A squeezed mechanical oscillator with millisecond quantum decoherence

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An enduring challenge in constructing mechanical-oscillator-based hybrid quantum systems is to ensure engineered coupling to an auxiliary degree of freedom and maintain good mechanical isolation from the environment, that is, low quantum decoherence, consisting of thermal decoherence and dephasing. Here we overcome this challenge by introducing a superconducting-circuit-based optomechanical platform that exhibits low quantum decoherence and has a large optomechanical coupling, which allows us to prepare the quantum ground and squeezed states of motion with high fidelity. We directly measure a thermal decoherence rate of 20.5 Hz (corresponding to $T_1 = 7.7$ ms) as well as a pure dephasing rate of 0.09 Hz, yielding a 100-fold improvement in the quantum state lifetime compared with prior optomechanical systems. This enables us to reach a motional ground-state occupation of 0.07 quanta (93% fidelity) and realize mechanical squeezing of $-2.7$ dB below the zero-point fluctuation. Furthermore, we observe the free evolution of the mechanical squeezed state, preserving its non-classical nature over millisecond timescales. Such ultralow quantum decoherence not only increases the fidelity of quantum control and measurement of macroscopic mechanical systems but may also benefit interfacing with qubits, and places the system in a parameter regime suitable for tests of quantum gravity.

Quantum control and measurement of mechanical oscillators has applications ranging from quantum metrology¹⁻³ and quantum computing⁴⁻⁵ to fundamental test of quantum mechanics itself⁶⁻⁷ or searches for dark matter⁸⁻¹⁰. This has been achieved by coupling mechanical oscillators to auxiliary degrees of freedom in the form of optical or microwave cavities¹⁰ or superconducting qubits¹¹⁻¹², allowing numerous advances such as mechanical squeezing¹³⁻¹⁵, quantum state transfer¹⁶⁻¹⁷, quantum transduction¹⁸⁻¹⁹ or teleportation²⁰. The decoherence of a mechanical oscillator induced by the interaction with its environment conceals macroscopic quantum phenomena and limits the realization of mechanical-oscillator-based quantum protocols¹⁻³,¹⁶⁻²¹. Quantum decoherence can be characterized with two independent rates: the thermal decoherence rate ($\Gamma_{\text{th}} = (n_{\text{th}} + 1)\Gamma$), where $n_{\text{th}}$ is the thermal bath occupation and $\Gamma$ is the bare damping rate), which describes the rate at which phonons are exchanged with the thermal bath, and the pure dephasing rate ($\Gamma_{\phi}$), caused by mechanical frequency fluctuations, that is, phonon-number-conserving interactions with the environment²². Even though the lowest thermal decoherence has been achieved in optomechanical crystals at millikelvin temperatures ($\Gamma_{\text{th}}/2\pi \approx 0.1$ Hz), such systems experience large dephasing ($\Gamma_{\phi}/2\pi \approx 4$ kHz), limiting their quantum coherence²². Soft-clamped dissipation-diluted Si$_3$N$_4$ membranes²¹ and levitated particles²² are other examples of optomechanical platforms interfacing with light, which achieved thermal decoherence rates of $\mathcal{O}(1$ kHz), but support limited optomechanical protocols as they operate in the non-resolved-sideband or cavity-free regimes, or suffer from optical

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observing free evolution of the prepared ground state as well as squeezed state, we report a thermal decoherence rate of 20.50 Hz (corresponding to 130 quanta s–1 motional heating rate) and a pure dephasing rate of 0.09 Hz, showing that quantum decoherence is dominated by thermal decoherence, comparable with motional decoherences achieved in trapped ion systems in a high vacuum 40, where the thermal decoherence is typically \( \mathcal{O}(10 \text{ Hz}) \), whereas the dephasing rate reaches \( \mathcal{O}(100 \text{ Hz}) \) (ref. 41).

**Ultracoherent circuit optomechanical platform**

We develop a nanofabrication process based on a silicon-etched trench, which enables us to substantially enhance the mechanical quality factor \( Q_m \) (Methods and Supplementary Information provide details about the fabrication process). Figure 1a shows a vacuum-gap capacitor with the top plate suspended on a circular trench with a gap size of 180 nm. The capacitor is shunted by a spiral inductor (Fig. 1b,c), forming a microwave LC resonator with a frequency of \( \omega_c/2\pi = 5.55 \text{ GHz} \) and a total decay rate of \( \kappa /2\pi = 250 \text{ kHz} \), which is inductively coupled to a waveguide (Fig. 1f). This superconducting circuit is operated in a dilution fridge with ~11 mK base temperature. The flat geometry of the top plate enhances the mechanical quality factor, resulting in an ultralow quantum decoherence, that is, both thermal decoherence and dephasing, and exhibit an efficient optomechanical coupling for quantum control and measurement. By observing free evolution of the prepared ground state as well as squeezed state, we report a thermal decoherence rate of 20.50 Hz (corresponding to 130 quanta s–1 motional heating rate) and a pure dephasing rate of 0.09 Hz, showing that quantum decoherence is dominated by thermal decoherence, comparable with motional decoherences achieved in trapped ion systems in a high vacuum 40, where the thermal decoherence is typically \( \mathcal{O}(10 \text{ Hz}) \), whereas the dephasing rate reaches \( \mathcal{O}(100 \text{ Hz}) \) (ref. 41).

**Fig. 1 | Ultracoherent circuit optomechanics.** a, False-coloured SEM image of a mechanically compliant parallel-plate capacitor. b, A microwave superconducting LC resonator consisting of the capacitor shunted by a spiral inductor. c, Magnified image showing a silicon-etched trench, inductor air bridges and galvanic connection. d, Focused-ion-beam cross section of a test capacitor—with a higher gap size than the main device—before removing a SiO\(_2\) sacrificial layer, where Pt is used as the focused-ion-beam protective layer. e, Schematic of the cross section of the suspended capacitor over the trench. f, Mode diagram of an optomechanical system. Here \( \kappa_m \) and \( \kappa_0 \) are the internal loss and external coupling rates of the cavity, respectively. g, Equivalent circuit diagram of the system. h, Ring-down trace, showing the energy decay of the mechanical oscillator with a rate of \( \Gamma_m/2\pi = 45 \text{ mHz} \). The red line is the exponential fit. The top inset shows the finite element method simulation of the fundamental mechanical mode of the drum. The blue line (bottom inset) shows the averaged PSD of the mechanics with 1 Hz measurement resolution bandwidth (grey line), indicating the frequency instability of less than 0.1 Hz.
plate (Fig. 1d,e) ensures minimal clamp and radiative mechanical losses, as well as stress relaxation in the aluminium thin film. A mechanical ring-down measurement (Fig. 1h) clearly exhibits the extremely low dissipation rate of $g_0/2\pi = 45$ mHz for the fundamental drumhead mode with a frequency of $\Omega_m/2\pi = 1.8$ MHz, corresponding to $Q_m = 4 \times 10^6$. This can be explained by the loss dilution factor $^{10}$ estimated to be $D_0 \approx 100$ from finite element method simulation for such a flat drumhead (Fig. 1h, top inset). The single-photon optomechanical coupling rate is measured to be $g_0/2\pi = 13.4 \pm 0.5$ Hz (Supplementary Information). We note that lower gap sizes lead to higher $g_0$ values, but not implemented in this work. Furthermore, the frequency fluctuation is observed below 0.1 Hz—inferred as an upper bound for dephasing—by measuring the power spectral density (PSD) of a thermomechanical sideband averaged over more than an hour and subtracting a measurement resolution bandwidth of 1 Hz (Fig. 1h (bottom inset) and Supplementary Information).

**High-fidelity optomechanical ground-state cooling**

The extremely high mechanical quality factor, together with the sufficient optomechanical coupling, enables us to perform an effective optomechanical sideband cooling$^{12}$ to prepare the mechanical oscillators in its quantum ground state with high fidelity. As shown by the schematic in Fig. 1g, in the resolved-sideband regime, where the quantum backaction does not influence the final phonon occupation$^{13}$, that is, $(\kappa/4\Omega_m)^2 = 0.001 \ll 1$ in our case, the phonon occupation of the mechanical oscillator in the presence of a cooling pump red-detuned by $B_m$ from the cavity frequency is given by

$$n_m = n_{m}^{th} + \frac{c}{1 + c} n_c,$$

where $c = 4 n_{0} g_{0}^{2}/(\kappa T_m)$ is the optomechanical co-operativity with the intracavity pump photon number $n_{0}$ and $n_c = \kappa_0 n_{0}^{th}/\kappa$ is the cavity thermal photon number induced by a finite loss rate of $\kappa_0$ to an intrinsic photon bath with $n_{0}^{th}$. The strong cooling pump may heat up the intrinsic photon bath occupation and consequently $n_c$, which normally imposes the minimum-achievable phonon occupancy in the large co-operativity limit, that is, $c \gg 1$ (ref. 12). We discovered that the thin native oxide layer in the galvanic connection between the top and bottom layers (Fig. 1c) is the dominant source of such cavity heating, which has been ubiquitous in all microwave optomechanical experiments. We substantially reduced the heating by removing the oxide to achieve $n_c = 0.05$ quanta at high co-operativities (Supplementary Information).

To reliably characterize the phonon occupation close to the ground state, we use optomechanical sideband asymmetry$^{14}$ as an out-of-loop calibration. As shown in Fig. 2a, we apply a strong cooling pump, and two weak, blue- and red-detuned probes with balanced powers to generate Stokes and anti-Stokes optomechanical sidebands, respectively, on the cavity resonance with a few kilohertz spacing to individually measure them. Figure 2b shows the simplified experimental setup, where a Josephson travelling-wave parametric amplifier$^{15}$ is
used to amplify the microwave signals with an added noise of $\frac{1}{2} + n_{\text{add}}^c = \frac{1}{2} + 0.3$ quanta and sufficient gain of $G_{\text{opt}} = 25\text{ dB}$ to suppress the classical noise dominated by the high-electron-mobility transistor amplifier ($n_{\text{add}}^c \approx 8$), enabling a nearly quantum-limited measurement of thermomechanical noise spectrum with an effective added noise of $\frac{1}{2} + n_{\text{add}}^c = \frac{1}{2} + 0.9$ (Fig. 2c shows the calibrated noise floor). Figure 2c,e shows the measured thermomechanical noise spectrum of the cavity thermal emission, as well as the Stokes and anti-Stokes sidebands, which are used for obtaining their powers, expressed as $P_c$, $P_b$, and $P_e$, respectively, by fitting a Lorentzian to the PSD of the sidebands. Although the sideband asymmetry may allow us to perform the cavity-free measurement of $n_m$, a finite cavity heating distorts the asymmetry, that is, $P_b \propto n_m + 1$ and $P_c \propto n_m$, preventing us from extracting $n_m$ without the prior knowledge of $n_c$ (ref. 43). Nevertheless, we are able to simultaneously extract both $n_c$ and $n_m$ without any calibration of the measurement chain by analytically obtaining them from the two sidebands powers normalized by the cavity thermal emission power, expressed by

$$\frac{P_b}{P_c} = \frac{\Gamma_b}{\kappa} \frac{n_m + 1 + 2n_c}{n_c} \quad \text{and} \quad \frac{P_c}{P_e} = \frac{\Gamma_c}{\kappa} \frac{n_m - 2n_c}{n_c}, \quad (2)$$

where $\Gamma_b$ ($\Gamma_e$) is the optomechanical (anti-)damping rate of the red (blue) probe. Importantly, this analysis allows us to calibrate the scaling factor between the actual occupations and the measured powers, which can be used to directly extract $n_c$ and $n_m$ independently from the cavity thermal emission and the sideband induced by the cooling pump, even when the two probes are off, thereby avoiding the quantum backaction induced by the blue probe and an additional cavity heating (Supplementary Information). Using the PSDs of the thermomechanical sideband from the cooling pump (Fig. 2d) and the cavity thermal emission when the two probes are off, we thus extract $n_c$ and $n_m$ as a function of the cooling pump co-operativity (Fig. 2f). The result shows a high-fidelity ground state cooling down to $n_m = 6.8(\pm 0.9) \times 10^{-9}$ quanta (93% ground-state occupation, which is $-8.7\text{ dB}$ of the zero-point energy), mainly limited by the cavity heating.

**Measurement of motional heating rate**

Next, we directly measure the thermal decoherence by recording the thermalization of the mechanical oscillator out of the ground state using a time-domain protocol (Fig. 3b), where we first prepare the ground state, and leave the system to freely evolve for a certain time $\tau_{\text{ev}}$. Using the optomechanical amplification technique with a blue-detuned pump (Supplementary Information), we intrinsically amplify the mechanical motion by $-50\text{ dB}$ with a minimal added noise, and measure both quadratures of motion encoded in the generated optomechanical sideband signal (Fig. 3a). Repeating this pulse sequence allows us to capture the quadrature distribution of the mechanical state, realizing quantum state tomography (Fig. 3d). As shown in Fig. 3c, we are able to precisely calibrate the amplification process using different phonon occupations as an input state, which are well calibrated by the sideband asymmetry measurement, resulting in $n_{\text{add}}^c = 0.80 \pm 0.09$ quanta.

Figure 3d shows examples of the measured quadrature distributions at different evolution times in units of $\sqrt{\text{quanta}}$. Figure 3e shows histogram for each quadrature is shown by purple bars. A Gaussian curve with the calculated variance is shown by the thick blue lines, reproducing each histogram. Each scatter plot contains 3,800 data points. Figure 3f shows examples of the measured quadrature distributions at different evolution times in units of $\sqrt{\text{quanta}}$. Figure 3e shows histogram for each quadrature is shown by purple bars. A Gaussian curve with the calculated variance is shown by the thick blue lines, reproducing each histogram. Each scatter plot contains 3,800 data points.
the free evolution from the ground state to the thermal equilibrium. The exponential fit results in a bare dissipation rate of $\Gamma_{\phi}/2\pi = 80$ mHz in the low-phonon-occupation regime, close to the value measured from the ring-down experiment with $n_m > 10^3$ (Fig. 1h). Figure 3e (right inset) shows the thermalization in shorter evolution times, where the thermal decoherence rate is directly measured as $\Gamma_{\phi}/2\pi = 20.5 \pm 0.6$ Hz, corresponding to a phonon lifetime of $T_1 = 7.7$ ms.

**Recording thermalization of squeezed mechanical state**
Finally, we generate a quantum squeezed state of our oscillator. Since the squeezed state is a phase-sensitive quantum state, its free evolution is subject to the dephasing in the system. Tracking its time evolution enables us to directly measure the quantum state lifetime and verify minimal dephasing in our mechanical oscillator. The ability to squeeze the mechanical oscillator critically relies on the residual thermal occupation on cooling, which in our case is below 0.1 quanta, implying that strong squeezing below the zero-point fluctuation is possible. We use the optomechanical dissipative squeezing technique by simultaneously applying two red- and blue-detuned pumps symmetrically with respect to the cavity frequency (Fig. 4b), and achieve $(X_{sq})^2/2 = -2.7^{+1.4}_{-2.3}$ dB squeezing in one quadrature of motion below the vacuum fluctuation and $(X_{sq})^2/2 = 8.1^{+0.3}_{-0.2}$ dB anti-squeezing in the other quadrature. These are obtained by subtracting the accurately calibrated $n_{\text{add}}$ in optomechanical amplification (Fig. 4a,c).

Figure 4d shows the measured quadrature scatter plots of the prepared squeezed state and its time evolution. We are able to record the free evolution of the prepared squeezed state (Fig. 4d) and observe the decoherence of both quadratures to thermal equilibrium (Fig. 4e). A slight difference is observed in the decoherence rates of the two quadratures, $(F_{sq} - F^{\text{meas}})/2\pi = 1.1(\pm 0.6)$ Hz. Comparing it with a numerical simulation allows us to characterize the dephasing rate of $\Gamma_{\phi}/2\pi = 0.09(\pm 0.05)$ Hz in our platform (Methods), in agreement with the measured frequency fluctuation discussed earlier.

We observe that the variance of the squeezed quadrature remains below the zero-point fluctuation up to 2 ms, demonstrating a remarkably long quantum-state storage time in a macroscopic mechanical oscillator.

**Conclusion and outlook**
The high-fidelity quantum control and measurement of mechanical oscillators with such extremely low thermal decoherence and pure dephasing rates may facilitate the implementation of qubit–mechanics interfaces and the generation of mechanical non-classical states, as well as realize long lifetime memories for quantum computation and communication. Furthermore, such a low quantum decoherence sets the stage to perform fundamental tests of quantum mechanics in macroscopic scales such as quantum gravity tests, as well as high-fidelity Bell tests and quantum teleportation.

**Online content**
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Methods

Fabrication of high-coherence electromechanical devices

Extended Data Fig. 1 shows an overview of the nanofabrication process we developed for circuit optomechanics. We define a trench in the substrate containing the bottom plate of the capacitor. The trench is then covered by a thick SiO₂ sacrificial layer, which inherits the same topography as the layer underneath. To remove this topography and obtain a flat surface, we use chemical mechanical polishing to flatten the SiO₂ surface. We then etch back the sacrificial layer down to the substrate layer and deposit the top aluminium plate of the capacitor. We release the structure by hydrofluoric acid vapour etching of SiO₂. At cryogenic temperatures, the high tensile stress ensures the flatness of the top plate. This will guarantee that the gap size is precisely defined by the depth of the trench and the thickness of the bottom plate and reduces the clamp losses in the mechanical oscillator. Supplementary Information provides a detailed description of each step.

Dephasing effect on thermalization of squeezed states

A mechanical squeezed state has a phase coherence between the Fock states, resulting in sensitivity to pure dephasing. By observing the free evolution of such a phase-sensitive state, we can characterize the pure dephasing rate of our mechanical oscillator. When the mechanical state is initialized in an isotropic Gaussian state, such as the vacuum or thermal states, the dephasing rates for the quadrature variances in all the phases are identical regardless of the amount of pure dephasing. However, when the mechanical state is initialized in the squeezed state, the dephasing rate of the quadrature variance in the squeezed axis is greater than that of the anti-squeezed axis if the pure dephasing rate is finite. Thus, we can use this property to extract the pure dephasing rate of the mechanical oscillator.

In the large-phonon-bath occupation limit (n_th ≫ 1), which is the case for our experiment, the free evolution of a mechanical oscillators with pure dephasing is described by a master equation:

\[ \frac{d\rho}{dt} = \Gamma_{\text{th}} [2\hat{a}^\dagger \hat{a} \rho + 2\Gamma_{\text{th}} (\hat{a}^\dagger \hat{a} \rho + \rho \hat{a}^\dagger \hat{a})/2] - \frac{1}{2} \left( \Gamma_{\text{th}} [2\hat{a}^\dagger \hat{a} \rho + 2\Gamma_{\text{th}} (\hat{a}^\dagger \hat{a} \rho + \rho \hat{a}^\dagger \hat{a})/2] + \Gamma_{\text{deph}} [2\hat{a}^\dagger \hat{a} \rho + 2\Gamma_{\text{deph}} (\hat{a}^\dagger \hat{a} \rho + \rho \hat{a}^\dagger \hat{a})/2] \right) \rho + 2\Gamma_{\text{th}} \rho \hat{a}^\dagger \hat{a} \rho + 2\Gamma_{\text{deph}} \rho \hat{a}^\dagger \hat{a} \rho - \frac{1}{2} \left( \Gamma_{\text{th}} [2\hat{a}^\dagger \hat{a} \rho + 2\Gamma_{\text{th}} (\hat{a}^\dagger \hat{a} \rho + \rho \hat{a}^\dagger \hat{a})/2] + \Gamma_{\text{deph}} [2\hat{a}^\dagger \hat{a} \rho + 2\Gamma_{\text{deph}} (\hat{a}^\dagger \hat{a} \rho + \rho \hat{a}^\dagger \hat{a})/2] \right) \rho, \]

where \( \Gamma_{\text{th}} \) is the thermal dephasing rate, \( \Gamma_{\text{deph}} \) is the pure dephasing rate, \( \rho \) is the density matrix and \( [2\hat{a}^\dagger \hat{a} \rho + 2\Gamma_{\text{th}} (\hat{a}^\dagger \hat{a} \rho + \rho \hat{a}^\dagger \hat{a})/2] \) is the Lindblad dissipator. Calculating the time evolution of the density matrix, we can numerically obtain the time evolution of the quadrature variances and phonon occupation as \( \langle X^2 \rangle = \text{Tr}[\rho X^2] \) and \( \langle b^2 \rangle_n = \text{Tr}[\rho b^2] \), respectively. Note that here we consider only a mechanical state with \( \text{Tr}[\rho X^1] = \text{Tr}[\rho b^1] = 0 \), for example, vacuum, thermal and squeezed states.

The blue and red circles in Extended Data Fig. 2a show the experimental results of the quadrature variances, namely, \( \langle X^2 \rangle \) in the squeezed axis and \( \langle X^2 \rangle \) in the anti-squeezed axis, as a function of the free evolution time, respectively, whereas the purple circles show the time evolution of the phonon occupation, that is, \( n_{\text{th}} = \langle X^2 \rangle + \langle X^2 \rangle / 2 - 1 / 2 \). Here we define the dephasing rate as a slope of the time evolution, that is, \( \Gamma_{\text{deph}} = \frac{\delta \langle X^2 \rangle}{\delta t} \). For linearly fitting the results (Extended Data Fig. 2a, black dotted lines), the dephasing rates are found to be \( \Gamma_{\text{deph}} = 0.09 \pm 0.03 \text{ Hz} \) and \( \Gamma_{\text{deph}} = 0.09 \pm 0.03 \text{ Hz} \), respectively. Therefore, to numerically determine the initial squeezed state using the experimentally obtained quadrature variances, that is, \( \langle X^2 \rangle = 0.27 \pm 0.10 \) and \( \langle X^2 \rangle = 3.27 \pm 0.20 \), we can consider the prepared squeezed state as a squeezed thermal state, which is defined as \( S \rho_{\text{th}} S^\dagger \), where \( S = \exp [\frac{1}{2} (\hat{b}^\dagger - \hat{b}^\dagger)] \) is a squeezing operator with a squeezing parameter of \( r \) and \( \rho_{\text{th}} \) is a thermal state with an average phonon number of \( n_{\text{th}} \). We can extract the thermal phonon number and the squeezing parameter as \( n_{\text{th}} = \langle X^2 \rangle / \langle X^2 \rangle - 1 / 2 \) and \( r = \frac{1}{4} \log \left( \frac{\langle X^2 \rangle}{\langle X^2 \rangle} \right) = 0.6 \pm 0.1 \), respectively.

Using the initial squeezed state, we calculate the time evolution of the quadrature variances for different pure dephasing rates and obtain the difference between the thermal decoherence rates of the squeezed and anti-squeezed variances. Extended Data Fig. 2b (green line) shows the numerical results of the differences in the thermal decoherence rates. As the dephasing rate increases, the difference becomes larger. We compare the numerical results with the experimentally obtained value (Extended Data Fig. 2b, blue line). From this, we find that the pure dephasing rate of our mechanical oscillator is found to be \( \Gamma_{\text{deph}} = 0.09 \pm 0.03 \text{ Hz} \).

Finally, we show the numerical results of the time evolution of the quadrature variances and the phonon occupation with a dephasing rate of \( \Gamma_{\text{deph}} = 0.09 \text{ Hz} \) (Extended Data Fig. 2a, green lines). The numerical results effectively reproduce the experimental ones.

Data availability

The data used to produce the plots within this paper are available via Zenodo at https://doi.org/10.5281/zenodo.7833893. All other data used in this study are available from the corresponding authors on reasonable request. Source data are provided with this paper.

Code availability

The code used to produce the plots within this paper is available via Zenodo at https://doi.org/10.5281/zenodo.7833893.

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Author contributions

A.Y. conceived the experiment. S.K., M.C. and A.Y. developed the theory. A.Y. designed and simulated the devices. A.Y. developed the fabrication process with assistance from M.C. M.C. and A.Y. fabricated the samples. A.Y. and M.C. developed the experimental setup. The measurement was performed by A.Y. and M.C., with assistance from S.K. The data analysis was performed by A.Y. with assistance from S.K. A.Y. introduced the phonon number calibration based on sideband asymmetry and conducted the numerical simulation for extracting the mechanical dephasing. The manuscript was written by A.Y., S.K., M.C. and T.J.K. T.J.K. supervised the project.

Competing interests

The authors declare no competing interests.

Additional information

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Extended Data Fig. 1 | Overview on the fabrication technique for highly coherent circuit optomechanics. The main steps of the process consist of etching a trench in the substrate followed by deposition of a sacrificial layer, planarization, top layer definition, release, and finally cool down. Due to the compressive stresses, the top plate may buckle up after the release. However, the drumhead shrinks and flattens at cryogenic temperatures, resulting in a controllable gap size.
Extended Data Fig. 2 | Free evolution of a mechanical squeezed state.
a, Quadrature variances and average phonon number of a squeezed state as a function of the free-evolution time. The blue, red and purple circles are the data for the quadrature variances in the squeezed and anti-squeezed axes, and the average phonon number, respectively. The black dotted lines are linear fits, while the green lines are the numerical simulation results. Error bars are corresponding to standard deviations. b, The difference of the thermal decoherence rates for the quadrature variances in the squeezed and anti-squeezed axes as a function of the pure dephasing rate. The green line shows the numerical simulation results, while the blue line is the experimentally obtained value. The shaded regions show the errors, respectively.