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An inhibited laser

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Traditional lasers function using resonant cavities, in which the round-trip optical path is exactly equal to an integer multiple of the intra-cavity wavelengths to constructively enhance the spontaneous emission rate. By taking advantage of the resonant cavity enhancement, the narrowest sub-10-mHz-linewidth laser [1] and a 10^{-16} -fractional-frequency-stability superradiant active optical clock (AOC) [2] have been achieved. However, never has a laser with atomic spontaneous radiation being destructively inhibited [3] in an anti-resonant cavity where the atomic resonance is exactly between two adjacent cavity resonances been proven. Herein, we present the first demonstration of the inhibited stimulated emission, which is termed an inhibited laser. Compared with traditional superradiant AOCs [4–9] exhibiting superiority for the high suppression of cavity noise in lasers, the effect of cavity pulling on the inhibited laser’s frequency can be further suppressed by a factor of $-(2\mathcal{F}/\pi)^2$. This study of the inhibited laser will guide further development of superradiant AOCs with better stability, and may lead to new searches in the cavity quantum electrodynamics (QED) field.

The significantly enhanced spontaneous decay rate of the spin in a resonant circuit, known as the Purcell effect [10], was first reported by Purcell in 1946. It was practically observed in the 1980s using atoms in resonant cavities both in the microwave [11] and optical [12, 13] domains. The enhanced spontaneous radiation has important application potential in cavity quantum electrodynamics (QED) [14–16], including for one-atom lasers [17], ion-trap lasers [18], and quantum logic gates [19] in quantum computers.

Essentially, the resonant cavity, whose cavity-mode frequency resonates with the peak of the emission line for atomic transition, enhances the strength of vacuum fluctuations, which promotes the atomic spontaneous radiation. Conversely, the spontaneous decay rate is suppressed when the cavity is off resonance, which was first proposed by Kleppner in 1981 and demonstrated through inhibited blackbody absorption [20] and inhibited spontaneous emission [3]. After, inhibited spontaneous emission was experimentally demonstrated in microwave and optical cavities in 1985 [21] and 1987 [12], respectively. Heinzen [12] pointed out that in an anti-resonant cavity, where the atomic frequency was exactly between two adjacent cavity resonances, the inhibition of the atomic spontaneous decay rate was the greatest. More strikingly,

through coupling with an anti-resonant cavity, the atomic radiative level shift vanished and the spectral linewidth decreased [22], which is potentially useful for precision measurements. Despite the experimental success of inhibited spontaneous emission, the working mechanism of inhibited laser is unknown.

Nevertheless, the demonstration of inhibited spontaneous emissions has provided credible evidence for the observation of inhibited stimulated emissions. The spontaneous emission can be viewed as a stimulated emission originating from the vacuum fluctuations, and the spontaneous emission below the threshold determines the spectrum of the laser above the threshold [23]. It has significant potential to achieve inhibited lasing, with the aid of a three- or four-level structure to increase the pumping efficiency and the multi-atom system to reach the strong-coupling regime [24].

Here, we report the first experimental demonstration of an inhibited laser. The general setup is depicted schematically in Fig. 1, sharing similarities with the proposed superradiant AOC based on thermal atoms [25]. $N \approx 1.8 \times 10^{11}$ pure Cs atoms are confined to the TEM₀₀ mode of a low-finesse optical cavity ($\mathcal{F} = 3.07$), whose dissipation rate is $\kappa = 2\pi \times 257$ MHz. Pumped by a 459 nm laser ($6S_{1/2}$ - $7P_{1/2}$), the atoms achieve stimulated emissions at a 1470 nm transition ($7S_{1/2}$ - $6P_{3/2}$). The relaxation rate of the atomic dipole $\Gamma = 2\pi \times 10.04$ MHz is much smaller than κ , which forms a bad-cavity regime [4, 5]. Unlike traditional resonant lasers, the inhibited laser is realized with a round-trip optical path equal to odd multiples of the half wavelength $2L = (2q + 1)\lambda/2$.

Suppose that the atom emitting the first photon by spontaneous radiation is located at the center of cavity, and the reflectivities of cavity mirrors are $R_1 = R_2 = R$, the ratio of the power of spontaneous radiation emitted into cavity P_c to the power into free space P_{free} is given by [22]

$$\frac{P_c}{P_{\text{free}}} = \frac{1 - R^2}{1 + R^2 - 2R \cos(\omega 2L/c)}, \quad (1)$$

where ω is the angular frequency of the radiation, and c the speed of light. $\Delta\phi = \omega 2L/c$ denotes the phase shift of the intracavity reflected field, and it also reflects the detuning of the cavity frequency ω_c from the atomic resonance ω_0 . A phase shift of 2π between two consecutive round trips of the radiation inside the cavity corresponds to the cavity-frequency detuning $\omega_c - \omega_0$ of one free spectral range (FSR). According to Eq. (1), the spontaneous decay rate from the atomic excited state is enhanced and

FIG. 1. **Working principle of the inhibited laser.** **a**, Level scheme for ^{133}Cs atom showing the 459 nm transition used for pumping and the 1470 nm transition as the inhibited lasing. Cs atoms are pumped by a 459 nm laser, which drives the $|1\rangle$ ($6S_{1/2}$) to $|2\rangle$ ($7P_{1/2}$) transition. Atoms are then transferred to the $|3\rangle$ ($7S_{1/2}$) state through spontaneous radiation. Utilizing weak cavity feedback, population inversion is built up between the $|3\rangle$ and $|4\rangle$ ($6P_{3/2}$) levels, and finally, 1470 nm lasing is realized. An anti-resonant inhibited laser is realized when the cavity-mode frequency is tuned to the center of two adjacent cavity modes centered around the frequency of the atomic transition. **b**, Sketch of the inhibited laser. Cs atoms (blue spheres) along the direction of the cavity mode are pumped by the 459 nm laser. Two cavity mirrors are coated with reflectivities of $R_1 = R_2 = 34.5\%$ at 1470 nm. To change the cavity-mode frequency, the location of the output mirror is tunable. Generally, the round-trip optical path is exactly equal to an integer multiple of the intracavity wavelengths, and the output mirror is located at M_2 . However, an inhibited laser (in red) is achieved when the optical path is equal to an odd multiple of the half wavelength $2L = (2q + 1)\lambda/2$ with output mirror at M_2 .

inhibited by a factor of $\frac{1+R}{1-R}$ compared with that in free space when the cavity is resonant ($\Delta\phi = q \cdot 2\pi$) and anti-resonant ($\Delta\phi = (2q + 1) \cdot \pi$), respectively. Accordingly, the suppression of the spontaneous emission rate under the anti-resonant state is weak in the low-reflectivity cavity, which is conducive to lasing.

The detuning, $\Delta = \omega - \omega_0$, of the radiation frequency ω from the atomic transition frequency ω_0 , can also be given by $\Delta = P(\omega_c - \omega_0)$, where $P \equiv d\omega/d\omega_c$ represents the cavity-pulling coefficient [4, 26]. In the bad-cavity limit, $P \approx \Gamma/\kappa \ll 1$ when the cavity is near resonant, and thus, $\Delta^2 \ll 4g^2(n+1)$. n is the intracavity photon number, and $g = \frac{\mu}{\hbar} \sqrt{\frac{\hbar\omega_0}{2\varepsilon_0 V_c}} = 1.99 \times 10^5 \text{ s}^{-1}$ is the atom-cavity coupling constant, where μ is the electric

dipole moment, ε_0 is the vacuum permittivity, and V_c is the mode volume. Consequently, the detuning Δ in the laser rate equation [27] is negligible (the exact calculations are given in the Methods section). Here, we modify the loss term in the classical laser rate equation to obtain a universal expression, which can be used to describe any cavity-frequency detuning condition, as follows:

$$\begin{aligned} \frac{dn}{dt} &= N_{\text{eff}} \frac{\rho_{33} - \rho_{44}}{\tau_{\text{cyc}}} \sin^2(\sqrt{n+1}gt_{\text{int}}) - \frac{n}{\tau} \\ &= N_{\text{eff}} \frac{\rho_{33} - \rho_{44}}{\tau_{\text{cyc}}} \sin^2(\sqrt{n+1}gt_{\text{int}}) - \eta\kappa n. \end{aligned} \quad (2)$$

The first term on the right side represents the gain, and the second term is the loss. For the gain term, the effective number of atoms that can be pumped to the $7P_{1/2}$ state is $N_{\text{eff}} = 5.71 \times 10^9$ with a pumping light intensity $I = 10 \text{ mW/mm}^2$ and vapor-cell temperature $T = 100^\circ\text{C}$. ρ_{ii} denotes the population probability at level $|i\rangle$ in Fig. 1a. τ_{cyc} is the cycle time for Cs atoms through a transition of $6S_{1/2} \rightarrow 7P_{1/2} \rightarrow 7S_{1/2} \rightarrow 6P_{3/2}$, and t_{int} is the interaction time between the atoms and the cavity mode.

The loss term in Eq. (2) is inversely proportional to the intracavity photon lifetime τ . Typically, for the laser output from a resonant cavity, $\tau = 1/\kappa$. However, if the cavity and the atomic-transition frequencies are not identical, $\tau < 1/\kappa$. τ is expressed exactly as $\tau = \frac{1}{\eta\kappa}$. The loss coefficient η , reflecting the destructive interference of the intracavity radiated fields, is defined as the ratio of the maximum power emitted into the cavity at the resonant condition P_{cmax} to the power at any cavity-frequency detuning P_c ,

$$\eta = \frac{P_{\text{cmax}}}{P_c} = \frac{1 + R^2 - 2R \cos(\omega 2L/c)}{(1 - R)^2}. \quad (3)$$

As for the resonant cavity, the loss coefficient exhibits a minimum of $\eta_{\text{min}} = 1$, and Eq.(2) is reduced to the traditional expression of the laser rate equation [27]. Instead, η reaches the maximum $\eta_{\text{max}} = \left(\frac{1+R}{1-R}\right)^2$ for the inhibited laser from an anti-resonant cavity. The black dots in Fig. 2a show the variation of τ with $\Delta\phi$.

Using Eqs. (2) and (3), we obtain the steady-state solution of the intracavity photon number n as a function of $\Delta\phi$. Utilizing $P_{\text{out}} = nh\nu\kappa$, the output laser power P_{out} with the change of $\Delta\phi$ is represented by a dark-blue line in Fig. 2a, and the experimental result is indicated by the light-blue line. This reflects that if the cavity is tuned to off-resonance, the intracavity radiated fields interfere destructively, resulting in a decreased laser power. We measured P_{out} as a function of $\Delta\phi$ at different pumping light intensities I and different vapor-cell temperatures T , as shown in Figs 2b and 2c, respectively. The power of the inhibited laser can be further improved with higher pumping light intensities and temperatures.

Theoretically, analogous to the inhibited spontaneous

FIG. 2. **Intracavity photon lifetime and primary laser power behavior.** **a**, Intracavity photon lifetime τ (black dots) is inversely proportional to the loss coefficient η described by Eq. (3). The experimental (light-blue line) and simulated (dark-blue line) results of the output laser power P_{out} match well under the condition of $g = 1.99 \times 10^5 \text{ s}^{-1}$, $\Omega = 4.30 \times 10^7 \text{ s}^{-1}$, and $N_{\text{eff}} = 5.71 \times 10^9$. **b**, Output laser power P_{out} as a function of the phase shift $\Delta\phi$ at different pumping light intensities I under vapor-cell temperature $T = 100^\circ\text{C}$. The blue and orange lines represent P_{out} vs. the pumping light intensity when the cavity is on-resonance and anti-resonance, respectively. The blue circles and the orange stars are the corresponding experimental results. The pumping light intensities reaching the laser threshold were around 0.45 and 5.20 mW/mm^2 when the cavity was resonant and anti-resonant, respectively. **c**, P_{out} as a function of $\Delta\phi$ at different vapor-cell temperatures T under $I = 10 \text{ mW}/\text{mm}^2$. Theoretical and experimental P_{out} vs. T when the cavity is resonant (green) and anti-resonant (purple) are separately depicted by the lines and dots, respectively. The corresponding temperatures for the laser threshold were 72.5 and 94.5 $^\circ\text{C}$, respectively.

TABLE I. Cavity-pulling coefficient in optical domain

	Stimulated emission		
	Theoretical	Experimental	Fitted
Resonant	$\frac{\Gamma}{\kappa} [4, 5]$ 0.039	0.038 ± 0.002	$T = 120^\circ\text{C}$ 0.040
Anti-resonant	$-\left(\frac{\pi}{2\mathcal{F}}\right)^2 \frac{\Gamma}{\kappa}$ -0.012	-0.019 ± 0.002	$T = 120^\circ\text{C}$ -0.0019

* In this work, $\mathcal{F} = 3.07$; $\Gamma = 2\pi \times 10 \text{ MHz}$; $\kappa = 2\pi \times 257 \text{ MHz}$.

emission, the linewidth of the inhibited laser has the potential to be narrower than that of the traditional resonant laser. Here, considering the cavity-modification effect, we obtain the general expression of the laser linewidth as follows:

$$\Delta\nu_L = \frac{\kappa}{4\pi n} N_{\text{sp}} \left(\frac{\Gamma}{\Gamma + \kappa} \right)^2 \left(1 + \frac{\Gamma_e}{\Gamma_e + \Gamma_g} \frac{n}{n_s} \right) \frac{1}{1 + \left(\frac{2\mathcal{F}}{\pi} \right)^2 \sin^2 \left(\frac{\omega L}{c} \right)}. \quad (4)$$

where $N_{\text{sp}} = \frac{N_e}{N_e - N_g}$ is the spontaneous-emission factor, N_e and N_g represent the populations of the excited and ground states, respectively; Γ_e , Γ_g , and Γ_{eg} are the decay rates of the atomic populations and polarization; and $n_s = \frac{\Gamma_{\text{eg}}}{2g^2} \frac{\Gamma_e \Gamma_g}{\Gamma_e + \Gamma_g}$ is the homogeneous saturation intensity in units of number of photons. Equation (4) shows four

extra features compared with the classical Schawlow-Townes equation [28]: (i) broadening of the linewidth due to the incomplete inversion, (ii) a bad-cavity effect [26] leading to linewidth narrowing, (iii) power broadening [29], and (iv) cavity-induced modification [12] following the absorption lineshape with the change of $\Delta\phi$. From the last term, the laser linewidth is expected to be narrowed by a factor of $\frac{1}{1 + (2\mathcal{F}/\pi)^2}$ for the inhibited laser compared with the resonant one, and simplify to the classical bad-cavity expression [29] under the resonant condition.

The laser linewidth was analyzed by measuring the beat-note spectrum between the tested laser and the reference laser, as shown in Fig. 3a. Limited to the intensity sensitivity of the photodetector, the cavity frequency of the reference laser should be coincident with the atomic resonance to improve the light intensity for beating. Although it is difficult to measure the beat-note spectrum between inhibited lasers due to the poor laser power, the

measured linewidth of the inhibited laser was comparable to that of the resonant laser.

FIG. 3. Linewidth and cavity-pulling characteristics. **a**, Beat-note spectrum between the reference laser and the tested laser. The cavity frequency of the reference laser coincides with the atomic transition frequency, while that of the tested laser is tunable by the cavity length. The beating spectrum (black circles) between the reference laser and the resonant tested laser was fitted by a Lorentzian function with a fitted linewidth of 1.2 kHz (red line). Moreover, the Lorentz fitting linewidth of the beat-note spectrum (black squares) between the reference laser and the inhibited laser was also 1.2 kHz (grey line). **b**, Frequency shift of the laser oscillation Δ as a function of cavity-frequency detuning from the atomic transition frequency $\omega_c - \omega_0$, whose adjustable range was around one FSR. We modify Eq. (5) with the laser rate equation to describe Δ . The fitting results of Δ at different temperatures for $R = 28\%$ are shown as the solid lines, and the black triangles represent the experimental results with $R = 34.5\%$ under $T = 100^\circ\text{C}$. The difference between the simulated data and experimental results is explained further in the main text.

Most importantly, the inhibited laser has the advantage of an enhanced suppression of the cavity-pulling effect. The relationship between the frequency shift of the oscillation frequency, i.e., Δ , and the cavity-frequency detuning from the atomic transition $\omega_c - \omega_0$ is analyzed comprehensively for spontaneous emission [30], which is

written as

$$\Delta = \frac{\Gamma}{4} \frac{\frac{4R}{(1-R)^2} \sin(2\omega L/c)}{1 + \frac{4R}{(1-R)^2} \sin^2(\omega L/c)}. \quad (5)$$

Therefore, for spontaneous radiation, the frequency shift caused by the cavity-frequency detuning is eliminated, not only when the atomic resonance coincides with one of the cavity resonances but also when the atomic resonance is halfway between two adjacent cavity resonances. Replacing $\frac{2\sqrt{R}}{1-R}$ by $\frac{2\mathcal{F}}{\pi}$, the cavity-pulling coefficients are equal to $\frac{2\mathcal{F}}{\pi} \frac{\Gamma}{\kappa}$ and $-\frac{2\mathcal{F}}{\pi} \frac{\Gamma}{\kappa} \frac{1}{1 + (2\mathcal{F}/\pi)^2}$ utilizing $\mathcal{F} = \frac{c}{2L} \frac{2\pi}{\kappa}$, when the cavity is resonant and anti-resonant, respectively. The ratio between the two coefficients is approximately equal to $-(\frac{2\mathcal{F}}{\pi})^2$.

Analogous to the spontaneous radiation, the ratio between the cavity-pulling coefficients when the cavity is resonant and anti-resonant is $-(\frac{2\mathcal{F}}{\pi})^2$ for the stimulated emission. The difference is that the pulling coefficient is $\frac{\Gamma}{\kappa}$ for the resonant bad-cavity laser [4, 5]. Accordingly, the cavity-pulling coefficient is around $-(\frac{\pi}{2\mathcal{F}})^2 \frac{\Gamma}{\kappa}$ when the cavity is anti-resonant. More specifically, for the stimulated emission, we should consider the atom-cavity interactions. Therefore, Eq. (5) is further modified by the laser rate equation to obtain the frequency shift of the stimulated emission. The fitted results are depicted by solid lines in Fig. 3b. In addition, we measured the frequency shift as a function of the cavity-frequency detuning from the atomic transition frequency. The experimental results (black triangles) agreed with the dispersive lineshape and were consistent with the fitted results. For comparison, the cavity-pulling coefficients discussed above are illustrated in Table I. The measured pulling coefficient of the inhibited laser was 1.71–2.35 times smaller than that of the resonant condition. Such an inhibited laser is characterized by an enhanced suppression of the cavity-pulling effect compared with the traditional AOCs.

The deviation between the experimental and theoretical cavity-pulling coefficients is analyzed. First, the general calculation for the cavity finesse based on the high-reflectivity approximation may not be preferably applicable to the low-reflectivity case, resulting in the deviation of the finesse calculation. Second, the non-linear hysteresis phenomenon of the piezoelectric ceramic used to adjust the cavity length could result in measurement errors of the cavity-frequency detuning. Third, the removal of the transverse mode degeneracy caused by the geometry change of cavity would disturb the single-mode requirement.

Differing from all existing types of lasers, the inhibited laser demonstrated here is inherently insensitive to cavity-length fluctuations, which leads to significantly smaller systematic perturbations compared to traditional AOCs [4–9]. In the future, with the cold atomic ensemble and the improved cavity finesse, we will expect a cavity-

pulling coefficient of the order 10^{-5} comparable with the results in reference [2], but solve its problem of pulsed operation. This novel superradiant AOC using the principle of inhibited laser might with better frequency sta-

bility, having wide range applications in precision measurements, such as test of variation of fundamental constants, gravitational potential of Earth, and search for dark matter.

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- [1] D. G. Matei, T. Legero, S. Häfner, C. Grebing, R. Weyrich, W. Zhang, L. Sonderhouse, J. M. Robinson, J. Ye, F. Riehle, U. Sterr, 1.5 μm lasers with sub-10 mHz linewidth, *Phys. Rev. Lett.* 118 (2017) 263202.
- [2] M. A. Norcia, J. R. K. Cline, J. A. Muniz, J. M. Robinson, R. B. Hutson, A. Goban, G. E. Marti, J. Ye, J. K. Thompson, Frequency measurements of superradiance from the strontium clock transition, *Phys. Rev. X* 8 (2018) 021036.
- [3] D. Kleppner, Inhibited spontaneous emission, *Phys. Rev. Lett.* 47 (1981) 233–236.
- [4] J. Chen, Active optical clock, *Chinese Sci. Bull.* 54 (3) (2009) 348–352.
- [5] J. G. Bohnet, Z. Chen, J. M. Weiner, D. Meiser, M. J. Holland, J. K. Thompson, A steady-state superradiant laser with less than one intracavity photon, *Nature* 484 (2012) 78–81.
- [6] G. A. Kazakov, T. Schumm, Active optical frequency standard using sequential coupling of atomic ensembles, *Phys. Rev. A* 87 (2013) 013821.
- [7] T. Laske, H. Winter, A. Hemmerich, Pulse delay time statistics in a superradiant laser with calcium atoms, *Phys. Rev. Lett.* 123 (2019) 103601.
- [8] S. A. Schäffer, M. Tang, M. R. Henriksen, A. A. Jørgensen, B. T. R. Christensen, J. W. Thomsen, Lasing on a narrow transition in a cold thermal strontium ensemble, *Phys. Rev. A* 101 (2020) 013819.
- [9] C. Chen, S. Bennetts, R. G. Escudero, B. Pasquiou, F. Schreck, Continuous guided strontium beam with high phase-space density, *Phys. Rev. Applied* 12 (2019) 044014.
- [10] E. M. Purcell, Spontaneous emission probabilities at radio frequencies, *Phys. Rev.* 69 (1946) 681.
- [11] P. Goy, J. M. Raimond, M. Gross, S. Haroche, Observation of cavity-enhanced single-atom spontaneous emission, *Phys. Rev. Lett.* 50 (1983) 1903–1906.
- [12] D. J. Heinzen, J. Childs, C. Monroe, J. E. Thomas, M. S. Feld, Enhanced and inhibited visible spontaneous emission by atoms in a confocal resonator, *Phys. Rev. Lett.* 58 (1987) 1320–1323.
- [13] F. D. Martini, G. Innocenti, G. R. Jacobovitz, P. Mataloni, Anomalous spontaneous emission time in a microscopic optical cavity, *Phys. Rev. Lett.* 59 (1987) 2955–2958.
- [14] H. Mabuchi, A. C. Doherty, Cavity quantum electrodynamics: Coherence in context, *Science* 298 (2002) 1372–1377.
- [15] A. Wallraff, D. I. Schuster, A. Blais, L. Frunzio, R.-S. Huang, J. Majer, S. Kumar, S. M. Girvin, R. J. Schoelkopf, Strong coupling of a single photon to a superconducting qubit using circuit quantum electrodynamics, *Nature* 431 (2004) 162–167.
- [16] T. Yoshie, A. Scherer, J. Hendrickson, G. Khitrova, H. Gibbs, G. Rupper, C. Ell, O. Shchekin, D. Deppe, Vacuum rabi splitting with a single quantum dot in a photonic crystal nanocavity, *Nature* 432 (2004) 200–203.
- [17] J. McKeever, A. Boca, A. D. Boozer, J. R. Buck, H. J. Kimble, Experimental realization of a one-atom laser in the regime of strong coupling, *Nature* 425 (2003) 268–271.
- [18] M. Keller, B. Lange, K. Hayasaka, W. Lange, H. Walther, Continuous generation of single photons with controlled waveform in an ion-trap cavity system, *Nature* 431 (2004) 1075–1078.
- [19] C. Monroe, D. M. Meekhof, B. King, W. M. Itano, D. Wineland, Demonstration of a fundamental quantum logic gate, *Phys. Rev. Lett.* 75 (1995) 4714–4717.
- [20] A. G. Vaidyanathan, W. P. Spencer, D. Kleppner, Inhibited absorption of blackbody radiation, *Phys. Rev. Lett.* 47 (1981) 1592–1595.
- [21] R. G. Hulet, E. S. Hilfer, D. Kleppner, Inhibited spontaneous emission by a rydberg atom, *Phys. Rev. Lett.* 55 (1985) 2137–2140.
- [22] D. J. Heinzen, M. S. Feld, Vacuum radiative level shift and spontaneous-emission linewidth of an atom in an optical resonator, *Phys. Rev. Lett.* 59 (1987) 2623–2626.
- [23] R. Loudon, The quantum theory of light, *Am. J. Phys.* 42 (2000).
- [24] M. Fox, J. Javanainen, Quantum optics: An introduction, *Phys. Today* 60 (9) (2007) 74–75.
- [25] T. Shi, D. Pan, J. Chen, Realization of phase locking in good-bad-cavity active optical clock, *Opt. Express* 27 (2019) 22040–22052.
- [26] S. J. M. Kuppens, M. P. Van Exter, J. P. Woerdman, Quantum-limited linewidth of a bad-cavity laser, *Phys. Rev. Lett.* 72 (24) (1994) 3815–3818.
- [27] C. M. Fang-Yen, Multiple thresholds and many-atom dynamics in the cavity qed microlaser, Ph.D. thesis, Massachusetts Institute of Technology (2002).
- [28] A. L. Schawlow, C. H. Townes, Infrared and optical masers, *Phys. Rev.* 112 (1958) 1940–1949.
- [29] A. Z. Khoury, M. I. Kolobov, L. Davidovich, Quantum-limited linewidth of a bad-cavity laser with inhomogeneous broadening, *Phys. Rev. A* 53 (1996) 1120–1125.
- [30] D. J. Heinzen, Radiative decay and level shift of an atom in an optical resonator, Ph.D. thesis, Massachusetts Institute of Technology (1988).

METHODS

Experimental details. To acquire sufficient gain, we take advantage of the multilevel structure of the Cs atom and multiple atoms interacting with a single mode of an optical cavity. As depicted in Fig. 1, a cloud of thermal Cs atoms collected in the low-finesse F-P cavity was pumped by the 459 nm continuous-wave laser. For typical lasers, the cavity length is exactly equal to an integral multiple of the half-wavelength, i.e., the cavity-

mode frequency resonates with the peak of the emission line for an atomic transition. However, the cavity length is equal to an odd multiple of the quarter wavelength, i.e., the atomic transition frequency is halfway between two adjacent cavity modes, for the inhibited laser. In this work, the cavity length was tunable through the PZT, of which the adjustable range was more than one-half wavelength of the laser oscillation.

Cavity-pulling coefficient. The integrated Invar F-P cavity consisted of a plane mirror M_1 and a plane-concave mirror M_2 (radius of curvature $r = 500$ mm) separated by a distance $L = 190$ mm. Therefore, the mode sustained by the cavity had Gaussian transverse profiles, of which the spot radii on the cavity mirrors M_1 and M_2 were $w_{s1} = 0.429$ mm and $w_{s2} = 0.337$ mm, respectively. The equivalent mode volume was $V_c = \frac{1}{4}L\pi\left(\frac{w_{s1}+w_{s2}}{2}\right)^2 = 21.89$ mm³. The cavity power decay rate was $\kappa = 2\pi \times 257$ MHz, and the free spectral range was FSR = 789 MHz. Therefore, the cavity finesse was $\mathcal{F} = 3.07$ [31].

The gain medium Cs atoms were pumped by the 459 nm laser through the velocity-selective mechanism. It was assumed that the pumping light intensity $I = 10$ mW/mm², while the corresponding saturation light intensity $I_s = \pi\hbar c\Gamma/3\lambda^3 = 1.27$ mW/cm². Therefore, the saturation broadening of state $|2\rangle$ in Fig. 1 caused by the pumping laser was

$$\frac{\Gamma_2}{2\pi} = \frac{\Gamma_{21} + \Gamma_{23} + \Gamma_{24}}{2\pi} \sqrt{1+s} = 26.38 \text{ MHz}, \quad (6)$$

where $\Gamma_{21} = 0.793 \times 10^6 \text{ s}^{-1}$, $\Gamma_{23} = 3.52 \times 10^6 \text{ s}^{-1}$, and $\Gamma_{24} = 1.59 \times 10^6 \text{ s}^{-1}$ are the decay rates of the $|2\rangle \rightarrow |1\rangle$, $|2\rangle \rightarrow |3\rangle$, and $|2\rangle \rightarrow |4\rangle$ transitions, respectively. s is the saturation factor represented by $s = I/I_s$. According to the velocity-selective scheme, only atoms in the direction of the cavity mode with a velocity less than $\Delta v = \frac{\Gamma_2}{2\pi} \times \lambda_{21}$ can be pumped to state $|2\rangle$ and then decay to state $|3\rangle$. Consequently, the Doppler broadening of $|3\rangle$ is $\frac{\Gamma_D}{2\pi} = \Delta v/\lambda_{34}$. Since the spontaneous decay rate of the 1470 nm transition $\Gamma_0 = 2\pi \times 1.81$ MHz [32], the atomic decay rate was $\Gamma = \Gamma_0 + \Gamma_D = 2\pi \times 10.04$ MHz, which was much smaller than κ . Accordingly, the cavity-pulling coefficient in the resonant cavity is $P \approx \Gamma/\kappa = 0.039$.

Intracavity photon number at steady state. For the atomic number density $n' = 1.57 \times 10^{13}/\text{cm}^3$ at a vapor-cell temperature of 100°C [33], the atomic number inside the cavity mode is $N = \frac{1}{4}n'\pi L_{\text{cell}}\left(\frac{w_{s1}+w_{s2}}{2}\right)^2 = 1.81 \times 10^{11}$, where the vapor-cell length is $L_{\text{cell}} = 10$ cm. Only the atoms with velocities between $-\Delta v/2 \sim \Delta v/2$ can be pumped to the $7P_{1/2}$ state. According to the Maxwell speed distribution, the effective atomic number N_{eff} is given by

$$N_{\text{eff}} = N \int_{-\frac{1}{2}\Delta v}^{\frac{1}{2}\Delta v} \sqrt{\frac{m}{2\pi k_B T}} \exp\left(\frac{-mv^2}{2k_B T}\right) dv \\ = 5.71 \times 10^9, \quad (7)$$

where m is the atomic mass, and k_B is the Boltzmann constant.

Utilizing the density matrix equations, the intracavity photon number at steady state as a function of the phase shift $\Delta\phi$ (or cavity-frequency detuning, $\nu_c - \nu_0$) is obtained. The atomic energy level is shown in Fig. 1a, where the energy states are labelled as $|i\rangle$. Using the rotating wave approximation (RWA) approximation, the density matrix equations for Cs atoms interacting with the 459 nm pumping laser are expressed as follows:

$$\begin{aligned} \frac{d\rho_{11}}{dt} &= -\Omega\rho_{12}^I + \Gamma_{21}\rho_{22} + \Gamma_{41}\rho_{44} + \Gamma_{61}\rho_{66}, \\ \frac{d\rho_{22}}{dt} &= \Omega\rho_{12}^I - (\Gamma_{21} + \Gamma_{23} + \Gamma_{25})\rho_{22}, \\ \frac{d\rho_{33}}{dt} &= \Gamma_{23}\rho_{22} - (\Gamma_{34} + \Gamma_{36})\rho_{33} - \frac{\rho_{33} - \rho_{44}}{\tau_{\text{cyc}}}, \\ &\quad \left(\frac{2g\sqrt{n+1}}{\sqrt{\Delta^2 + 4g^2(n+1)}}\right)^2 \sin^2\left(\frac{\sqrt{\Delta^2 + 4g^2(n+1)}}{2}\tau_{\text{cyc}}\right), \\ \frac{d\rho_{44}}{dt} &= \Gamma_{34}\rho_{33} + \Gamma_{54}\rho_{55} - \Gamma_{41}\rho_{44} + \frac{\rho_{33} - \rho_{44}}{\tau_{\text{cyc}}}, \\ &\quad \left(\frac{2g\sqrt{n+1}}{\sqrt{\Delta^2 + 4g^2(n+1)}}\right)^2 \sin^2\left(\frac{\sqrt{\Delta^2 + 4g^2(n+1)}}{2}\tau_{\text{cyc}}\right), \\ \frac{d\rho_{55}}{dt} &= \Gamma_{25}\rho_{22} - (\Gamma_{54} + \Gamma_{56})\rho_{55}, \\ \frac{d\rho_{66}}{dt} &= \Gamma_{36}\rho_{33} + \Gamma_{56}\rho_{55} - \Gamma_{61}\rho_{66}, \\ \frac{d\rho_{12}^I}{dt} &= \frac{1}{2}\Omega(\rho_{11} - \rho_{12}) + \rho_{12}^R\Delta' - \frac{1}{2}\Gamma_{21}\rho_{12}^I, \\ \frac{d\rho_{12}^R}{dt} &= -\rho_{12}^I\Delta' - \frac{1}{2}\Gamma_{21}\rho_{12}^R, \\ \frac{dn}{dt} &= N_{\text{eff}}\frac{\rho_{33} - \rho_{44}}{\tau_{\text{cyc}}}\left(\frac{2g\sqrt{n+1}}{\sqrt{\Delta^2 + 4g^2(n+1)}}\right)^2 \\ &\quad \sin^2\left(\frac{\sqrt{\Delta^2 + 4g^2(n+1)}}{2}t_{\text{int}}\right) - \eta\kappa n. \end{aligned} \quad (8)$$

Ω is the Rabi frequency, and Γ_{ij} represents the rate of decay from $|i\rangle$ to $|j\rangle$. $\tau_{\text{cyc}} = \frac{1}{\Omega} + \frac{1}{\Gamma_{23} + \Gamma_{25}} + \frac{1}{\Gamma_{34} + \Gamma_{36}} + \frac{1}{\Gamma_{41}}$ is the cycle time for Cs atoms through a complete transition of $6S_{1/2} \rightarrow 7P_{1/2} \rightarrow 7S_{1/2} \rightarrow 6P_{3/2}$. The interaction time between the atoms and the cavity mode is given by $t_{\text{int}} \approx \frac{1}{\Gamma_{34} + \Gamma_{36} + \Gamma_{41}}$. Ideally, we would simplify the equations by setting the frequency detuning between the pumping laser and the atomic transition of $|1\rangle$ to $|2\rangle$ to be zero. ρ_{12}^I and ρ_{12}^R represent the energy shift and the power broadening, respectively. ρ_{ii} denotes the population probability of atoms in the corresponding state, and the result is shown in Fig. 4. $\Delta = \omega - \omega_0$ is the frequency detuning of the laser oscillation from the atomic transition. Since $\Delta^2 \ll 4g^2(n+1)$, we assume that $\Delta = 0$ in the main text. To verify the correctness of this assumption, we give the most accurate description of the detuning Δ in Eq. (8). The photon number at steady

state is calculated by inserting the fitted result of Δ in Fig. 3b into Eq. (8).

The results of the intracavity photon number at steady state with and without considering the detuning Δ are shown as the green dotted line and the red solid line in Fig. 5. This shows that there was little difference between the photon number obtained by Eq. (2) and the last equation of Eq. (8). The difference between photon numbers obtained by the two equations is shown in the inset of Fig. 5, which illustrates that the differences are both zero when the cavity is resonant and anti-resonant. The difference is eliminated when the mode frequency exactly coincides with the center frequency of the gain profile. In addition, when the mode frequency is tuned to the center of two adjacent cavity resonances, the effects of cavity-pulling of the two adjacent cavity modes on the laser frequency are equal and opposite. Hence, the difference is also zero for the inhibited laser. This result demonstrates that the approximation used in the main text is reasonable.

To further characterize the output laser power as a function of the phase shift at different pumping efficiencies and atomic densities, the pumping light intensity and the vapor-cell temperature are adjustable, as depicted in Fig. 2a. According to Eq. (7), the intracavity effective atomic number is influenced by both the pumping light intensity and the vapor-cell temperature, while the Rabi frequency of the pumping laser is relative to the pumping light intensity, which is described as $\Omega = \sqrt{\frac{3\lambda_{21}^3\Gamma_{21}I}{2\pi\hbar c}}$. N_{eff} and Ω as functions of I and N_{eff} vs. T are shown in Fig. 6a and b, respectively.

Photon number as function of cavity decay rate. According to Eq. (2), the function of the intracavity photon number n with phase shift $\Delta\phi$ varies with the cavity-mirror reflectivity R , namely, the cavity power loss rate κ . When the cavity is anti-resonant ($\Delta\phi = (2q+1)\pi$), the intracavity photon number decreased with the increase in the reflectivity, which is shown in Fig. 7. The intracavity photon number for the inhibited laser was smaller than 1 when the reflectivity increased to 80% with $g = 1.99 \times 10^5 \text{ s}^{-1}$, $\Omega = 4.30 \times 10^7 \text{ s}^{-1}$, and $N_{\text{eff}} = 5.71 \times 10^9$. Nevertheless, n could be further improved with a higher pumping light intensity and a higher atomic number density.

END NOTES

Data availability

The data represented in Figs. 1–7 are available as Source Data. All other data that support the plots within this paper and other findings of this study are available from the corresponding author upon reasonable request.

References

31. Riehle, F. Frequency Standards: Basics and Applications. Hoboken, NJ, USA: Wiley. 2006.

32. O. S. Heavens, Radiative Transition Probabilities of the Lower Excited States of the Alkali Metals, J. Opt. Soc. Am. **51**, 1058-1061 (1961).

33. Steck, Daniel A. Cesium D line data. available online at <http://steck.us/alkalidata>.

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Author contributions

J.C. conceived the idea to use an anti-resonant cavity to realize the inhibited laser as a stable active optical clock. T.S. performed the experiments and carried out the theoretical calculations. T. S. wrote the manuscript. D. P. and J. C. provided revisions. Both authors contributed equally to the discussions of the results.

Competing interests

The authors declare no competing interests.

Additional information

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FIG. 4. Numerical results of the population probability of each state ρ_{ii} in Fig. 1a, under $g = 1.99 \times 10^5 \text{ s}^{-1}$, $\Omega = 4.30 \times 10^7 \text{ s}^{-1}$, and $N_{\text{eff}} = 5.71 \times 10^9$.

FIG. 5. Photon number at steady state as a function of the phase shift for $g = 1.99 \times 10^5 \text{ s}^{-1}$, $\Omega = 4.30 \times 10^7 \text{ s}^{-1}$, and $N_{\text{eff}} = 5.71 \times 10^9$. The red line and the green dotted line represent the photon number when $\Delta = 0$ and $\Delta \neq 0$, respectively. The inset shows the difference between the photon numbers when $\Delta = 0$ and $\Delta \neq 0$.

FIG. 6. **a**, Effective atomic number and Rabi frequency vs. the pumping light intensity at a vapor-cell temperature of 100°C . **b**, Effective atomic number as a function of the vapor-cell temperature under a pumping light intensity of 10 mW/mm^2 .

FIG. 7. Intracavity photon number n as a function of the phase shift $\Delta\phi$ at different cavity-mirror reflectivity R values for $g = 1.99 \times 10^5 \text{ s}^{-1}$, $\Omega = 4.30 \times 10^7 \text{ s}^{-1}$, and $N_{\text{eff}} = 5.71 \times 10^9$. The blue and red lines represent the photon number when the cavity is resonant and anti-resonant, respectively. The orange and the purple lines show the results at $R = 34.5\%$ and $R = 80\%$ ($n \leq 1$), respectively.