

Cavity dark mode mediated by atom array without atomic scattering loss

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We realize a ring cavity strongly interacting with an atom array with configurable spatial structures. By preparing the atom array with a maximized structure factor, we observe the emergence of a cavity dark mode, where the standing-wave nodes are dynamically locked to the positions of the atoms. The dark mode is decoupled from the atoms, protecting the system from dissipation through atomic scattering, but still mediates strong coupling and enables efficient conversion between two optical modes. Moreover, we impart an arbitrary large phase shift on the converted optical fields by translating the atom array. This strongly interacting ring cavity system with single-atom addressability opens ways to quantum optical engineering and the generation of photonic quantum states based on the geometrical structure of atom arrays.

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Strong atom-light interactions are crucial for advancing quantum information technology [1,2]. By using atoms strongly coupled to optical cavities [3–6] or waveguides [7–12], efficient coupling between different optical modes can be achieved. This enables photon switching in various propagation directions [7], atomic Bragg mirrors [8,9], all-optical circulators and routers [10–12], and large optical phase shifts [13–15]. These capabilities are vital for optical quantum engineering [16–25] and the advancement of quantum networks [26–29].

To achieve maximum atom-light coupling strength, the most direct approach is to position atoms where the optical field is strongest. In a Fabry-Pérot cavity, this means placing atoms at the standing-wave antinodes [30–34]. However, at these positions, the strong atom-light coupling is often accompanied by strong spontaneous scattering of light by the atoms into free space, causing photon loss and atomic heating [35,36]. Here, we demonstrate a cavity dark mode that, despite arising from strong atom-light interactions, enables efficient coupling between two optical modes while protecting the atom-cavity system from spontaneous atomic scattering.

The cavity dark mode is analogous to the well-known atomic dark state in electromagnetically induced transparency (EIT) [37], where atomic excitation is eliminated through the destructive interference of two coupling fields, enabling dissipation-free atomic state transfer as in stimulated Raman adiabatic passage. Photonic dark states have been realized in various systems. For instance, in a coupled-cavity system [38,39] where two Fabry-Pérot fiber cavities, each containing an atomic ensemble, are connected by a fiber, a cavity dark mode is formed, separating photonic and atomic

excitations into different cavities so that the atoms are not exposed to light. Similarly, optomechanical dark modes [40,41] have been demonstrated, where a superposition of two cavity modes with different frequencies decouples from the mechanical oscillator, thereby suppressing mechanical dissipation.

In this Letter, we realize a cavity dark mode inside a ring cavity that strongly interacts with an atom array [42,43] with configurable spatial structures. The ring cavity supports two counterpropagating traveling-wave modes [44–48]. Each mode is strongly coupled to the atoms, with the coupling strength determined by the atom array structure factor. When the atom array is configured with a maximized structure factor, with the atomic separations being integer multiples of half the cavity wavelength, a cavity dark mode emerges as a coherent superposition of the two traveling-wave modes, dynamically forming a standing wave with the nodes precisely locked to the atoms. This alignment decouples the atom array from the dark mode, eliminating dissipation from atomic scattering loss. Despite this decoupling, the atom array still mediates strong coupling between the two cavity modes, leading to optical mode conversion. This cavity dark mode is robust against the relative movement between the atom array and the cavity mirrors, as the dark mode can follow the motion of the atom array. By displacing the atom array, we can impart an arbitrarily large phase shift on the converted optical mode.

In our experiment, an in-vacuum optical ring cavity comprising four mirrors has a small waist of $\sim 7 \mu\text{m}$ and finesse of $\mathcal{F} = 4.4(1) \times 10^4$. Laser-cooled ^{87}Rb atoms are loaded into a one-dimensional tweezer array at the center of the cavity and further cooled to $5.2 \mu\text{K}$ by Raman sideband cooling. The Hamiltonian is given by

$$\hat{H}/\hbar = -\delta(\hat{a}_+^\dagger \hat{a}_+ + \hat{a}_-^\dagger \hat{a}_-) - \Delta \sum_{j=1}^N \hat{\sigma}_j^+ \hat{\sigma}_j^- - \left[ig \sum_{j=1}^N (e^{ikx_j} \hat{a}_+ + e^{-ikx_j} \hat{a}_-) \hat{\sigma}_j^+ + \text{H.c.} \right]. \quad (1)$$

Here, $\hat{a}_\pm$ corresponds to the annihilation operator of the linearly polarized cavity traveling-wave modes $e^{\pm ikx}$,

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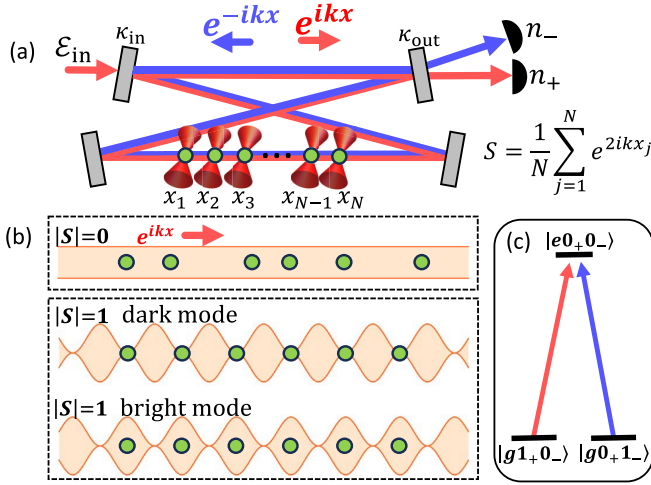


FIG. 1. Experimental schemes. (a) An atom array of ^{87}Rb atoms with the configurable structure factor S is coupled to an optical ring cavity. We drive the cavity by an input probe beam $\mathcal{E}_{\text{in}}$ in the forward $+\hat{x}$ direction. Two single-photon counting modules detect the output photons $n_{\pm}$ of the two cavity traveling-wave $\hat{a}_{\pm}$ modes $e^{\pm ikx}$. κ_{in} (κ_{out}) is the coupling rate of the input (output) mirror. (b) When $S = 0$, only the traveling-wave $\hat{a}_{+}$ mode is excited in the cavity. When $|S| = 1$, meaning the atomic separations are integer multiples of half the cavity wavelength $\lambda/2$, the dark/bright mode can be dynamically formed such that all the atoms are located at the nodes/antinodes. (c) The energy level structure illustrates the cavity dark mode. The destructive interference of a coherent superposition of the cavity $\hat{a}_{\pm}$ modes eliminates atomic excitation.

respectively. The coupling between the atom and each of the cavity traveling-wave modes is g . We probe the cavity with an input probe beam $\mathcal{E}_{\text{in}}$ in the forward $+\hat{x}$ direction and measure the output photon number $n_{\pm}$ in the $\hat{a}_{\pm}$ modes [Fig. 1(a)]. The probe and cavity resonance frequencies are tuned near the $5S_{1/2}, F = 2 \rightarrow 5P_{3/2}, F' = 3$ transition, with the probe-cavity detuning δ and the probe-atom detuning Δ . $\hat{\sigma}_j^{\pm}$ are atomic raising and lowering operators. Atoms are prepared in $|F = 2, m_F = 2\rangle$ by optical pumping.

Under the weak excitation approximation and adiabatically eliminating the atoms [49], the equations of motion for the $\hat{a}_{\pm}$ modes are

$$\begin{aligned} \frac{d\hat{a}_{+}}{dt} &= \left(i\delta - \frac{\kappa}{2} + \frac{Ng^2}{i\Delta - \gamma/2} \right) \hat{a}_{+} \\ &\quad + \frac{Ng^2S^{*}}{i\Delta - \gamma/2} \hat{a}_{-} - i\sqrt{\kappa_{\text{in}}}\mathcal{E}_{\text{in}}, \\ \frac{d\hat{a}_{-}}{dt} &= \left(i\delta - \frac{\kappa}{2} + \frac{Ng^2}{i\Delta - \gamma/2} \right) \hat{a}_{-} + \frac{Ng^2S}{i\Delta - \gamma/2} \hat{a}_{+}. \end{aligned} \quad (2)$$

Here, we characterize the spatial structure of the atom array by the structure factor defined by $S = \sum_{i=1}^N e^{2ikx_i}/N$, which plays a crucial role in mediating the coupling between the two cavity modes $\hat{a}_{\pm}$. We deterministically prepare atom arrays with a variable atom number N and fully controlled positions with a precision of 5 nm by arranging tweezers using acousto-optic deflectors. Our cavity decay rate $\kappa/(2\pi) = 33.6(8)$ kHz, the atomic decay rate $\gamma/(2\pi) = 6.07$ MHz, and the

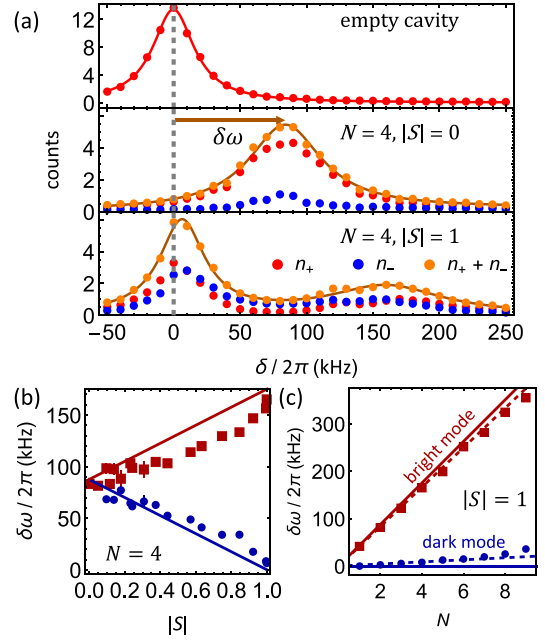


FIG. 2. Cavity dynamical modes tuned by the atom array spatial structures. The cavity forward (red), backward (blue), and total (orange) transmissions are measured for (a) empty cavity and atom arrays with $N = 4$ and $|S| = 0$, $|S| = 1$. The solid lines are Lorentzian fits to total transmission. (b) The cavity shifts of the two modes $\hat{c}_{1,2}$ are measured when varying the atom array structure factor $|S|$. The bright and dark modes emerge when $|S| = 1$. The solid lines are the theoretical calculations under ideal conditions with $C = 12.5(1)$. (c) At $|S| = 1$, the cavity shifts of the bright (square) and dark (circle) modes for atom arrays containing up to nine atoms. $\Delta/(2\pi)$ is chosen to be 30 MHz. The solid lines are the theoretical calculations with $C = 12.5(1)$ and $|S| = 1$. The dashed lines are $|S| = 0.9$ considering the finite temperature of atoms.

single-atom cooperativity $C = 4g^2/(\kappa\gamma) = 12.5(1)$ for each of the $\hat{a}_{\pm}$ modes.

In Eq. (2), when the structure of the atom array $|S| = 0$, the $\hat{a}_{\pm}$ modes are decoupled and hence are the eigenmodes of the system. Since the cavity is driven in the forward $+\hat{x}$ direction, only the $\hat{a}_{+}$ mode is excited. The cavity forward transmission displays a resonance shifted by $\delta\omega = NC\kappa\gamma\Delta/(4\Delta^2 + \gamma^2)$ relative to the empty cavity and with minimal transmission in the cavity backward direction [Fig. 2(a)]. Along with the cavity resonance shift, the traveling-wave mode $\hat{a}_{+}$ also suffers atomic scattering loss, resulting in an increase of the cavity linewidth from κ to $\kappa + \delta\kappa = \kappa + NC\kappa\gamma^2/(4\Delta^2 + \gamma^2)$.

When the structure factor $|S| \neq 0$, the $\hat{a}_{\pm}$ modes are coupled and no longer the eigenmodes. The cavity system is best described in a new basis by performing the transformations $\hat{c}_1 = \frac{1}{\sqrt{2}}(\frac{S}{|S|}\hat{a}_{+} + \hat{a}_{-})$, $\hat{c}_2 = \frac{1}{\sqrt{2}}(\frac{S}{|S|}\hat{a}_{+} - \hat{a}_{-})$. We obtain a new set of decoupled equations after substituting the transformations into Eq. (2):

$$\begin{aligned} \frac{d\hat{c}_1}{dt} &= \left(i\delta - \frac{\kappa}{2} + \frac{Ng^2(1 + |S|)}{i\Delta - \gamma/2} \right) \hat{c}_1 - i\frac{S}{\sqrt{2}|S|}\sqrt{\kappa_{\text{in}}}\mathcal{E}_{\text{in}}, \\ \frac{d\hat{c}_2}{dt} &= \left(i\delta - \frac{\kappa}{2} + \frac{Ng^2(1 - |S|)}{i\Delta - \gamma/2} \right) \hat{c}_2 - i\frac{S}{\sqrt{2}|S|}\sqrt{\kappa_{\text{in}}}\mathcal{E}_{\text{in}}. \end{aligned} \quad (3)$$

From Eqs. (3), the $\hat{c}_{1,2}$ modes have the resonance frequencies $\delta\omega_{1,2} = NC(1 \pm |S|)\kappa\gamma\Delta/(4\Delta^2 + \gamma^2)$ and the linewidth broadening $\delta\kappa_{1,2}/\kappa = NC(1 \pm |S|)\gamma^2/(4\Delta^2 + \gamma^2)$. When $NC|S|$ is large, the cavity exhibits two distinct resonances that are well separated as shown in Fig. 2(a). By fitting the total transmission $n_+ + n_-$ with a sum of two Lorentzian line shapes, we can determine the photon number $n_{1,2}$ in the $\hat{c}_{1,2}$ modes, respectively. We vary the structure factor $|S|$ and the atom number N , and measure the resonance frequencies of the $\hat{c}_{1,2}$ modes for atom arrays containing up to nine atoms [Figs. 2(b) and 2(c)].

A special case occurs when the atom array is arranged with $|S| = 1$, meaning the atomic separations are integer multiples of $\lambda/2$. Under this condition, a cavity dark mode emerges, when the resonance of the $\hat{c}_2$ mode becomes the same as that of the empty cavity, exhibiting zero cavity shift $\delta\omega = 0$ and zero cavity broadening $\delta\kappa/\kappa = 0$. The dark mode is analogous to EIT. It is an equal superposition of the $\hat{a}_{\pm}$ modes, with the phase controlled by the atom array structure factor S , such that all the atoms are located at the nodes of the standing wave as shown in Fig. 1(b). As a result, atoms are shielded from light and introduce no photon loss into free space through atomic scattering, allowing the dark mode to maintain the bare cavity linewidth. Nevertheless, the two counterpropagating traveling waves $\hat{a}_{\pm}$ are still strongly coupled by the interactions mediated with the atoms. In contrast, the $\hat{c}_1$ mode has its antinodes aligned with the atoms, exhibiting a cavity shift and atomic scattering loss twice that of the $|S| = 0$ case, with the single-atom cooperativity of $2C = 25.0(2)$, and is referred to as the bright mode. To characterize the quality of the dark mode that we experimentally generate, we measure the dark mode purity $\mathcal{D}$, and the relative linewidth broadening $\delta\kappa/\kappa$ due to atomic scattering loss. The first quantity, the purity $\mathcal{D}$, describes how well the dark mode can be selectively generated while suppressing the excitation of the unwanted lossy bright mode. It is defined as $\mathcal{D} = n_2/(n_2 + n_1)$, where $n_{1,2}$ are the photon numbers in the $\hat{c}_{1,2}$ modes, respectively, when the probe beam is tuned to the dark mode resonance frequency $\delta\omega_2$. The $n_{1,2}$ values are extracted from the Lorentzian fit of the cavity total transmission in Fig. 2(a). In order to suppress the photon number n_1 from bright mode leakage, the dark mode needs to be spectrally well separated from the bright mode. In Fig. 3(a) we calculate the purity $\mathcal{D}$ using Eqs. (3). At a given atomic detuning Δ , $\mathcal{D}$ increases with the collective cooperativity NC . The second quantity, the linewidth broadening $\delta\kappa/\kappa$, describes the level of atomic scattering loss that is achieved in practice. Ideally, the dark mode has $\delta\kappa/\kappa = 0$. However, due to the finite spatial spread in traps, atoms do not perfectly remain at the dark mode nodes, resulting in a slight reduction of the structure factor $|S| < 1$ so that $\delta\kappa/\kappa \neq 0$. Through Raman sideband cooling of atoms in our tweezer traps with a radial trapping frequency of $\omega_m = 2\pi \times 120$ kHz, we obtain the atomic spatial spread $\sigma = 31$ nm, much less than the dark mode standing-wave period of $\lambda/2 = 390$ nm. This leads to a reduction of the structure factor to $|S| = 0.9$. We compute $\delta\kappa/\kappa$ in Fig. 3(a) from Eqs. (3). At a fixed detuning, $\delta\kappa/\kappa$ increases with NC [Fig. 3(a)] because the atomic scattering of photons increases. This shows the competing effects of increasing NC at a fixed Δ when we try to optimize both the dark mode purity $\mathcal{D}$ and

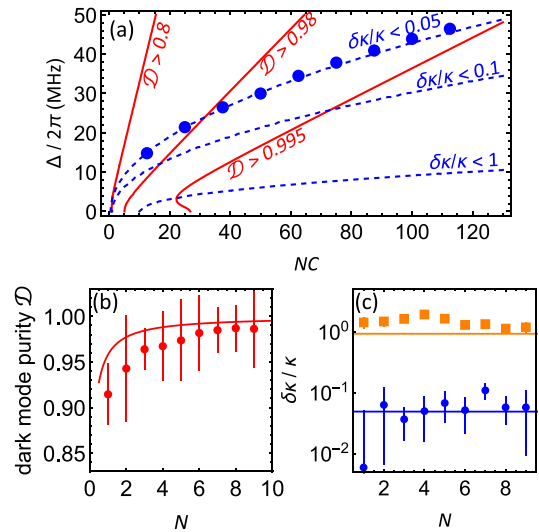


FIG. 3. Characterization of the cavity dark mode by the dark mode purity $\mathcal{D}$ and the atomic scattering broadening $\delta\kappa/\kappa$. (a) Contour plots of $\mathcal{D}$ and $\delta\kappa/\kappa$ with the collective cooperativity NC and the atomic detuning Δ at $|S| = 0.9$. With increasing NC , the area to the right of the red contour lines corresponds to larger purity $\mathcal{D}$. With increasing Δ , the region above the blue contour lines corresponds to smaller cavity broadening $\delta\kappa/\kappa$. We vary the atom number N and adjust the corresponding Δ along the contour line of $\delta\kappa/\kappa = 0.05$ indicated by blue circles in (a). The measured dark mode purity $\mathcal{D}$ increases with N (b) while the measured $\delta\kappa/\kappa$ (blue circles) remains constant (c). In comparison, the cavity bright mode [orange squares in (c)] has much larger loss than the dark mode (blue circles). All solid lines are theoretical calculations according to Eqs. (3).

the loss $\delta\kappa/\kappa$. However, as long as NC is large enough, we can always adjust Δ to simultaneously achieve large $\mathcal{D}$ and small $\delta\kappa/\kappa$. Therefore, the collective cooperativity NC is the only fundamental physical quantity that determines the dark mode quality.

In Fig. 3(c), we vary N and correspondingly adjust Δ to obtain $\delta\kappa/\kappa \sim 0.05$, which is already close to our cavity linewidth measurement resolution limit. Due to our large value of $C = 12.5(1)$, we measure $\mathcal{D} > 0.98$ with atom arrays containing more than five atoms [Fig. 3(b)]. In comparison, the loss of the bright mode is much larger than the dark mode, with $\delta\kappa/\kappa > 1$.

A high-quality dark mode enables low-loss conversion between the cavity modes. We characterize the overall conversion efficiency as the ratio between the output photon number n_- of the backward-traveling mode $\hat{a}_-$ and the input photon number n_{in} in the forward direction, $\chi = n_-/n_{\text{in}}$. Using Eqs. (2), when $NC \gg 1$, the efficiency takes a simple form as $\chi = \eta_{\text{in}}\eta_{\text{out}}/(1 + \delta\kappa/\kappa)^2$, where $\eta_{\text{out}} = \kappa_{\text{out}}/\kappa = 0.28$ and $\eta_{\text{in}} = \kappa_{\text{in}}/\kappa = 0.03$ accounting for the output and input photon coupling ratio [49]. In Fig. 4 we show the normalized conversion $\chi/\eta_{\text{in}}\eta_{\text{out}}$, excluding the effects of the finite input and output coupling ratios, as a function of N . This normalized $\chi/\eta_{\text{in}}\eta_{\text{out}}$ describes the conversion for the intracavity $\hat{a}_-$ mode. It increases with N , reaching $\sim 80\%$ for atom number $N > 1$, and is limited by the finite $\delta\kappa/\kappa \sim 0.05$. In contrast, the optical mode conversion operated with the bright mode is much lower due to greater photon loss from atomic scattering.

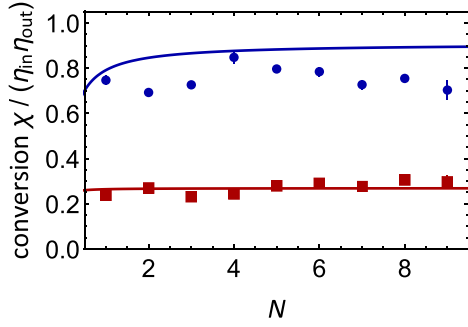


FIG. 4. Conversion $\chi/\eta_{\text{in}}\eta_{\text{out}}$ between two optical modes. The cavity dark mode (blue circles) has a much higher conversion efficiency than the bright mode (red squares). For the dark mode, the conversion efficiency quickly increases to near unity with the atom number $N > 1$, owing to the large value of $C = 12.5(1)$. The conversion is limited by the small linewidth broadening $\delta\kappa/\kappa$, which is due to the finite spatial spread of the atoms. The solid lines are theoretical calculations according to Eqs. (2).

This shows that, in addition to $\delta\kappa/\kappa$, the conversion efficiency χ provides a complementary measure of the atomic scattering loss.

The dark mode can also realize arbitrary large optical phase shifts on the converted optical mode. This behavior is formally attributed to the continuous $U(1)$ symmetry inherent in the ring cavity system. Equations (2) are invariant under a spatial displacement of the atom array by $x \rightarrow x + X$, accompanied by the cavity phase shifts as $\hat{a}_+ \rightarrow \hat{a}_+$, $\hat{a}_- \rightarrow \hat{a}_- e^{2ikX}$. This means that the intracavity field, given by $\hat{E}(x) \propto \hat{a}_+ e^{ikx} + \hat{a}_- e^{-ikx}$, is locked to the atoms and translates together with the displacement of the atoms.

To measure the phase ϕ of the converted output field of the $\hat{a}_-$ mode, we perform an interference experiment by directing the cavity forward and backward transmitted fields to a beam splitter and measuring the output photon numbers [Fig. 5(a)]. Displacing the atom array of four atoms by a distance X from any initial position, we plot the correlations of $\cos \phi_1$ and $\cos \phi_2$ in Fig. 5(c), where ϕ_1 and ϕ_2 are the phases of the cavity output field before and after displacing the atom array, respectively. We obtain the phase change $\Delta\phi = \phi_2 - \phi_1$ from the correlations and confirm that an arbitrary large optical phase shift can be achieved as $\Delta\phi = 2kX$ [Fig. 5(b)].

There are several directions for future research with this system. First, while the overall conversion efficiency is currently limited by the coupling ratio of the input and output mirrors, near-unity conversion can be realized by using a one-sided cavity with the same input and output mirror of dominant coupling [50] such that $\eta_{\text{in}} = \eta_{\text{out}} = 1$. This will enable efficient generation and routing of directional photons,

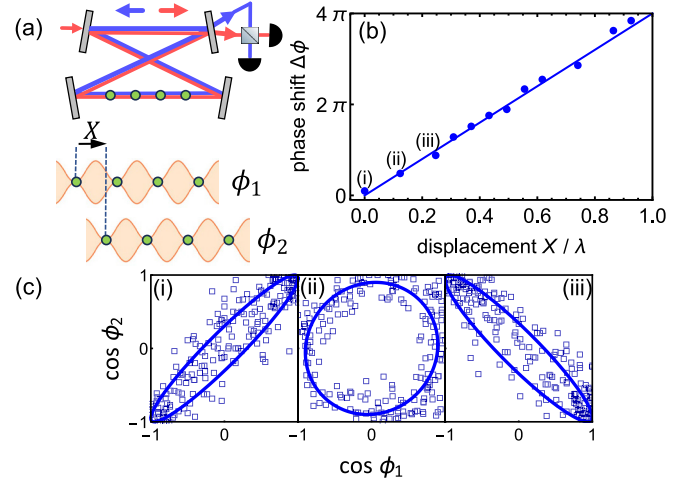


FIG. 5. Large phase shift of the cavity backward field imparted by the atom array. (a) Displace an atom array by distance X , and measure the phase shift $\Delta\phi$ of the output of the cavity backward field $\hat{a}_-$ through the interference experiment. (b) Arbitrarily large phase shifts can be realized by the atom array displacement. (c) The correlations of $\cos \phi_2$ and $\cos \phi_1$ before and after the atom array movement that are used to extract $\Delta\phi = \phi_2 - \phi_1$. The solid lines are fit to the data.

essential for applications in quantum networks. Second, optical cavities can be used for the midcircuit readout of atomic states [51]. However, changes in the atomic state and atom loss limit the readout fidelity. The cavity dark mode, resulting from strong interactions with atoms, carries information of the atomic state while eliminating atomic excitation, and offers another method to improve readout fidelity. Third, the ability of atom arrays to impart arbitrary large phase shifts suggests promising future directions. We plan to utilize this capability to store single photons in the array and realize photon phase gates [14,15] with arbitrarily large phases. Lastly, since the ring cavity offers two counterpropagating cavity modes and multiple photon output ports, coupling photons to atom arrays with arbitrary structures will enable the generation of entangled states of many photons [52] in different spatial directions. This integrated system of an optical ring cavity with atom arrays introduces possibilities for photon-photon interactions and all-optical quantum information processing based on the tuning geometrical structures of atom arrays.

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