

Collective strong coupling of thermal atoms to integrated microring resonators

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Strong coupling between atomic ensembles and high-quality optical cavities enables collective and nonlinear phenomena that are central to cavity quantum electrodynamics (cQED). Many experiments have been performed on this topic, with most of them focusing on cold atoms. Here, we experimentally demonstrate collective strong coupling between thermal rubidium vapor and high-quality silicon nitride microring resonators (MRRs) on an integrated photonic chip. We observe cavity-mode splitting, with a measured collective coupling strength of $g_N/2\pi \approx 1$ GHz and a collective cooperativity of $C_N \approx 2$ at 110 °C, indicating coherent energy exchange between the atomic ensemble and the cavity mode despite rapid decoherence in the thermal vapor system. We infer an average of 20 atoms effectively participating in the collective interaction, yielding a single-atom cooperativity of $C_0 = 0.1$ and approaching the single-atom strong-coupling regime. Our results establish the integrated thermal vapor MRR platform as a robust, compact, and scalable system for studying collective and nonlinear phenomena in cQED.

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I. INTRODUCTION

Research in cavity quantum electrodynamics (cQED) dates back to the 1940s, when Purcell showed that an emitter's radiative properties can be modified by its environment [1], particularly by an optical cavity. Theoretical descriptions of single-atom and many-atom cavity coupling followed in the 1960s [2,3], and strong-coupling signatures such as mode splitting were later observed with atomic beams [4–6], cold atoms, and trapped ions [7–13]. The cold atoms, in particular, offer long coherence time, low motional broadening, and single-particle control, which are essential for demonstrating extremely clean cQED. Recent advances in microfabrication have enabled high-finesse microcavities with small mode volumes, further enhancing atom-cavity interactions, with many strong-coupling demonstrations across different platforms such as miniaturized Fabry-Pérot cavities [14], microphotonic-crystal cavities [15], microbottle resonators [16,17], microtoroidal resonators [18], and microring resonators [19].

Despite these major successes of cold atoms interfaced with either macroscopic cavities or nanophotonic cavities,

cooling and trapping systems for cold atoms are necessary, which makes these experimental platforms often mechanically delicate, experimentally demanding, and bulky. These limitations could cause problems for applications where strong light-matter interaction is important, but individual-atom control is not essential and the footprint of the system matters. Therefore, we are motivated in combining nanophotonic microcavities with thermal atomic vapors. It takes the opposite philosophy of cold atoms and trades the single-particle control for scalability, compact footprint, robustness, and manufacturability. However, the key question is whether these advantages can be harnessed despite the associated thermal effects, without sacrificing essential coherent atom-light phenomena such as strong coupling.

Recent works have explored the feasibility of achieving strong coupling between thermal atoms and microcavities [20–22]. Despite the transit time decay loss and Doppler broadening associated with thermal atoms, experiments have nevertheless demonstrated collective strong coupling between thermal atomic vapor and high-quality microcavities, such as microspheres in vacuum chamber [23] or microring resonators (MRRs) in glued vapor cells [24]. However, microspheres in vacuum are not naturally suited for reproducible, lithographic, multidevice integration, while the glued vapor cell has limiting operation condition due to the performance of the glue.

In this work, we demonstrate collective strong coupling between thermal rubidium (Rb) vapor and integrated MRRs in a fully anodically bonded vapor cell, evidenced by a mode splitting of 2 GHz, corresponding to collective coupling strength $g_N/2\pi \approx 1$ GHz and a collective cooperativity of $C_N \approx 2$. We

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estimate that approximately 20 Rb atoms contribute to the collective coupling, corresponding to a single-atom cooperativity of $C_0 \approx 0.1$. By measuring the splitting systematically at different Rb vapor densities, we find good agreement with the theoretical prediction for collective coupling strength $g_N = \sqrt{N}g_0$ [3], where N is the average number of Rb atoms involved in the interaction and g_0 is the single-atom coupling strength. Furthermore, we also theoretically simulate the system with a simplified Tavis-Cummings model with input-output formalism that includes Doppler broadening and transient effects. The result of the simulation agrees well with the measurement.

The anodically bonded Rb-vapor-cell platform provides strong mechanical bonding, reliable hermetic vacuum sealing of 1×10^{-8} mbar over months, and endurance of operating temperatures above 200°C . The physical device used in this work occupies a small volume of $20 \times 30 \times 5 \text{ mm}^3$, less than a matchbox. Previous works have demonstrated that, with proper designs, hundreds of MRRs can be fabricated on a photonic chip and coupled to each other [25–27]. Such scalability enables the system to be extended to larger resonator networks and integrated into the small Rb-vapor-cell platform, without rebuilding the whole experiment. The only active control required for the atomic medium is regulation of the Rb vapor density through thermal control of the Rb reservoir. Compared with trapped single atom cQED systems, this platform sees challenges from thermal atomic motion, but it possesses low experimental complexity, compact footprint, scalability, and potential for fast, localized reconfiguration, making it a promising architecture for applications such as compact quantum sensing, cavity-enhanced spectroscopy, and nonlinear optics.

II. EXPERIMENTAL SETUP

We design and fabricate ring resonators from a 250-nm-thick Si_3N_4 layer deposited on thermally grown SiO_2 substrate on a Si wafer by low-pressure chemical vapor deposition (LPCVD). The ring radius is varied from 13 to $24 \mu\text{m}$. The ring width is set to $2 \mu\text{m}$ to confine the mode in the radial direction and reduce the scattering loss from sidewalls of the ring. A 1- μm -thick SiO_2 cladding layer is then deposited on top of the structured Si_3N_4 layer. Finally, the cladding above the ring resonator and the waveguide coupling region is locally removed [“opened,” Fig. 1(a)] to enable interaction between Rb atoms and the MRRs. After all fabrication steps, we dice the chip out of the wafer. The typical chip size is about $20 \times 30 \text{ mm}^2$.

To load Rb vapor into the vicinity of the MRR, we anodically bond the photonic chip to a borosilicate glass cell with a size of $15 \times 24 \times 5 \text{ mm}^3$, following the technical methods of previous works [22,28,29]. The size is not a fixed limitation, and is set by the current chip design and cell fabrication convention, which could be further miniaturized. An additional tube, serving as the rubidium reservoir, is attached to the glass cell and heated, separately [Fig. 1(b)]. The size of the Rb reservoir tube can be further reduced to reach even smaller footprint. A high-purity metallic rubidium source with natural isotopic composition ($\approx 72\%$ ^{85}Rb , $\approx 28\%$ ^{87}Rb) is then transferred and hermetically sealed into the reservoir

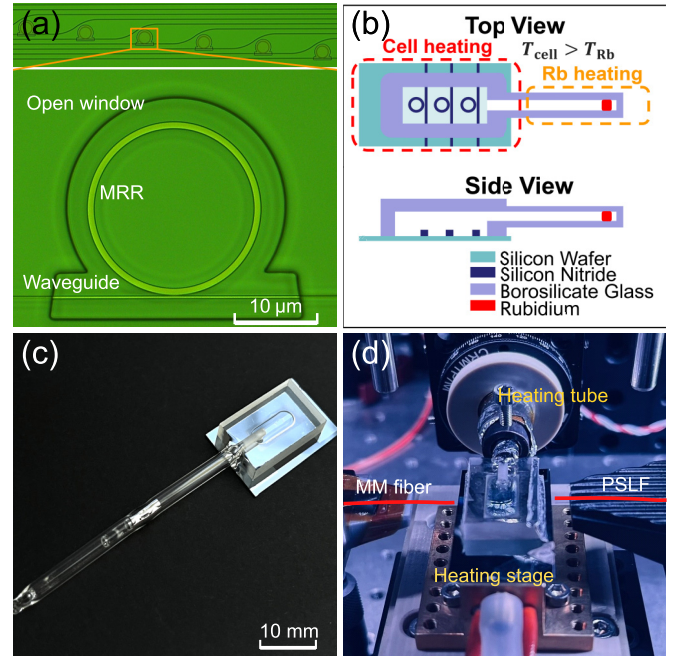


FIG. 1. (a) Optical microscope image of a group of MRRs on a chip. The enlarged view shows an MRR with radius of $10 \mu\text{m}$ coupled to a bus waveguide. The dark edge marks the opened window (image courtesy of IMS CHIPS). (b) Top and side-view schematics of the Rb vapor cell anodically bonded to an integrated photonic chip. The cell and Rb reservoir are independently heated, with $T_{\text{cell}} > T_{\text{Rb}}$. (c) Photograph of the anodically bonded Rb vapor cell, with Rb in the reservoir tube. A 10 mm scale bar is provided for reference. (d) Photograph of the vapor cell mounted in the heating systems. A polarization-maintaining single-mode lensed fiber (PSLF) couples light into the waveguide and a multimode (MM) fiber collects the transmission light, both fibers highlighted in red.

tube of the bonded cell under vacuum conditions (pressure below 1×10^{-7} mbar). This makes the complete integrated vapor cell platform for this work, which is shown in Fig. 1(c). To prevent Rb condensation on the device, the vapor cell is heated by two independent heaters, with the cell temperature maintained at least 20°C above the Rb reservoir temperature. Figure 1(d) shows the separate heating systems together with the fiber-coupling setup of the photonic chip. Details of the measurement scheme are provided in Fig. 5 of Appendix B.

III. MEASUREMENT AND RESULTS

We repeatedly scan the probe laser every 10 ms, while the MRR resonance is slowly swept across the Rb D_2 transitions at a rate of 0.4 GHz/s . To achieve this, a tightly focused IR laser is directed onto the MRR and intensity modulated, providing localized tuning of the cavity resonance through thermo-optic effect. Continuous acquisition of the probe transmission during this sweep yields the two-dimensional transmission map shown in Fig. 2. At $T_{\text{Rb}} = 30^\circ\text{C}$, the Rb vapor density near the MRR is negligible [Fig. 2(a)]. From the bare-cavity spectrum, we extract a loaded quality factor $Q_{\text{load}} = \omega_c/\kappa = 3 \times 10^5$ and a finesse $\mathcal{F} = 1700$, with $\kappa/2\pi = 1.2 \text{ GHz}$. We further infer an intrinsic loss rate $\kappa_i/2\pi = 1 \text{ GHz}$ from $Q_{\text{int}} = \omega_c/\kappa_i = 2Q_{\text{load}}/(1 + \sqrt{\mathcal{F}}_0)$,

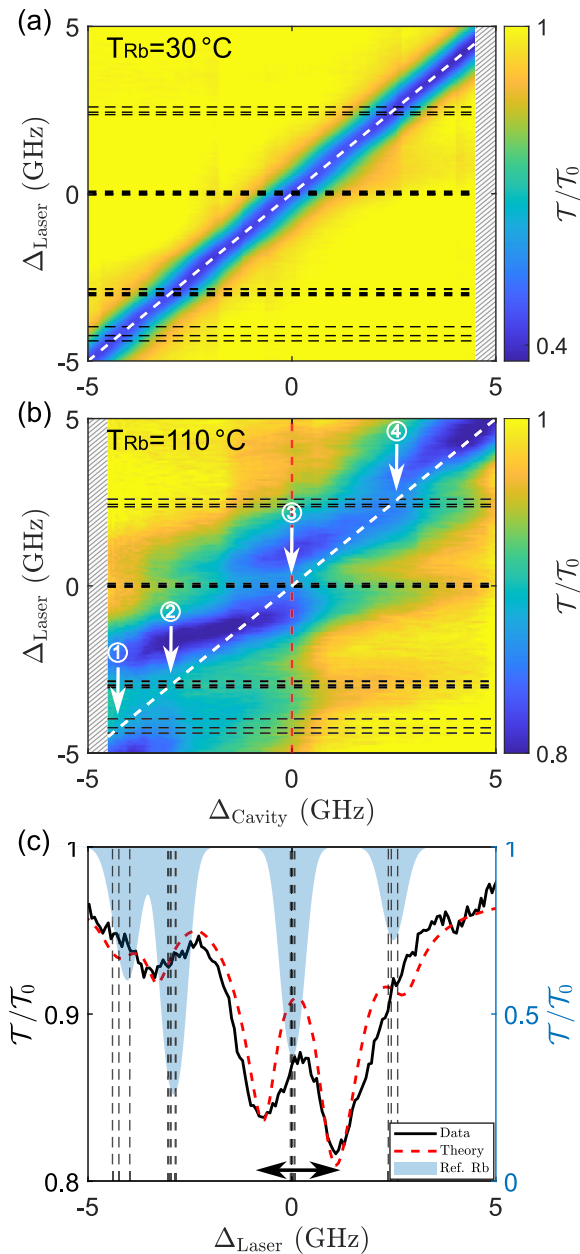


FIG. 2. (a), (b) Two-dimensional colormaps of the transmission spectrum at $T_{\text{Rb}} = 30$ and 110°C . Colorbars show the normalized transmission T/T_0 . The cavity-detuning axes are aligned; hatched regions indicate missing data. Horizontal black dashed lines denote the hyperfine transitions of ^{85}Rb and ^{87}Rb ; the white dashed line marks the MRR resonance. In panel (b), four white arrows mark the center-of-mass positions of the expected avoided crossings, and the red dashed line indicates the selected transmission slice near the ^{85}Rb $5^2S_{1/2}$ $F_g = 2 \rightarrow 5^2P_{3/2}$ $F_e = 2$ transition, presented in panel (c). In panel (c), the red dashed curve shows the theoretical prediction; the blue shaded region indicates the reference Rb vapor cell spectrum at the same T_{Rb} ; black dashed lines denote the same hyperfine transitions as in panels (a) and (b) and black double arrow indicates the mode splitting.

T_0 being the resonance transmission of the MRR, and an external coupling rate $\kappa_e/2\pi = 230\text{ MHz}$ [30].

With higher Rb density at $T_{\text{Rb}} \geq 80^\circ\text{C}$, avoided crossings emerge. The result at $T_{\text{Rb}} = 110^\circ\text{C}$ is shown in Fig. 2(b), in

which we expect four avoided crossings indicated by numbered white arrows. Based on the hyperfine structure of ^{85}Rb and ^{87}Rb (Fig. 6 in Appendix C), the excited-state splittings of the $5^2P_{3/2}$ manifold are smaller than the Doppler ($\sim 580\text{ MHz}$) and transit-time ($\sim 280\text{ MHz}$) broadenings at this temperature. As a result, transitions from the same ground hyperfine state merge into a single broadened avoided crossing, yielding one feature per ground state. However, only three distinct avoided crossings are observed. This is because the ground states ^{85}Rb ($5^2S_{1/2}$, $F = 3$) (arrow ①) and ^{87}Rb ($5^2S_{1/2}$, $F = 2$) (arrow ②) are separated by only $\sim 1.2\text{ GHz}$, which is smaller than 2 GHz separation of the other ground states, causing their corresponding avoided crossings to merge.

We obtain the mode-splitting spectrum at $T_{\text{Rb}} = 110^\circ\text{C}$ by slicing Fig. 2(b) along the red dashed line at arrow ③, defined as zero cavity detuning. The measured splitting is $\approx 2\text{ GHz}$, corresponding to a collective coupling strength of $g_N/2\pi \approx 1\text{ GHz}$. From the evanescent mode volume and vapor density, we estimate $N \approx 20$ participating ^{85}Rb atoms at 110°C (see Appendix C). The splitted spectrum is asymmetric about the atomic transition due to off-resonant collective coupling to the ^{85}Rb ($5^2S_{1/2}$, $F_g = 3$) $\rightarrow$ ($5^2P_{3/2}$, $F_e = 2, 3, 4$) transitions. This is supported by our theoretical simulations including all hyperfine transitions of the Rb D_2 manifold (Fig. 8 in Appendix D). The simulated transmission spectrum at zero cavity detuning at $T_{\text{Rb}} = 110^\circ\text{C}$ is shown as the red dashed curve in Fig. 2(c).

When the MRR is detuned a few gigahertz from the Rb transitions, it approaches the bare-cavity regime. We observe that the resonance linewidth broadens while the contrast decreases with increasing T_{Rb} , as shown in Fig. 3(a). The cavity photon decay rate κ and contrast A_0 are extracted by fitting a single Lorentzian function to the normalized bare-cavity transmission spectrum; the error bars in Fig. 3(a) represent the fitting uncertainties. Consequently, measurements are limited to $T_{\text{Rb}} \leq 110^\circ\text{C}$, as at higher temperatures the resonance feature becomes barely discernible. The increase in Rb vapor density introduces additional loss, likely due to scattering from Rb condensation or Rb diffusion on the MRR. This loss is reversible, and the original resonance line shape is recovered upon reducing the vapor density by lowering the reservoir temperature.

Figures 3(b) and 3(c) summarize the measured coupling strength g_N and cooperativity C_N as functions of the average atom number N at zero cavity detuning. Each data point is obtained from repeated measurements, with vertical error bars indicating statistic errors and horizontal error bars corresponding to the uncertainty in N due to temperature fluctuations of $\pm 0.5^\circ\text{C}$. Figure 3(b) is fitted with $g_N = \sqrt{N}g_0$, yielding $g_0/2\pi \approx 225\text{ MHz}$, consistent with the theoretical value at the maximum evanescent field (Fig. 7 in Appendix C). Figure 3(c) shows the expected cooperativity $C = g_N^2/(\kappa\gamma)$ (black), calculated using $\kappa/2\pi = 1.2\text{ GHz}$ and $\gamma/2\pi \approx 240\text{ MHz}$ at $T_{\text{Rb}} = 30^\circ\text{C}$, together with a linear fit (black dashed line). In practice, κ varies with T_{Rb} [Fig. 3(a)], and $\gamma = \langle v_{\text{Rb}} \rangle / \delta_{1/e}$ scales with $\sqrt{T_{\text{Rb}}}$. Using $\kappa(T_{\text{Rb}})$ and $\gamma(T_{\text{Rb}})$ leads to the measured cooperativity shown in orange in Fig. 3(c), giving $C_N \approx 2$ for ~ 20 ^{85}Rb atoms.

The mode splitting at a given T_{Rb} depends on the probe power and follows the saturation behavior, as shown in

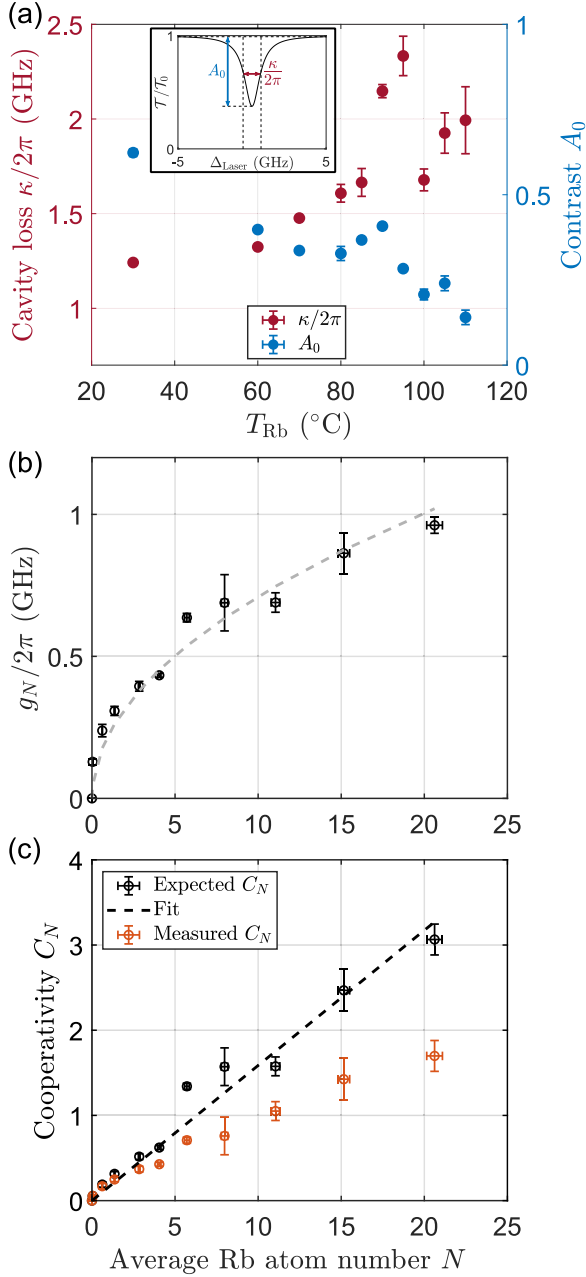


FIG. 3. (a) The temperature dependence of cavity-photon loss rate κ (in red) and the resonance contrast A_0 (in blue). Inset shows explicitly how κ and A_0 are extracted from the bare-cavity resonance profile. (b) Measured collective coupling strength g_N . The gray dashed line shows a fit with $g_N = g_0\sqrt{N}$. This panel shares a common x axis with panel (c). (c) Calculated cooperativity with $C_N = \frac{g_N^2}{\kappa\gamma}$. Black data correspond to C_N with κ at $T_{\text{Rb}} = 30^\circ\text{C}$, black dashed line is a linear fit to the black data with $C = NC_0$, and orange data show C_N with $\kappa = \kappa(T_{\text{Rb}})$.

Fig. 4(a). This reflects a reduction in the number of atoms that respond linearly to the probe field. Within our theory model presented in Appendix D, the collective coupling scales with the population difference $\langle\sigma^z\rangle$ [see Eq. (D7)]. For weak excitation, $\langle\sigma^z\rangle \neq 0$, and atoms couple linearly to the cavity, producing mode splitting. As the probe power exceeds the saturation threshold, $\langle\sigma^z\rangle \rightarrow 0$, reducing the effective

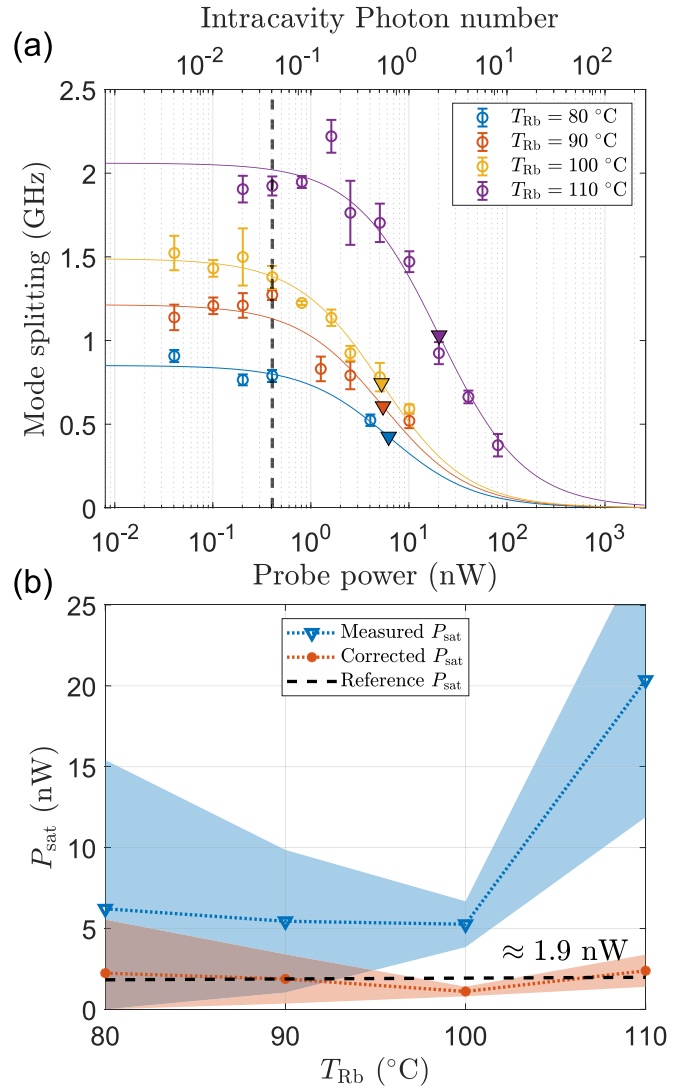


FIG. 4. (a) Measured mode splitting depending on probe power inside the waveguide at different T_{Rb} . Solid lines are fits using the saturation relation $S = S_0/(1 + P/P_{\text{sat}})$, where S_0 denotes the weak-probe splitting and S is the power-dependent splitting [24]. Triangles denote the fitted saturation power for each T_{Rb} . Black dashed line marks the probe power for measuring all cooperativity data. (b) Extracted saturation power from panel (a) (in blue) and corrected saturation power (in orange), with dotted lines as guidance to the eye and shaded region as the uncertainty of the fitting in panel (a). Black dashed line is the reference P_{sat} power inferred with $\kappa(T_{\text{Rb}} = 30^\circ\text{C})$ and $\langle n \rangle_{\text{sat}} \approx 0.74$.

number of participating atoms; the system approaches the bare-cavity regime and the splitting vanishes. The extracted saturation power at each T_{Rb} (triangles) is compared with the probe power of 0.45 nW (black dashed line) used for the cooperativity measurements, confirming operation well below saturation.

The fitted saturation power P_{sat} (blue) in Fig. 4(b) increases from 6.5 nW at $T_{\text{Rb}} = 80^\circ\text{C}$ to above 20 nW at 110°C . To explain this trend, we estimate the intracavity photon number $\langle n \rangle_{\text{sat}} = \gamma^2/(2g^2)$ [31] required to saturate an atom at the maximum evanescent field, with $g/2\pi = 220$ MHz and $\gamma = \langle v_{\text{Rb}} \rangle / \delta_{1/e} \propto \sqrt{T_{\text{Rb}}}$. Over $80^\circ\text{C} \leq T_{\text{Rb}} \leq 110^\circ\text{C}$, we find

a weak increase of $\langle n \rangle_{\text{sat}}$ from 0.71 to 0.77, allowing us to approximate $\langle n \rangle_{\text{sat}} \approx 0.74$.

The corresponding input power follows from the steady-state solution of Eq. (D5),

$$P_{\text{sat}} = \langle n \rangle_{\text{sat}} \hbar \omega \frac{\kappa^2}{4\kappa_e}, \quad (1)$$

revealing a dependence on $\kappa^2/(4\kappa_e)$, which increases with T_{Rb} [Fig. 3(a)] and accounts for the observed rise in P_{sat} . To isolate this effect, we normalize each measured P_{sat} by $\kappa^2/(4\kappa_e)$ relative to its initial-state value at $T_{\text{Rb}} = 30^\circ\text{C}$. The normalized values (orange) collapse to $P_{\text{sat}} \approx 1.9\text{ nW}$, consistent with $\langle n \rangle_{\text{sat}} \approx 0.74$.

IV. CONCLUSION

In this work, we establish thermal-vapor integrated photonic MRR cell as a promising platform for cQED. While this platform trades single-atom control for a compact footprint and scalability, collective interactions still enable access to the strong-coupling regime. We demonstrate this by observing collective strong coupling in this platform using a high- Q ($\approx 3 \times 10^5$) Si_3N_4 microring resonator coupled to thermal atoms in the anodically bonded vapor cell. Cavity-mode splitting confirms the strong-coupling condition $g_N^2/(\kappa\gamma) \geq 1$, with the extracted coupling strength and cooperativity following the expected scaling with atom number. At $C_N \approx 2$, corresponding to ~ 20 ^{85}Rb atoms, we infer a single-atom cooperativity of $C_0 \approx 0.1$. We also develop a theoretical model that reproduces the experimental results (see Appendix D).

These results mark progress toward the single-atom strong-coupling regime ($C_0 \geq 1$) with this thermal-vapor-cell platform. Reaching this limit will require reduced mode volume, enhanced evanescent-field overlap, and higher Q , achievable with photonic-crystal cavities [15,20], suspended microring geometries [32], and reduced surface scattering loss [33].

Beyond the single MRR system demonstrated in this work, the compact footprint and inherent scalability of this platform make it particularly attractive for extension to coupled multicavity architectures [25–27] and enable on-chip quantum networks with localized reconfiguration [34]. This platform may enable applications in on-chip quantum networks [35–37], nonlinear optics [38–40], chip-scale quantum photonics [18,41], and precision spectroscopy [42,43].

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The authors declare no conflicts of interest.

DATA AVAILABILITY

The data that support the findings of this article are openly available [44].

APPENDIX A: NUMERICAL DESIGN AND NANOFABRICATION OF MRR CHIPS

The design of the MRR structures requires consideration of several key properties, including the Q factor, critical coupling condition, free spectral range, and resonance wavelength. These properties depend on physical parameters such as the propagation loss α_{loss} , the external coupling rate κ_e between the bus waveguide and the MRR, and the ring size, which are in turn determined by the geometrical parameters: ring width W , gap distance G , and ring radius R .

Rather than searching numerically for a single optimum geometry, which may be shifted by fabrication errors, we use simulations to evaluate the sensitivity of the MRR performance to variations in these geometrical parameters and to determine appropriate parameter sweep ranges and step sizes. The ring width W is fixed at $2\text{ }\mu\text{m}$ for all MRRs to reduce bending-induced propagation loss associated with the small ring radius. The gap distance G is varied from 70 to 290 nm in steps of 20 nm to cover a broad range of coupling conditions. Three groups of rings with nominal radii of 13, 19, and 24 μm are implemented, with the radius in each group varied by $\pm 40\text{ nm}$ in steps of 10 nm. This design accounts for fabrication uncertainties and increases the probability of obtaining MRRs with the desired resonance frequency.

The MRR structures are fabricated on a 150 mm silicon wafer. We apply a 3.3- μm -thick buried oxide layer of SiO_2 grown by wet thermal oxidation. This is followed by depositing 250-nm-thick Si_3N_4 as a device layer using low-pressure chemical vapor deposition. The device patterns are first defined in a photoresist layer by electron beam lithography and subsequently transferred into the underlying 250-nm-thick Si_3N_4 layer by reactive ion etching. The patterned Si_3N_4 layer is then covered with a SiO_2 cladding layer with an LPCVD process based on tetraethyl orthosilicate. Afterward, the wafer undergoes chemical mechanical polishing and is coated with a polycrystalline silicon layer to obtain an optically flat and anodically bondable surface for integration with the vapor cell. Finally, to enable interaction between the MRR mode and the atomic vapor, the upper cladding is locally removed around each MRR by a selective etching process using a polysilicon hardmask and buffered hydrofluoric acid.

APPENDIX B: MEASUREMENT SETUP AND LOCALIZED CAVITY TUNING

Figure 5(a) shows the schematic of the measurement setup. A tunable laser at 780 nm is split into two beams: One probes the transmission spectrum of a reference Rb cell, providing a frequency reference for the measurement, while the other, together with a tracking laser at 980 nm, is coupled into a PSLF. The PSLF is mounted on a piezo-stabilized stage and couples light into the waveguide via edge coupling. A MM fiber collects the transmitted light, which is then separated by a dichroic mirror into the 780 nm signal and the 980 nm tracking signal, the latter serving as a

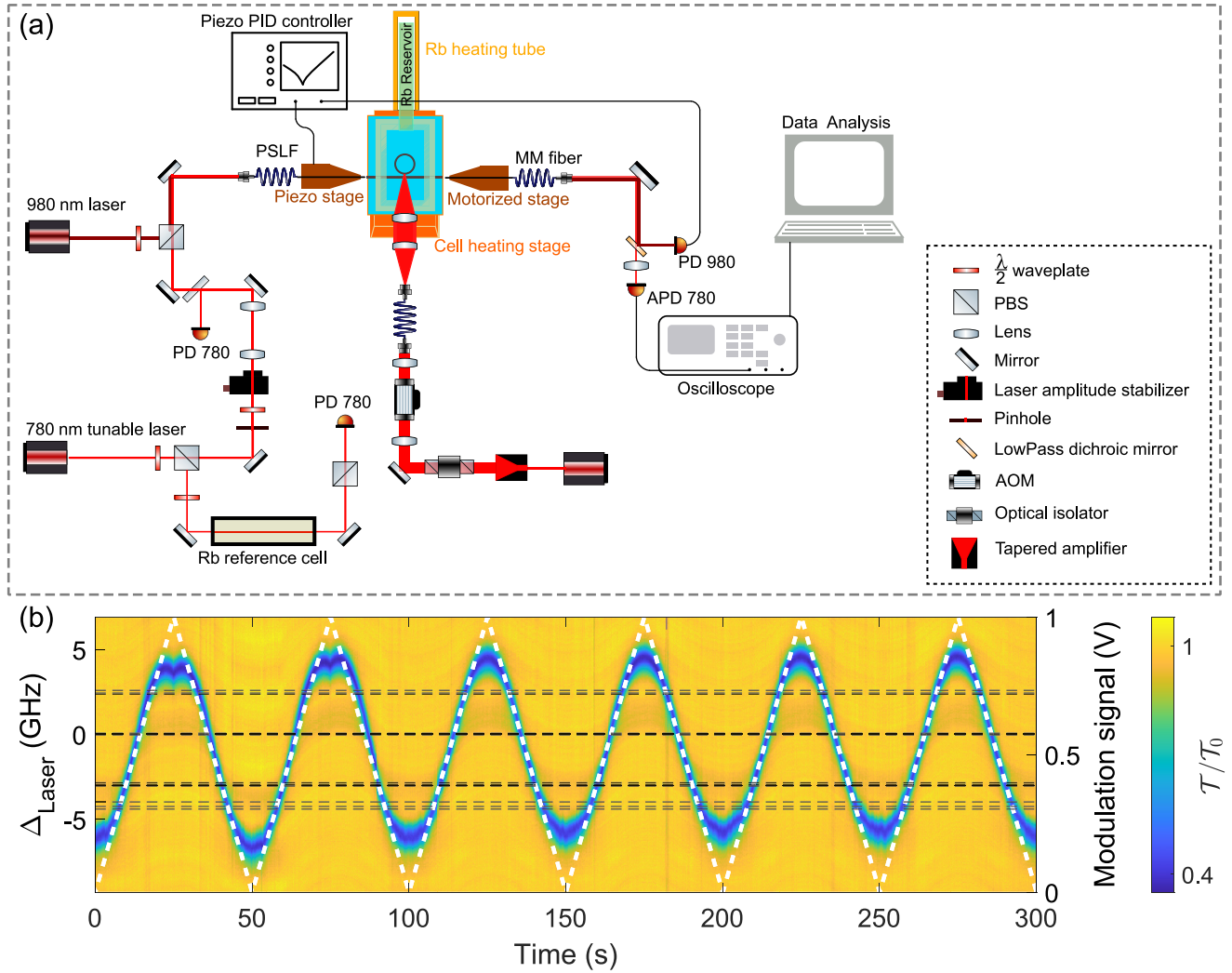


FIG. 5. (a) Schematic of the measurement setup. A 780 nm tunable laser sends in a weak probe laser, while a 980 nm laser sends in a strong laser to stabilize the fiber coupling efficiency. An infrared laser locally tunes the resonance of the MRR. (b) Time-resolved transmission map of the bare MRR resonance tuning under the illumination of the AOM modulated infrared laser (white dashed lines) with a period of 50 s. The probe detuning Δ_{Laser} is referenced to the ^{85}Rb $5S_{1/2}, F_g = 2 \rightarrow 5P_{3/2}, F_e = 3$ transition, while black dashed lines mark the relevant Rb D_2 hyperfine transitions.

feedback signal for the piezo-stabilized stage. The piezo-stabilized stage compensates the thermal drift and vibration, stabilizing the PSLF coupling efficiency at an optimum level during the measurement.

An infrared laser is amplified and tightly focused onto the MRR to locally tune the resonance via the thermo-optic effect, with its intensity modulated by an acousto-optic modulator (AOM). Figure 5(b) shows the recorded transmission map of the bare MRR resonance tuning under the illumination of the AOM modulated infrared laser. The resulting resonance tuning range is approximately 10 GHz. For an AOM modulation period of 50 s, the MRR responds sufficiently rapidly to follow the IR-induced tuning with negligible hysteresis. The MRR resonance shift does not strictly follow the triangular modulation signal due to the fact that the AOM amplitude modulation of the laser is not perfectly linear. However, one could easily linearize the MRR resonance position with the bare MRR transmission map at low Rb temperature, such as $T_{\text{Rb}} = 30^\circ\text{C}$.

The linearized transmission map is shown in Fig. 2(a). The same linearization data at $T_{\text{Rb}} = 30^\circ\text{C}$ can provide calibration for transmission maps measured at higher T_{Rb} .

The local tuning rate of the MRR is chosen to be sufficiently slow (0.4 GHz/s) compared with the probe-laser scan speed (1 THz/s), so that during each 10 ms probe-laser scan, the cavity resonance shift is negligible. In this way, one could treat each probe scan as a time bin and construct a full transmission map of the MRR resonance at different cavity detuning positions. The local infrared laser tuning allows fast and local reconfiguration of MRR devices in the present integrated thermal vapor platform, avoiding additional on-chip heater structures. Further improvements in tuning speed could be achieved by integrating phase-change materials [45,46] or by exploiting the electro-optic response of the photonic material platform [47], thereby enabling rapid local tuning and dynamic reconfiguration of MRR-based systems when needed.

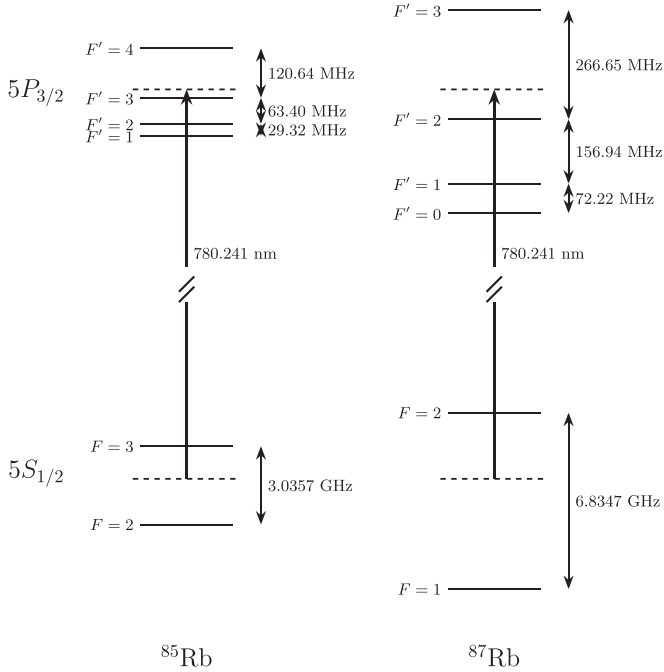


FIG. 6. Hyperfine level structure of the D_2 transition manifolds of ^{85}Rb and ^{87}Rb [51]. The hyperfine excited-state splittings are smaller than the Doppler broadening and transient decoherence rate, resulting in the smearing and merging of avoided crossings in Fig. 2(b).

APPENDIX C: COUPLING STRENGTH $g_0(\mathbf{r})$ AND AVERAGE Rb ATOM NUMBER IN THE MODE VOLUME

We simulate the electric field distribution $E(\mathbf{r})$ of the fundamental TE_{00} cavity mode of the MRR. From the simulated mode distribution, we extract the effective refractive index $n_{\text{eff}} = 1.81$, the evanescent-field decay length $\delta_{1/e} \approx 90$ nm, and the cavity mode volume $V_{\text{mode}} = 10.1 \mu\text{m}^3$.

For a single atom in the evanescent field of an MRR mode containing a single photon, we calculate its position-dependent single-atom-single-photon coupling strength $g_0(\mathbf{r})$ using the standard definition [48]:

$$g_0(\mathbf{r}) = \frac{d}{\hbar} \cdot E_{\text{ph}}(\mathbf{r}), \quad (\text{C1})$$

where d denotes the Rb D_2 transition dipole moment, with the contribution from all hyperfine transitions (Fig. 6) weighted according to their relative strengths and the natural abundances of the two Rb isotopes. $E_{\text{ph}}(\mathbf{r})$ is the cross-sectional evanescent electric field distribution of the MRR mode, defined as

$$E_{\text{ph}}(\mathbf{r}) = \sqrt{\frac{\hbar\omega}{2\varepsilon_0 V_{\text{mode}}}} \cdot \psi(\mathbf{r}), \quad (\text{C2})$$

where $\hbar\omega$ is the single-photon energy at a wavelength of 780 nm and ε_0 is the vacuum permittivity. The cavity mode volume V_{mode} tells how tightly the electromagnetic energy of the cavity mode is confined in space, and $\psi(\mathbf{r})$ is the normalized mode function that captures the location dependence of the electromagnetic field of the mode. In practice, we could numerically simulate the mode shape of the MRR and

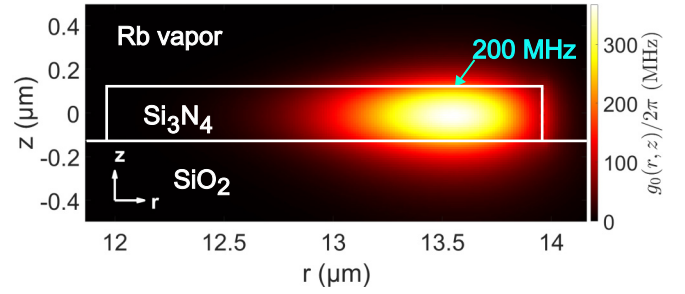


FIG. 7. Calculated single-atom coupling strength $g_0(\mathbf{r})$. White solid lines indicate the waveguide cross-sectional boundaries and the substrate boundary, with the Si_3N_4 core and the SiO_2 substrate beneath it, while the surrounding region above is filled with Rb vapor. The blue arrow indicates the location of the maximum single-atom coupling strength on the surface of the MRR.

calculate the single-photon mode-field distribution E_{ph} . Then, we could infer the coupling strength $g_0(\mathbf{r})$, as shown in Fig. 7.

The definition of the cavity mode volume V_{mode} and the relation with the mode function $\psi(\mathbf{r})$ is given by [23]

$$V_{\text{mode}} = \frac{\int \varepsilon(\mathbf{r}) |E(\mathbf{r})|^2 d\mathbf{r}}{\max(\varepsilon(\mathbf{r}) |E(\mathbf{r})|^2)} = \int d^3r |\psi(\mathbf{r})|^2. \quad (\text{C3})$$

The denominator of the definition of the cavity mode volume usually takes the global maximum of the electromagnetic energy density, including the value inside the Si_3N_4 core. The mode function is normalized to the cavity mode volume such that its value at the maximum electromagnetic energy density is approximately equal to 1.

On the other hand, the thermal Rb atoms are randomly distributed in the space surrounding the MRR. Only the evanescent mode field, which is outside the Si_3N_4 material of the MRR cavity, is accessible to the atoms. Therefore, the actual volume of the cavity mode interacting with atoms is not the cavity mode volume, but a small fraction of it. We estimate the evanescent-field mode volume within the vapor by taking the denominator of Eq. (C3) to be the local maximum of the electromagnetic energy density in the vapor surrounding the MRR. To account for the portion of the cavity mode accessible to the atoms, we likewise restrict the integration in the numerator to the vapor region:

$$V_{\text{evan.}} = \frac{\int_{\text{vapor}} \varepsilon(\mathbf{r}) |E(\mathbf{r})|^2 d\mathbf{r}}{\max(\varepsilon(\mathbf{r}) |E(\mathbf{r})|^2)_{\text{vapor}}}, \quad (\text{C4})$$

where $\varepsilon(\mathbf{r})$ and $E(\mathbf{r})$ are the position-dependent permittivity and electric-field amplitude. The denominator takes the maximum evanescent-field energy density in the vapor as the reference, and the evanescent mode volume is $V_{\text{evan.}} = 2.54 \mu\text{m}^3$.

With the Rb vapor number density of $\rho = 11.23 \mu\text{m}^{-3}$ at $T_{\text{Rb}} = 110^\circ\text{C}$ [49,50], and the natural abundance of ^{85}Rb ($\eta_{^{85}\text{Rb}} = 72.17\%$) that contributes to the measured mode splitting in Fig. 2(c), the average number of ^{85}Rb atoms coupled to the cavity evanescent field at zero cavity detuning is estimated to be

$$N_{\text{eff}} = V_{\text{evan.}} \times \rho \times \eta_{^{85}\text{Rb}} \approx 20. \quad (\text{C5})$$

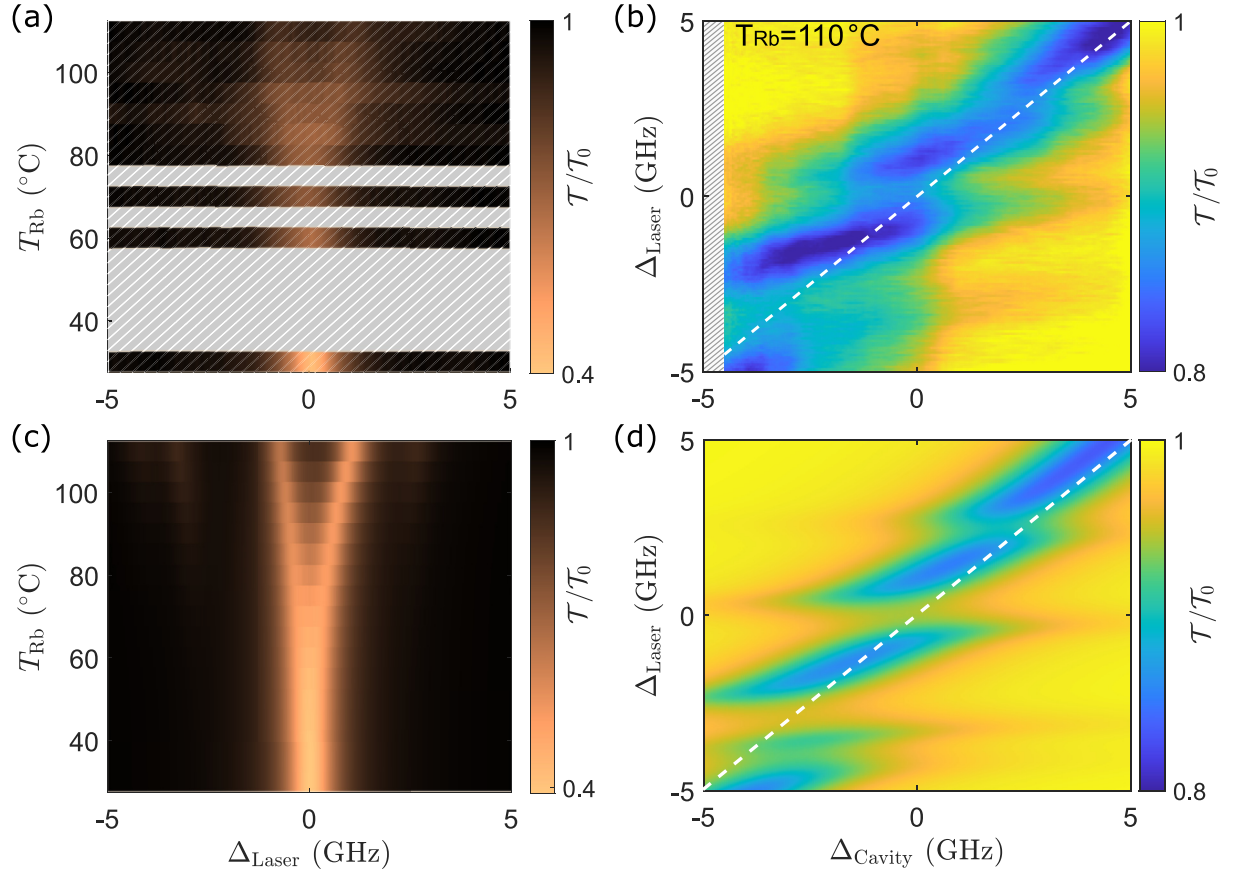


FIG. 8. (a) Measured transmission spectra for different T_{Rb} at zero cavity detuning. The colorbar represents the normalized transmission T/T_0 . For $T_{\text{Rb}} \leq 80$ °C, hatched region denotes missing data. (b) Measured transmission map at $T_{\text{Rb}} = 110$ °C. The white dashed line tracks the MRR resonance detuning. (c), (d) Theoretically simulated results of panels (a) and (b), sharing the same x axis with panels (a) and (b), respectively.

With this estimated Rb atom number, we calculate a single-atom coupling strength of $\bar{g}_0/2\pi \approx 220$ MHz, consistent with the calculated coupling strength $g_0(\mathbf{r})$ at the location of maximum evanescent field in Fig. 7.

APPENDIX D: THEORETICAL MODEL

We model the atom-cavity interaction using the Tavis-Cummings Hamiltonian, and treat dissipation and coupling to external fields within the Gardiner-Collett input-output formalism [52], in which the total Hamiltonian of our system is described by

$$H = H_{\text{sys}} + H_{\text{bath}} + H_{\text{int}}. \quad (\text{D1})$$

The closed cavity-atom system is described by the Tavis-Cummings Hamiltonian:

$$H_{\text{sys}} = \hbar\omega_c a^\dagger a + \sum_{n=1}^N \hbar\omega_{a,n} \sigma_n^+ \sigma_n^- + \sum_{n=1}^N \hbar g_n (a^\dagger \sigma_n^- + a \sigma_n^+), \quad (\text{D2})$$

where ω_c is the cavity frequency, $\omega_{a,n}$ is the atomic transition frequency of atom n , and g_n is the corresponding atom-cavity coupling strength. Here, a is the cavity-mode annihilation operator, and σ_n^+ and σ_n^- are the raising and

lowering operators of atom n . The external baths consist of the bus-waveguide channel and the intrinsic cavity-loss channel:

$$H_{\text{bath}} = \hbar \int \omega c^\dagger(\omega) c(\omega) d\omega + \hbar \int \omega d^\dagger(\omega) d(\omega) d\omega, \quad (\text{D3})$$

where $c(\omega)$ and $d(\omega)$ are annihilation operators at angular frequency ω for the bus-waveguide continuum mode and intrinsic-loss bath mode, respectively.

The interaction Hamiltonian between the closed system and the baths is

$$H_{\text{int}} = i\hbar \int \kappa_e(\omega) [c^\dagger(\omega) a - a^\dagger c(\omega)] d\omega + i\hbar \int \kappa_i(\omega) [d^\dagger(\omega) a - a^\dagger d(\omega)] d\omega. \quad (\text{D4})$$

Here, $\kappa_e(\omega)$ is the amplitude coupling between the cavity and the bus continuum, and $\kappa_i(\omega)$ is the amplitude coupling between the cavity and the intrinsic-loss bath. Under the Heisenberg picture, we derive the equations of motion from the Heisenberg-Langevin equations. With the rotating-wave approximation, Markovian approximation, and mean-field

approximation (valid in our strongly decoherent system), the equations in the laboratory frame reduce to

$$\dot{a} = -\left(\frac{\kappa}{2} + i\omega_c\right)a - i \sum_{n=1}^N g_n \sigma_n^- + \sqrt{\kappa_e} a_{\text{in}} + \sqrt{\kappa_i} a_{0,\text{in}}, \quad (\text{D5})$$

$$\kappa \equiv \kappa_e + \kappa_i, \quad \kappa_e = 2\pi |\kappa_e(\omega_c)|^2 \rho(\omega_c), \quad (\text{D6})$$

$$\dot{\sigma}_n^- = -\left(\frac{\gamma_n}{2} + i\omega_{a,n}\right)\sigma_n^- + ig_n a \sigma_n^z, \quad (\text{D7})$$

$$a_{\text{out}} = a_{\text{in}} - \sqrt{\kappa_e} a. \quad (\text{D8})$$

Here, κ_e is the observable cavity energy-decay rate into the bus waveguide, and κ_i is the intrinsic cavity-loss rate. $\rho(\omega_c)$ is the bath density of states. a_{in} is the input field operator oscillating at the probe frequency ω_p in the bus waveguide, and $a_{0,\text{in}}$ is the vacuum-noise input operator (which vanishes in expectation-value calculations). γ_n is the atomic polarization decay rate of atom n , including spontaneous decay and transient effects. σ_n^z is the Pauli- z operator, whose expectation value gives the population difference $\langle \sigma^z \rangle = \rho_{ee} - \rho_{gg}$, where ρ_{ee} and ρ_{gg} are the excited- and ground-state populations, respectively. Under weak excitation, $\sigma_n^z \approx -1$. a_{out} is the output field operator in the bus waveguide. The probe laser is introduced as a coherent input in the waveguide channel within the input-output formalism, such that the input operator is replaced by its expectation value $a_{\text{in}} \rightarrow \alpha_{\text{in}} e^{-i\omega_p t}$ in the Heisenberg-Langevin equations, with ω_p as the probe laser frequency. Here, we select a single cavity mode by choosing the helicity set by the propagation direction of light in the waveguide.

By solving the steady-state expectation values of these operator equations in the rotating frame, we obtain the bus-waveguide transmission function:

$$T(\omega_p) = \left| \frac{\alpha_{\text{out}}}{\alpha_{\text{in}}} \right|^2 = \left| 1 - \frac{\kappa_e}{\frac{\kappa}{2} - i\Delta_c + \sum_{n=1}^N \frac{g_n^2}{\frac{\gamma_n}{2} - i\Delta_{a,n}}} \right|^2, \quad (\text{D9})$$

with

$$\Delta_c \equiv \omega_p - \omega_c, \quad \Delta_{a,n} \equiv \omega_p - \omega_{a,n}. \quad (\text{D10})$$

Notably, a similar transmission function can be derived from a classical analog model of two linearly coupled damped oscillators. In the weak-excitation limit, the Tavis-Cummings model becomes linear; therefore, a coherently driven quantum system evolves equivalently to a system of two classically coupled linear oscillators. As a result, transmission spectroscopy of this system does not require treatment beyond a linear dispersive model, as has been studied in the 1990s with atomic beams [4–6]. We simulate the transmission spectrum with the derived transmission function [Eq. (D9)] using Monte Carlo sampling of atomic positions and thermal velocities to include Doppler shifts and transient decoherence, together with the position-dependent single-atom coupling strength $g_0(\mathbf{r})$ for each atom. We plot the simulated transmission spectrum at the same cavity-atom detuning as in Fig. 2(c) for each T_{Rb} , and compare with the measured spectra, as shown in Fig. 8. The simulations are in good agreement with the measurements and reveal the asymmetry of the mode splitting due to the off-resonance collective coupling from other hyperfine transitions.

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