

Interfacing of an optical nanofiber with tunably spaced atoms in an optical lattice

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We experimentally demonstrate efficient interfacing of a large number of atoms to an optical nanofiber using an optical lattice with tunable spacing ($0.88 - 1.5 \mu\text{m}$) projected onto the nanofiber. The lattice beam and reflections from the nanofiber yield trap potentials that provide tight confinement in all motional degrees of freedom $\approx 220 \text{ nm}$ above the nanofiber surface, enabling efficient atom-photon coupling. We achieve trapping of ≈ 1300 atoms in periodic trap sites with a trap lifetime of $\approx 15 \text{ ms}$. We also observe the effect of varied lattice periods on the atomic motional frequencies. Our new scheme is adaptable to other nanophotonic cold-atom systems and provides a versatile and scalable platform for studying photon-mediated long-range collective interactions.

I. INTRODUCTION

Efficient coupling between quantum emitters and an electromagnetic (EM) mode [1, 2] is an essential element for photon-based quantum information processing. A one-dimensional (1D) periodic array of atoms interfaced with an optical waveguide [3–6] has been of particular interest [7, 8], due to the collective enhancement of atom-photon coupling. One promising platform to achieve this is optical nanofibers [9–13] for their lossless photon guiding, which enables long-range collective interactions between a large number of atoms [14–17]. Moreover, the suspended structure allows for easy integration with cold-atom systems, and its inherent connectivity to conventional fibers is advantageous for distributed quantum systems [18–20].

A widely used method for creating a periodic array along the nanofiber is to utilize the evanescent field modes of the trapping light propagating through the fiber [11, 21–25]. In this setup, the lattice period (d_{lat}) is determined by the wavelength of the counter-propagating trapping beams, which is constrained by the conditions for stable trapping. An interesting regime for the collective dynamics of an atomic array arises when d_{lat} satisfies the one-dimensional Bragg condition:

$$d_{\text{lat}} = \frac{q \lambda_0}{2 n_{\text{eff}}}, \quad (1)$$

where the integer q is the diffraction order, λ_0 is the free-space wavelength corresponding to the atomic resonance, and n_{eff} is the effective refractive index of the propagating mode. However, to meet this along the nanofiber, the trap laser wavelength $\lambda_{\text{laser}} \approx q \lambda_0$ has to be either near-resonant ($q = 1$) [26, 27], which is prone to decoherence due to a large photon scattering, or extremely far off-resonant ($q \geq 2$) [28], requiring large laser power.

An alternative approach is to use optical tweezers [29] to interface the atoms with the nanofiber [30, 31], controlling the array period independent of the wavelength of the trapping laser. However, interfacing a large number of atoms ($\sim 10^3$) with a nanofiber remains a challenge.

In this work, we use a side-illuminated optical lattice, featuring a wide range of quasi-continuously adjustable lattice period ($0.88 - 1.5 \mu\text{m}$), to couple a large number of atoms (1270 ± 35) to an optical nanofiber. Our tunable lattice period is set by the interference between two converging beams, similar to an optical accordion [32–35]. Instead of opto-mechanically shifting the optical paths or actively steering the beams, we employ a stationary $4f$ imaging system that passively converts the phase pattern of our custom grating into a sinusoidal intensity modulation at the nanofiber [36]. The straightforward $4f$ imaging scheme enables the generation of an optical lattice within a highly elongated envelope profile, allowing an efficient and stable interface with the nanofiber. Our method leverages the advantages of both an evanescent-mode optical lattice (encompassing a large number of sites and high periodicity) and optical tweezers (offering high tunability), with modest experimental complexity. Our system provides a versatile tool for controlling the positions of a large number of atoms coupled to an optical mode to study various photon-mediated collective dynamics [37–41]. The technique is applicable to more general nanophotonic waveguide systems [5, 42–44], and with its compatibility with optical tweezers and conventional fiber systems, our system has potential application to quantum networking [19, 31].

II. SETUP

A schematic diagram of the experiment is shown in Fig. 1. We deliver 280 mW of lattice laser power to the $4f$ imaging system with a single-mode fiber (SMF). The cylindrical lens (CL) forms a highly elongated beam profile on the grating plate (GP), overlapping with one of the compactly stacked 1D phase gratings, as depicted in the upper-right inset. We manually adjust the position of the

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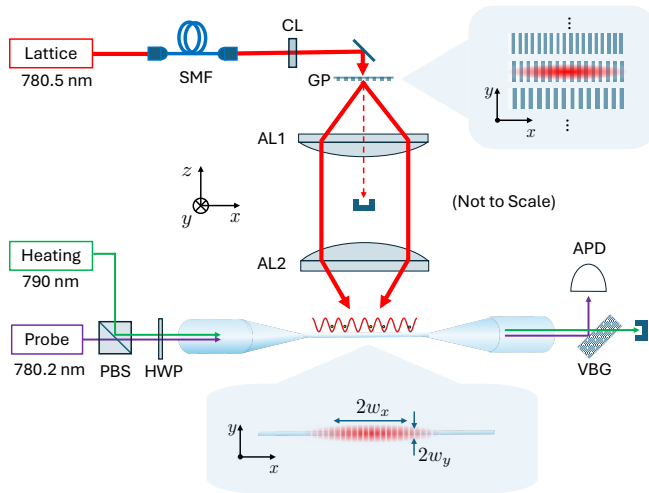


FIG. 1. Schematic diagram of the experiment: The optical lattice is formed by illuminating the grating plate with the lattice beam (see upper-right inset), which is projected onto the nanofiber via the $4f$ imaging system composed of two aspherical lenses (ALs). The probe and heating beams co-propagate through the nanofiber with orthogonal polarizations. The avalanche photodiode (APD) measures the transmission of the probe beam. SMF: single-mode fiber, CL: cylindrical lens, GP: grating plate, PBS: polarizing beam splitter, HWP: half-wave plate, VBG: volume Bragg grating.

GP along the y direction to overlap the desired grating with the beam precisely. The $4f$ imaging system, consisting of a pair of aspherical lenses (ALs) with focal lengths 200 and 100 mm, collects the ± 1 st-order diffractions from the gratings and combines them onto the nanofiber surface. Both the grating plate and the $4f$ imaging system are situated outside the vacuum chamber, enabling easy access and adjustment. Near the nanofiber surface, an optical lattice is formed along the fiber axis (x direction), enveloped by a highly elongated Gaussian profile with a waist radius of 1.5 mm in the x direction and $8\ \mu\text{m}$ in the y direction, denoted as w_x and w_y , respectively. We achieve three-dimensional confinement via reflection of the lattice beams from the nanofiber surface. Further technical details regarding our adjustable lattice system can be found in Reference [36].

The lattice laser is detuned by $-130\ \text{GHz}$ ($0.26\ \text{nm}$) from the ^{85}Rb D_2 transition, creating an attractive trapping potential for the atoms in the ground state at the local intensity maxima. For our typical trap depth of $U_0 = k_B \times 0.5\ \text{mK}$, where k_B is the Boltzmann constant, the off-resonant photon scattering rate is $4.9 \times 10^2/\text{s}$, which is negligible on the time scales of interest ($\sim 100\ \text{ns}$). The corresponding recoil heating rate $0.2\ \text{mK/s}$ is also a negligible rate. Background atomic collisions and a hot nanofiber surface limit the lifetime of the trap on a much shorter time scale ($\sim 10\ \text{ms}$) [45]. We note that our choice of the lattice detuning was due to the laser available at the time, and we plan to use a larger detuning with higher power in future experiments.

We fabricate the optical nanofiber by heating and stretching a standard single-mode fiber (Thorlabs SM800) using a hydrogen-oxygen flame and a computer-controlled motor stage [46]. The tapered nanofiber has a radius $R = 240\ \text{nm}$ over a length of 7 mm. To prevent background atoms from adhering to the dielectric surface, we deliver a constant $100\ \mu\text{W}$, 790 nm heating laser through the nanofiber. The heating laser is tuned to the tune-out wavelength between the D_1 and D_2 transition lines to eliminate the scalar light shift on the atomic ground state. The 780 nm probe beam, tuned to the D_2 cycling transition between $|F = 3\rangle$ and $|F = 4\rangle$, co-propagates with the heating laser through the nanofiber and is quasi-linearly polarized [47] in the z direction.

As the lattice beams encounter the nanofiber surface, interference patterns are formed due to the surface reflection. To study the local intensity structure, we employ a finite element method (FEM) simulation in the electromagnetic wave frequency domain, as illustrated in Fig. 2. In this simulation, the nanofiber is illuminated by two lattice beams from the top at an angle of 46° relative to each other, resulting in a lattice period $d_{\text{lat}} = 1.0\ \mu\text{m}$. Both beams are polarized in the y direction, perpendicular to the nanofiber axis. We assume a plane-wave incidence for the lattice beams, rather than the finite Gaussian profile used in the experiment, for computational efficiency. The intensity scale is displayed in arbitrary units. Fig. 2 (a) illustrates the intensity profile rendered in the $x - z$ plane ($y = 0$), depicting the standing wave pattern in both axial (x) and radial (z) direction above the nanofiber (shaded rectangular region). The lattice contrast along the nanofiber axis remains unaltered by the presence of the nanofiber, providing nearly sinusoidal potential.

Fig. 2 (b) illustrates the intensity profile rendered in the $y - z$ plane ($x = 0$), containing the radial cross-section of the nanofiber (the bottom circle) at an antinode of the lattice. The localized interference pattern provides both radial and azimuthal trap confinement near the nanofiber. A similar structure has been studied in Reference [30] for an optical tweezer interfaced with a nanofiber. We find a local intensity maximum approximately $220\ \text{nm}$ above the nanofiber surface (marked with "x"). The corresponding coupling efficiency of the spontaneous emission of the atom into the nanofiber-guided mode is predicted to be $\approx 1.8\%$. The subsequent nearest trapping site lies $650\ \text{nm}$ away from the nanofiber surface, and the corresponding coupling efficiency drops significantly to 0.014% . Because of the significant disparity in coupling efficiencies, we focus only on the trapping sites closest to the fiber surface.

Fig. 2 (c) describes the radial trapping potentials above the nanofiber surface, incorporating both lattice-induced potential (solid red line) and the surface-induced van der Waals potential (solid black line) for the ^{85}Rb ground state [48, 49]. Here, we select the intensity of the lattice beam to provide a ground-state energy shift equivalent to the trap depth $U_0 = k_B \times 0.5\ \text{mK}$ at the local

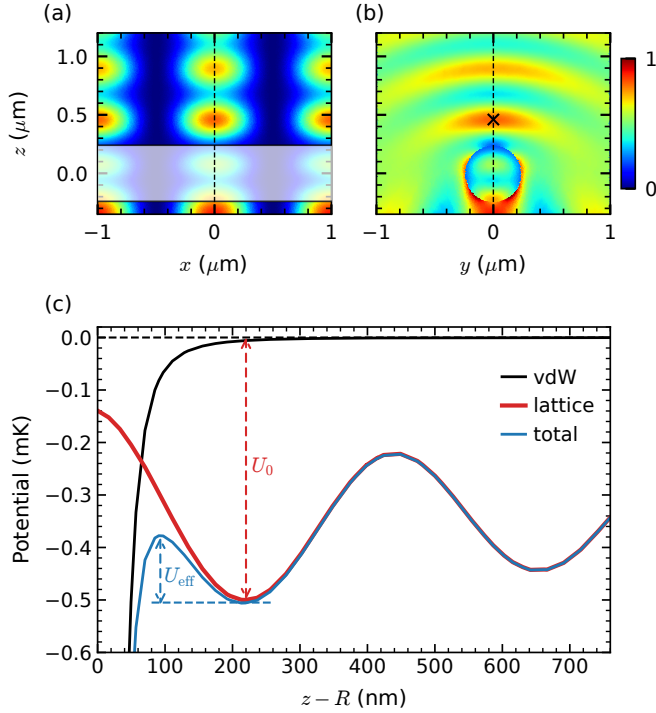


FIG. 2. Simulated intensity of the lattice light interfaced with the nanofiber rendered in the (a) x - z and (b) y - z plane, residing at $y = 0$ and $x = 0$, respectively. The geometric convention follows that of Fig. 1. The shaded rectangle in (a) and the circle in (b) represent the cross-section of the nanofiber in the respective plane. The color-map scale is displayed in arbitrary units. The local intensity maximum nearest to the nanofiber surface is marked with "x". (c) The estimated radial trapping potentials along the z axis (dashed lines in (a) and (b)), corresponding to the van der Waals (vdW) potential (solid black line), lattice (solid red line), and total potential (solid blue line) are plotted against the distance from the nanofiber surface $z - R$, where $R = 240$ nm is the radius of the nanofiber.

intensity maximum. The total potential is represented as a solid blue line, showing the reduced effective trap depth (U_{eff}) in proximity to the fiber surface. We note that due to the van der Waals potential, the lattice potential must be $U_0 \gtrsim k_B \times 0.1$ mK to provide a stable radial confinement at the nearest intensity maximum. The intensity profile formed on the other side of the nanofiber does not provide a stable trapping potential when combined with the van der Waals potential. Consequently, this enables the formation of a 1D atomic array exclusively on one side of the nanofiber, which is advantageous for studying chiral waveguide quantum optics [41, 50–53].

III. EXPERIMENT

Our experimental sequence begins with a 2.5-second loading of the magneto-optical trap (MOT) surrounding the nanofiber, followed by a 30-ms sub-Doppler cooling

phase. The lattice beams remain on throughout the experiment. Once we turn off the MOT lasers, we apply a “push-out” beam driving both the $|5^2S_{1/2} F = 2\rangle \rightarrow |5^2P_{1/2} F = 3\rangle$ and $|5^2S_{1/2} F = 3\rangle \rightarrow |5^2P_{3/2} F = 3\rangle$ transitions for 4 ms along the y direction, removing the ambient atoms from the MOT, and leaving only the atoms trapped in the lattice. During the lattice holding time, the lattice beam power provides a trap depth of $U_0 \approx k_B \times 0.5$ mK. After a holding time of 5 to 60 ms, we gradually reduce the trap depth to $U_0 \approx k_B \times 0.2$ mK over a duration of 1 ms, prior to interrogation. This reduces the Stark-induced light shift and inhomogeneous spectral broadening during our optical depth measurement. We verified that reducing the trap depth for less than 2 ms does not cause a significant loss of trapped atoms. We turn on a free-space repumping beam driving the D_1 transition 0.2 ms before the interrogation stage, so the atoms are optically active for the probe beam. Following this, we send a 1-ms probe pulse through the fiber, driving the $|F = 3\rangle \rightarrow |F = 4\rangle$ atomic D_2 cycling transition.

We measure the transmitted probe beam power P_T in the presence of atoms, the incident probe power P_{in} without atoms, and the background power P_{bg} without probe beam sequentially measured within the same experimental shot to calculate the transmission $\mathcal{T}$:

$$\mathcal{T} = \frac{P_T - P_{\text{bg}}}{P_{\text{in}} - P_{\text{bg}}}. \quad (2)$$

Fig. 3 (a) displays the transmission spectrum of atoms after a lattice holding time of 5 ms, for a lattice period of $1.2 \mu\text{m}$. For comparison, we acquire data both with (blue circles) and without (red squares) the lattice beams present. Error bars are smaller than the marker size. The solid lines represent the fits to the transmission spectra using:

$$\mathcal{T}(\Delta) = \exp\left(\frac{-\text{OD}}{1 + 4\left(\frac{\Delta - \Delta_{\text{LS}}}{\Gamma_{\text{eff}}}\right)^2}\right) + y_0, \quad (3)$$

where OD is the optical depth, Δ is the detuning of the probe laser, Δ_{LS} is the lattice-induced light shift, Γ_{eff} is the effective linewidth accounting for the inhomogeneous broadening, and y_0 is an offset [11, 54, 55]. For fitting, OD, Δ_{LS} , Γ_{eff} , and y_0 are taken as free parameters. The fitted ODs are 10.9 ± 1.3 (blue trace) and 0.05 ± 0.01 (red trace). The stark contrast between the two spectra confirms that the absorption due to the untrapped atoms is negligible.

Fig. 3 (b) shows an extension of the OD measurements using transmission spectra for different lattice holding times, ranging from 5 to 60 ms. The empty circles represent the measured OD values, and the solid line is a fitting exponential decay curve, resulting in the trap lifetime of 14.7 ± 1.4 ms. Here, we should note that Eq. 3, used to deduce the OD values, does not fully account for the inhomogeneous spectral broadening caused by the

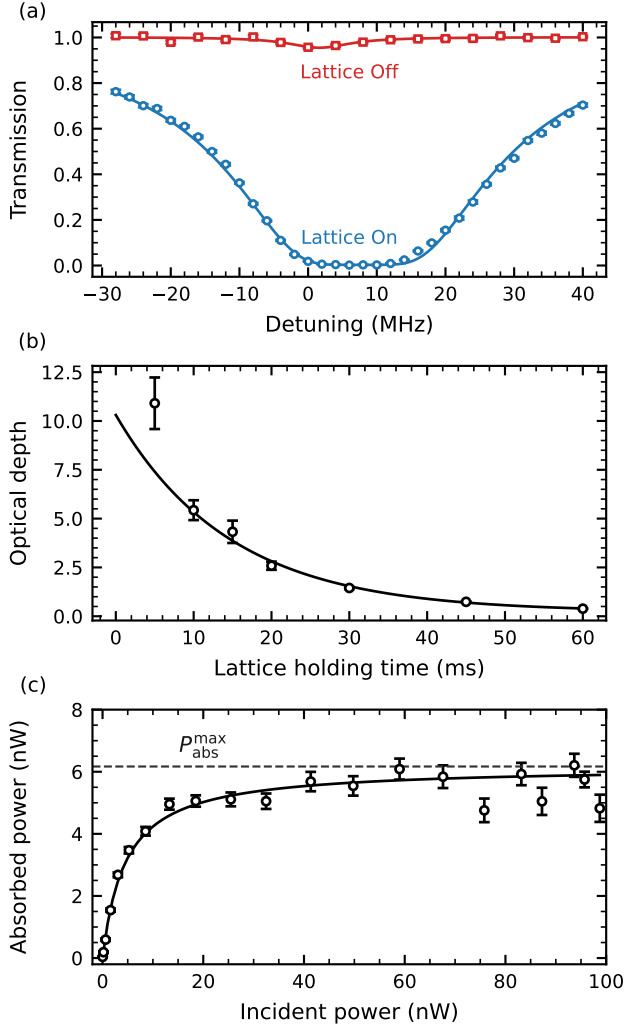


FIG. 3. (a) Transmission data with a 5-ms holding time while the lattice beams are on (blue circles) and off (red squares), fitted by Eq. 3. (b) The optical depth is plotted against various lattice holding times. The solid line represents an exponentially decaying function that fits the data, with a time constant of 14.7 ms. (c) The absorbed probe power P_{abs} is plotted for the range of incident power P_{in} . The solid line represents the data fit according to Eq. 5, and the horizontal dashed line represents the asymptotic limit of the absorption, $P_{\text{abs}}^{\text{max}}$, obtained from the fit.

thermal distribution of the atomic motion, intensity variations across different lattice sites, and the differential tensor polarizability between magnetic sublevels of the excited state [55].

To determine the number of trapped atoms in the lattice, we conduct a saturation measurement [11], with a lattice holding time of 5 ms. With the frequency fixed to the light-shifted resonance of the D_2 cycling transition, we scan the intensity of the probe beam that extends beyond the atomic saturation. For highly saturated ensemble of N atoms, the absorbed probe power

$P_{\text{abs}} = P_{\text{in}} - P_{\text{T}}$ asymptotically approaches:

$$P_{\text{abs}}^{\text{max}} = N P_{\text{atom}}^{\text{max}} = N E_{\text{ph}} \Gamma_0 / 2, \quad (4)$$

where $P_{\text{atom}}^{\text{max}} = 4.85$ pW is the power radiated by a fully saturated single Rb atom, $E_{\text{ph}} = 2.55 \times 10^{-19}$ J is the energy of a single photon, and $\Gamma_0 = 38.1 \times 10^6/\text{s}$ is the decay rate of the excited state [55].

We plot the absorbed probe power P_{abs} against the incident probe power P_{in} in Fig. 3 (c) as empty circles. The solid line fits the data using:

$$P_{\text{abs}} = P_{\text{abs}}^{\text{max}} \frac{P_{\text{in}}/P_{\text{sat}}}{1 + P_{\text{in}}/P_{\text{sat}}} \quad (5)$$

where $P_{\text{abs}}^{\text{max}}$ and P_{sat} (saturation power of the ensemble) are free parameters [54]. $P_{\text{abs}}^{\text{max}} = 6.2(1)$ nW is obtained from the fit, and using Eq. 4, we estimate the number of atoms to be $N = 1270 \pm 35$. We extract the optical depth per single atom, $\text{OD}_{\text{atom}} \approx 0.86\%$, based on the total OD value of 10.9 ± 1.3 from Fig. 3 (a).

Finally, to investigate the effect of the variation in the lattice period on the atomic motional degree of freedom, we perform parametric heating-induced atomic loss spectroscopy [56]. We modulate the lattice beam power, targeting a frequency twice the lattice's axial trap frequency, f_{ax} . The relationship between f_{ax} and the lattice period d_{lat} is given by:

$$f_{\text{ax}} = \frac{1}{d_{\text{lat}}} \sqrt{\frac{U_0}{2M}}, \quad (6)$$

where U_0 is the lattice trap depth and M is the mass of the ^{85}Rb atom. To mitigate the axial inhomogeneity of trap depth due to the Gaussian beam profile, we restrict the spatial extent of the lattice to approximately 1 mm around its center by placing a slit on the grating plate to mask the beam intensity outside the targeted extent. Following the 4-ms push-out stage, we modulate the lattice beam power with an amplitude of $\approx 3\%$ of the total power for 20 ms. For these trap-frequency measurements, we reduce the time-averaged trap depth to $k_{\text{B}} \times 0.4$ mK (80% of the maximum depth) to stay in the linear regime of the acousto-optic modulator (AOM) that controls the lattice beam power. We ramp down the lattice depth to $U_0 \approx k_{\text{B}} \times 0.1$ mK for 1 ms before the interrogation to remove the parametrically heated atoms more effectively. We fix the probe laser frequency at the light-shifted resonance of the D_2 cycling transition, and keep the power below 1 pW to avoid atomic saturation.

The results of the loss spectroscopy for different lattice periods d_{lat} are presented in Fig. 4 (a). As the modulation frequency approaches $2f_{\text{ax}}$, more atoms escape the lattice, resulting in a prominent main transmission peak in each spectrum. On the left side of these main peaks, a smaller hump at f_{ax} emerges due to the anharmonicity in the trapping potential and the imperfections in the intensity modulation that introduce the second-harmonic spectral component. We analyze these

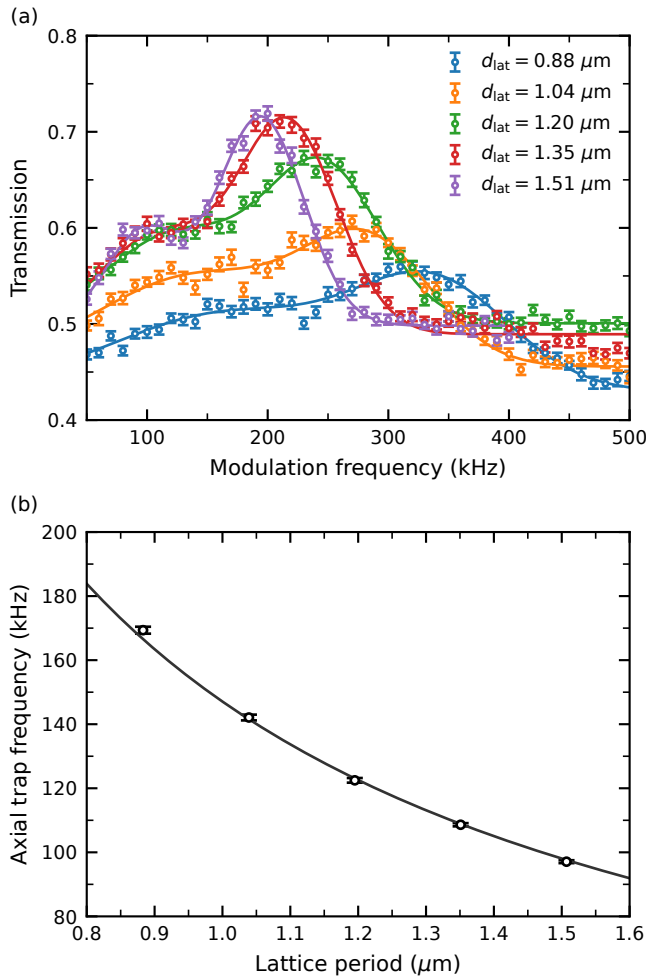


FIG. 4. (a) The results of the atom loss spectroscopy induced by the intensity modulation are presented for various lattice periods (d_{lat}). The solid curves fit the data with double Gaussian curves to estimate the axial trap frequency f_{ax} . (b) The axial trap frequencies f_{ax} from the fittings from (a) are plotted against the varied lattice periods. The solid line fits the data by Eq. 6.

data using a double Gaussian curve (solid lines), which captures both the main peak and the secondary hump at $2f_{\text{ax}}$ and f_{ax} , respectively. From these fits, we extract f_{ax} corresponding to various lattice periods, as shown in Fig. 4 (b) as empty circles. The solid line fits the data according to Eq. 6, from which we estimate the trap depth of $U_0 = k_B \times 0.443(4)$ mK, which is consistent with the predicted value, demonstrating our control over the lattice period.

From the FEM simulations (see Fig. 2), applying a harmonic approximation around the intensity maximum, the trapping frequencies for radial and azimuthal degrees of freedom are estimated to be ≈ 300 kHz and ≈ 70 kHz, respectively. However, despite our extended scan beyond the presented modulation frequency range, we were unable to identify those trapping frequencies through this

measurement. We hypothesize that this may be due to reduced visibility caused by the anharmonicity of the trapping potentials in those degrees of freedom and the inseparability between them. Additionally, the signal corresponding to the azimuthal trap frequency may have been dominated by the axial frequencies, hindering clear identification.

With the fixed lattice extent (≈ 1 mm), the measured number of atoms remained nearly constant ($N \approx 1000$) throughout the range of lattice periods. This results in the atomic filling factor of approximately 0.9 to 1.5 per site, which is significantly higher than half-filling in typical photon-assisted collisional blockade regime [57, 58]. We attribute this difference to the larger trap volume and shorter lattice holding time (5 ms) employed in our experiment. We anticipate that employing near-ground-state cooling of the atomic motion and a larger trap depth would bring the atoms into the blockade regime and realize a half-filled sub-Poissonian number distribution. It is worth noting that the Poissonian distribution of the atom numbers in the individual sites does not pose a significant hindrance to the application to study collective radiation. The 1D nature of the infinite-range interaction [14] effectively overlooks the specific atomic number distribution.

Our ongoing work aims to enhance the trap depth and achieve near ground-state cooling [59] to provide better axial localization. To establish a highly reflective atomic Bragg mirror, we need to increase the Debye-Waller factor $f_{\text{DW}} = e^{-4k^2\sigma_{\text{ax}}^2}$, where k and σ_{ax} denote the wave number of the probe light and the atomic spread in the axial direction, respectively [26, 60, 61]. The axial atomic spread $\sigma_{\text{ax}} = d_{\text{lat}}\sqrt{k_B T/(2U_0)}/\pi$ depends both on the temperature T and the trap depth U_0 . Our current temperature $T \approx 30 \mu\text{K}$, trap depth $U_0 = k_B \times 0.44$ mK, and $d_{\text{lat}} = 1.0 \mu\text{m}$ yields $f_{\text{DW}} \approx 0.3$. With a higher power laser system and enhanced cooling to $10 \mu\text{K}$, we can achieve $f_{\text{DW}} = 0.9$.

IV. SUMMARY AND OUTLOOK

In summary, we have demonstrated a new method for interfacing a one-dimensional ordered array of over a thousand atoms, with tunable lattice periods, to an optical nanofiber using a simple optical imaging technique. We have verified the maximum number of trapped atoms to be 1270 ± 35 through a saturation measurement and estimated the single-atom optical depth to be $\approx 0.86\%$. To confirm control of the lattice period, we conducted parametric heating loss spectroscopy and found a good agreement between the experimental results and the model. We highlight that our method holds potential for application in various cold-atom nanophotonic systems [5, 42–44] beyond the nanofiber platforms.

Our approach provides lattice constant tunability irrespective of the trap beam frequency, suitable for exploring collective radiative behavior of the atomic array

coupled to nanophotonic waveguides. This method allows for the exact Bragg conditions to create an atomic mirror [7, 26, 27, 62], without dissipation from the trapping potential. The wide range of lattice tunability allows study of externally driven collective radiative behavior, providing both symmetric (even q) and antisymmetric (odd q) coupling conditions. These options enable us to explore super- and sub-radiant behavior, respectively [37–40]. The single-sided trap geometry provides a convenient platform to study chiral quantum optics, which typically requires additional experimental complexity to trap atoms only on one side of the waveguide [41, 50, 53]. The combination of chirality and precisely ordered collective coupling may offer a promising opportunity to advance the quantum optics field.

Our work is an important stepping stone towards realizing quantum networking architectures [31], providing ability to interface thousands of atoms to the waveguide with great flexibility. Our platform can be combined with optical tweezers to interface free-space tweezer arrays for quantum computation with nanophotonic waveguides for quantum networking. While the static lattice can host

thousands of atoms near the waveguide and facilitate their collective interaction, the dynamic tweezers could dedicate their finite resources to individually transporting [30, 63, 64] or addressing [65–68] single atoms. The nanofiber’s inherent connectivity to conventional fibers further promotes this interface as a promising building block for a functional quantum network, providing seamless optical linkage between the local quantum processors [19, 31].

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