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Dissipative Bose-Hubbard system with intrinsic two-body loss
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We coherently transfer the MI in the three-dimensional (3D) optical lattice. Then the ground state of bosonic ytterbium atoms is recently reported in the loss behavior of a system of the dissipative metastable state. First we create a BEC in the ground state of bosonic ytterbium atoms, and then the generation of a highly entangled Dicke state. Three-body loss was implemented by introducing several types of dissipation, the influence of the dissipation on the quantum systems has been revealed. The bosonic atoms in the three-dimensional optical lattice in the metastable state. Nevertheless, we have an interesting possibility of the quantum many-body physics. The bosonic atoms in the three-dimensional optical lattice in the metastable state. With this MI as an initial state, we investigate the stability of the system and find that the atom loss is suppressed by the strong correlation between atoms, which is attributed to both the on-site interaction and the inelastic loss. Also, as we decrease the potential depth of the lattice, we observe the growth of the phase coherence and find its suppression owing to the dissipation.

We report an experimental study of dynamics of the metastable state of bosonic ytterbium atoms in an optical lattice. The dissipative Bose-Hubbard system with on-site two-body atom loss is realized via its intrinsic strong inelastic collision of the metastable atoms. We investigate the atom loss behavior with the unit-filling Mott insulator as the initial state and find that the atom loss is suppressed by the strong correlation between atoms, which is attributed to both the on-site interaction and the inelastic loss. Also, this novel scheme of the initial state preparation enables us to observe the growth of the phase coherence as we decrease the lattice depth, otherwise impossible to create, and we quantitatively reveal that the formation of a sizable phase coherence is suppressed by the dissipation.

The bosonic atoms in the optical lattice can be regarded as the dissipative Bose-Hubbard system described by a master equation in Lindblad form [32]:

$$\frac{d\hat{\rho}}{dt} = -i\hbar[H, \hat{\rho}] + \mathcal{L}(\hat{\rho}),$$

where $H$ is the Bose-Hubbard Hamiltonian

$$\hat{H} = \frac{U_{ee}}{2} \sum_j \hat{n}_j(\hat{n}_j - 1) - J \sum_{\langle j,k \rangle} (\hat{a}_j^\dagger \hat{a}_k + \text{h.c.}) + \sum_j (\epsilon_j - \mu)\hat{n}_j,$$

and $\mathcal{L}(\hat{\rho})$ represents the dissipation due to the inelastic collision between two atoms in the state

$$\mathcal{L}(\hat{\rho}) = \frac{\Gamma_{\epsilon\epsilon}}{4} \sum_j (\hat{a}_j^\dagger \hat{a}_j \hat{\rho} \hat{\rho}^\dagger - \hat{\rho} \hat{\rho}^\dagger \hat{a}_j^\dagger \hat{a}_j - 2\hat{a}_j \hat{\rho} \hat{a}_j^\dagger + \hat{a}_j \hat{a}_j^\dagger \hat{\rho} \hat{\rho}^\dagger).$$

$U_{ee}$, $J$, and $\Gamma_{\epsilon\epsilon}$ represent the on-site interaction, the tunneling amplitude, and the inelastic collision rate, respectively. $\epsilon_j$ is the confining potential of the site $j$ and $\mu$ is the chemical potential. Index $e$ ($g$) denotes the state. $\hat{a}_j$ is the annihilation operator of the state at a site $j$ and $\hat{n}_j = \hat{a}_j^\dagger \hat{a}_j$. The nearest-neighboring pairs of lattice sites. We note that there exists the one-body loss process due to the photon scattering and the spontaneous emission, the loss rate of which is $1 \sim 2$ order of magnitude smaller than the two-body loss rate.

For the full characterization of the system, it is necessary to measure the strength of the on-site interaction...
between the $^3P_2$ state atoms. This has never been done because of the difficulty associated with the rapid loss of atoms in the $^3P_2$ state due to the large inelastic collision. We determine the scattering length by establishing a new spectroscopic technique with double-excitation process by utilizing the inelastic loss property.

We start with a preparation of the MI state of the $^1S_0$ atoms with singly- and doubly-occupied sites at the lattice depth of $V_0 = 18 E_R$ for the $^1S_0$ state. Here, $E_R = h^2/(2mL^2)$ is a recoil energy, where $m$ is the mass of the $^{174}$Yb atom, $h$ is the Planck’s constant and $L = 532 \text{ nm}$ is the wavelength of the lattice beam. Because the polarizability of the $^3P_2$ state for the 532 nm lattice beam is different from that of the $^1S_0$ state, the lattice depth depends on the atomic state, which is taken into account in the determination of the lattice depth and the calculation of the interaction (see Supplementary material).

We then excite a single $^1S_0$ state atom in the doubly-occupied sites into the $^3P_2$ state by adiabatic rapid passage (ARP) with a frequency-swept pulse with a 507 nm laser under a bias magnetic field of 200 mG. We perform the experiment with the atoms in the magnetic sublevel of $m_I = -2$. The atoms in the singly-occupied sites are not excited because of the well-separated resonance frequencies between the singly- and doubly-occupied sites due to the interaction (see Fig. 1 (c)). Subsequently, we apply the second excitation pulse with a variable frequency. If the second pulse successfully excites a remaining $^1S_0$ state atom in the doubly-occupied sites, two $^3P_2$ state atoms occupy the same site, resulting in the strong atom loss due to the inelastic collision with the rate $\Gamma_{ee}$ (Fig. 1 (a)). In the optical lattice, $\Gamma_{ee}$ is determined by the inelastic collision coefficient $\beta_{ee}$ and the confinement of the lattice potential through the relation $\Gamma_{ee} = \beta_{ee} \int |w(r)|^2 dr$, where $w(r)$ is the Wannier function of the lowest band. $\beta_{ee}$ is expected to a half of the inelastic collision coefficient with a thermal gas $\beta_{ee}^\text{thermal} = 5.1(6) \times 10^{-11} \text{cm}^3 / \text{s}$ [16, 33].

For the detection, atoms in the $^3P_2$ state are repumped back to the $^1S_0$ state using repumping lasers of 770 nm and 649 nm which are resonant to the $^3P_2 - ^3S_1$ and $^3P_0 - ^3S_1$ transitions, respectively. The $^3P_2$ atoms absorbing a 770 nm photon is excited to the $^3S_1$ state. Then the $^3S_1$ atoms decay into the $^3P_J$ states ($J = 0, 1, 2$). The atoms which decay to the $^3P_1$ state return to the $^1S_0$ state emitting 556 nm photon. The atoms which decay to the $^3P_0$ and $^3P_2$ state are again excited to the $^3S_1$ state absorbing 649 nm and 770 nm photon, respectively. The repumped $^1S_0$ atoms are captured by a magneto-optical trap (MOT) with the $^1S_0 - ^1P_1$ transition. The fluorescence from the MOT is detected by an electron-multiplying charge-coupled-device camera.

Figure 1 (b) shows the spectrum of the above-mentioned double-excitation spectroscopy. We observe a large dip around $+10 \text{ kHz}$ detuning from the $^1S_0 - ^3P_2$ transition of the singly-occupied atoms, which does not have the counterpart in the spectrum of the low-intensity single pulse spectroscopy (Fig. 1 (c)). We determine the interaction shifts as $(U_{ee} - U_{eg})/h = +10.7(3) \text{ kHz}$ and $(U_{eg} - U_{gg})/h = -9.70(5) \text{ kHz}$. From these results and the known scattering length $a_{gg} = +104.9(1.5) a_0$ [34], we obtain $a_{eg} = -201.5(1.5) a_0$ and $a_{ee} = +110(8) a_0$, where $a_0$ is the Bohr radius. This means that the on-site interaction between the $^3P_2$ atoms is repulsive and comparable to the dissipation strength: the dimensionless dissipation strength is $\hbar \Gamma_{ee}/U_{ee} = 0.94(13)$, which does not depend on the lattice depth. We note that the on-site interaction which can be measured is intrinsically affected by the coupling to the environment [35].

As the basic property of the dissipative quantum many-body system, we first study the stability of the unit-filling MI state in the presence of the two-body dissipation. Here, we measure the loss rate which varies as a function of the lattice depth because $J$, $U_{ee}$, and $\Gamma_{ee}$ depend on the lattice depth.

We first prepare the unit-filling MI state of the $^3P_2$ state in almost the same manner as in the double-excitation spectroscopy, except that the lower atom number is loaded so that the doubly-occupied sites are not created. After ramping up the lattice, we excite the atoms to the $^3P_2$ state by the ARP. The remaining $^1S_0$ atoms are blasted by applying 399 nm resonant light.

FIG. 1. (Color online). (a) Schematic of the double-excitation spectroscopy. The on-site interaction strength manifests itself in the shift of the excitation frequency at the doubly-occupied sites, which differs from the excitation frequency of the atoms at the singly-occupied site. (b) The spectrum of the double-excitation spectroscopy. The dashed line represents the fluorescence count measured after the first excitation, which corresponds to the atoms excited by the first pulse. The horizontal axis represents the detuning of the second-excitation pulse frequency from the transition of the atoms in the singly-occupied sites. The resonance observed in the negative detuning corresponds the de-excitation from the $^1S_0 + ^3P_2$ state to the $^1S_0$ state and peak around +35 kHz represents the excitation for the blue-sideband. (c) The spectrum of the single-excitation spectroscopy for comparison.
The atom number in the $^3P_2$ state $N(t)$ decreases as

$$
\dot{N}(t) = -\frac{n_0 \kappa}{N(0)} N(t)^2 - \xi N(t),
$$

where $\kappa$ is the two-body loss rate and $\xi$ is the one-body loss rate. $n_0$ is the initial filling factor estimated by the ARP excitation efficiency, which is typically 90%. The one-body loss is mainly induced by the photon scattering with the $^3P_2 - ^3S_1$ transition at 770 nm due to the 532 nm lattice beam, the rate $\gamma_{sc}$ of which depends on the intensity of the lattice beam. The spontaneous emission rate is $\gamma_{sp} = 67(7)$ mHz [36]. The one-body loss rate is given by $\xi = \gamma_{sc} + \gamma_{sp}$, which is calculated up to $\sim 0.3$ Hz.

Figure 2 (a) shows the typical decay of the atom number of the $^3P_2$ state. By fitting Eq. (4) to the data with the calculated one-body loss rate $\xi$, we extract the two-body loss rate $\kappa$ (see Fig. 2 (b)). The loss rate $\kappa$ is suppressed compared to the tunneling rate $6J/\hbar$ (see Fig. 2 (b)), which naively characterizes the time scale of the creation of the double occupancy. In the sufficiently deep lattice, $\kappa$ can be suppressed to the order of Hz.

This suppression is attributed to the formation of the strong correlation. When the tunneling is much smaller than the other energy scales ($J \ll \hbar \Gamma_{ee}$), $\xi$ is given by [10, 37]

$$
\kappa = \frac{16z(J/\hbar)^2}{\Gamma_{ee}} \left[ 1 + \left( \frac{2U_{ee}}{\hbar \Gamma_{ee}} \right)^2 \right]^{-1}.
$$

Here, $z = 6$ is the coordination number. We fit Eq. (5) to the data with the fitting parameter of $\beta_{ee}$. The best-fit value is $\beta_{ee} = 2.5(6) \times 10^{-11}$ cm$^3$/s, which is well agree with the half of $\beta_{ee}^{thermal} = 5.1(6) \times 10^{-11}$ cm$^3$/s [16, 33].

The correlation is characterized by the pair correlation function $g^{(2)} \equiv \langle n_i(n_i - 1) \rangle / \langle n_i \rangle^2$, which can be estimated from the experimental result according to the relation $g^{(2)} = \kappa / \Gamma_{ee}$ [10, 37]. Figure 2 (c) shows that $g^{(2)}$ is much smaller than 1, which means strong anti-bunching correlation, namely, the creation of the double occupation is strongly suppressed. Since $g^{(2)} = 4zJ^2/(\hbar \Gamma_{ee}/2)^2 + U_{ee}^2$ from Eq. (5), the reduction of the $g^{(2)}$ is attributed to both of the on-site elastic interaction $U_{ee}$ and the inelastic loss $\hbar \Gamma_{ee}$. In our experimental parameter of $\hbar \Gamma_{ee}/U_{ee} = 0.94(13)$, the inelastic interaction contributes to the formation of the correlation in addition to the elastic interaction, although the strength of the inelastic collision does not achieve the quantum Zeno region as the previous experiments [10–12] in which $\kappa$ decrease as the strength of the dissipation increases.

In order to confirm that the suppression of the doubly-occupied sites is not due to the reduction of the tunneling amplitude itself but due to the correlation effect as a result of the occupation of the atoms in the nearest neighboring site, we observe the tunneling dynamics from the initial state in which there is no atom in the nearest neighboring sites along one direction. After preparing the MI state with the $^1S_0$ atoms, we form the optical superlattice by adding the long lattice with 1064 nm laser along the $x$ axis with the relative phase between two lattice beam adjusted to make potential difference between A and B layers, which separates the excitation frequency (Fig. 3 (a)). We selectively excite the atoms to the $^3P_2$ state only in the A layer with ARP and blast the remaining $^1S_0$ atoms. Then we remove the additional lattice and monitor the atom number. The detection is also selectively performed with the coherent transfer to the $^1S_0$ state using ARP. We observe fast decrease of the atom number in the A layer and increase of the atom number in the B layer (Fig. 3 (b)), which indicates the tunneling of the atoms along the $x$ axis. We simply describe the
From the fitting, we obtain the tunneling rate of $a_{cc}$ and $a_{gg}$ are almost the same within the error: $a_{cc}/a_{gg} = 1.05(7)$. The lattice depth is adjusted for each state.

Without dissipation (the $^1S_0$ state), around $V_0 \sim 10 E_R$ we observe the transition from a MI with a broad distribution to the SF with a clear interference pattern characterizing the presence of the phase coherence, which is consistent with the theoretical value of the critical lattice depth of $V_0 = 11.29(16) E_R$. On the other hand, the atom distribution of the dissipative $^3P_2$ state is modified. Although we still observe the anisotropic interference pattern in the shallow lattice region, the interference pattern is unclear.

For the quantitative analysis, we evaluate the atom number, the visibility of the interference peaks, and the width of the atom distribution obtained by the TOF images (Fig. 4 (b-d)). For the $^3P_2$ system, the number of atoms starts to decrease around $V_0 = 10 E_R$. This significant atom loss reflects the start point of the melting of the MI, which creates the double occupation. The visibility of the interference peaks is defined as $v = (N_{\text{max}} - N_{\text{min}})/(N_{\text{max}} + N_{\text{min}})$ [39], where $N_{\text{max}}$ is the sum of the atom number in the regions of first-order interference peaks, and $N_{\text{min}}$ is that in the regions at the same distance from the central peak along the diagonals. In both cases, $v$ increases with the ramp-down of the lattice. This increase in the $^3P_2$ state is more moderate compared to that in the $^1S_0$ state (Fig. 4 (c)). In addition, the narrowing of the width of the density distribution is also more moderate in the case of the $^3P_2$ state (Fig. 4 (d)). These results suggest that the growth of the phase coherence is suppressed by the intrinsic on-site two-body dissipation. Similar behavior is observed in the previous experiment [12], where the two-body loss is artificially introduced using the photo-association technique and the effect of the dissipation on the coherence and the phase transition is discussed both in theory and experiment.

We estimate the effect of the momentum kick in the repumping process by the deconvolution analysis of the atom distribution. Because the repumping laser is irradiated along the imaging axis, the effect of the recoil due to the absorbing process of photon is not observed. On the other hand, the expansion of the distribution of the atoms due to the recoil of the photon emission is observed because the direction of the photon emission is random and isotropic. After repumping process, the repumped $^1S_0$ state atoms expand in the accordance with the sum of the original momentum and the recoil momentum obtained by the photon emission. We estimate the width of the expansion of the atom cloud with calculating the average number of the emitted photons $N_{\text{ph}}$ through the tunneling behavior as follows:

$$\dot{N}_A(t) = -RN_A(t) + RN_B(t), \quad \dot{N}_B(t) = RN_A(t) - RN_B(t). \tag{6}$$

Here, $N_A (B)$ is the atom number in the A (B) layers, and $R$ represents the tunneling rate between the A and B layers. With the initial condition that all atoms are placed in the A layer, Eq. (6) yields $N_B(t) = [1 - \exp(-2Rt)]N_0/2$, where $N_0$ is the initial atom number. We fit this function to the initial 0.1 sec data of the atom number in the B layer, as shown in the inset of Fig. 3 (b). From the fitting, we obtain the tunneling rate of $R = 42(6)$ Hz, which is much larger than the observed $\kappa$ in the case of the unit-filling MI (see Fig. 2 (b)) and is consistent with the relaxation time scale $4J/h = 50$ Hz discussed in Ref [38].

We also investigate the effect of the dissipation on the quantum phase transition from the MI to the SF state. We first prepare the unit-filling MI of the $^3P_2$ state in the same manner as the preparation of the initial state of the loss rate measurement. The lattice depth is $V_0 = 20 E_R$ for the $^3P_2$ state. Then we ramp down the lattice, in which the lattice ramp-down speed is $-2 E_R/\text{ms}$. The atom number and the momentum distribution during the ramp-down dynamics are obtained from the density distribution of the time-of-flight (TOF) absorption image. After ramping down the lattice to the final lattice depth, we suddenly turn off all the trap and take the image after 6-ms expansion of the atom cloud (Fig. 4 (a)). The atoms in the $^3P_2$ state are repumped back to the $^1S_0$ state 1 ms before taking the absorption image. For comparison, we observe the atoms in the dissipationless $^1S_0$ state. We compare the two results as a function of the lattice depths because the scattering lengths $a_{cc}$ and $a_{gg}$ are consistent with the relaxation time scale $4E_R$.
repumping process (Table I) with the assumption that the repumping process is instantaneously finished. In the numerical calculation, we obtain the momentum distribution due to the recoil in the repumping process, which is well approximated by the gaussian function with a half width at half maximum of $1.2 \times 2.7 \times 3.7 \times 10^7$ \text{Hz}. Here, $k_L = 2\pi/\lambda L$ is the wave number of the lattice beam. After turning off all the trap, the $^3P_2$ atoms expands in 5 ms. Then the atoms get the recoil momenta through the repumping process and expand in 1 ms. We reconstruct the original atom distribution by deconvoluting the recoil momentum distribution from the atom distribution obtained from the TOF image, and estimate original visibility and width, as shown in the yellow triangles in the Fig. 4 (c) and (d), which shows that the effect on the TOF image is limited qualitatively and does not change the whole behavior of these values qualitatively.

We note that, in the case of the $^3P_2$ state, the formation of the interference pattern is still observed, which suggests the growth of the phase coherence in the metastable state. Because of the strong inelastic collision, it is difficult to create the BEC in the metastable state and load it into the optical lattice. On the other hand, in our method with the slow ramp-down of the lattice, we can load the metastable atoms into the shallow optical lattice with suppressing the inelastic collision between atoms.

In conclusion, we have realized the dissipative Bose-Hubbard system with the metastable $^3P_2$ state of $^{174}$Yb by first creating a MI state in the ground state and the subsequent coherent transfer of the atoms into the $^3P_2$ state, evading the large inelastic loss process in the state preparation. We fully characterize the system by measuring the scattering length between two $^3P_2$ atoms by developing the double-excitation method. In the 3D optical lattice, we investigate the atom loss behavior with the unit-filling MI as the initial state and find that the atom loss is suppressed by the strong correlation between atoms, which is attributed to both the on-site interaction and the inelastic loss. Also, as we decrease the potential depth of the lattice, we observe the growth of the phase coherence and find that the formation of a sizable phase coherence is suppressed by the dissipation.

It is expected that similar behaviors will be observed with the $^3P_0$ state of Yb [18, 19] and other two-electron atomic species [13, 14, 17]. The strong suppression of the inelastic collision between atoms in the metastable state is beneficial to avoid unwanted losses in the experiment for two-component many-body physics [30] and the manipulation of the $^3P_2$ atoms exploiting the magnetic dipole moment [20, 23]. Also, our system will be a candidate for the investigation of the phase of the non-Hermitian Bose-Hubbard system [40], which is realized by combination with a quantum gas microscope to post-select the events that no atomic losses occur. Our spectroscopic method for measuring the on-site interaction enables one to search anisotropy-induced Feshbach resonances [41] between the $^3P_2$ atoms, which allows for the control of the sign and the amplitude of the parameter $h\Gamma_{ee}/U_{ee}$ and systematic investigation of the effect of the dissipation.

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