An ytterbium quantum gas microscope with narrow-line laser cooling

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Abstract

We demonstrate site-resolved imaging of individual bosonic $^{174}$Yb atoms in a Hubbard-regime two-dimensional optical lattice with a short lattice constant of 266 nm. To suppress the heating by probe light with the $^{1}S_0^0-^{1}P_1$ transition of the wavelength $\lambda = 399$ nm for high-resolution imaging and preserve atoms at the same lattice sites during the fluorescence imaging, we simultaneously cool atoms by additionally applying narrow-line optical molasses with the $^{1}S_0^0-^{3}P_1$ transition of the wavelength $\lambda = 556$ nm. We achieve a low temperature of $T = 7.4(13) \mu K$, corresponding to a mean oscillation quantum number along the horizontal axes of 0.22(4) during the imaging process. We detect, on average, 200 fluorescence photons from a single atom within a 400 ms exposure time, and estimate a detection fidelity of 87(2)%. The realization of a quantum gas microscope with enough fidelity for Yb atoms in a Hubbard-regime optical lattice opens up the possibilities for studying various kinds of quantum many-body systems such as Bose and Fermi gases, and their mixtures, and also long-range-interacting systems such as Rydberg states.

1. Introduction

Ultracold quantum gases in optical lattices have proven extremely useful for the study of quantum phases and the dynamical evolutions of strongly correlated many-body system described by a Hubbard model [1]. Well-known examples include a quantum phase transition from a superfluid (SF) to a Mott insulator (MI) for bosonic species [2--4], and a crossover from a metal to a MI for fermionic species [5, 6]. To fully exploit the potential of ultracold atoms in an optical lattice as a quantum simulator, it is a great advantage to have access to the in-trap atom distribution with single-atom resolution. In 2009, a quantum gas microscope (QGM) was realized with bosonic $^{87}$Rb for the first time [7]. Site-resolved imaging has been used to study the SF to MI transition [8, 9], and strongly correlated dynamics in optical lattices [10--12]. Site-resolved imaging systems have been realized for other atomic species such as fermionic $^{40}$K [13--15] and $^6$Li [16--18] very recently.

Extending the applicability of a QGM technique to atomic species beyond alkali-metal atoms is an important step for a further revolution. In particular, a successful application of a QGM technique to two-electron atoms such as alkaline-earth metal and ytterbium (Yb) atoms is remarkable because it offers many unique possibilities for quantum simulation and quantum information research. In fact, recent studies demonstrate that a system of two-electron atoms in an optical lattice is an ideal platform for the study of SU($N$) physics [19--22], two-orbital SU($N$) physics [23--26], and topological physics [27]. In addition, a variety of stable isotopes, 5 bosonic and 2 fermionic isotopes in the case of Yb atoms for example, enable us to study various kinds of many-body systems such as ultracold Bose and Fermi gases and Bose--Bose [28, 29], Bose--Fermi [29--31], and Fermi--Fermi [20] mixtures in an optical lattice. The existence of nuclear spin degrees of freedom in the ground state $^1S_0$ and long-lived metastable states $^3P_0$ and $^3P_2$ offers unique possibilities for quantum memory and quantum computation [32--35]. Additionally we can tune interatomic interactions between the $^1S_0$ and $^3P_2$ states by an anisotropy-induced magnetic Feshbach resonance [36]. Furthermore, a high-resolution laser spectroscopy of atoms in an optical lattice using the ultranarrow $^1S_0^0-^3P_0$ and $^1S_0^0-^3P_2$ optical transitions is also demonstrated both for bosons and fermions, revealing the novel behaviour of the atomic interaction of the system [24--26, 36, 37]. There has been also considerable interest in high-lying Rydberg states of two-electron atoms [38, 39] in an optical lattice.
because of an additional degrees of freedom for probing and manipulation provided by the remaining valence electron of a singly excited Rydberg state. The successful application of a QGM technique to these systems will determine electron of a singly excited Rydberg state. The successful application of a QGM technique to these systems will de

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Our experiment starts from loading \( ^{174}\text{Yb} \) atoms into a magneto-optical trap (MOT) in a metal chamber and transferring the atoms into a glass cell (Schott AG BOROFLOAT) using an optical tweezer (OT). The detail of the MOT and OT setup is described in [43]. The position of atoms in the glass cell is about 5.5 nm below the surface of the glass cell, and just 6.23 mm under a high-resolution objective with numerical aperture of NA = 0.75 (Mitutoyo G Plan Apo HR50x (custom)), which is schematically shown in figure 1(b). After creating a Bose–Einstein condensate (BEC) of \( 5 \times 10^4 \) atoms after 10 s evaporative cooling in a crossed optical trap formed by the OT beam and another 532 nm beam, we load the BEC into a vertical lattice generated by the interference of two laser beams with the wavelength of \( \lambda = 532 \text{ nm} \) propagating at a relative angle of \( \alpha = 62.9^\circ \). The vertical lattice has a spacing of \( \lambda / (2 \sin (\alpha / 2)) = 4.9 \mu\text{m} \) and the trap frequency along the vertical axis (z-axis) of \( \omega_z = 2\pi \times 2 \text{ kHz} \) at this loading stage, as explained in detail in our previous work [44].

The atoms, just after being loaded into the vertical lattice, spread over several, typically three, layers, as shown in the left panel of figure 2(a). In this situation, although we can focus on the atoms in one selected layer with an objective depth of less than 1 \( \mu\text{m} \), we always have contributions from the atoms in other layers which considerably blur the image. To observe a clear image of the atoms in only one layer, we blow away the atoms trapped in the 1\textsuperscript{st} and 3\textsuperscript{rd} layers by alternately exciting the atoms into \( ^3\text{P}_2 \) state followed by the rapid inelastic collision decay and repumping back them into \( ^1\text{S}_0 \) state. The blast time is typically 250 ms. (c) Low-lying energy levels of \( ^{174}\text{Yb} \) atoms relevant for blast.

2. Experimental setup and atom preparation

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Finally, we load the 2D atom cloud in a single layer into horizontal 2D optical lattices (x- and y-axes) by simultaneously ramping up the potential of the vertical and horizontal optical lattices, where the wavelength of horizontal lattices is \( \lambda = 532 \text{ nm} \) and the lattice spacing is 266 nm. The vibrational frequencies along the three
axes at the fluorescence imaging stage are \( (\omega_x, \omega_y, \omega_z) = 2\pi \times (300, 300, 15.7) \, \text{kHz} \), corresponding to the lattice depths \( (U_x, U_y, U_z) = k_B \times (300, 300, 250) \, \mu\text{K} \), respectively.

3. Narrow-line laser cooling in an optical lattice

An important prerequisite for realizing a QGM is to preserve the atoms at their sites during fluorescence imaging. The \(^1\text{S}_0 \rightarrow ^3\text{P}_1\) transition provides high-resolution imaging with the diffraction limited resolution of a FWHM of 274 nm in our system. However, the high Doppler cooling limit temperature \( T_D = 690 \, \mu\text{K} \) makes quite difficult to preserve atoms at their sites. In addition, the lack of hyperfine structure in the ground state \(^1\text{S}_0\) of bosonic Yb atom makes impossible to apply sub-Doppler cooling techniques such as polarization-gradient cooling and Raman sideband cooling. To resolve this difficulty, we simultaneously cool the atoms with Doppler and sideband cooling using the \(^1\text{S}_0 \rightarrow ^3\text{P}_1\) narrow-line transition.

To efficiently cool all the atoms in an optical lattice with the narrow-line transition \(^1\text{S}_0 \rightarrow ^3\text{P}_1\), we need to suppress the inhomogeneity of light shift between the \(^1\text{S}_0\) and \(^3\text{P}_1\) states, otherwise the detuning for cooling is not optimized simultaneously for all atoms. Here the light shift in the ground state \(^1\text{S}_0\) is given by \( \Delta E_L = -(1/4)\alpha \rightangle I \), where \( I \) is the laser intensity of the wavelength \( \lambda = 532 \, \text{nm} \) and \( \alpha = 37.9 \, \text{Hz/(Wcm}^{-2}\text{)} \) is the calculated scalar polarizability in the \(^1\text{S}_0\) state. The light shift in the magnetic sublevel \( m_f \) of the \(^3\text{P}_1\) state is given as \(^{[45]}\)

\[
\Delta E_L^m = -\frac{1}{4} \alpha_s(m_f, \theta) I,
\]

where

\[
\alpha_s(m_f, \theta) = \alpha_s^s - \alpha_s^\perp \frac{1 - 3 \cos^2 \theta}{2} (3m_f^2 - 2).
\]

Here \( \theta \) is the angle between the quantization axis and the polarization of laser beams, \( \alpha_s^s \) and \( \alpha_s^\perp \) are the scalar and tensor polarizabilities in the \(^3\text{P}_1\) state, respectively. Importantly, equation (2) provides the possibility of tuning the polarizability \( \alpha_s(m_f, \theta) \) to coincide with \( \alpha_s \) by choosing an appropriate angle \( \theta \) for \( m_f \), thus canceling the light-shift effects of the \(^1\text{S}_0 \rightarrow ^3\text{P}_1\) transition. For this possibility we perform a laser spectroscopy with the \(^1\text{S}_0 \rightarrow ^3\text{P}_1\) transition for various laser intensities and angles \( \theta \), and accurately determine \( \alpha_s^s \) and \( \alpha_s^\perp \) as \( 22.4(2) \) and \(-7.6(1) \, \text{Hz/(Wcm}^{-2}\text{)} \), respectively. With these values, our current setup of the polarizations of all the lattice beams parallel to the vertical axis provides \( \alpha_s(m_f = 0, \theta = 0)/\alpha_s = 0.99 \). In our experiment, however, we slightly tilt a magnetic field from the vertical direction by an angle \( 6.1^\circ \) which gives \( \alpha_s(m_f = 0, \theta = 6.1^\circ)/\alpha_s = 0.98 \). This setup enables us to excite the atoms into the \(^3\text{P}_1(m_f = 0)\) state, when the polarizations of the 556 nm cooling light along the horizontal axes are set to vertical, and those along the vertical axis horizontal (see figure 1(b)). Note that the light shift of the \(^1\text{S}_0 \rightarrow ^3\text{P}_1\) transition for probing is smaller than the natural linewidth of this transition of 29 MHz, and so it is not a problem. The total intensities of 399 nm and 556 nm beams correspond to the saturation parameters of \( s_{399} \sim 1 \times 10^{-3} \) and \( s_{556} \sim 1 \), respectively. With this dual molasses, Moiré patterns of about 6 \( \mu\text{m} \) pitch are observed as a result of the interference between the cooling molasses beam of 556 nm and the optical lattice of 532 nm. To erase this unwanted Moiré pattern, we modulate the phase of the standing wave of the 556 nm optical molasses by modulating retro-reflecting mirrors via the attached Piezo transducers, as explained in detail in our previous work \(^{[44]}\).

The fine-tuning of the relative angle between a magnetic field and lattice laser polarizations indeed gives us a reasonably narrow resonance of the \(^1\text{S}_0 \rightarrow ^3\text{P}_1(m_f = 0)\) transition for atoms in the optical lattice during fluorescence imaging. Figure 3(a) shows the spectra of atoms in our deepest horizontal optical lattices of \( U_x = U_y = 1500 \, E_B \) where \( E_B = h^2/(2 \, m \lambda^2) = k_B \times 200 \, \text{nK} \) is recoil energy of lattice beam. The top panel shows the fluorescence counts of 399 nm probe molasses light as a function of a frequency of 556 nm cooling molasses beams along the horizontal axes, in which we simultaneously apply the probe light and weak cooling molasses lights, \( s_{399} \sim 1 \times 10^{-3} \) and \( s_{556} \sim 1 \), and we can observe many fluorescence counts during a 400 ms exposure time only when the cooling is efficient at favorable detunings. We obtain the optimal frequency \( f_{bg} = -337(18) \, \text{kHz} \) with the width of 318(12) kHz (FWHM). The bottom panel shows the optical density measured by absorption imaging with a 556 nm beam irradiated along the horizontal axis as a function of a frequency of the 556 nm probe light. In this measurement we do not apply 399 and 556 nm molasses beams. Note that we set a zero frequency detuning as the resonance frequency of this spectrum. We determine the optimal detuning of cooling beam along the horizontal axes \( s_{556}/2\pi = f_{bg} = -337(18) \, \text{kHz} \). The same measurements are done at several horizontal lattice depths, as shown in figure 3(b). In our lattice system, Lamb-Dicke parameters are \( n_f = \sqrt{h \lambda^2/(2 \, m \omega_p)} = (0.11, 0.11, 0.48) \), where \( \lambda \) is a wavevector of 556 nm light, and \( \mu = x, y, z \). Although the frequency separation between the cooling sideband \( f_{bg} \) and the carrier transition \( f = 0 \), corresponding to the trap frequencies \( \omega_x \) and \( \omega_z \), is not large enough compared with the natural linewidth of 184 kHz for the 556 nm cooling transition, as shown in figure 3(a), the responsible cooling mechanism along the
x- and y-axes should be sideband cooling, because the optimal detuning of cooling beam along the horizontal axis depends on the horizontal lattice trap depth and is consistent with the trap frequency along the horizontal axis, as shown in figure 3(b). We note that the shape of the resonance does not change for all the horizontal lattice potential depths of $E_x = 250$ to $1500 E_R$ studied in this measurement.

The temperature of the atoms during the fluorescence imaging is accurately measured by laser spectroscopy using the ultranarrow transition $^1S_0 - ^3P_0 j = 3/2$ transition after cooling by sideband (horizontal axis) and Doppler (vertical axis) cooling of 556 nm. The line shows fit to the data. The ratio of the red to the blue sideband peaks $S_y/S_x = 0.32(6)$, and the mean oscillation quantum number along horizontal axis $\langle n \rangle = 0.22(4)$, corresponding to the temperature along horizontal axis $T_H = 7.4(13)$ $\mu$K. The error in the determination of the mean vibrational occupation number $\langle n \rangle$ comes from a fitting error of the peak heights of red- and blue-sidebands ($S_y$ and $S_x$).

The unique feature of our scheme is the separation of the cooling and probing processes during the fluorescence imaging. We can therefore study the effect of the cooling beams alone. Here we study the temporal evolution of the temperature with narrow-line laser cooling, especially at the early stage of cooling, to investigate what happens during the cooling process. This is especially interesting because the atom loss rate for the light-assisted collision associated with the $^1S_0 - ^3P_1 (m_J = 0)$ states is small [49] and the atoms in multiply-occupied sites would be heated without loss, differently from the case of alkali atoms. In this measurement, the temperature is measured by a time-of-flight method with absorption imaging. To correctly estimate the temperature along the horizontal axis, we numerically calculate the size of the atom cloud after a time-of-flight, assuming an initial...
Boltzmann occupation of the each vibrational level and a ballistic expansion of the cloud. Figure 4(a) shows the results of the measurements. The temperatures rapidly increase within several milliseconds followed by the rather slow decrease towards the steady-state value obtained by the ultranarrow line laser spectroscopy of figure 3(c). This behaviour is explained as an effect of a light-assisted collision due to the near-resonant cooling light. Namely, atoms in multiply-occupied sites should be heated by release of the kinetic energy subsequent on a light-assisted collision. This is confirmed by further measurements with applying a photoassociation (PA) pulse for removal of multiply-occupied sites before imaging, shown in figure 4(b). In spite of the fluctuation of data, it is clear that the behaviour of the temperatures of figure 4(b) is different from that of figure 4(a). Although all the following single-site resolved imaging data presented in this paper are measured without the application of PA light, this initial heating effect is negligible because the multiply-occupied sites are almost absent in sparse atomic samples used for our QGM measurement.

4. Site-resolved imaging

We image the atomic fluorescence onto the EMCCD camera (Andor iXon EM Blue). In figure 5(a) we show one illustrative example of the obtained images. Note that, just before the fluorescence imaging, we intentionally select only about 2% of the atoms for easier evaluation of the performance of the QGM. Such dilution of the atoms is done by performing a weak excitation with the 1S0 → 3P2 transition, and then returning the atoms back into the ground state 1S0. Figure 5(b) shows our measured point spread function (PSF), obtained by averaging over 10^4 fluorescence images of individual atoms. We find that our PSF can be well approximated by a double Gaussian:

$$\text{PSF}(r) = A \left[ \exp \left( -\frac{r^2}{2\sigma_1^2} \right) + B \exp \left( -\frac{1}{2} \left( \frac{r - r_0}{\sigma_2} \right)^2 \right) \right] + C$$

with widths \(\sigma_1, \sigma_2\), main and relative amplitudes \(A, B\), a spatial offset \(r_0\), and an overall count offset \(C\). The fit result shows our PSF is well described with \(\sigma_1 = 154(1)\) nm and \(\sigma_2 = 153(10)\) nm, and also has a FWHM of 364 nm, and we detect on average 200 photons per atom within 400 ms fluorescence time. In our imaging system, the spherical aberration remains, making the resolution of PSF worse than ideal one. Our system has a total fluorescence collection efficiency of 6.0%, given by the objective’s solid angle of \(\Omega/4\pi = 17\%\), 51% total transmission through the imaging optics, and quantum efficiency of 70% of our EMCCD camera. The corresponding atomic fluorescence rate of \(\sim 8300\) photons/s is large enough to unambiguously identify the presence or absence of an atom for each lattice site. Note that the maximum number of detected photons is limited by the cooling rate of the narrow transition [44]. Although the resolution of measured PSF is about 1.4 times larger than the lattice spacing, we successfully determine the atomic distribution by deconvolution of images.
For reconstructing an atom distribution from our obtained images, we first determine a lattice angle and spacing from the isolated, single-site resolved signals [8]. Our lattice axes are oriented approximately along the vertical and horizontal axes with respect to the images. The histogram of the mutual distances in the coordination rotated by a small angle $\phi$ is shown in figure 5(c) and (d). We fit a periodic array of Gaussians to the observed histogram. Figure 5(e) shows the width of the Gaussians in the histogram as a function of a coordinate rotation angle. The red dotted line is a fit with a double Gaussian of equation (3) and yields a minimum width at rotation angle of $\phi = 0.482 \pm 0.01$ degrees. The histogram is obtained at the coordinate rotation angle of 0.482(3)$^\circ$ and the lattice constant of 2.66(1) pixels on CCD plane corresponding to the lattice constant of 266 nm. From the same analysis of the other lattice axis, we also obtain the coordinate rotation angle of $-0.664(4)^\circ$ and lattice constant of 2.65(1) pixels on CCD plane corresponding to the lattice constant of 266 nm. Thus, the magnification ratio of our imaging system is 159.7(4) and one pixel of our CCD camera corresponds to 100.2(3) nm on the objective plane. These values are used for the deconvolution analysis of our images. Figure 5(f) shows a raw image of the limited region of

![Figure 5](image-url)
Figure 5. (a) with grid lines showing lattice separation and orientation. In figure 5(g), we show a reconstructed atom distribution where red circles and black dots represent the atoms and the lattice sites, respectively. In figure 5(h), we also show the reconstructed atom distribution convoluted with the model PSF of equation (3), which is compared with the raw image of figure 5(f).

An important aspect of QGM is the high-fidelity of the imaging process characterized by loss and hopping rates during the fluorescence imaging. For this purpose, we take two successive images of the same atoms with 400 ms exposure time and 300 ms delay between the two images, and observe the change in the distribution. We precisely tune the detuning $\Delta \nu$ of cooling molasses along the horizontal axes, and evaluate the loss and hopping rates during the fluorescence image from the two successive images. Figure 6(a) and (b) show the results of reconstructed atom distributions at $\Delta \nu/2\pi = -394$ and $-206$ kHz, respectively. Red circles and blue squares in the panels show the lattice sites occupation of the first and second image, respectively. (c) The pinned, lost, and hopping fractions. The fraction of pinned atoms (blue circles) shows the number of atoms preserved at the same lattice sites in the two successive fluorescence images (400 ms exposure time, 300 ms delay between the two images). The fraction of lost atoms (red squares) shows difference of the number of atoms between the two images. The fraction of hopping atoms (green triangles) shows the number of atoms appearing on a previously empty site in the second image. All fractions are normalized to the number of atoms in the first image.

5. Conclusion

In conclusion, we demonstrate a bosonic $^{174}$Yb QGM for a 2D optical lattice with a short lattice constant of 266 nm. The atoms are preserved in the lattice sites during fluorescence imaging by narrow-line laser cooling which successfully combines Doppler cooling and sideband cooling. The resulting temperature is $T = 7.4(13)$ $\mu$K, corresponding to a mean oscillation quantum number along the horizontal axes 0.22(4). The PSF has a reasonably small width comparable to the ideal value, enabling the identification of the presence and absence of atoms by the deconvolution analysis. The high fidelity of the imaging process, which is an important aspect of QGM, is confirmed by the measurement of loss rate of 6.5(18)% and hopping rate of 6.7(15)% for 400 ms exposure time.
While we perform the experiment with bosonic Yb$^{174}$ atoms, our method is applicable for a QGM for other Yb isotopes, including fermionic Yb$^{171}$ and Yb$^{173}$. In fact, our preliminary result shows that we can successfully obtain site-resolved images for fermionic Yb$^{171}$ atoms. In addition, the sideband cooling demonstrated in this work can be straightforwardly applied to other alkaline-earth atoms such as strontium, especially for an optical lattice with magic wavelength. The realization of a QGM with enough fidelity for Yb atoms in a Hubbard-regime optical lattice opens up the possibilities for studying various kinds of quantum many-body systems, and also long-range-interacting systems such as Rydberg states.

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Appendix. Deconvolution procedure

Here we describe the deconvolution procedure of our fluorescence image. First step is to determine a PSF, a lattice spacing, and a coordinate rotation angle from the isolated, single-site resolved signals, as is explained in detail in section 4. From the determined PSF of equation (3) and the location of lattice sites, we next calculate a trial image

$$I_{\text{est}}(r, \alpha) = \sum_n \alpha_n \text{PSF}(|r - r_n|), \quad (A.1)$$

where $r$ represents a position in an image plane, $r_n$ a position of a lattice site $n$, $\alpha_n$ a fitting parameter corresponding to a probability of finding an atom at a lattice site $n$, and $\alpha = (\alpha_1, \alpha_2, \alpha_3, \ldots, \alpha_n)$. Note that $\alpha_n$ represents a fitted value of a peak fluorescence count at the $n$-th lattice site. Finally we determine the parameters $\alpha$ by minimizing a quantity $\Delta(\alpha) = \sum_n (I_{\text{raw}}(r) - I_{\text{est}}(r, \alpha))^2$, where $I_{\text{raw}}(r)$ is a raw fluorescence image. As a typical example of this deconvolution procedure, we show in figure A1 a histogram of $\alpha_n$ determined in several regions of $9 \times 9$ sites picked up from raw images.

References

[1] Jaksh D and Zoller P 2005 Annals Phys. 315 52–79
[2] Greiner M, Mandel O, Esslinger T, Hänsch T W and Bloch I 2002 Nature 415 39–44
[3] Stöferle T, Moritz H, Schori C, Kohl M and Esslinger T 2004 Phys. Rev. Lett. 92 130403
[4] Spielman I B, Phillips W D and Porto I V 2007 Phys. Rev. Lett. 98 080404
[5] Schneider U, Hackermüller L, Will S, Best T, Bloch I, Costi T A, Helmes R W, Rasch D and Rosch A 2008 Science 322 1520–5
[6] Jördens R, Strohmaier N, Günter K, Moritz H and Esslinger T 2008 Nature 445 204–7
[7] Bakr W S, Gillen J L, Peng A, Fölling S and Greiner M 2009 Nature 462 74–7
[8] Sherson J F, Weitenberg C, Endres M, Cheneau M, Bloch I and Kuhr S 2010 Nature 467 68–72
[9] Bakr W S, Peng A, Tai M E, Ma R, Simon J, Gillen J L, Foellng S, Pollet J L and Greiner M 2010 Science 329 547–50
[10] Preiss P M, Ma R, Tai M E, Lukin A, Rispoli M, Zupancic P, Lahini Y, Islam R and Greiner M 2015 Science 347 1229–33
[11] Fukuhara T et al 2013 Nat. Phys. 9 235–41
[12] Cheneau M, Barmettler P, Poletti D, Endres M, Scaufler P, Fukuhara T, Gross C, Bloch I, Kollath C and Kuhr S 2012 Nature 481 484–7
[13] Cheuk I W, Nichols M A, Okan M, Gersdorf T, Ramasseh V V, Bakr W S, Lompe T and Zweifel M W 2015 Phys. Rev. Lett. 114 193001
[14] Haller E, Hudson J, Kelly A, Cotta D A, Peaudecour B, Bruce G D and Kuhr S 2015 Nat. Phys. 11 738–42
[15] Edge J A, Anderson R, Jervis D, McKay D C, Day R, Trotzky S and Thywissen J H 2015 Phys. Rev. A 92 063406
[16] Parsons M F, Huber F, Mazureenko A, Chiu C S, Setiawan W, Wooley-Brown K, Blatt S and Greiner M 2015 Phys. Rev. Lett. 114 213002
[17] Omran A, Boll M, Hilker T, Kleinlein K, Salomon G, Bloch I and Gross C 2015 arXiv:1510.04599
[18] Greif D, Parsons M F, Mazureenko A, Chiu C S, Blatt S, Huber F, Ji G and Greiner M 2015 arXiv:1511.06366
[19] Cazalilla M, Ho A and Ueda M 2009 New J. Phys. 11 103033
[20] Taie S, Yamazaki R, Sugawa S and Takahashi Y 2012 Nat. Phys. 8 825–30
[21] Hazzard K R A, Gurarie V, Hermele M and Rey A M 2012 Phys. Rev. A 85 041604
[22] Pagano G et al 2014 Nat. Phys. 10 198–201
[23] Gorshkov A, Hermele M, Gurarie V, Xu C, Julienne P, Ye J, Zoller P, Demler E, Lukin M and Rey A 2010 Nat. Phys. 6 289–95
[24] Sczara F, Hoofrachter C, Hfer M, De Groot P, Bloch I and Folling S 2014 Nat. Phys. 10 779–84
[25] Zhang X, Bishof M, Bronley S, Kraus C, Safroynova M, Zoller P, Rey A and Ye J 2014 Science 345 1467–73
[26] Cappellini G et al 2014 Phys. Rev. Lett. 113 120402
[27] Mancini M et al 2015 Science 349 1510–3
[28] Sugawa S, Yamazaki R, Taie S and Takahashi Y 2011 Phys. Rev. A 84 011610
[29] Stellmer S, Grimm R and Schreck F 2013 Phys. Rev. A 87 031611
[30] Tey M K, Stellmer S, Grimm R and Schreck F 2010 Phys. Rev. A 82 011608
[31] Sugawa S, Inaba K, Taie S, Yamazaki R, Yamashita M and Takahashi Y 2011 Nat. Phys. 7 642–8
[32] Derevianko A and Cannon C C 2004 Phys. Rev. A 70 062319
[33] Daley A J, Boyd M M, Ye J and Zoller P 2008 Phys. Rev. Lett. 101 170504
[34] Gorshkov A V, Rey A M, Daley A J, Boyd M M, Ye J, Zoller P and Lukin M D 2009 Phys. Rev. Lett. 102 110503
[35] Shiibata K, Kato S, Yamaguchi A, Uetake S and Takahashi Y 2009 Appl. Phys. B 97 753–8
[36] Kato S, Sugawa S, Shiibata K, Yamamoto R and Takahashi Y 2013 Phys. Rev. Lett. 110 173201
[37] Yamaguchi A, Uetake S, Kato S, Ito H and Takahashi Y 2010 New J. Phys. 12 103001
[38] Millen J, Lochhead G and Jones M P A 2010 Phys. Rev. Lett. 105 213004
[39] McQuillen P, Zhang X, Strickler T, Dunning F B and Killian T C 2013 Phys. Rev. A 87 013407
[40] Mukherjee R, Millen J, Nath R, Jones M and Pohl T 2011 J. Phys. B: At. Mol. Opt. Phys. 44 184010
[41] Miranda M, Inoue R, Okuyama Y, Nakamoto A and Kozuma M 2015 Phys. Rev. A 91 063414
[42] Fukuhara T, Sugawa S, Sugimoto M, Taie S and Takahashi Y 2009 Phys. Rev. A 79 041604
[43] Kato S, Shiibata K, Yamamoto R, Yoshikawa Y and Takahashi Y 2012 Appl. Phys. B 108 31–8
[44] Shiibata K, Yamamoto R and Takahashi Y 2014 J. Phys. Soc. Jpn. 83 014301
[45] Le Kien F, Schneeweiss P and Rauschenbeutel A 2013 Eur. Phys. J. D 67 92
[46] Wineland D J and Itano W M 1979 Phys. Rev. A 20 1521–40
[47] Wineland D J, Itano W M, Bergquist J C and Hulet R G 1987 Phys. Rev. A 36 2220–32
[48] Diedrich F, Bergquist J C, Itano W M and Wineland D J 1989 Phys. Rev. Lett. 62 403–6
[49] Lee J, Lee H, Noh J and Mun J 2015 Phys. Rev. A 91 053405
[50] Honda K, Takahashi Y, Kuwamoto T, Fujimoto M, Toyoda K, Ishikawa K and Yabuzaki T 1999 Phys. Rev. A 59 R934–7