Deterministic creation of entangled atom-light Schrödinger-cat states

Bastian Hacker,* Stephan Welte, Severin Daiss, Armin Shaukat, Stephan Ritter,† Lin Li,‡ and Gerhard Rempe
Max-Planck-Institut für Quantenoptik, Hans-Kopfermann-Straße 1, 85748 Garching, Germany
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Coherent superpositions of classical objects entangled with quantum objects – dubbed Schrödinger cats – illustrate the weirdness of quantum physics. Several laboratory experiments with ions or microwaves explored such cats in a setting where the classical objects were confined to boxes that protect the superposition state. New applications in a distributed quantum network, however, require unconfined cats that propagate through optical channels between network nodes. Approximate cat states of light have been produced, but only probabilistically. Here we deterministically create atom-light cat states by reflecting laser pulses from an optical cavity containing a single atom in a superposition of two internal states. In contrast to single-photon experiments, each trial establishes entanglement. Suitable measurements on the atom produce a plethora of optical cat states with nonclassical characteristics. We furthermore implement a deterministic quantum gate between the atom and an optical cat, thus demonstrating its potential for continuous-variable quantum information processing with loss correction and for fault-tolerant quantum computing.

As early as 1935, Schrödinger formulated a Gedankenexperiment with a live cat and a radioactive atom, both inside a box. The decaying atom triggers a death mechanism that kills the cat. However, without detailed knowledge of what happens inside the box, the atom and the cat are in an entangled superposition state, with the weird consequence that the cat might simultaneously be dead and alive. Our experiment resembles this scenario as it produces entanglement between a (mesoscopic) pulse of light and a single atom inside a cavity. However, our experiment differs in the fact that the light propagates outside the cavity in a well-defined spatial and temporal mode, unconfined by any box. The light exists in an odd or even superposition of “alive”, “dead”, and entangled with the state of the atom, respectively. Performing a state detection even when atom and light are far apart. In fact, the entanglement with an atom is a unique resource not shared by previous experiments generating approximate optical cat states.

These schemes relied on the extraction of a photon from a squeezed vacuum beam and are therefore inherently probabilistic. Our experiment follows a different route and implements a deterministic protocol proposed some decade ago.

We employ a single $^{87}$Rb atom trapped at the centre of a high-finesse (6 × 10$^5$) optical cavity (see Methods). The atom acts as a three-level system consisting of the two ground states, $|\downarrow\rangle$ and $|\uparrow\rangle$, separated by a microwave transition, and one excited state $|e\rangle$. The atom in state $|\uparrow\rangle$ is strongly coupled to the fundamental mode of the cavity with parameters $(g, \kappa, \gamma) = 2\pi \times (7.8, 2.5, 3.0)$ MHz. Here, $g$ denotes the atom-photon coupling constant, $\kappa$ is the cavity field decay rate, and $2\gamma$ is the spontaneous atomic decay rate on the transition $|\uparrow\rangle \leftrightarrow |e\rangle$. The cavity is actively stabilised on this transition with a wavelength of 780 nm. It is furthermore single-sided in the sense that the optical losses $\kappa$ are dominated by the transmission of the mirror through which light is coupled into and out of the cavity with rate $\kappa_c = 2\pi \times 2.3$ MHz.

The protocol to create cat states is outlined in Fig. 1a. Starting with an atom in the coupling state $|\uparrow\rangle$, we apply a $\pi/2$ pulse to generate an equal superposition state $(|\uparrow\rangle + |\downarrow\rangle)/\sqrt{2}$. Next, we reflect a coherent pulse $|\alpha\rangle = e^{-|\alpha|^2/2} \sum_n \alpha^n/\sqrt{n!} |n\rangle$ with an arbitrary amplitude $\alpha$ (and the photonic Fock states $|n\rangle$) from the cavity. This creates a phase shift that depends on the state of the atom:

$$
\begin{align*}
|\uparrow\rangle |\alpha\rangle &\rightarrow \left|\uparrow\right\rangle |\alpha\rangle, \\
|\downarrow\rangle |\alpha\rangle &\rightarrow \left|\downarrow\right\rangle -|\alpha\rangle.
\end{align*}
$$

Thus, after the reflection, the state is $(|\uparrow\rangle |\alpha\rangle + |\downarrow\rangle -|\alpha\rangle)/\sqrt{2}$, an entangled state of the atom and the light field. With the initial state of the light field, $|\alpha\rangle$, being classical and potentially macroscopic, the entangled atom-light state is our “Schrödinger cat”. A subsequent $\pi/2$ spin rotation on the atom produces the state

$$
\frac{1}{2} \left[ |\uparrow\rangle (|\alpha\rangle -|-\alpha\rangle) + |\downarrow\rangle (|\alpha\rangle +|-\alpha\rangle) \right],
$$

where the light is in an odd or even superposition of “alive”, “dead”, and entangled with the $|\uparrow\rangle$ and $|\downarrow\rangle$ eigenstates of the atom, respectively. Performing a state detection on the atom, the light part of the state (2) is projected onto either an odd cat state $(|\alpha\rangle -|-\alpha\rangle)/N^\pm$ if the atom is observed in $|\uparrow\rangle$ or an even cat state $(|\alpha\rangle +|-\alpha\rangle)/N^{-}$ if the atom is found in $|\downarrow\rangle$. Here even and odd refers to the photon number parity of the cat states in Fock space, i.e., a state containing only even and odd photon numbers. $N^\pm$ are the respective normalisation factors. More complex cat states with arbitrary superposition coefficients can be generated by suitable rotations of the atomic state before the measurement. The whole protocol described above lasts 45 $\mu$s.

The created cat states propagate in a well-defined optical mode and can be sent through an optical fibre that may be part of a larger quantum network. In our experiment, we send the light to a homodyne detector for characterisation (see Methods). Each measurement returns an amplitude of the optical field for a given phase value. The full optical state is defined by its distribution in phase space, the Wigner function $W(q,p)$ that we reconstruct from a larger set of measurements. Here, $q$ and $p$ are the field quadratures spanning phase space.
Fig. 1 | Cat state generation protocol and Wigner functions of experimentally measured cat states. a, Steps required to generate a cat state. Initially, the atom is prepared (1) in the state $|\uparrow\rangle$ before a $\pi/2$ rotation is applied (2) to bring it into the coherent superposition $(|\uparrow\rangle + |\downarrow\rangle)/\sqrt{2}$. Then, a coherent state $|\alpha\rangle$ is reflected from the cavity and thus an entangled state $(|\uparrow\rangle|\alpha\rangle - |\downarrow\rangle|\alpha\rangle)/\sqrt{2}$ is created (3). A subsequent $\pi/2$ rotation prepares $1/2[(|\alpha\rangle - |\alpha\rangle) + (|\alpha\rangle + |\alpha\rangle)]$ (4). The last step in the protocol is a state detection on the atom that projects the optical part onto the even or the odd cat state (5). b, Wigner functions of an even (left) and an odd (right) cat state with $\alpha = 1.4$, obtained when the atom is measured in $|\downarrow\rangle$ or $|\uparrow\rangle$, respectively. The fringes between the two Gaussian distributions are shifted by $2\pi$ between the two cases. The upper row is reconstructed without loss-correction, whereas the lower row was corrected for 25% propagation and detection losses, which affect the state after its creation at the cavity. c, Three-dimensional plot of the measured Wigner function of an odd cat state as it emerges from the cavity.

Our reconstruction basis is the density matrix $\rho$ in truncated Fock space. We observe the characteristic Wigner functions (Fig. 1b-c) with two Gaussian peaks and interference fringes in the centre that encode the coherent nature of the superposition state. The even cat state displays a local maximum and the odd cat state a local minimum at the centre of the Wigner distribution.

A characteristic feature of cat states, which has no classical explanation, is the existence of negative regions in the Wigner function. We observe a minimum value of $-0.016 \pm 0.004$ for the odd cat state with $\alpha = 1.4$, shown in Fig. 1b. If we correct for propagation losses that occur after the creation, the minimum value becomes $-0.083 \pm 0.005$ (Fig. 1b-c). Interesting for applications is that even cat states are “squeezed”, i.e., the uncertainty of one of the field quadratures, $\Delta p$, is smaller than for any classical state. Experimental results are displayed in Fig. 2a that shows the minimal width of the measured quadrature distributions as a function of $\alpha^2$. We observe squeezing with a maximal amount of 1.18(3) dB at $\alpha^2 = 0.5$ (for details see Methods). The measured values follow the theoretical prediction $\Delta p/\Delta p_{\text{vac}} = \sqrt{1 - 4(1 - L)\alpha^2}/(1 + \exp 2\alpha^2)$ with relative optical losses $L$ and the uncertainty of the vacuum state $\Delta p_{\text{vac}}$. With increasing $\alpha^2$ the width $\Delta p$ increases due to experimental noise in the optical phase $\phi$ (angle between the two Gaussian distributions in phase space), attributed to residual fluctuations of the cavity and laser frequencies.

It is now worthwhile to discuss the optical losses. They reduce the visibility $V$ of the nonclassical coherences which appear as fringes in the Wigner function. To quantify this effect, we express the visibility as the difference between even- and odd-cat Wigner functions at the centre of the cat state, $V = \frac{2}{3}(W_{\text{even}}(0,0) - W_{\text{odd}}(0,0))$. The theoretical calculation (see Methods) gives $V = \frac{\sinh(2(1 - L)\alpha^2)}{\sinh(2\alpha^2)}$ for an initially ideal cat state. Measured visibilities are presented in Fig. 2b.

In our experiment, losses stem from the finite cooperativity of the atom-cavity system, $C = g^2/(2\kappa\gamma) = 4.1$, and the reduced escape efficiency $\eta_{\text{esc}} = \kappa/\kappa = 0.92$ of the cavity that does not reflect all of the incoming light. For a coherent input of size $\alpha$ the effective size of the output cat state $\alpha_{\text{out}} = |\alpha_+ - \alpha_-|/2$, i.e. the phase-space distance between the two coherent contributions, is obtained from input-output theory (see Methods) as $\alpha_{\text{out}} = \eta \alpha$ with

$$\eta = \frac{\gamma_+ g^2}{g^2 + \kappa\gamma} = \frac{\eta_{\text{esc}} C}{C + 1/2} = 0.81.$$  \hspace{1cm} (3)
Fig. 2 | Nonclassical properties of measured cat states. a, Squeezing. Even cat states are squeezed with reduced $p$-quadrature fluctuations for low $\alpha^2$ (light red area), with a maximum of 1.18(3) dB. The dotted and dashed lines show the theoretical prediction for even cat states with losses $L_{\text{eff}} = 0.19$ and $L_{\text{det}} = 0.46$ for the states right after the cavity and at the homodyne setup, respectively. The full model (solid lines) for even (red), odd (yellow) and combined (gray) cat states is affected by atomic state detection errors of 1.3% at small $\alpha^2$ and random noise in the phase angle $\phi$ at larger $\alpha^2$. The latter limits the observable squeezing to states with $\alpha^2 < 1.3$. b, Interference fringe visibility $V$ of measured Wigner functions. The dotted and dashed lines show the effect of experimental losses $L_{\text{eff}}$ and $L_{\text{det}}$, respectively, which reduce the visibility with increasing $\alpha^2$. For $\alpha^2$ close to 0, errors in the atomic state detection dominate and cause the visibility to drop (solid line, fitted to data). All data points are without correction for optical losses. Error bars depict statistical standard errors.

This corresponds to a loss of $L_{\text{cav}} = 1 - \eta^2 = 34\%$.

Remarkably, not all these losses reduce the coherence of the created cat state like a classical absorber would do. This is caused by some losses that effectively occur before the cat state is created. Theoretically, we find that the effective coherence-reducing losses $L_{\text{eff}}$ are (see Methods)

$$L_{\text{eff}} = 1 - \eta = 19\%, \quad (4)$$

much less than $L_{\text{cav}}$. $L_{\text{eff}}$ are the fundamental losses in the setup that determine the quality of the created cat states.

Further losses occur along the optical beam path from the cavity to the detection setup. These are mainly optical absorption losses $L_o = 14\%$. The mode matching between the signal and the local oscillator used in the homodyne setup introduces additional losses $L_m = 6.0\%$. The detector there comprises two photodiodes with $L_d = 1.5\%$. The subtracted photodiode signals are fed into a field-programmable gate array (FPGA) for further processing. We identified several more noise sources which add effective losses of $L_n = 5.5\%$ in total. Combining these loss channels as $1 - L_{\text{pred}} = \prod_i (1 - L_i)$ with $i \in \{o, m, d, n\}$, we obtain the propagation and detection losses $L_{\text{pred}} = 25\%$. Together with $L_{\text{eff}}$, this sets a lower limit to

the observed losses in our reconstructed states which we find as $L_{\text{det}} = 46(1)\%$ as obtained from the fit (solid line) in Fig. 2b, well below the threshold of 50\%, where the Wigner function becomes positive. While the protocol puts no intrinsic constraint on the size of produced cat states, in our experiment the size is limited by these losses to $\alpha^2 \lesssim 4$.

The even and odd cat states shown so far are special cases of more general coherent state superpositions of the form

$$|\psi_{\text{cat}}\rangle = \frac{1}{\mathcal{N}} \left( \cos(\xi/2) |\alpha\rangle + e^{i\theta} \sin(\xi/2) |e^{i\phi} \alpha\rangle \right) \quad (5)$$

with the parameters $\alpha$, $\phi$, $\theta$ and $\xi$. Our experiment allows to control all these parameters, the phase and modulus of $\alpha$, the optical phase $\phi$ between the coherent contributions, the superposition phase $\theta$ that determines even or odd cat states and the population fraction of the two coherent contributions $\xi$. This opens the possibility to create more complex states that could be useful for continuous-variable error correction codes.

Specifically, the angle of the last spin rotation (step 4 in Fig. 1a) controls the relative amplitude of the two coherent state contributions in the generated cat state. Scanning the Raman-pulse area $\xi$ from 0 to $\pi/2$ continuously transforms the optical state (after measuring the atom) from a coherent state into a cat state with equal amplitudes of coherent contributions, see Fig. 3a.
The cat state in equation (5) can be continuously tunned from an even cat state into an odd cat state via a change of the phase θ of the final π/2 rotation (step 4 in Fig. 1a). This phase is imprinted onto the observed interference fringes in the Wigner function. Fig. 3b shows the continuous transition from an even into an odd cat state and vice versa for 0 ≤ θ ≤ π. The optical phase φ is varied between 0 and 2π via a detuning between the impinging light and the cavity resonance. On resonance, φ = π.

So far we have used the light-matter entanglement provided by the Schrödinger cat state to produce an optical cat state by making a measurement on the atom. We now show that light and matter are indeed entangled by observing an entanglement witness. To this end, we measure the atomic state in three different detection bases via the application of spin rotation pulses. Representing the spin in terms of Pauli matrices σ = {σx, σy, σz}, and the corresponding optical state as joint Wigner functions W = {Wx, Wy, Wz}, which play the role of Stokes parameters, allows us to express the full density matrix of the entire atom-cat system as ρac = 1/2 ∑σ Wσ. Here Wx is the average photonic density matrix with the atom traced out. Wx,y,z are sums of photonic density matrices weighted with the outcome of the atomic spin measurement of ±1 for outcomes |↑⟩ and |↓⟩, respectively, in the measurement bases x, y and z. Wigner representations of our measured Wσ are displayed in Fig. 4. The obtained density matrix ρac has a negativity of N = 0.005 ± 0.005 for an input state with α = 1.4. The statistically significant positive value witnesses the entanglement.

The reflection mechanism (equation 1) that was so far used to create entanglement is essentially a quantum-logic gate between an atom and a light field. Here, the coherent states |α⟩ and |−α⟩ form an orthogonal qubit basis as long as α is large enough. For our value, α = 1.4, the overlap of these basis states is |⟨−α|α⟩|2 = exp(−4|α|2) = 3.9 × 10⁻⁴, small enough for a good qubit. To probe our mechanism as a gate, we employ a basis set of input states |↑, α⟩, |↑, −α⟩, |↓, α⟩, |↓, −α⟩. In this basis, the gate acts as a CNOT with the truth table

\[
\begin{array}{c|cccc}
\alpha & |↑, +α⟩ & |↑, −α⟩ & |↓, +α⟩ & |↓, −α⟩ \\
\hline
0 & 1 & 0 & 0 & 0 \\
1 & 1 & 0 & 0 & 0 \\
\end{array}
\]

Here the atom serves as the control qubit that can flip the optical target qubit. We characterise this gate with coherent pulses containing a mean photon number of |α|² = π ≈ 2, the same value as in the measurements confirming the entanglement between the atom and the light field. The quantum nature of the gate has thus been demonstrated in the context of Fig. 4. The classical truth table is shown in Fig. 5 and exhibits the expected CNOT behaviour. Here, the optical losses change the input states |±α⟩ to output states |±αdet⟩ with αdet = α√1 - Tdet = 1.0. This change is well-characterised and thus predictable, and the overlap between the two output states is still small (1.8 × 10⁻²). We calculate the mean overlap of the observed output states with the expected output states and find values around 86%. The reduction from 100% comes mainly from phase noise due to cavity and laser frequency fluctuations that broaden the coherent-state Wigner function of the reflected field in the azimuthal direction. Note that the gate acts as a CPHASE in the |↑⟩/|↓⟩ ⊗ |cat⟩ basis.
We emphasise that in contrast to all protocols that have so far been realised with single optical photons as qubits, the creation of Schrödinger cat states as demonstrated here prepares an entangled light-matter state in each trial, without any post-selection. All relevant parameters are controlled, giving access to all degrees of freedom required for continuous-variable quantum information processing with coherent-state superpositions as qubits\cite{17,18}. The size of these qubits is only limited by optical losses, elimination of which is the key challenge for future improvements, e.g., by utilising loss-correcton codes that are possibly similar to those which have already been tested in the microwave domain\cite{29}. Optical cats as qubits could therefore be a promising alternative to single photons for quantum communication in a future quantum internet\cite{30}.

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\section*{Author Contributions}

Experimental data were taken and analysed by B.H., S.W., S.D. and L.L.. The paper was written by B.H., S.W. and G.R., with input from all authors.
Methods

Experimental setup. A schematic of the setup is shown in Extended Data Fig. 1. We employ an optical Fabry-Pérot cavity with a mirror separation of 0.5mm and finesse of $F = 6 \times 10^4$. The relevant cavity QED parameters are $(g, \kappa, \gamma) = 2\pi(7.8, 2.5, 3.0) \text{MHz}$ which place our atom-cavity system in the strong coupling regime. In our notation, $g$ is the atom-photon coupling rate, $\kappa$ the cavity field decay rate and $2\gamma$ the spontaneous decay rate of the excited atomic state $|e\rangle$. Initially, atoms are trapped in a magneto-optical trap and transferred into the cavity. Once an atom arrives at the centre of the cavity, a three dimensional standing wave configuration is switched on. We employ a red detuned dipole trap ($1064\text{nm}$) perpendicular to the cavity axis along the $x$ direction in Extended Data Fig. 1 and two blue detuned (771 nm) repulsive standing waves in the $y$ and $z$ direction. The atom is held in this configuration on average for $10$s. We monitor the atom with a camera and a high numerical aperture (NA = 0.4) objective. Conditioned on the presence of an atom in the cavity, we start the experimental protocol. Initially, the atom is pumped to the state $|\uparrow\rangle = \{5^2S_{1/2}, F = 2, m_F = 2\}$ via a right-circularly polarised pump laser along the cavity axis on the $F = 2 \leftrightarrow F' = 3$ transition. The pump light is applied for $120 \mu$s, and successful pumping of the atom is heralded by a reduction of the cavity transmission$^{31}$. After this step, the atom can effectively be considered as a three level system consisting of the states $|\uparrow\rangle$, $|\downarrow\rangle = \{5^2S_{1/2}, F = 1, m_F = 1\}$ and $|e\rangle = \{5^2P_{3/2}, F' = 3, m_F = 3\}$. Spin rotations and the generation of superposition states of $|\uparrow\rangle$ and $|\downarrow\rangle$ are realized with a pair of Raman lasers (green arrow in Fig. 1). The frequency difference between these two orthogonally polarised laser beams corresponds to the frequency difference between the states $|\downarrow\rangle$ and $|\uparrow\rangle$ of $6.8\text{GHz}$. The Raman lasers are $131\text{GHz}$ red detuned from the $^{87}\text{Rb}$ D1 transition at 795 nm and impinge onto the atoms from the side. Each application of the cat state protocol is followed by a Sisyphus cooling scheme that is applied to the atom for $1700\mu$s. The entire protocol is repeated at a rate of $500\text{Hz}$.

Homodyne detection. We characterise the optical states via balanced homodyne detection$^{10}$. The signal is mixed with a $1.8\text{mW}$ local oscillator (LO) beam of a continuous-wave laser on a 50/50 non-polarising beam splitter (NPBS). The intensities of the two output beams are measured on separate photodiodes (Hamamatsu S3883) and subtracted electronically. This difference signal is proportional to the signal beam amplitude and amplified to a measurable level. We record a time trace of each signal pulse with a field-programmable gate array (FPGA), and average the amplitude over the known temporal mode of the pulse. This results in one electric field amplitude (quadrature) in each homodyne measurement. The quadrature is a projection of the amplitude in phase space under a projection angle that is given by the relative phase of the LO with respect to the reflected beam. We do not interferometrically stabilise this phase in our homodyne measurements. The phase thus changes due to thermal drifts of the employed optics and a $1\text{Hz}$ detuning of the LO that ensures uniform sampling over all projection angles. For each homodyne measurement the relative phase between signal and LO is determined separately. To this end, the power of the reflected beam is increased by a factor of 30 for a sufficient signal-to-noise ratio and the LO phase is swept for $100\mu$s over $2\pi$ to record a full interference oscillation. In order not to lose the atom from the intra-cavity dipole-traps during this measurement, it is pumped to the non-coupling state $|\downarrow\rangle$, which is almost completely decoupled from the cavity. In each shot, the phase of the reflected beam with respect to the local oscillator can be determined with an accuracy of $2$ degrees. Phase drifts within a few $100\mu$s between the experiment and the phase determination are negligible.

For the state reconstruction, we perform many repeated measurements, each resulting in one quadrature value and a corresponding LO phase which determines the projection angle in phase space. We acquire on the order of $10^6$ samples for each measured optical state, and each outcome can be considered as one projective measurement of the density matrix $\rho$. The density matrix in truncated Fock space is obtained via maximum likelihood estimation using the iterative $R\rho R$ algorithm$^{32}$. This ensures a physical density matrix and allows for the correction of optical losses. We obtain an estimate for the uncertainty (covariance matrix) of each reconstructed density matrix using the Hessian of the likelihood function$^{33}$. The density matrix fully characterises the quantum state and therefore allows us to compute all derived quantities, such as the Wigner function, photon statistics or fringe visibility. Some quantities, such as squeezing, can directly be derived from the raw data and do not depend on the density matrix reconstruction.

Analytic treatment of cavity losses. Our atom-cavity system is well described by input-output theory$^{34}$ in the limit of slowly varying light intensities and low atomic excitation probability. For simplicity, we assume the case where cavity, atom and coherent input light $|\alpha\rangle$ are on resonance, and all amplitudes are real. We consider four modes: cavity reflection $|r\rangle$, cavity transmission $|t\rangle$, mirror losses $|m\rangle$ and scatter-
The total losses can be expressed as \( r_{l_0} = \frac{Ng^2 + (\kappa - 2\kappa)\gamma}{Ng^2 + \kappa\gamma} \alpha \) (7), \( t_{l_0} = \frac{2\sqrt{\kappa}\kappa\gamma}{Ng^2 + \kappa\gamma} \alpha \) (8), \( m_{l_0} = \frac{2\sqrt{\kappa}\kappa\gamma}{Ng^2 + \kappa\gamma} \alpha \) (9), and \( a_{l_0} = \frac{2\sqrt{\kappa}\gamma\sqrt{Ng^2 + \kappa\gamma}}{\alpha} \) (10), respectively. Here, \( N \) is the number of coupling atoms, \( N = 0 \) for \( |\downarrow\rangle \) and \( N = 1 \) for \( |\uparrow\rangle \). \( \kappa \) are optical field decay rates into the respective modes and \( \kappa = \kappa + \kappa_m \).

With a coherent input \(|\alpha\rangle\) of amplitude \( \alpha \), the output will be a superposition of two coherent fields of amplitude \( r_l \) and \( r_r \), which we call \( |r_l\rangle \) and \( |r_r\rangle \), respectively. If the moduli of the two amplitudes differ, the resulting cat state in phase space will be off-centred, which is easily correctable with a displacement. In our experiment however, the two amplitudes are nearly identical. The total size of the cat state is given by the peak separation

\[
\alpha_{\text{out}} = \frac{1}{2} |r_l - r_r| = \frac{\kappa_r g^2}{\kappa g^2 + \kappa\gamma} \alpha = \eta \alpha ,
\]

where we use the definition

\[
\eta := \frac{\kappa_r g^2}{\kappa g^2 + \kappa\gamma} (= 0.81) .
\]

The total losses can be expressed as \( L_{\text{cav}} = 1 - \alpha_{\text{out}}^2/\alpha^2 = 1 - \eta^2 = 0.34 \).

Henceforth we consider the reflected light modes \(|r_l\rangle \), \(|r_r\rangle \) and the loss modes \(|l_l\rangle := |t_l\rangle |m_1\rangle |a_l\rangle \) and \(|l_r\rangle := |t_r\rangle |m_1\rangle |a_r\rangle \), which have overlaps of

\[
\langle r_l | r_l \rangle = e^{-2\eta^2 a^2} ,
\]

\[
\langle l_l | l_l \rangle = \langle t_l | t_l \rangle \langle m_1 | m_1 \rangle \langle a_l | a_l \rangle = e^{-2(1-\eta)\alpha^2} ,
\]

\[
\langle l_l | l_r \rangle = e^{-2\eta^2 a^2} .
\]

We reflect the light when the atom is in an equal superposition of \( |\uparrow\rangle \) and \( |\downarrow\rangle \) and probe the atom after a consecutive \( \pi/2 \) rotation with phase \( \theta \). The produced optical state right after the reflection and state detection of the atom is

\[
|\psi_{\text{out}}\rangle = \frac{|r_l\rangle |t_l\rangle + e^{i\theta} |r_r\rangle |l_l\rangle}{\sqrt{2(1 + e^{-2\eta^2 a^2} \cos \theta)}}
\]

for the atom measured in \( |\downarrow\rangle \), and similar with \( \theta \to \theta + \pi \) for the atom in \( |\uparrow\rangle \). Light in the loss modes will be dissipated in the environment. The remaining optical state \( \rho \) is obtained by tracing out the losses,

\[
\rho = \text{tr}_{\text{in}} |\psi_{\text{out}}\rangle \langle \psi_{\text{out}}| .
\]

Here, the coherence terms are reduced due to optical losses by a factor \( \exp(-2(1 - \eta)\alpha_0^2) \). The state \( \rho \) is equivalent to a cat state of original amplitude \( \alpha_0 = \sqrt{\eta} \alpha \) that has undergone coherence-reducing intensity losses\(^2\) of \( L_{\text{eff}} = 1 - \eta = 0.19 \).

Thus, only a part of the total cavity losses \( L_{\text{cav}} = 1 - \eta^2 \) affects the coherences. In terms of optical depth, the coherence-reducing losses are exactly half the total losses.

The Wigner function \( W \), given for a generic lossy cat state\(^2\), is:

\[
W(q,p) = \frac{1}{2\pi} \left( e^{-p^2 - (q - \sqrt{2}r_l)^2} + e^{-p^2 - (q - \sqrt{2}r_r)^2} + 2e^{-2(1-\eta)\alpha^2}e^{-q^2 - (q - (r_l + r_r)/\sqrt{2})^2} \cdot \cos(\theta + \sqrt{8\eta \alpha p}) \right)/(1 + e^{-2\eta^2 a^2} \cos \theta) .
\]

It consists of two Gaussian peaks of separation \( \sqrt{8\eta \alpha} \) and a fringe term in the centre whose amplitude is reduced by \( \exp(-2(1 - \eta)\alpha_0^2) \) through the losses. The fringe centre \( (q_0, p_0) = ((r_l + r_r)/\sqrt{2}, 0) \), is equal to \((0, 0)\) in our experiment, because the reflection amplitudes \( r_l \) and \( r_r \) have the same magnitude. The interference fringe visibility defined by

\[
V = \frac{\sinh(2\eta \alpha_0^2)}{\sinh(2\alpha_0^2)} = \frac{\sinh(2(1 - L_{\text{eff}})\alpha_0^2)}{\sinh(2\alpha_0^2)} ,
\]

taking into account the normalization of \( W \). If the coherent contributions in the final cat states have little overlap, \( \eta \alpha_0^2 \gg 1 \), the visibility decays exponentially with respect to the losses and to the cat size \( \alpha_0^2 \):

\[
V \approx \exp(-2L_{\text{eff}} \alpha_0^2) .
\]

The fringes essentially vanish when the number of lost photons \( L_{\text{eff}} \alpha_0^2 \) exceeds one half. This effectively limits the achievable size of a cat-state.

Finally we compute the fidelity of \( \rho \) with respect to an ideal cat state of the same size

\[
|\psi_{\text{cat}}\rangle = |r_l\rangle + e^{i\theta} |r_r\rangle \sqrt{2(1 + e^{-2\eta^2 a^2} \cos \theta)}
\]

which yields

\[
F = \langle \psi_{\text{cat}} | \rho | \psi_{\text{cat}} \rangle = 1 - \frac{(1 - e^{-4\eta^2})}{(1 - e^{-2(1-\eta)\alpha_0^2})} \frac{1 - e^{-2(1-\eta)\alpha_0^2}}{(1 + e^{-2\eta^2 a^2} \cos \theta)} .
\]
Squeezing of cat states. Extended Data Fig. 2 shows an example of an even cat state of $\alpha = 0.7$ and the respective histogram along the $p$-axis. The data show squeezing of 1.18(3) dB. Squeezing can only be observed for even cat states. The odd cat states have a minimum at the centre of phase space which necessarily broadens the distribution. Fig. 2a in the main text shows a scan of the cat size and the corresponding width of the respective distributions for the even and the odd cat states.

**Extended Data Figure 2 | Squeezed cat state at $\alpha = 0.7$.** a, Reconstructed Wigner distribution from $\approx 7 \times 10^4$ measured quadrature values by means of the maximum likelihood estimation in Fock space. No loss-correction was applied. b, Histogram of measured quadrature values for projection angles in the interval $90^\circ \pm 5^\circ$. The dashed curve shows the respective histogram of a vacuum state. c, Standard deviation of quadrature values vs. projection angle given by the relative phase of the local oscillator. Fitted values range from 0.62 (squeezed) to 0.99 (anti-squeezed). The vacuum noise level is at $1/\sqrt{2}$ (dashed line).

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