Strongly inhibited transport of a 1D Bose gas in a lattice

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We report the observation of strongly damped dipole oscillations of a quantum degenerate 1D atomic Bose gas in a combined harmonic and optical lattice potential. Damping is significant for very shallow axial lattices (0.25 photon recoil energies), and increases dramatically with increasing lattice depth, such that the gas becomes nearly immobile for times an order of magnitude longer than the single-particle tunneling time. We emphasize, and discuss further below, that the inhibited transport is not due to Bloch oscillations [1, 2], where transport is frustrated by Bragg reflection at the Brillouin zone (BZ) boundary, as has been seen in previous experiments [1, 2, 6]. Our method to realize an ensemble of independent 1D Bose gases is similar to earlier work [12]. We produce a nearly pure 87Rb condensate of \( N = (0.8 \pm 1.6) \times 10^5 \) atoms in the \( |F = 1, m_F = -1\) state in a Ioffe-Pritchard magnetic trap \( (\nu_x = 29 \text{ Hz}, \nu_y = 8 \text{ Hz}) \). We next partition the BEC into an array of independent, vertical 1D “tubes” by adiabatically applying a transverse (in the \( xy \) plane) 2D confining lattice [12]. The confining lattice is ramped on during 200 ms to a depth of approximately \( 30E_R \) (where \( E_R = \hbar^2/2m\lambda^2 \) is the photon recoil energy, and \( \lambda \) is the laser wavelength). The combined magnetic and optical potential results in approximately 5000 occupied tubes, each with an axial frequency of \( \omega_0/2\pi \approx 60 \text{ Hz} \). We observe a Thomas-Fermi density envelope in the combined magnetic and optical potential, and calculate [12] cloud radii of \( r_x = 14(1) \mu\text{m} \), \( r_y = 20(1) \mu\text{m} \), and \( r_z = 10.6(5) \mu\text{m} \) for \( N = 1.4 \times 10^5 \). From this we estimate a peak 1D density of \( 4.8(4) \times 10^4 \text{ cm}^{-1} \) in the central tube, and a peak 3D density of \( 4.7(4) \times 10^{14} \text{ cm}^{-3} \). Subsequently, we corrugate the tubes by adiabatically applying, over 20 ms, an axial (vertically along \( z \) ) 1D lattice. The Rayleigh length of the axial lattice beams is large enough that they do not significantly modify the axial harmonic potential. All lattice beams derive from a single Ti:Sapphire laser operating at \( \lambda = 810 \text{ nm} \), far detuned from the atomic resonances at 780 nm and 795 nm. The pairs of lattice beams are detuned from each other by 6 MHz, making them effectively independent [22]. The final configuration consists of three independent standing waves, each formed from a pair of counter-propagating beams.

We excite dipole oscillations of the center-of-mass of the atoms in all the tubes by suddenly \( (\lesssim 150 \mu\text{s}) \) applying a linear magnetic field gradient, thus displacing...
the total harmonic trap (but not the lattice) axially by $z_0 \approx 3 \mu m$. This displacement is less than 30% of $r_z$, and corresponds to approximately eight axial lattice sites \(24\). The waists of the Gaussian transverse lattice beams \(w_0 \approx 210 \mu m\) are much larger than both $z_0$ and the size of the trapped cloud.

The oscillation in the position of the atoms is too small for our imaging system to clearly resolve. We therefore observe oscillation in velocity by waiting a variable time \(t_w\) after the initial displacement, then suddenly turning off all trapping potentials (with time constants of $\approx 250 \mu s$ and $\approx 150 \mu s$ for the optical and magnetic potentials, respectively), and imaging the atoms after a time-of-flight \(t_{TOF} = 18.4 \text{ ms}\). The turn-off of the optical lattice is fast compared to the oscillation period, but slow enough to avoid diffraction of the atoms (i.e., adiabatic with respect to band excitations).

We observe damped dipole oscillations for axial lattice depths from $V = 0E_R$ to $2E_R$, as seen in Fig. 1a. In the absence of an axial lattice, we observe oscillations (period \(T = 15.4 \text{ ms}\)) consistent with no damping (Fig. 1a), indicating that tube-to-tube dephasing and trap anharmonicities are not significant on the timescale of our experiments. However, the oscillations are noticeably damped in a lattice only \(0.25E_R\) deep. Such a shallow lattice modulates the atomic density by only 6%, and modifies the single-particle energy-quasimomentum dispersion relation \(E(q)\) from that of a free-particle around only the last few percent of the BZ. (We note here, and discuss further below, that the amplitude of motion is kept well within the quadratic part of \(E(q)\) for shallow lattices).

Beyond a lattice depth of $3E_R$ the motion is overdamped, and there are no oscillations. In this case, the atoms’ velocity can be quite small, so we use a technique that maps the atoms’ position in the trap to the cloud position after TOF. The experiment proceeds as before, except that after the trap is displaced by \(z_0\), the atoms are allowed to relax toward their equilibrium position at \(z = 0\) for a fixed time \(t_w = 90 \text{ ms}\). We then rapidly (with time constant $\approx 250 \mu s$) turn off only the axial lattice. The remaining transverse lattice and magnetic potentials are left on for 3.75 ms (approximately a quarter-period of undamped axial harmonic motion), then turned off simultaneously (as in the underdamped experiment). This converts the axial displacement \(z(t_w)\) into a velocity, which we measure by TOF.

Figure 2a shows \(z(t_w = 90 \text{ ms})\) as a function of axial lattice depth. For the shallowest lattices this wait time is sufficient for the atoms to damp to the equilibrium position \(z = 0\). For the deepest lattices the motion is so overdamped that there is negligible motion during this time, and the position remains \(z \approx z_0\). We note that in the absence of damping, atoms would tunnel through the lattice to the equilibrium position in a time \((T/4)\sqrt{m^*/m}\), where \(m^*\) is the effective mass. This time is only 8 ms for non-interacting particles in a \(10E_R\) lattice.

To quantify the damping we model the motion as damped simple harmonic, \(m^* \ddot{z} = -b \ddot{z} - k \dot{z}\), and extract a damping constant $b = b(V)$ for different axial lattice depths $V$. For underdamped motion, we simultaneously fit the oscillation data for 8 depths to the expression

$$\dot{z}(t) = A \frac{k}{\omega m^*} e^{-bt/2m^*} \sin(\omega t),$$

where $\omega \equiv \sqrt{k/m^* - (b/2m^*)^2}$, and $A$ and $k$ are fit parameters common across all $V$. For overdamped motion, we determine $b = b(V)$ from the overdamped solution for $z(t_w)$, which in the limit of strong damping simplifies to

$$z(t_w) \approx z_0 \left( \frac{e^{-bt_w/b}}{1-km^*/b^2} + \frac{e^{-bt_w/m^*}}{1-b^2/km^*} \right),$$

where $z_0$ and $b$ are inputs derived from measurements of undamped oscillations. In our analysis we use a single-particle calculation of the effective mass $m^*$ $\text{22}$. Figure 2b shows a plot of $(b(V))/b_0$ versus lattice depth $V$, where $b_0 \equiv 2m\omega_0$ corresponds to critically damped harmonic motion for $\omega_0 \equiv \sqrt{k/m}$. We show data for both the under- and overdamped regimes, and note that the damping constant increases by at least a factor 1000 for a 30-fold increase in lattice depth.

The axial width of the cloud after TOF can provide information about the distribution of atomic quasimomenta in the lattice. The lattice turn-off time constant of 250 $\mu s$ is long enough to avoid diffraction, but short
enough to be non-adiabatic with respect to inter-well tunneling and interactions. (Related experiments in 2D [21] and 3D [24] lattices support this conclusion.) In the absence of interactions, and neglecting the initial size of the cloud, the turn-off maps the single-particle quasimomentum distribution of atoms in the lattice to free particle (i.e., plane-wave) momentum states that can be directly observed in TOF. In the presence of interactions, the mapping is complicated by mean-field repulsion during TOF. A variational calculation [27] indicates that mean-field repulsion is in fact the dominant contributor to the axial TOF width in our system. Therefore, the extracted TOF width greatly overestimates the width of a narrow initial quasimomentum distribution.

An example TOF image is shown in Fig. 4, together with cross-sectional profiles of the optical depth along the axial and transverse directions. The first BZ of the transverse lattice is uniformly filled, producing a uniform, square spatial distribution (in the \( xy \) plane) after TOF. Our imaging system views this square distribution along the diagonal in the \( xy \) plane, resulting in a triangular profile, from which we extract the width of the square (open symbols, Fig. 3b). In the axial direction the distribution is narrower and reasonably well fit to a Gaussian, from which we extract the axial (along \( z \)) TOF 1/e half-width (filled symbols, Fig. 3b). Even for the strongly overdamped data, the axial TOF width (which, we recall, overstates the width of narrow quasimomentum distribu-

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**FIG. 2:** Overdamped motion of a 1D Bose gas in an optical lattice. (a) Plot of the atoms’ position 90 ms after shifting the trap, as a function of lattice depth. Inset depicts relaxation of the atoms toward equilibrium. Immediately after the trap is displaced to \( z_0 \), the atoms (open circle) begin to move toward equilibrium, reaching a displacement \( z \) from equilibrium after 90 ms (solid circle). Also shown is the initial position of the atoms (dash-dot line). (b) The 1/e half-widths of Gaussian fits to axial TOF distributions (closed symbols), and the half-widths of square transverse TOF distributions (open symbols) resulting from a uniformly filled BZ (see text, and Fig. 2b). Also shown (dashed line) is the BZ calculated from lattice parameters.

**FIG. 3:** Plot of the reduced damping constant \( b/b_0 \) for various depths \( V \) of the axial lattice, as determined using the underdamped (squares) and overdamped (circles) experimental techniques. Near critical damping (\( V \approx 3E_R \)), the analysis cannot distinguish between underdamped and overdamped motion, so only upper and lower bounds are shown at \( V = 2.25 \) and 2.50.

**FIG. 4:** Cross-sections (a) of a TOF absorption image (b) of the expanded atom cloud. The plot of the transverse cross-section (along \( x \)) is offset vertically by 0.5 units for clarity. The solid lines are Gaussian (triangular) fits to the axial (transverse) profiles, respectively. The dashed ovals in (b) indicate the peak-to-peak range of motion of undamped dipole oscillations.
tions), is much narrower than the BZ. This implies that the inhibition of transport is not due to effects related to Bloch oscillations of a filled BZ, as observed in Refs. [4, 6]. Furthermore, we do not see a significant difference in TOF width between atoms that undergo damped harmonic motion and those that are unexcited but held for an equal time [28]. This is in stark contrast to earlier experiments on 3D BECs [8, 11], where strong damping was accompanied by a pronounced broadening and fragmentation of the quasimomentum distribution.

Large amplitude dipole oscillations in a lattice can damp due to dynamical instabilities caused by particle interactions. For a 3D BEC moving in an optical lattice, such an instability point occurs at \( q_\pi/2 \) (\( q_\pi \equiv 2\pi/\lambda \) is the BZ boundary), where the dispersion relation has an inflection point, as predicted in [10, 11], and observed in [11]. This effect is manifested as a large increase in the width of the quasimomentum distribution. Here, in contrast, we keep the maximum (single-particle) quasimomentum \( q_{\text{max}} \) of the oscillation small by limiting the initial energy of displacement \( E(z_0) = \hbar \omega_0^2 z_0^2/2 \). For \( V < 2E_R \), our choice of \( z_0 \) corresponds to \( q_{\text{max}} \approx q_\pi/5 \). For deeper lattices our fixed \( z_0 \) corresponds to a larger \( q_{\text{max}} \), but is always less than \( q_\pi/2 \) for \( V < 9E_R \). (In a separate experiment, we excited oscillations in our system with twice the usual amplitude, and saw stronger damping that was accompanied by a broadening of the axial TOF width by nearly a factor of two.)

For small amplitude dipole oscillations (\( q_{\text{max}} \ll q_\pi/2 \)) of a 3D BEC in an optical lattice, the effect of the lattice is merely to increase the effective mass, leading to undamped motion at a lower frequency [6, 11]. In the reduced dimensionality system of our 1D Bose gas, we have seen that the optical lattice has a qualitatively different effect. To highlight the difference between these two situations, we excited dipole oscillations in a 3D BEC (i.e., no transverse confining lattice) in a \( 4E_R \) axial lattice, and saw no damping. This is in contrast to the results of the same experiment in a 1D Bose gas (i.e., with a transverse confining lattice), shown in Fig. 3 where \( b(V = 4E_R)/b_0 \approx 50 \) corresponds to extremely overdamped motion.

After we performed these experiments, theoretical treatments appeared which suggested that zero-temperature quantum fluctuations can lead to substantial damping of transport in a 1D atomic Bose gas [29, 30]. Our observations, including the significant damping in lattices too shallow to support a Mott-Insulator (MI) phase [41], can be explained by the mechanisms of Refs. [29, 30], but appear to be inconsistent with a mechanism involving incompressibility [13].

It is possible that there is a temperature dependence to the damping; unfortunately, we can derive little information on temperature from the TOF widths. Future experiments to investigate the temperature dependence could shed light on the relative importance of quantum and thermal fluctuations to dissipation, a question of interest in, for example, the development of ultrathin superconducting wires [17]. We also look forward to testing other explicit predictions of these theories, such as the dependence of damping on displacement and dimensionality. We note that the periodic potential of an optical lattice is free from defects, and that 1D atomic Bose gases are well isolated from the environment, yielding a relatively clean system in which to compare experiment with theory. The ability to continuously and dynamically vary the confining potentials makes optical lattice experiments attractive for future studies of superfluidity in low dimensional quantum systems.

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