Large array of Schrödinger cat states facilitated by an optical waveguide

Wui Seng Leong, Mingjie Xin, Zilong Chen, Shijie Chai, Yu Wang & Shau-Yu Lan

Quantum engineering using photonic structures offer new capabilities for atom-photon interactions for quantum optics and atomic physics, which could eventually lead to integrated quantum devices. Despite the rapid progress in the variety of structures, coherent excitation of the motional states of atoms in a photonic waveguide using guided modes has yet to be demonstrated. Here, we use the waveguide mode of a hollow-core photonic crystal fibre to manipulate the mechanical Fock states of single atoms in a harmonic potential inside the fibre. We create a large array of Schrödinger cat states, a quintessential feature of quantum physics and a key element in quantum information processing and metrology, of approximately 15000 atoms along the fibre by entangling the electronic state with the coherent harmonic oscillator state of each individual atom. Our results provide a useful step for quantum information and simulation with a wide range of photonic waveguide systems.
The quantum engineering toolbox developed from free space atom–light interaction experiments over the decades has led to many exciting breakthroughs in the areas of quantum information, metrology and simulation. Departing from free space interactions towards photonic system interfaces could provide new paradigms\textsuperscript{1–17}, such as increased interaction strength, scalability of the devices and novel design of functionalities. For instance, atoms trapped by the evanescent fields of nanofibres and photonic crystal slabs have been used to study long-range atom–atom interactions mediated by the coupling of superradiant emission into the structure\textsuperscript{4,14}. In quantum communication applications, single collective excitation\textsuperscript{6} and light storage\textsuperscript{8–11} have been demonstrated in nanofibres and hollow-core fibres. Moreover, tight and diffraction-free confinement of atoms in photonic structures was used for optical switching\textsuperscript{7,17}. In metrology, cold atoms in hollow-core fibres have also shown potential applications in timekeeping and sensing\textsuperscript{8,12,13}. Regardless of these achievements, quantum control and manipulation of both external and internal degrees of freedom of atoms using photonic waveguide modes have not been realised.

Here, we demonstrate coherent excitation of the Fock states of atoms in an optical harmonic potential formed by the fundamental LP\textsubscript{01} mode of a hollow-core photonic crystal fibre. We implement an anti-Jaynes–Cummings–Hamiltonian to excite the coherent harmonic oscillator states\textsuperscript{18} (denoted as a coherent state here) and create a Schrödinger cat (SC) state\textsuperscript{18,19} using the LP\textsubscript{02} mode. Our realisation of the SC states is based on a 3-mm array of atoms trapped in an optical lattice potential inside a hollow-core fibre, as shown in Fig. 1. The diffraction-free optical waveguide allows us to prepare an array of harmonic potentials with nearly identical axial trapping frequency. The axial vibrational energy levels of the lattice form the Fock state basis $|n\rangle$ for our experiments.

**Results**

**Ground-state cooling to the Zeeman insensitive state.** We first collect an ensemble of $^{85}\text{Rb}$ atoms by Doppler and sub-Doppler cooling 5 mm above a 4-cm-long open-end hollow-core photonic crystal fibre. The fibre is a hypocycloid-shaped photonic crystal fibre with $1/e^2$ mode field radius of 22 μm\textsuperscript{11–13}. Atoms are then transported into the fibre at a velocity of 2 cm s$^{-1}$ by an optical conveyor belt using a moving optical lattice formed by a pair of counterpropagating beams\textsuperscript{15}. The optical lattice has a period of 410.5 nm, with a power of 110 mW per beam. When atoms are in the fibre, we hold the atoms in a stationary lattice and optically pump them into the magnetic-field-insensitive $|F = 2, m = 0\rangle$ state to avoid the influence of inhomogeneous magnetic fields, where $F$ denotes the hyperfine ground state of $^{85}\text{Rb}$ and $m$ is the Zeeman state. The vibrational frequency of the harmonic-like trap formed by each lattice site is $\omega = 2\pi \times 400$ kHz in the axial direction and $\omega_\pi = 2\pi \times 3.5$ kHz in the radial direction.

To prepare atoms in the axial vibrational ground state of the optical harmonic potential, we implement Raman sideband cooling (RSC) to cool atoms to the $|F = 2, m = 0, n = 0\rangle$ state\textsuperscript{20}. The cooling cycle starts with exciting atoms from $|F = 2, m = 0, n > 0\rangle$ to $|F = 3, m = 0, n = 1\rangle$ with a pair of linearly and orthogonally polarised Raman lasers (RB1 and RB2) at 821 nm, as shown in Fig. 2a, where a magnetic field of 2 G is applied along the fibre axis to define the quantisation axis and break the Zeeman degeneracy. A linearly polarised depump beam is used to bring atoms back to the $|F = 2, m, n = 1\rangle$ state in the Lamb–Dicke regime, where the coupling between the internal spin states and motional states is strongly suppressed during spontaneous emission, to preserve the vibrational quantum number. Finally, a $\pi$-polarised optical pump beam is applied to accumulate atoms in the $|F = 2, m = 0, n = 1\rangle$ state to complete the cooling cycle. The cooling process continues until all the atoms are in the $|F = 2, m = 0, n = 0\rangle$ state, which is the dark state for all the laser beams. Figure 2b shows the vibrational spectroscopy before and after cooling. The mean vibrational quantum number $<n\rangle$ is determined by taking the ratio of the areas of the first red sideband $A_{rb}$ and first blue sideband $A_{bb}$ as $<n\rangle = (A_{rb}/A_{bb})/(1 - (A_{rb}/A_{bb}))$. We achieve $<n\rangle = 0.25$ after cooling from $<n\rangle = 3.3$ before cooling. An optical depth (OD) of one corresponds to approximately $1.5 \times 10^4$ atoms, which gives an average of two atoms per lattice site\textsuperscript{12}.

**Fig. 1 Schematic of the experimental setup.** Two 150-mm achromatic lenses are used to couple all the beams into the fibre with 70% coupling efficiency. The $\pi$-polarised optical pump beam incident from the side of the fibre can efficiently prepare atoms in the $|F = 2, m = 0\rangle$ state with >90% efficiency. Each lattice site in the fibre represents a harmonic oscillator potential for atoms (red filled circles). The coloured Gaussian wave packets in the potential illustrate the coherent harmonic oscillator states that are entangled with the spin states. The inset shows a near-field image of the waveguide mode. PBS: polarising beam splitter. NPBS non-polarising beam splitter, NF notch filter, APD avalanche photodiode.
Study of the coherence between Fock states. To study the coherence of the spin Fock states in our system, we use RB1 and RB2 detuned by 3 GHz–ω/2π to excite a two-photon transition between |F = 2, m = 0⟩ and |F = 3, m = 0⟩ to |F = 3, m = 0⟩ and |F = 2, m = 0⟩. The coherence time of the two spin states |⟩⟩ and |⟩⟩ measured by a microwave π/2–π–π/2 spin-echo sequence. We observe 37% and 22% contrasts at 0.1 ms and 0.9 ms separation times between the through a driving lasers, respectively, as shown in Fig. 3a. In our differential light shift compensated optical lattice potential, the spin coherence of the spin Fock states in our system, we use RB1 and RB2 to excite atoms to the |F = 3, m = 0⟩ state as a function of the detuning with and without RSC. Most of the atoms are accumulated in the n = 0 state after RSC; therefore, the blue sidebands remain (mostly n = 0 to n = 1 transition), and the red sidebands are suppressed. Smaller OD after RSC is mainly due to atom loss during RSC. Error bars are the standard error of the mean of four experimental runs.

Preparation of a quantum coherent harmonic oscillator state. The oscillation amplitude of a classical harmonic oscillator increases when the oscillator is subject to an external sinusoidal force at its natural frequency. Similarly, for a quantum harmonic oscillator, resonant driving can be achieved when the oscillation frequency of the force is at the vibrational frequency of the quantum harmonic oscillator. The oscillatory force in our experiment is created by the time-dependent dipole force from a superconducting circuit. The coherent state created by a harmonic oscillator can be written in the Fock state basis with a Poissonian distribution as \( \alpha \exp(-|\alpha|^2/2) \sum_{n=0}^{\infty} \frac{|\alpha|^n}{n!} n^{|n>} \). Therefore, its amplitude can be determined by the mean phonon number \( \langle n > \rangle \) obtained from vibrational spectroscopy measurements. Figure 3b shows the amplitude of the coherent state \( |\alpha> \) versus the laser driving time \( t \), where atoms are initially prepared in \( |⟩⟩ \) state. The pair of driving lasers have wavelengths of 795 nm and are 3.78 GHz red-detuned from the \( |F = 2> \) to \( |F' = 2> \) transition, with \(-13 \mu W\) per beam. They are linearly polarised and relatively detuned by frequency \( \omega' \). To avoid the effect of the radial motion while creating a coherent state, we operate the driving lasers in the short pulse regime \( t << 1/\omega' \). When \( \delta t \) is small, the amplitude of the coherent state is linear with the driving time \( t \) and independent of the detuning \( \delta = \alpha = \eta \Omega t/2 \). The data show a linear trend confirming the approximation of our system.

However, this amplitude measurement does not reveal information about the coherence between the superposition of Fock states in the coherent state. We demonstrate the coherence in the coherent state by applying two consecutive driving lasers pulses with phase difference \( \phi \). The mean phonon number after the application of two successive displacement operators \( D(\alpha) \) and \( \bar{D}(\alpha e^{i\phi}) \) is \( \langle n> = n \eta > + 2|\alpha|^2 + 2|\alpha|^2 \cos \phi \), as shown in Fig. 3c. The mean phonon number returning to the initial value \( \langle n> = 0 \) when the phase \( \phi = \pi \) indicates the coherence between different Fock states in the coherent state. The fitted data show the coherent state with an amplitude \(|\alpha|=0.36\).

Creation of an array of the SC states. The SC state is a superposition of two spatially separated but localised classical states entangled with an auxiliary superposition state, where the projection measurement on the auxiliary state (or atomic decay) collapses the SC state into one or the other widely separated classical states (or live or dead cat). The SC state can be created by entangling a superposition of coherent states with the internal spin states of an atom. For trapped ions SC states, the coherent state \( |\alpha> \) has been created in ions’ motional Fock states basis, and superconducting circuits, the coherent states have been formed by the photonic Fock states in resonators with \(|\alpha|=12\). In SC states with neutral atoms, and superconducting circuits, the coherent states have been formed in harmonic potentials. Here, we create the SC states in an optical lattice by entangling the coherent states and spin states in an optical waveguide.

The protocol of the SC state preparation is shown in Fig. 4a. A microwave π/2 pulse of 7.5 μs brings the atoms from the
Fig. 3 Preparation of a quantum coherent harmonic oscillator state. a Coherence time measurements between $|\uparrow, n\rangle$ and $|\uparrow, n-1\rangle$ states. A $\pi/2-\pi-\pi/2$ spin-echo sequence is implemented on the atoms before sideband cooling. The duration of the $\pi/2$ laser pulses is 7.5 $\mu$s. The OD of $|\uparrow, n-1\rangle$ is measured with the varying phase of the final $\pi/2$ laser pulse. Error bars are the standard error of the mean of four experimental runs. The curves are sinusoidal fits to the data. b Measurements of the amplitude of the coherent state as a function of driving lasers pulse duration, where atoms are initially prepared in $|\uparrow\rangle$ state. We determine the amplitude of the coherent state by measuring $<n|$ with vibrational spectroscopy, similar to Fig. 2b. Each data point corresponds to one vibrational spectroscopy measurement. Error bars are the standard error of mean of four experimental runs. The straight line is a linear fit to the data with a fixed intercept at zero. c Measurements of $<n|$ of a coherent state after two consecutive driving lasers pulses with phase difference $\phi$. The coherence of Fock states components in the coherent state is examined by applying two consecutive driving lasers pulses with phase difference $\phi$. After the two driving lasers pulses, we measure $<n|$ from the vibrational spectra and plot it for different phases between the first and second displacement pulses. Error bars are the standard error of the mean of four experimental runs. The curve is a sinusoidal fit to the data.

The population of atoms in $|\uparrow\rangle$ is normalised to 0.5 when the driving lasers are off, and the relative phase is set at $\delta_M = \pi/2$. When the microwave phase is set to $\delta_M = 0$ ($\pi$), the SC state is termed the odd (even) cat state. The odd (even) cat states only contain odd (even) Fock states, similarly to squeezed vacuum states. We use this approach to verify the validity of the SC state created in our experiment. Figure 4c shows a comparison of the SC states when $\delta_M = 0$ and $\pi$, and it clearly confirms the trend of population change of Eq. (3) for odd (even) cat states.

Our implementation of the generation of SC states deviates from the ideal case of Eq. (3) due to the imperfect conditions of our experiment. We modify Eq. (3) to fit the data in Fig. 4b by introducing the effects of the residual excitation of the driving lasers on the $|\uparrow\rangle$ state, contrast $C$, and imperfect initial ground-state preparation as

$$P_1 = \frac{1}{2} \left( 1 - C e^{-(2|\alpha|^2+1)|\beta|^2(1-\cos\theta)} \cos(\delta_M + |\beta|^2\sin\theta) \right), \quad (4)$$

where $\beta = (1 + \epsilon e^{i\phi})|\alpha|$, $\epsilon|^2 = (e^{i\phi'} + \epsilon)/(1 + e^{i\phi'})$, $\phi' = \phi - \phi_t$, and $\phi_t = \omega^2 \Delta t$ is the phase advance of the driving lasers due to the increase in the pulse time $\Delta t$. For the residual driving, the parameter $\epsilon = 0.19$ is used to consider the off-resonant dipole force on $|\uparrow\rangle$ (see "Methods"). The imperfect ground-state preparation is taken into account by adding a factor of $2|\alpha|^2 n+1 = 1.5$ into the exponent of Eq. (3). The fitting parameters $|\alpha|$, $C$ and $\phi_t$ are displayed in Fig. 4d. The phase advance $\phi_t$ is linear in time, as expected, due to the free evolution of the phase of the beating of the driving lasers. For $t = 0.45 \mu$s, the amplitude of the coherent state is $\alpha = 1.42$, which

\[\text{Equation (1) is an entangled state between a superposition of spin-up $|\uparrow\rangle$ and spin-down $|\downarrow\rangle$ states. To confirm this, we measure the interference of the two separated coherent states by applying another microwave $\pi/2$ pulse with a phase shift $\delta_M$ relative to the first two microwave pulses, and the state becomes}

$$|\downarrow\rangle \left( a e^{i\delta_M} |\alpha\rangle \right) - |\downarrow\rangle \left( a e^{i\delta_M} |\alpha\rangle \right). \quad (2)$$

The population of atoms in state $|\uparrow\rangle$ is then determined by the overlap between the two coherent states in phase space. We detect the population of the atoms in state $|\uparrow\rangle$ by absorption detection (OD) with varying phase $\phi$. The probability of finding the atoms in $|\uparrow\rangle$ is

$$P_1 = \frac{1}{2} \left( 1 - |\alpha|^2 (1 - \cos\phi) \cos(\delta_M + |\alpha|^2\sin\phi) \right). \quad (3)$$

Figure 4b shows the experimental data of SC state interference fringes for different driving lasers pulse durations. The population of the atoms in $|\uparrow\rangle$ is normalised to 0.5 when the driving lasers are off, and the relative phase is set at $\delta_M = \pi/2$. When the microwave phase is set to $\delta_M = 0$ ($\pi$), the SC state is termed the odd (even) cat state. The odd (even) cat states only contain odd (even) Fock states, similarly to squeezed vacuum states. We use this approach to verify the validity of the SC state created in our experiment. Figure 4c shows a comparison of the SC states when $\delta_M = 0$ and $\pi$, and it clearly confirms the trend of population change of Eq. (3) for odd (even) cat states.

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corresponds to a maximum separation of the two coherent wave packets of $2\sqrt{2}\alpha z_0 = 54$ nm, where $z_0^2 = (\hbar/2m\omega)$ is the variance in the ground-state wave packet size, $\hbar$ is the reduced Planck constant and $m$ is the mass of $^{85}$Rb. The fidelity of the state can be estimated from the contrast as $C^{\alpha\beta}$. The loss of contrast suggests phase noise during the operation, which also widens the narrow interference feature. The increase of the interference width leads to the saturation of $|\alpha|$ after $t = 0.45 \mu s$ for the fitting of Eq. (4).

We study the effect of radial motion during the free evolution between driving lasers pulses by inserting an additional free evolution time $\tau$, as shown in Fig. 4e. The sequence follows $\pi/2 - D(\alpha) - \tau - \pi - D(\alpha e^{i\phi}) - \pi/2$. The data show that the radial motion does not significantly degrade the contrast at our experimental timescale. We also numerically integrate Eq. (4) over the radial distribution of atoms with radial position-dependent parameters of $\eta_1$, $\Omega$ and $\delta$, as shown in Fig. 4f.

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**Fig. 4 SC state interferometer.** a Protocol for creating the SC state in phase space representation. The top and bottom columns represent $|\uparrow\rangle$ and $|\downarrow\rangle$, respectively. The circles in red indicate the uncertainty in the ground state and the coherent state. The microwave $\pi/2$ and $\pi$ pulses are used to manipulate the internal spin states. The displacement operator $D(\alpha)$ and $D(\alpha e^{i\phi})$ are created by driving lasers pulses. b Interference fringes of the SC state interferometer. The population of atoms in $|\uparrow\rangle$ state $P_\uparrow$ is measured with the relative phase $\phi$ of the two driving lasers pulses for different pulse durations. Error bars are the standard error of the mean of four experimental runs. c Even and odd SC states. The population of atoms in $|\uparrow\rangle$ state $P_\uparrow$ is measured with the relative phase $\phi$ of the two driving lasers pulses for $\delta_m = 0$ (odd SC state) and $\delta_m = \pi$ (even SC state). The driving lasers pulse duration is $0.45 \mu s$. Error bars are the standard error of the mean of four experimental runs. d Fitting parameters of Eq. (4) for the experimental data in panel b is plotted against different driving lasers pulse duration. e Measurements of the $|\uparrow\rangle$ state population $P_\uparrow$ in cat interferometer operation versus phase $\phi$ with additional free evolution time $\tau$. The total free evolution time between the two driving lasers pulses for $\tau = 1 \mu s$ is $16 \mu s$. Error bars are the standard error of the mean of four experimental runs. f Simulation of the temperature effect on the SC state interferometer. The simulation uses the fitted parameters for $t = 0.45 \mu s$ in panel d. The $1/e$ Gaussian radius of the atomic cloud in the radial direction is $0.9 \mu m$ and $8 \mu m$ for $2 \mu K$ and $150 \mu K$, respectively.
Derivation of the displacement operator. The Hamiltonian of a forced harmonic oscillator in the interaction picture has a form similar to the classical harmonic oscillator 

\[ H(t) = \hbar f(t) a e^{-i \omega t} + f(t) a^\dagger e^{i \omega t}, \]

where \( f(t) \) is a time-dependent force, and \( a \) and \( a^\dagger \) are the annihilation and creation operators, respectively. The state after some interaction time \( t \) can be characterised by the time-evolution operator

\[ \mathcal{U}(t) = \exp \left( -\frac{i}{\hbar} \int_0^t H(t') dt' \right) = e^{\mathcal{H} \alpha^2 - \alpha^2}, \]

where we only consider the first term in the exponent and \( \alpha \) is defined as

\[ \alpha(t) = -\frac{i}{\hbar} \int_0^t f(t') e^{i \omega t'} dt'. \]

The time-dependent force \( f(t) \) in the operator is provided by a moving standing wave resulting from a pair of counterpropagating lasers whose frequency difference \( \omega \) is set to be near the vibrational frequency \( \omega \) of the harmonic potential. The interaction Hamiltonian between the spin states \( |\uparrow\rangle \) of atoms and the lasers can be written as

\[ \hat{H}(t) = \frac{\hbar}{2} \left[ \Omega |\alpha e^{i(\phi - \delta)} + \alpha^* e^{i(\phi + \delta)}\rangle\langle\uparrow| \right], \]

where \( \Omega \) is the Rabi frequency of the driving lasers, \( |\alpha\rangle \) is the effective wave number of the lasers, \( \omega \) is the relative frequency of the driving lasers, and \( \phi, \delta \) are the relative phase between the two lasers. In the Lamb–Dicke regime, where \( k_0 |\alpha| < \hbar \omega / \Omega |\alpha| \), the driving lasers from the vibrational frequency, the corresponding Hamiltonian for the first blue sideband is

\[ \hat{H}(\tau) = \frac{\hbar}{2} \Omega \left[ \alpha e^{i(\phi - \delta)} + \alpha^* e^{i(\phi + \delta)} \right]|\uparrow\rangle\langle\uparrow|. \]

Schrodinger cat-state interferometer. In the preparation of the coherent harmonic state, ideally, the driving lasers should only excite atoms in \( |\uparrow\rangle \), \( |\downarrow\rangle \), \( |\alpha\rangle \), \( |\alpha^*\rangle \), respectively, where \( \Phi = \phi + \omega \Delta t, \Delta \) is the relative phase of the driving lasers and \( \omega \Delta t \) is the phase advance during the extra waiting time between driving lasers pulses, the state then becomes

\[ \alpha(\tau) \sin(\delta / 2) e^{i \Phi} \]

The global phase term \( |\phi\rangle \sin(\delta / 2) \) can be ignored in our measurements. The final \( \pi / 2 \) microwave pulse with phase \( \delta_P \) relative to the first two microwave pulses creates the state

\[ \frac{1}{2} \left[ (e^{i \Phi} |\uparrow\rangle + |\downarrow\rangle) \right] \frac{1}{2} \left[ (|\uparrow\rangle + e^{i \Phi} |\downarrow\rangle) \right] \]

where we define \( \beta \equiv (e^{i \Phi} |\uparrow\rangle + |\downarrow\rangle) / (1 + e^{i \Phi}) \). The population of atoms in the \( |\uparrow\rangle \) state is

\[ P_1 = \frac{1}{2} + \frac{1}{2} e^{\Phi} \sin(\theta) \]

Temperature effect in the Schrödinger cat-state interferometer. Due to the finite temperature of the atoms in the radial direction, atoms at different locations of the radial trap experience different axial trapping frequencies and therefore different Lamb–Dicke parameters and Rabi frequencies of the driving lasers. Assuming that the above parameters have a Gaussian distribution radially, we numerically integrate Eq. (3) over the position \( r \) to obtain the coherent state

\[ \frac{1}{2} \left[ (1 - C_\alpha (\gamma^{(1)}(\bar{r}) |\alpha\rangle \sin(\theta) \sin(\theta) \right) dr, \]

where the centre of the fibre is defined as \( r = 0, w_t^2 = W^2 k_0 T^2 / 2U \) is the 1/e radius of the spatial distribution of the atomic cloud in the radial direction, \( W = 22 \mu \text{m} \) is the 1/e^2 mode field radius, \( k_0 \) is the Boltzmann constant, \( T_i \) is the radial atom.
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Author contributions

W.S.L., M.X., Z.C., S.C. and S.-Y.L. contributed to designing the experiment, taking and analysing the data and preparing the paper.

Competing interests

The authors declare no competing interests.

Additional information

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Correspondence and requests for materials should be addressed to S.-Y.L.

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