ν -photon scattering process the matrix element scales like $\sim \hbar^2 a_s n_0 J_0 \mathcal{J}_{\nu}(\alpha) / (\nu\hbar\omega M)$ [45], with a prefactor that for odd ν is about five times smaller than for even ν at the measured momentum spread of $0.2\pi/a$ and where n_0 is the condensate density.

For the larger driving frequency, ω_h , *interband* scattering dominates. In this regime, single-photon single-particle interband coupling is strong (with matrix elements $\sim \alpha E_r$ [45]) but off resonant. This leads to a perturbative admixture of states from the first excited band ($b = 1$) with $m = -1$ to the lowest band ($b = 0$) with $m = 0$ and vice versa. As a result of this coupling, already ordinary zero-photon collisions, which are stronger than ν photon scattering processes, give rise to excitations by scattering atoms between these dressed bands. Another consequence of this admixture of the highly dispersive first excited band to the rather narrow lowest band is the formation of a double-well structure within the lowest band for sufficiently large driving strengths α [4, 6, 45]. As a result, the condensate reforms at the new minima of the dispersion at finite quasimomenta, see insets in Fig. 1 (b). We compute the matrix elements for resonant interband excitations using degenerate perturbation theory. In leading order, we encounter three different single-photon ($\nu = 1$) processes, such as the one depicted in Fig. 3 (b), involving a single-particle single-photon interband transition and a

zero-photon two-particle scattering event. Their matrix elements scale like $\sim \hbar^2 a_s n_0 E_r \alpha / (\hbar \omega M)$. The leading correction stems from two-photon ($\nu = 2$) processes, an example of which is shown in Fig. 3(c), giving rise to matrix elements that are a factor of $\alpha E_r / (\hbar \omega)$ smaller [45].

Applying Fermi's golden rule and integrating over the time-dependent Thomas-Fermi profile in a local density approximation, we derive the rates Γ_ν of atoms scattered out of the condensate via ν -photon processes [45] and find that they are proportional to $(a_s N_0)^{7/5}$. For the lower driving frequency ω_l , the scattered particles will not have enough energy to leave the trap and will dissipate their entire energy into the system via ordinary (zero-photon) collisions. This excites additional atoms out of the condensate, leading to a decay rate for condensed atoms $\dot{N}_0 = -f\beta\hbar\omega \sum_\nu \nu \Gamma_\nu / 2$ with inverse temperature β and a numerical factor $f \sim O(1)$ that depends on the details of the system [45]. Due to the finite momentum of these newly created thermal atoms, $\nu = 1$ scattering now becomes dominant and they will absorb photons at an even faster rate than condensed atoms. In a truly closed system, this form of heating would scale linearly with the photon energy $\hbar\omega$. Due to the finite trap depth, however, the system is effectively open and, for the higher driving frequency ω_h , scattered particles typically have sufficient energy to quickly leave the trap without dissipating the absorbed energy. In this regime we expect $\dot{N}_0 = -\sum_\nu \Gamma_\nu$.

Comparison between theory and experiment.—

Due to the thermalization of the absorbed photon energies described above, the measured loss rates of condensed atoms at low driving frequency ω_l will be larger than the total scattering rate $\sum_\nu \Gamma_\nu$ (Fig. 2(a-c)). We observe a factor of $f\beta\hbar\omega \approx 10$, which provides a lower bound for the temperature of the condensate. Assuming for simplicity an ideal homogeneous gas results in a realistic lower bound of 15 nK. While typical temperatures of the BEC will likely be higher [45], the differences are most likely due to resonant scattering of thermal atoms, which is not included in the theory. In contrast, for a driving frequency of ω_h , shown in Fig. 2(d-f), the loss rate of condensed atoms coincides with the total scattering rate since the absorbed photon energy is carried away with the scattered particles leaving the trap. This highlights the advantage of working at larger driving frequencies.

Fig. 2(c) and (f) show a rather good general agreement with the expected scaling with scattering length of $\dot{N}_0 \propto a_s^{7/5}$, demonstrating that the dominant loss mechanisms are indeed interaction driven and that the Thomas-Fermi local-density approximation is consistent with our data. While the data at large driving frequency ω_h follows the theory rather well for not too strong interactions and driving, we observe increasing discrepancies both for larger scattering lengths and larger driving strengths. When the mean free path of excited atoms ($\propto 1/a_s^2$) becomes on the order of the size of the BEC, excited atoms will undergo additional collisions while leav-

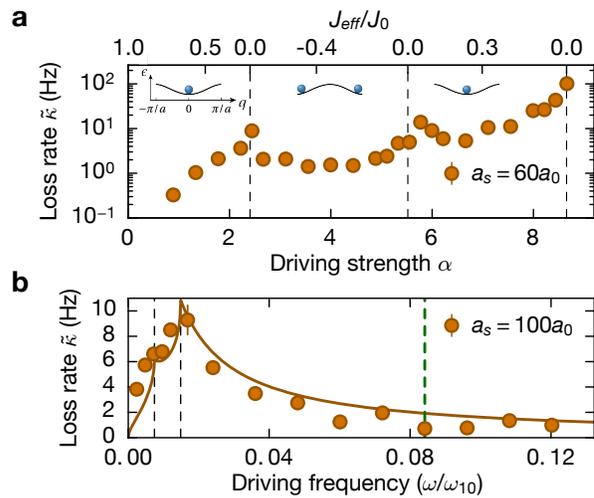


FIG. 4. **Loss rates for large driving strengths and small frequencies** (a) When scanning the driving strength α at frequency ω_l , we observe peaks in the loss rate whenever the effective tunneling $J_{\text{eff}} = J_0 \cdot \mathcal{J}_0(\alpha)$ goes through zero (dashed lines). The insets sketch the lowest band for positive and negative tunneling. (b) For a fixed driving strength $\alpha = 1.1$, we observe a peak in the loss rate for driving frequencies close to the bandwidth of the lowest band. The solid line shows the theory scaled by $f\beta\hbar\omega$ with a temperature of 15 nK (see text). The dashed black lines indicate one (two) times the effective bandwidth, below which the number of accessible states for two (single) photon scattering becomes reduced [45]. The dashed green line marks the frequency ω_l . Error bars indicate fit errors.

ing the cloud, thereby giving rise to an additional loss of condensate atoms similarly to the low frequency case [45]. This is most clearly visible when plotting the data vs scattering length, see Fig. 2(f), where for weak driving a discrepancy to the $a_s^{7/5}$ scaling can be observed for scattering lengths larger than $\simeq 100 a_0$. For higher driving strengths we expect the onset of additional scattering channels with $\nu > 2$, which are not included in the theory for ω_h [45]. We note that two degenerate minima appear in the lowest band for $\alpha > 0.7$, giving rise to the small kinks in the expected loss rates in Fig. 2(e).

We also measured the loss rates for large driving amplitudes and low frequency ω_l , see Fig. 4(a). We can observe clear maxima in the loss rate whenever the effective tunneling matrix element J_{eff} is close to zero and attribute them to zero-photon scattering in the effectively flat band. Interestingly, the loss rate decreases again once the sign of the effective tunneling matrix element J_{eff} changes. Fig. 4(b) shows the loss rate for various frequencies close to the bandwidth. Since there are fewer modes available for frequencies below the bandwidth of the lowest band, a clear decrease in the loss rates can be observed [45].

Conclusion and Outlook.— We have measured the loss rates of an interacting BEC in a driven one-dimensional optical lattice. They approximately scale

with the interaction as $a_s^{7/5}$, in agreement with a theoretical description based on a Thomas-Fermi approximation and Fermi's golden rule. For large driving frequencies, scattered particles can leave the trap and carry away the absorbed energy quanta $\hbar\omega_h$. This mechanism of continuous Floquet evaporative cooling can act as a powerful general strategy to reduce heating rates in Floquet engineered quantum gases. Furthermore, the two-particle scattering processes considered here rely on exciting transverse motion and might therefore be absent in a three-dimensional lattice. Another intriguing possibility is the use of non-ergodic or many-body localized systems, where the dynamics can be immune to these

heating processes.

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