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Statement & Readings

Towards quantum superpositions of a mirror

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Abstract

Theoretical and experimental aspects of testing the quantum superposition principle for relatively massive objects will be presented. In collaboration with Prof. R. Penrose a specific experiment for achieving this was proposed [1]. The motivation for Prof. R. Penrose to consider such an experiment stems from his prediction that gravitational effects play a role in the reduction of a quantum superposition to a single "branch" of the superposition. There are two levels of reasoning that led Penrose to his prediction [2]. The highest level is related to the evolution of the universe; starting from the Big Bang the universe appears to evolve into a collection of black holes that evaporate into zero-mass fields. In order to reconcile this observation with the notion that entropy increases, Penrose concludes that there must be a physical process by which state-space reduces (in order to counterbalance the increase of entropy). He then identifies this with quantum measurements in which a state-space reduction (collapse of the wave-function) seems to take place. The second level of reasoning is based on the observation that a quantum wave function $\Psi(x,y,z,t)$ describing a spatial superposition of a massive object is simply inconsistent with the theory of relativity; a spatial superposition of a massive object implies two different space time structures (space-time curvature is determined by the position of the mass) and can therefore not be described by a single set of coordinates *x, y, z, t*. Penrose has estimated the energy associated with a superposition of two space-time structures resulting from a spatial superposition of a massive object, and argues that this amount of excess energy is allowed for a certain amount of time, invoking the Heisenberg uncertainty relation for time and energy. Based on this argument one can estimate that objects observed in quantum superpositions to date (such as atoms, BECs, superconducting currents and C60 molecules, and even the mechanical resonator recently investigated by A, Cleland [Nature 2010]) had much too small mass to observe the reduction of the wave function within the time scale of the performed experiments. It will, according to Prof. R. Penrose, require objects of about 10^{-12} kg (approximately corresponding to the objects considered below) to witness the collapse of the wave function on a time scale of the order of a second. The designed experiment [1] should bring a tiny mirror in a quantum superposition for several microseconds, thus approaching the regime of interest to test Penrose's predictions.

The experiment envisioned in 2003 has received a lot of attention and the challenge to design, fabricate and test such opto-mechanical systems that can display quantum-mechanical behavior is currently a strongly emerging field of science and technology. The reason for this huge interest is partly the fundamental importance, partly the application of optical cooling (an active cooling scheme will be presented [3]), partly the interest in quantum information processing where quantum mechanical resonators can couple to other quantum systems, and partly the potential for metrology applications in weak force and momentum detection. A good overview of state-of-the-art research in this field can be found in [4].

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Towards Quantum Superpositions of a Mirror

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We propose an experiment for creating quantum superposition states involving of the order of 10^{14} atoms via the interaction of a single photon with a tiny mirror. This mirror, mounted on a high-quality mechanical oscillator, is part of a high-finesse optical cavity which forms one arm of a Michelson interferometer. By observing the interference of the photon only, one can study the creation and decoherence of superpositions involving the mirror. A detailed analysis of the requirements shows that the experiment is within reach using a combination of state-of-the-art technologies.

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Introduction.—In 1935 Schrödinger pointed out that according to quantum mechanics even macroscopic systems can be in superposition states [1]. The associated quantum interference effects are expected to be hard to detect due to environment induced decoherence [2]. Nevertheless, there have been proposals on how to create and observe macroscopic superpositions in various systems [3–7], as well as experiments demonstrating superposition states of superconducting devices [8] and large molecules [9]. One long-term motivation for this kind of experiment is the search for unconventional decoherence processes [5,10].

In several of the above proposals a small quantum system (e.g., a photon [4–6] or a superconducting island [7]) is reversibly coupled to a large system (e.g., a moveable mirror [4–6] or a cantilever [7]) in order to create a macroscopic superposition. The existence of the quantum superposition of the large system is verified by observing the disappearance and reappearance of interference for the small system, as the large system is driven into a superposition and then returns to its initial state. The challenge is to find a feasible implementation of this idea.

Our proposal develops on the ideas in Refs. [4,5]. We also use results from Ref. [6], which relies on coupling between atoms and photons in a microcavity to create and detect superposition states of a moveable mirror. In particular, the formalism used in Ref. [6], based on Refs. [11,12], is applicable to our case. The main purpose here is to show that our purely optical proposal has the potential to be performed with current technology.

Principle.—The proposed setup, shown in Fig. 1, consists of a Michelson interferometer which has a high-finesse cavity in each arm. The cavity in arm (A) contains a tiny mirror attached to a micromechanical oscillator, similar to the cantilevers in atomic force microscopes. The cavity is used to enhance the radiation pressure of the photon on the mirror. The initial superposition of the photon being in either arm causes the system to evolve into a superposition of states corresponding to two distinct locations of the mirror. The observed interference of the photon allows one to study the creation of coherent superposition states of the mirror.

The system can be described by a Hamiltonian [6,11]

$$H = \hbar\omega_c a^\dagger a + \hbar\omega_m b^\dagger b - \hbar G a^\dagger a(b + b^\dagger), \qquad (1)$$

where ω_c and a are the frequency and creation operator for the photon in the cavity, ω_m and b are the frequency and phonon creation operator for the center of mass motion of the mirror, and $G = (\omega_c/L)\sqrt{(\hbar/2M\omega_m)}$ is the coupling constant, where L is the cavity length and M is the mass of the mirror.

Let us suppose that initially the photon is in a superposition of being in either arm A or B, and the mirror is in

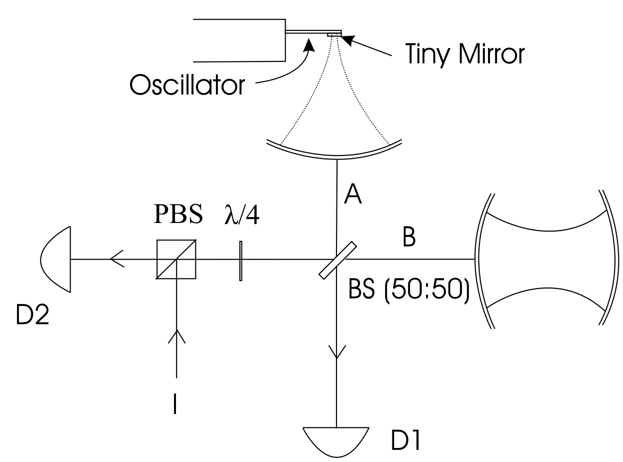

FIG. 1. The proposed setup: a Michelson interferometer for a single photon, where in each arm there is a high-finesse cavity. The cavity in arm A has a very small end mirror mounted on a micromechanical oscillator. The single photon comes in through I. If the photon is in arm A, the motion of the small mirror is affected by its radiation pressure. The photon later leaks out of either cavity and is detected at D1 or D2.

its ground state $|0\rangle_m$. Then the initial state is $|\psi(0)\rangle = (1/\sqrt{2})(|0\rangle_A|1\rangle_B + |1\rangle_A|0\rangle_B)|0\rangle_m$. After a time t the state of the system will be given by [6,12]

$$|\psi(t)\rangle = \frac{1}{\sqrt{2}}e^{-i\omega_c t}[|0\rangle_A|1\rangle_B|0\rangle_m + e^{i\kappa^2(\omega_m t - \sin\omega_m t)}|1\rangle_A|0\rangle_B \times |\kappa(1-e^{-i\omega_m t})\rangle_m], \tag{2}$$

where $\kappa = G/\omega_m$, and $|\kappa(1-e^{-i\omega_m t})\rangle_m$ denotes a coherent state with amplitude $\kappa(1-e^{-i\omega_m t})$. In the second term on the right-hand side the mirror moves under the influence of the radiation pressure of the photon in cavity A. The mirror oscillates around a new equilibrium position determined by the driving force. The parameter κ quantifies the displacement of the mirror in units of the size of the ground state wave packet.

The maximum interference visibility for the photon is given by twice the modulus of the off-diagonal element of the photon's reduced density matrix. By tracing over the mirror one finds from Eq. (2) that the off-diagonal element has the form $\frac{1}{2}e^{-\kappa^2(1-\cos\omega_m t)}e^{i\kappa^2(\omega_m t-\sin\omega_m t)}$. The first factor is the modulus, reaching a minimum after half a period at $t = \pi/\omega_m$, when the mirror is at its maximum displacement. The second factor gives the phase, which is identical to that obtained classically due to the varying length of the cavity.

In the absence of decoherence, after a full period, the system is in the state $(1/\sqrt{2})(|0\rangle_A|1\rangle_B + e^{i\kappa^2 2\pi}|1\rangle_A|0\rangle_B) \times |0\rangle_m$, such that the mirror is again disentangled from the photon. Full interference can be observed if the photon is detected at this time, provided that the phase factor $e^{i\kappa^2 2\pi}$ is taken into account. This revival, shown in Fig. 2, demonstrates the coherence of the superposition state that exists at intermediate times. For $\kappa^2 \gtrsim 1$ the superposition involves two distinct mirror positions. If the environment of the mirror "remembers" that the mirror has moved, then, even after a full period, the photon will still be entangled with the mirror's environment, and thus the revival will not be complete. Therefore the setup can be used to measure the decoherence of the mirror.

Here we have assumed that the mirror starts out in its ground state. We will argue below that optical cooling close to the ground state should be possible. However, in Ref. [6] it was shown that this is not necessary for observing the revival, although for a thermal mirror state with an average phonon number $\bar{n} = 1/(e^{\hbar\omega_m/kT} - 1)$ the revival peak is narrowed by a factor of $\sqrt{\bar{n}}$, leading to stricter requirements on the stability; see Fig. 2 and the discussion below. We now discuss the experimental requirements for achieving a superposition of distinct mirror positions and for observing the revival at $t = 2\pi/\omega_m$.

Conditions for displacement by ground state size.—We require $\kappa^2 \gtrsim 1$, which implies the momentum imparted by the photon has to be larger than the initial quantum uncertainty of the mirror's momentum. Let N denote the number of round-trips of the photon in the cavity during one period of the mirror's motion, such that $2NL/c = 2\pi/\omega_m$. The condition $\kappa^2 \gtrsim 1$ can be written

$$\frac{2\hbar N^3 L}{\pi c M \lambda^2} \gtrsim 1, \tag{3}$$

where λ is the wavelength of the light. The factors entering Eq. (3) are not all independent. The achievable N, determined by the quality of the mirrors, and the minimum mirror size (and hence M) both depend on λ. The mirror's lateral dimension should be an order of magnitude larger than λ to limit diffraction losses. The thickness required in order to achieve sufficiently high reflectivity depends on λ as well.

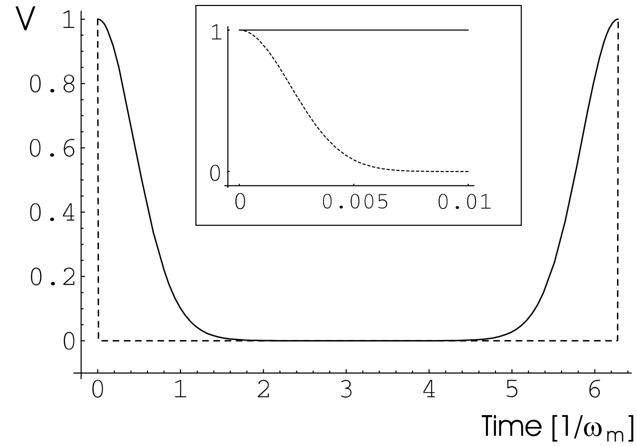

FIG. 2. Time evolution of the interference visibility V of the photon over one period of the mirror's motion for the case where the mirror has been optically cooled close to its ground state ($\bar{n} = 2$, solid line) and for $T = 2$ mK, which corresponds to $\bar{n} = 100\,000$ (dashed line—see also inset). The visibility decays after $t = 0$, but in the absence of decoherence there is a revival of the visibility after a full period. The width of the revival peak scales like $1/\sqrt{\bar{n}}$.

Equation (3) allows one to compare the viability of different wavelength ranges. While the highest values for N are achievable for microwaves using superconducting mirrors (up to 10^{10}), this is counteracted by their longer wavelengths. On the other hand, there are no good mirrors for highly energetic photons. The optical regime is optimal, given current mirror technology. We propose an experiment with λ around 630 nm.

The cavity mode needs to have a sharp focus on the tiny mirror, which requires the other cavity end mirror to be large due to beam divergence. The maximum cavity length is therefore limited by the difficulty of making large high-quality mirrors. We propose a cavity length of 5 cm, and a small mirror size of $10 \times 10 \times 10$ μm, leading to a mass of order 5×10^{-12} kg.

Such a mirror on a mechanical oscillator can be fabricated by coating a silicon cantilever with alternating layers of SiO_2 and a metal oxide. The best current optical mirrors are made in this way. A larger silicon oscillator has been coated with SiO_2/Ta_2O_5 and used as part of a high-finesse cavity in Ref. [13].

For the above dimensions the condition Eq. (3) is satisfied for $N = 5.6 \times 10^6$. Correspondingly, photon loss per reflection must be smaller than 3×10^{-7}, about a factor of 4 below reported values for such mirrors [14] and for a transmission of 10^{-7}, consistent with a 10 μm mirror thickness. For these values, about 1% of the photons are still left in the cavity after a full period of the mirror. For the above values of N and L one obtains a frequency $\omega_m = 2\pi \times 500$ Hz. This corresponds to a spread of the mirror's ground state wave function of order 10^{-13} m.

The fact that a relatively large L is needed to satisfy Eq. (3) implies that the creation of superpositions following the microcavity based proposal of Ref. [6] imposes requirements beyond current technology. A large L is helpful because, for a given N, it allows a lower frequency ω_m, and thus a more weakly bound mirror that is easier to displace by the photon.

Decoherence.—The requirement of observing the revival puts a bound on the acceptable environmental decoherence. To estimate the expected decoherence we model the mirror's environment by an (Ohmic) bath of harmonic oscillators. The effect of this can approximately be described by a decoherence rate $\gamma_D = \gamma_m k T_E M (\Delta x)^2/\hbar^2$ governing the decay of off-diagonal elements between different mirror positions [2]. Here γ_m is the damping rate for the mechanical oscillator, T_E is the temperature of the environment, which is constituted mainly by the internal degrees of freedom of the mirror and cantilever, and Δx is the separation of two coherent states that are originally in a superposition. This approximation is strictly valid only for times much longer than $2\pi/\omega_m$ and for Δx large compared to the width of the individual wave packets. Here we assume that the order of magnitude of the decoherence is well captured by γ_D. If the experiment achieves $\kappa^2 \gtrsim 1$, i.e., a separation by the size of a coherent state wave packet, $\Delta x \sim \sqrt{(\hbar/M\omega_m)}$, the condition $\gamma_D \lesssim \omega_m$ can be cast in the form

$$Q \gtrsim \frac{kT_E}{\hbar\omega_m}, \tag{4}$$

where $Q = \omega_m/\gamma_m$ is the quality factor of the mechanical oscillator. For $Q \gtrsim 10^5$, which has been achieved [15] for silicon cantilevers of approximately the right dimensions and frequency, this implies that the temperature of the environment has to be of the order of 2 mK, which is achievable with state-of-the-art dilution refrigerators.

Optical cooling.—Cooling the mirror's center of mass motion significantly eases the stability requirements for the proposed experiment. A method for optical cooling of a mirror via feedback was first proposed in Ref. [16]. By observing the phase of the output field of a cavity, its length can be measured with high precision. This can be used to implement a feedback mechanism that cools the center of mass motion of the mirror far below the temperature of its environment. A variation of the original scheme was experimentally implemented in Ref. [17], where a vibrational mode of a macroscopic mirror was cooled using a feedback force proportional to the natural damping force, but larger by a gain factor g. The size of g determines the achievable final temperature for a given T_E. For a tiny mirror, large gain values are realistic using the radiation pressure of a second laser beam to implement the feedback force. To analyze cooling to the quantum regime, one has to take into account the fact that measurement and feedback introduce noise, Ref. [18].

For our proposed experiment the constant component of the feedback laser has to balance the force from the measurement field, since otherwise the mirror would start to oscillate when the light is turned off. Adapting Ref. [19], the final energy of the cooled mirror is given by

$$E_c = \frac{\hbar\omega_m}{2}\frac{1}{2(1+g)}\left[\frac{4k_B T_E}{\hbar\omega_m} + 2\zeta + \frac{g^2}{\eta\zeta}\right], \tag{5}$$

where T_E is the temperature of the mirror's environment, $\zeta = (64\pi c P/M\gamma_m\omega_m\lambda\gamma_c^2 L^2)$, with P the light intensity incident on the measurement cavity and γ_c the cavity decay rate, and η the detection efficiency. The first term in Eq. (5) comes from the original thermal fluctuations, which are suppressed by the feedback. The second term is the back action noise from the measurement and feedback light. It differs from the formula of Ref. [19] by a factor of 2 to include the noise from the feedback laser. The third term is the noise due to imperfect measurement. Increasing the light intensity in the cavity improves the measurement precision, but also increases the back action noise.

The energy of the mirror can be made very close to its ground state energy choosing realistic parameter values; $E_c = \hbar\omega_m$ can be achieved with $g = 6 \times 10^5$, $T_E = 2$ mK, $P = 10^{-8}$ W, $\gamma_c = 3 \times 10^7$ s^{-1}, $\lambda = 800$ nm, $\eta = 0.8$, $\gamma_m = 0.03$ s^{-1}, and M, ω_m, L as before. The necessary feedback force for such a high value of g can be achieved with a feedback laser intensity modulation of $\Delta P_{fb} = 10^{-6}$ W. To balance the measurement field, the constant component of the feedback laser should be $\bar{P}_{fb} = 4 \times 10^{-6}$ W. The relatively large value of γ_c can be achieved in the cavity used in the superposition experiment by working at a wavelength away from where the mirrors are optimal.

Once the mirror has been cooled close to its ground state, which is reached in a time of order $1/(\gamma_m g)$ [20], the measurement and feedback laser fields should be turned off simultaneously. Then the experiment proceeds as described above. Reheating of the mirror happens at a time scale of $1/\gamma_m$ [20] and thus is not a problem for a high-Q oscillator. After every run of the experiment, the mirror has to be reset to its initial state by the optical cooling procedure.

Stability.—The distance between the large cavity end mirror and the equilibrium position of the small mirror has to be stable to of order $\lambda/20N = 0.6 \times 10^{-14}$ m over the whole measurement time, which is determined as follows. A single run of the experiment starts by sending a weak pulse into the interferometer, such that on average 0.1 photons go into either cavity. This probabilistically prepares a single-photon state as required to a good approximation. The two-photon contribution has to be kept low because it causes noise in the interferometer. Considering the required low value of ω_m and the fact that approximately 1% of the photons remain after a full period for the assumed loss, this implies a detection rate of approximately 10 photons per minute in the revival interval. Thus we demand stability to of order 10^{-14} m over a few minutes. Stability of order 10^{-13} m/min for an STM at 8 K was achieved with a rather simple suspension [21]. Gravitational wave observatories using interferometers also require very high stability in order to have a length sensitivity of 10^{-19} m over time scales of a ms or greater, for arm lengths of order 1 km [22]. If the mirror is in a thermal state, the revival peak is narrowed by a factor $\sqrt{\bar{n}}$ [6], leading to lower count rates in the revival interval and thus making the stability requirements stricter by the same factor, cf. Fig. 2.

The experiment also requires ultrahigh vacuum conditions in order to ensure that events where an atom hits the cantilever are sufficiently rare not to cause significant errors, which is at the level of about 5/s. Background gas particle densities of order $100/\mathrm{cm}^3$ have been achieved [23] and are sufficient for our purposes.

Outlook and conclusions.—In principle the proposed setup has the potential to test wave function reduction models, in particular, the one of Ref. [5]. We estimate that the ratio Q/T needs to be improved by about 6 orders of magnitude from the values discussed in this Letter ($Q = 10^5$ and $T = 2$ mK) to make the predicted wave function decoherence rate comparable to the environmental decoherence rate. However, temperatures as low as 60 μK have been achieved with adiabatic demagnetization [24], while Q is known to increase with decreasing temperature [15] and through annealing [25].

We have performed a detailed study of the experimental requirements for the creation and observation of quantum superposition states of a mirror consisting of 10^{14} atoms, approximately 9 orders of magnitude more massive than any superposition observed to date. Our analysis shows that, while very demanding, this goal appears to be within reach of current technology.

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LETTERS

Sub-kelvin optical cooling of a micromechanical resonator

Dustin Kleckner[1] & Dirk Bouwmeester[1]

Micromechanical resonators, when cooled down to near their ground state, can be used to explore quantum effects such as superposition and entanglement at a macroscopic scale[1–3]. Previously, it has been proposed to use electronic feedback to cool a high frequency (10 MHz) resonator to near its ground state[4]. In other work, a low frequency resonator was cooled from room temperature to 18 K by passive optical feedback[5]. Additionally, active optical feedback of atomic force microscope cantilevers has been used to modify their response characteristics[6], and cooling to approximately 2 K has been measured[7]. Here we demonstrate active optical feedback cooling to 135 ± 15 mK of a micromechanical resonator integrated with a high-quality optical resonator. Additionally, we show that the scheme should be applicable at cryogenic base temperatures, allowing cooling to near the ground state that is required for quantum experiments—near 100 nK for a kHz oscillator.

Using a laser tuned to the resonance fringe of a high finesse optical cavity, it is possible to observe very small fluctuations in the length of the cavity due to brownian motion of one or both of the end mirrors. We have developed an optical cavity with one rigid large mirror, 6 mm in diameter and with a 25 mm radius of curvature, and one tiny plane mirror, 30 μm in diameter, attached to a commercial atomic force microscope cantilever of dimensions $450 \times 50 \times 2\,\mu m$ with a fundamental resonance of 12.5 kHz (Fig. 1b). An optical finesse of 2,100 and a mechanical quality factor of 137,000 have been achieved with the system[8]. The motion of the tiny mirror/cantilever is monitored by measuring the transmission of the cavity at a frequency on the side of an optical resonance peak. To do this, we use about 1 mW from a 780 nm tunable diode laser which is locked to the resonance fringe using the integrated signal from a photo-multiplier tube which monitors the light transmitted through the cavity (Fig. 1a). The time derivative of this signal is proportional to the velocity of the cantilever tip and is used to modulate the amplitude of a second, 980 nm, diode laser focused on the cantilever less than 100 μm away from the tiny mirror. The radiation pressure exerted by this feedback laser counteracts the motion of the mirror and effectively provides cooling of the fundamental mode.

The effective feedback gain can be varied over several orders of magnitude by sending the feedback laser through a variable neutral density filter. The average power in the feedback beam when it reaches the cantilever is of the order of 1 mW at the highest gain settings and proportionally lower otherwise. The mean modulation depth of the feedback beam varies from nearly 100% to less than 5% as gain is increased. The vibration spectrum of the cantilever as a function of gain is shown in Fig. 2. The r.m.s. thermal amplitude of the cantilever without feedback is 1.2 ± 0.1 Å. From this value, one can calculate that the spring constant of the cantilever is $0.15 \pm 0.01\,N\,m^{-1}$, in agreement with the manufacturer-specified range, and the effective mass of the cantilever fundamental mode is $(2.4 \pm 0.2) \times 10^{-11}$ kg.

To determine the effective gain of the feedback loop and the temperature of the fundamental mode, we fit a lorentzian plus a constant background to the vibration spectrum of the cantilever for each value of feedback gain. The temperature is determined from the area under the lorentzian without the background, while the gain is determined by the width of the resonance. The linewidth provides a good measure of gain because it is directly determined by the damping rate whereas the cantilever amplitude may be affected by other sources of noise in the feedback loop. Cooling is observed over more than three orders of magnitude. The lowest temperature we are able to measure is 135 ± 15 mK, or a cantilever r.m.s. amplitude of 0.023 ± 0.002 Å, with a gain (the ratio of feedback to mechanical damping) of $g = 2{,}490 \pm 90$ (Fig. 2b). The lowest trace in Fig. 2b, indicating an even lower temperature, cannot be reliably fitted owing to the laser noise floor. Since the optical finesse is not the current limiting factor, we operate the opto-mechanical system at a finesse of only 200,

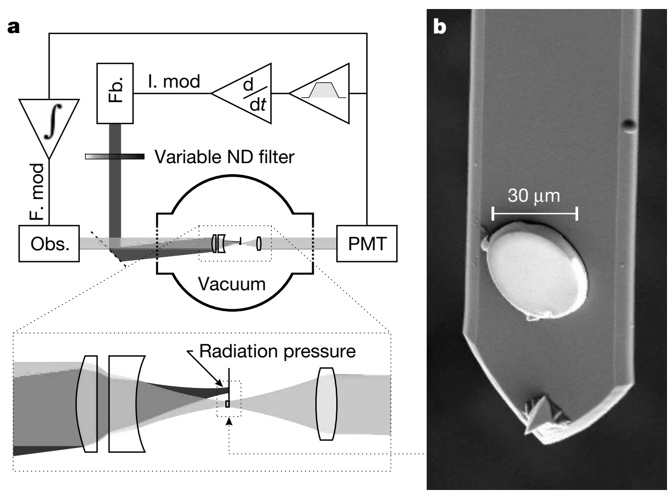

Figure 1 | The experimental system. a, Diagram of the feedback mechanism: a 780 nm observation laser (Obs.) is frequency locked to the optical cavity (shown magnified at bottom) with an integrating circuit (via the laser frequency modulation input, f. mod), using the signal from a photomultiplier tube (PMT). This signal is also sent through a 1.25 kHz bandpass filter at 12.5 kHz and a derivative circuit (d/dt) to provide an intensity-modulating signal (I. mod.) for the 980 nm feedback laser (Fb.). The feedback laser is attenuated with a variable neutral density (ND) filter to adjust the gain of the feedback. The feedback force is exerted on the cantilever via this laser's radiation pressure. **b**, Scanning electron microscope image of the tip of the cantilever with attached mirror.

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produced by slight cavity misalignment, which makes the system less sensitive to transient vibrations.

The amplitude of the mirror motion can be calculated in the presence of feedback by assuming that the Langevin force—the effective thermal force that maintains brownian motion—remains constant while the mechanical susceptibility of the mirror is reduced by the dissipation due to the radiation feedback pressure. It suffices to consider only the fundamental mode of the mirror motion, represented by a damped harmonic oscillator. In this approximation, the power spectrum of the mirror's motion in the presence of feedback becomes[9]:

$$S_x^{\mathrm{fb}}[\Omega] = \frac{2\Gamma_0 k_B T_0}{M} \frac{1}{(\omega^2 - \Omega^2)^2 + (1+g)^2 \Gamma_0^2 \Omega^2} \quad (1)$$

where Ω is the observation frequency, ω is the resonator frequency, Γ_0 is the mechanical damping factor, M is the effective mass of the resonator mode, k_B is Boltzmann's constant, T_0 is the bulk temperature of the resonator and g is the gain. $g = 0$ corresponds to the vibration spectrum in the absence of feedback. The motion of the oscillator in the presence of feedback is the same as that of an oscillator with lower temperature and a higher damping constant:

$$T_{\mathrm{fb}} = (1+g)^{-1} T_0 \quad (2)$$

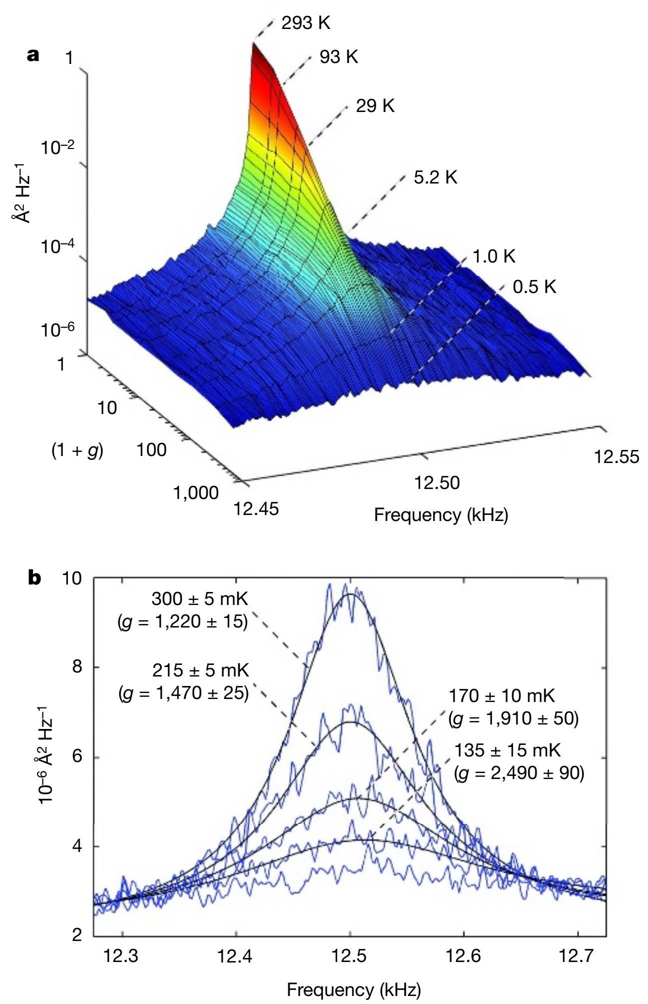

Figure 2 | Single-sided thermal vibration spectrum of the cantilever as it is cooled. g is the dimensionless gain factor, which is the ratio of feedback to mechanical damping. **a**, Spectrum at low to moderate gains. **b**, Spectrum near the background noise level for large gains. The blue curves correspond to experimental data, and the black curves to fits of a gaussian function plus a background. The lowest trace cannot be reliably fitted.

$$\Gamma_{\mathrm{fb}} = (1+g)\Gamma_0 \quad (3)$$

The optical feedback scheme, when analysed in terms of noiseless classical light fields, can be seen as a virtual viscous force, which unlike a real viscous force creates dissipation without introducing fluctuations. As discussed below, the cooling temperature as demonstrated here is limited by laser frequency fluctuations. Ultimately, optical cooling should be limited by the balance of residual heating and quantum noise in the observation and feedback laser signals.

For a signal-to-noise ratio of one in spectral density at the peak of the mechanical resonance, the temperature of the cantilever would be (as can be derived from equation (1)):

$$T_{\min} \cong \sqrt{\frac{T_0 M \omega^3 S_{\mathrm{noise}}}{2 k_B Q}} \quad (4)$$

where S_{noise} is the equivalent position noise in the interferometer measurement and $Q = \omega/\Gamma_0$ is the mechanical quality factor. For higher values of gain, the feedback signal is mostly noise and lower temperatures can not be conclusively demonstrated. For our experiment, the equivalent noise level is $\sqrt{S_{\mathrm{noise}}} \approx 10^{-3}\,\text{Å}\,\mathrm{Hz}^{-1/2}$. This corresponds to the expected noise due to the frequency fluctuations of a free running tunable laser diode, which are of order $10^3\,\mathrm{Hz}\,\mathrm{Hz}^{-1/2}$ at the resonance frequency of 12.5 kHz (ref. 10). With the system in vacuum at pressures of 10^{-6} mbar, so as to maximize the mechanical quality factor of the cantilever, this noise level corresponds to a minimum temperature of the order of 100 mK, in good agreement with the experimental data.

An alternative approach to study the cooling is to analyse the temporal response of the system by gating the signal to the feedback laser. The characteristic time constant for the system to reach equilibrium after the cooling is turned on is given by:

$$\tau_{\mathrm{fb}} = \Gamma_{\mathrm{fb}}^{-1} = (1+g)^{-1} \Gamma_0^{-1} \quad (5)$$

To observe this behaviour, we monitor the cantilever over many 10 s periods during each of which the cooling is on for 3 s. Data for cooling to 1.8 ± 0.2, 4.0 ± 0.2 and 6.4 ± 0.1 K and returning to thermal equilibrium are shown in Fig. 3. The cooling times are measured to be 9.0 ± 0.5, 19 ± 1 and 27 ± 1 ms, respectively. The reheating time is found to be indistinguishable for all three gains with an average of $\tau_0 = 1.30 \pm 0.05$ s. This is in agreement with the linewidth of the cantilever measured without feedback, $\Gamma_0 = 680 \pm 50$ mHz. In

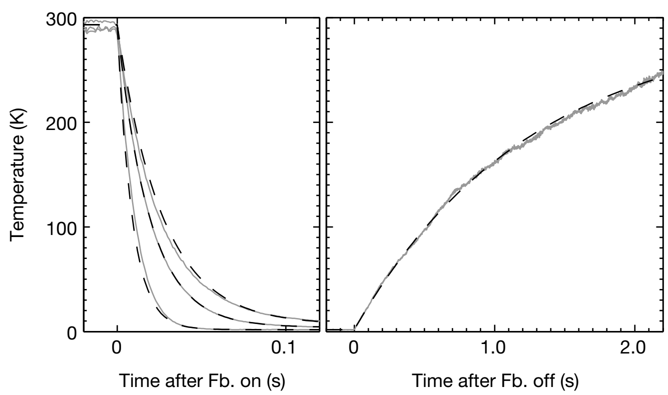

Figure 3 | Temporal response of the cantilever to cooling pulses. The temperature is determined by calculating the total vibrational amplitude of the cantilever between 12 and 13 kHz in 1 ms bins and subtracting the background. Each data set is the average of 1,000 samples. The three sets in the left panel correspond to cooling to 6.4, 4.0 and 1.8 K (solid lines, top to bottom). Heating is shown (right panel) for only one data set (1.8 K), as all three are nearly coincident. The dashed lines are fits to exponential decays, used to determine the cooled temperature and the cooling and reheating times. Fb. refers to the feedback system.

accordance with theory, the ratio of the reheating to the cooling times, τ_0/τ_{fb}, and the corresponding ratio of the spectral linewidths from the earlier measurements, Γ_{fb}/Γ_0, are found to be the same as the cooling factor, T_0/T_{fb}, within statistical uncertainties.

In experiments where optical feedback is used on cantilevers with non-uniform composition, radiation pressure is typically overwhelmed by the photothermal force, which is an effective force due to thermally induced bending[5,6]. Although this is not the case for single-crystal silicon cantilevers, the addition of a tiny mirror on the tip of our cantilever should produce a weak photothermal force. This force can be distinguished from radiation pressure by its dependence on the intensity modulation frequency of the feedback laser. Whereas radiation pressure is independent of modulation frequency, the photothermal force is not, because it has a characteristic response time, τ, related to the thermal relaxation time of the cantilever. A simple model for the frequency dependence of the photothermal force, $F_{pt}(\Omega)$, gives:

$$F_{pt}(\Omega) \cong \int_0^{\infty} \frac{F_{pt}(0)}{\tau} e^{-\frac{t}{\tau}} e^{-i\Omega t} dt = \frac{F_{pt}(0)}{1+i\Omega\tau} \quad (6)$$

where $e^{-i\Omega t}$ corresponds to the input power modulation, and $e^{-\frac{t}{\tau}}$ is due to the thermal relaxation. This is consistent with the frequency dependence of the photothermal force as described in previous work[6]. To test for the presence of photothermal force in our resonator, the feedback laser was modulated at a range of frequencies from 100 Hz to 20 kHz and the mechanical response of the cantilever was measured as before (Fig. 4). The power in the feedback laser reflected from the cantilever was determined to have a mean of 2.7 ± 0.5 mW and a modulation amplitude of 1.0 ± 0.2 mW, independent of the modulation frequency. This results in a radiation pressure force of $F_{rad} = 2P_{mod}/c = 6.7 \pm 1.3$ pN(where P_{mod} is the amplitude of the power modulation and c is the speed of light) at the modulation frequency.

If the driving frequency is sufficiently far from the cantilever resonance, the mechanical damping constant can be ignored and the amplitude of the cantilever's motion should be of the form:

$$A(\Omega) = \left| \frac{\frac{A_{pt}}{1+i\Omega\tau} + A_{rad}}{1 - \left(\Omega/\omega\right)^2} \right| \quad (7)$$

where Ω is the driving frequency, ω is the resonance frequency, τ is the photothermal characteristic time, and A_{rad} and A_{pt} are the magnitudes of the motion due to the radiation pressure and photothermal force alone, at zero frequency. The term in the denominator is due to mechanical amplification by the cantilever resonance. This equation fits well to the measured response (Fig. 4), resulting in $A_{rad} = 0.470 \pm 0.005$ Å, $A_{pt} = -6.3 \pm 0.2$ Å and $\tau = 30 \pm 2$ ms. At frequencies greater than 5 kHz, radiation pressure is observed to be the dominant force mechanism, whereas the photothermal force is relevant only at lower frequencies.

Assuming the constant force background described by A_{rad} is entirely due to radiation pressure, one can calculate the spring constant of the cantilever at the position where the feedback laser is focused to be $k = F_{rad}/A_{rad} = 0.14 \pm 0.03$ N m^{-1}, in agreement with the value for the spring constant obtained earlier. Near the fundamental resonance of the cantilever, the radiation pressure is calculated to be almost 5 times larger than the photothermal force. Additionally, the two forces should be nearly 90° out of phase at this frequency, given that the time constant of the photothermal force is found to be 30 ± 2 ms. Thus the radiation pressure is responsible for almost all of the demonstrated feedback cooling; in the absence of photothermal force, the total feedback force would be reduced by less than 3%.

When optical cooling is active, the cantilever's motion is strongly damped, making it undesirable for many types of measurements. In some cases this problem can be overcome with a stroboscopic cooling scheme, where measurements are only made in the periods when the cooling is off. In addition to being of direct importance for the aforementioned massive superposition experiment, this scheme has already been theoretically shown to be useful for high sensitivity measurements of position and weak impulse forces[11]. Because the cooling is faster than the heating by a factor $(1+g)$, a low temperature can be maintained even when the cooling is off the majority of the time. However, maintaining low temperatures requires that the measurement window be short; if it is, for example, one oscillation period long, the temperature of the oscillator will have increased by $\Delta T \approx 2\pi T_0/Q$ by the end of each measurement window, meaning that cooling past this point results in marginal improvement.

We now evaluate the potential for reaching even lower temperatures for the purpose of studying quantum effects in similar systems. Reference 3 proposes an experiment: putting a mechanical oscillator in a quantum superposition of vibrating and not-vibrating by interaction with the light pressure of a single photon in an optical cavity of which one end mirror is attached to the oscillator. Appropriate for such a scheme would be a 250-μm-long silicon cantilever with a 20-μm-diameter dielectric mirror on the tip and a resonance frequency of 1 kHz. Because of the constraints of environmentally induced decoherence[12,13], the bulk temperature must be less than $T_{EID} = Q\hbar\omega/k_B = 8$ mK for the cantilever to remain coherent over one period, given $Q \approx 150{,}000$. This temperature is achievable by conventional means; nuclear adiabatic demagnetization of $PrNi_5$ (ref. 14) could be employed as the final, vibration-free, cooling stage, as it is able to be started from temperatures previously demonstrated

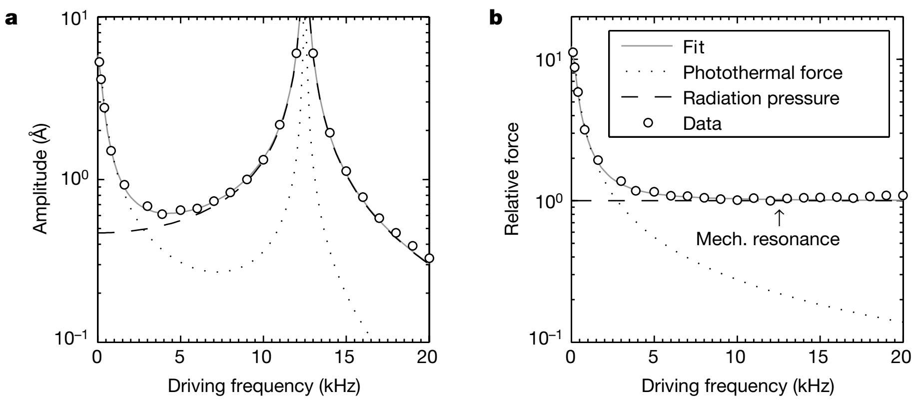

Figure 4 | Response of the cantilever to an external intensity-modulated laser. a, The amplitude of the cantilever's motion at the driving frequency. **b,** The force on the cantilever, calculated by dividing the amplitude by the mechanical amplification of the cantilever. In both graphs the magnitude of the contributions (ignoring phase differences) of the photothermal force and radiation pressure are shown as dashed and dotted lines, respectively. The slight deviation of the fit from the data at higher frequencies is due to higher-order flexural modes.

for vibration-isolated cold stages ($\sim$100 mK)[15,16]. The observation period for a massive superposition experiment is one oscillation long, thus the maximum useful cooling factor is $Q/2\pi \approx 25{,}000$ as discussed above. This corresponds with a temperature of 300 nK or a mean oscillator quantum number of only 2π.

It has been shown theoretically that optical feedback still works in the quantum regime, allowing cooling to the ground state[17]. Experimentally, cooling to the quantum regime requires the capability to accurately monitor the position of the cantilever without introducing significant heating. With an optical finesse $F = 5 \times 10^5$, which should be technologically achievable[8], an observation beam power of 1 aW, or about 5,000 photons per second, is enough to reduce shot noise to the appropriate level. Assuming the thermal conductivity of the cantilever is reduced to the one-dimensional quantum limit, the cantilever's thermal resistivity will be roughly 30 mK aW^{-1} (ref. 18). The heating from the feedback laser can be reduced by use of a sufficiently long wavelength laser so that absorption is negligible; this is not possible for the readout beam, which must be resonant with the optical cavity. Thus as long as the observation laser has relatively low absorption in the cantilever/mirror, it should not significantly affect the bulk temperature. This implies that cooling a kHz oscillator to near its ground state should be possible, drastically simplifying the experimental requirements to observe quantum phenomena in this system.

We have demonstrated active laser feedback cooling of a micromechanical oscillator, using only radiation pressure, from room temperature to 135 mK. Furthermore, we have shown that this cooling method could be used in addition to traditional cryogenics to reach much lower temperatures, even near the ground state of a kHz oscillator. This in turn would significantly aid the realization of proposals to create and investigate massive quantum superpositions.

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Cavity Optomechanics: Back-Action at the Mesoscale
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Cavity Optomechanics: Back-Action at the Mesoscale

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The coupling of optical and mechanical degrees of freedom is the underlying principle of many techniques to measure mechanical displacement, from macroscale gravitational wave detectors to microscale cantilevers used in scanning probe microscopy. Recent experiments have reached a regime where the back-action of photons caused by radiation pressure can influence the optomechanical dynamics, giving rise to a host of long-anticipated phenomena. Here we review these developments and discuss the opportunities for innovative technology as well as for fundamental science.

The reflection of a photon entails momentum transfer, generally referred to as "radiation pressure," with the resulting force called the scattering force. Besides this scattering force, the spatial variation of an intensity distribution can give rise to a gradient or dipole force. Interest in radiation pressure was first generated by the trapping of dielectric particles using laser radiation (*1*). This technique is widely adapted today in the biological and biophysical sciences and is known as the "optical tweezer." In atomic physics, the ability to cool atoms with the use of radiation pressure (*2*, *3*) has enabled many advances (*4*), including the realization of exotic quantum states such as Bose-Einstein condensates.

Radiation pressure can also have an effect on macroscale mechanical masses (such as on an optical interferometer's mirror) and has been considered theoretically for decades (*5*, *6*). The mutual coupling of optical and mechanical degrees of freedom in an optical resonator (or optical cavity) has been explored in laser-based gravitational wave interferometers, in which radiation pressure imposes limits on continuous position detection. Beyond setting detection limits, radiation pressure can also influence the dynamics of a harmonically bound mirror. A discernable effect on mirror motion was first demonstrated in the optical bistability resulting from the static elongation of cavity length caused by radiation pressure (*7*), and later, in work demonstrating the optical spring effect (a radiation-pressure–induced change in stiffness of the "mirror spring") (*8*). These phenomena, however, do not rely on the cavity delay; rather, each results from an adiabatic response of the cavity field to mechanical motion. Phenomena of a purely dynamical nature were predicted (*5*, *9*) to arise when the decay time of the photons inside the cavity is comparable to or longer than the mechanical oscillator period. Creating such delays through an electro-optic hybrid system was later proposed and demonstrated to induce radiation-pressure "feedback cooling" of a cavity mirror (*10*, *11*), also known as cold damping. Whereas in subsequent attempts dynamic radiation-pressure phenomena were masked by thermal effects (*12*), recent advances in micro- and nanofabrication made it possible to access the regime where the effects of cavity-enhanced radiation pressure alone dominate the mechanical dynamics. Demonstrations of mechanical amplification (*13*, *14*) and cooling (*14*–*16*) via dynamical back-action signal that a paradigm shift (*17*) in the ability to manipulate mechanical degrees of freedom is now under way, which has long been anticipated (*18*, *19*). Central to all current work is the role of back-action in setting dynamical control and performance limits. This review is intended to provide context for these recent accomplishments and also to present an overview of possible and anticipated future research directions.

Dynamical Back-Action Versus Quantum Back-Action

Photons at optical frequencies are uniquely suited to measure mechanical displacement for several reasons. First, because of the high energy of optical photons (~1 eV), thermal occupation is negligible at room temperature. Moreover, present-day laser sources are available that offer noise performance that is limited only by quantum noise. To measure displacement, a commonly used experimental apparatus is a Fabry-Perot interferometer,

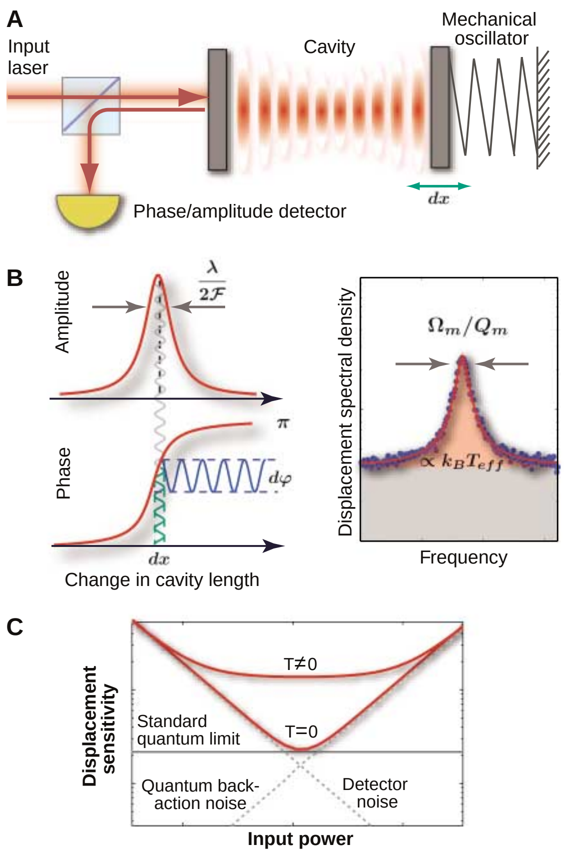

Fig. 1. (**A**) Schematic of the cavity optomechanical interaction of a cavity field (red) and a moveable mirror. (**B**) Transduction mechanism for the laser resonantly probing the cavity. The mechanical motion (green) causes the reflected field to be phase modulated around its steady-state value. This occurs because the mirror motion changes the total cavity length and thereby changes the resonance frequency of the cavity by $\omega_0 \frac{dx}{L}$, where L is the separation between the two mirrors and dx is the mirror displacement. Owing to the high Finesse of the cavity ($\mathcal{F}_\pi$, which describes the number of reflections a photon undergoes on average before escaping the cavity), the conversion of mechanical amplitude to the phase of the field is enhanced (i.e., $d\varphi \approx \frac{F}{\lambda} \cdot dx$, where $d\varphi$ is the change in the phase of the reflected laser field and λ is the incident wavelength of the laser), allowing minute mirror displacements to be detected. The reflected amplitude is left unchanged. (Right) Fourier analysis of the reflected phase reveals the mechanical spectrum of the mirror motion. Mechanical resonance frequency (Ω_m), quality factor (Q_m), and temperature (T_{eff}) can be determined using this spectrum. (**C**) Sensitivity of the interferometer measurement process for the case of a zero-temperature mechanical oscillator mirror and for finite temperature T. For low-input laser power, detector noise due to the quantum shot noise of the laser field dominates, whereas at higher laser power the quantum fluctuations of the light field cause the mirror to undergo random fluctuations (quantum back-action). At the optimum power, the two sources of fluctuation contribute equally to the measurement imprecision, constituting the SQL. At finite temperature, the mechanical zero-point motion is masked by the presence of thermal noise.

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whose purpose is to determine differential changes in distance between the two end mirrors (Fig. 1A). To account for the mirror suspension or the internal mechanical modes of a mirror, it is assumed that the end mirror is free to oscillate. This harmonic confinement can be either intentional or intrinsic, as we will discuss later. The high-reflectivity end mirrors enhance the number of roundtrips photons undergo (by a factor $\mathcal{F}/\pi$, where $\mathcal{F}$ is the cavity Finesse) and enable very sensitive measurement of the end mirror position (Fig. 1B). For a laser resonant with the cavity, small changes in cavity length shift the cavity resonance frequency and, enhanced by the cavity Finesse, imprint large changes in the reflected phase of the laser field. To date, the best displacement sensitivities attained with optical interferometers [such as those at the Laser Interferometer Gravitational Wave Observatory (LIGO) or Fabry-Perot cavities (*20*)] are already exceeding 10^{-19} m/$\sqrt{\text{Hz}}$, which implies that a displacement equivalent to 1/1000 of the radius of a proton can be measured in 1 s.

This extremely high sensitivity, however, also requires that the disturbances of the measurement process itself must be taken into account. The ultimate sensitivity of an interferometer depends on the back-action that photons exert onto the mechanically compliant mirror, caused by radiation pressure. In terms of mirror-displacement measurement, two fundamental sources of imprecision exist (Fig. 1C). First, there is the detector noise that, for an ideal laser source (emitting a coherent state) and an ideal detector, is given by the random arrival of photons at the detector; i.e., shot noise. The detector signal-to-noise ratio increases with laser power, thereby improving the measurement precision. Increasing power, however, comes at the expense of increased intracavity optical power, causing a back-action onto the mirror. This leads to a second source of imprecision: The resulting random momentum kicks of reflected photons create a mirror-displacement noise. This random force causes the mechanical oscillator to be driven and thus effectively heated. Although this noise can also contain a contribution due to classical sources of noise (excess phase or amplitude noise), it is ultimately, under ideal circumstances, bound by the quantum nature of light and is termed quantum back-action (*21*, *22*). Taking into account both contributions, the optimum sensitivity of an interferometer is achieved at the standard quantum limit (SQL). At the SQL, detector noise and quantum back-action noise contribute each a position uncertainty equal to half of the zero-point motion of the mirror, where the latter is given by $x_0 = \sqrt{\hbar/2m\Omega_m}$ [$\hbar$ is Planck's constant divided by 2π, m is the effective mass (*23*) of the mirror, and Ω_m is the mirror's harmonic frequency]. Much research in the past decade has also focused on ways of circumventing this limit. For example, the use of squeezed light (*24*) can enable surpassing this limit. So far, however, experiments with mechanical mirrors have not observed the radiation-pressure quantum back-action because it is masked by the random, thermal motion of the mirror (Fig. 1C). Fluctuations of the radiation-pressure force have been observed in the field of atomic laser cooling (*25*), where they are responsible for a temperature limit (the Doppler limit).

The optical cavity mode not only measures the position of the mechanical mode, but the dynamics of these two modes can also be mutually coupled. This coupling arises when the mechanical motion changes the intracavity field amplitude, which thereby changes the radiation-pressure force experienced by the mirror. For small displacements, this occurs when the laser is detuned with respect to the cavity resonance (Fig. 2A). This mutual coupling of optical and mechanical degrees of freedom can produce an effect called dynamic back-action that arises from the finite cavity delay. This delay leads to a component of the radiation-pressure force that is in quadrature (out of phase) with respect to the mechanical motion. The component is substantial when the cavity photon lifetime is comparable to, or larger than, the mechanical oscillator period and creates an effective mechanical damping of electromagnetic origin. This is the essence of dynamic back-action (*5*), which, like quantum back-action, modifies the motion of the object being measured (the mirror). Unlike quantum back-action, which effectively sets a measurement precision (by causing the mirror to be subjected to a stochastic force resulting from quantum fluctuations of the field), the effect of dynamic back-action is to modify the dynamical behavior of the mirror in a predictable manner. Two consequences of this form of back-action in the context of gravitational wave detection have been identified. With a laser field blue-detuned relative to the optical cavity mode, the mirror motion can be destabilized (*5*) as a result of mechanical amplification (*13*). Similar to the operation of a laser, the onset of this instability occurs when the mechanical gain equals the mechanical loss rate and could thus create an effective limit to boosting detection sen-

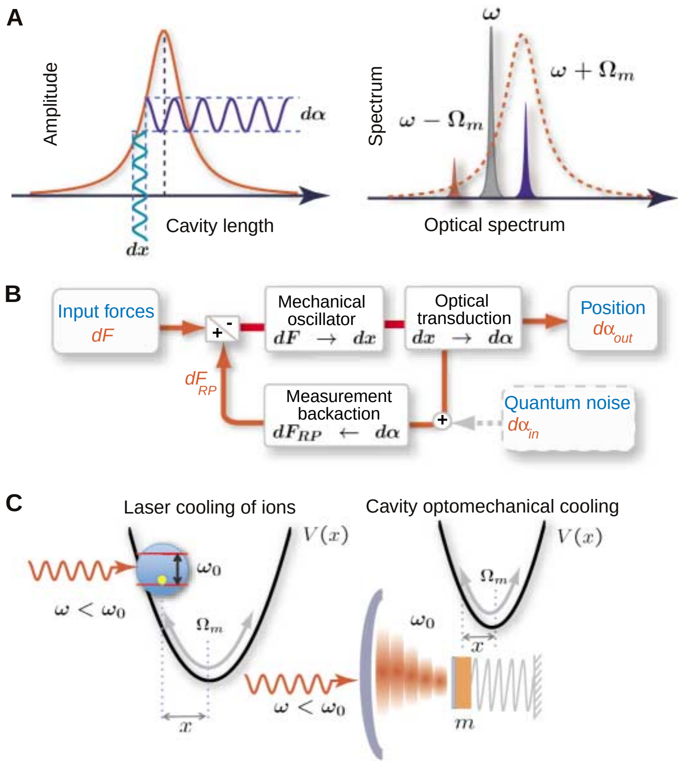

Fig. 2. (**A**) Dynamic back-action results from the coupling of the mechanical motion to the fluctuations of the intracavity field amplitude ($d\alpha$), which occurs when the cavity is excited in a detuned manner. In the frequency domain, the amplitude modulation at Ω_m can be interpreted as sidebands around the optical laser frequency (shown at right). The sideband amplitudes are asymmetric because of the density of states of the cavity. This photon imbalance results in work on the mechanical oscillator (either amplification or cooling), as is further detailed in Fig. 4. (**B**) Basic elements of a feedback loop describing the measurement process and its back-action on the mirror. The mechanical oscillator is subject to a force dF (e.g., because of the thermal force or an externally applied signal force) that induces a mechanical response (dx). The latter causes a change in the optical field (either in amplitude $d\alpha$ or in phase, depending on the detuning), allowing measurement of mechanical position. This transduction is not instantaneous on account of the finite cavity lifetime. For a detuned laser, the amplitude change caused by this measurement process feeds back to the mechanical oscillator through the radiation-pressure force, closing the feedback loop. The sign of the feedback depends on the cavity detuning and can produce either damping (red-detuned pump) or amplification (blue-detuned pump). In a quantum description, this feedback branch is not noiseless but is subjected to quantum noise of the optical field ($d\alpha_{in}$), which yields a random force due to the quantum fluctuations of the field (i.e., the quantum back-action). Although dynamic back-action can be prevented by probing the cavity on resonance (causing $d\alpha = 0$ and thereby preventing feedback), the quantum back-action nevertheless feeds into the mechanical oscillators' input (and thereby reinforces the SQL. $d\alpha_{out}$ are the amplitude fluctuations of the reflected laser field; dF_{RP} are the fluctuations in the radiation pressure force. (**C**) Analogy of dynamical back-action cooling to the laser cooling of harmonically bound ions. In both the case of a harmonically trapped ion and a harmonically oscillating end mirror of a cavity, a dissipative force arises because of the Doppler effect. $V(x)$ denotes the trapping potential of the mirror and ion.

sitivity by increasing optical power in interferometers. On the other hand, a red-detuned pump wave can create a radiation component of mechanical damping that leads to cooling of the mechanical mode; i.e., a reduction of the mechanical mode's Brownian motion (*9*, *26*).

One description of this process is given in Fig. 2B, wherein a feedback loop that is inherent to the cavity optomechanical system is described. The elements of this loop include the mechanical and optical oscillators coupled through two distinct paths. Along the upper path, a force acting on the mechanical oscillator (for instance, the thermal Langevin force or a signal force) causes a mechanical displacement, which (for a detuned laser) changes the cavity field due to the optomechanical coupling (the interferometric measurement process). However, the amplitude fluctuations, which contain information on the mirror position, are also coupled back to the mechanical oscillator via radiation pressure (lower path), resulting in a back-action. A blue-detuned pump wave sets up positive feedback (the instability), whereas red detuning introduces negative feedback. Resonant optical probing (where the excitation frequency equals the cavity resonance frequency, $\omega = \omega_0$) interrupts the feedback loop because changes in position only change the phase, not the amplitude, of the field. As described below, this feedback circuit also clarifies the relation between "feedback cooling" and cooling by dynamic back-action.

Experimental Systems

Systems that exhibit radiation-pressure dynamic back-action must address a range of design considerations, including physical size as well as dissipation. Dynamic back-action relies on optical retardation; i.e., is most prominent for photon lifetimes comparable to or exceeding the mechanical oscillation period. Very low optical dissipation also means that photons are recycled many times, thereby enhancing the weak photon pressure on the mirror. On the other hand, the mechanical dissipation rate governs the rate of heating of the mechanical mirror mode by the environment, limiting the effectiveness of optomechanical cooling. It also sets the required amplification level necessary to induce regenerative oscillations. These considerations illustrate the importance of high optical Finesse and mechanical Q in system design.

It is only in the past 3 years that a series of innovative geometries (shown in Fig. 3) has reached a regime where the observation of radiation-pressure dynamic back-action could be observed. These advances have relied on the availability and improvements in high-Finesse mirror coatings (as used in gravity wave detectors) and also on micro- and nanofabrication techniques [which are the underlying enabling technology for nano- and microelectromechanical systems (*27*)]. A commonly used hybrid system consists of a conventional-input mirror made with a high-reflectivity coating and an end mirror whose dimensions are mesoscopic and which is harmonically suspended. This end mirror has been realized in multiple ways, such as from an etched, high-reflectivity mirror substrate (*14*, *15*), a miniaturized and harmonically suspended gram-scale mirror (*28*), or an atomic force cantilever on which a high-reflectivity and micron-sized mirror coating has been transferred (*29*). A natural optomechanical coupling can occur in optical microcavities, such as microtoroidal cavities (*13*) or microspheres, which contain co-existing high-Q, optical whispering gallery modes, and radio-frequency mechanical modes. This coupling can also be optimized for high optical and mechanical Q (*30*). In the case of hybrid systems, yet another approach has separated optical and mechanical degrees of freedom by using a miniature high-Finesse optical cavity and a separate nanometric membrane (*31*). Whereas the aforementioned embodiments have been in the optical domain, devices in the micro- and radiowave domain have also been fabricated (*22*, *32*), such as a nanomechanical resonator coupled to a superconducting microwave resonator (*33*).

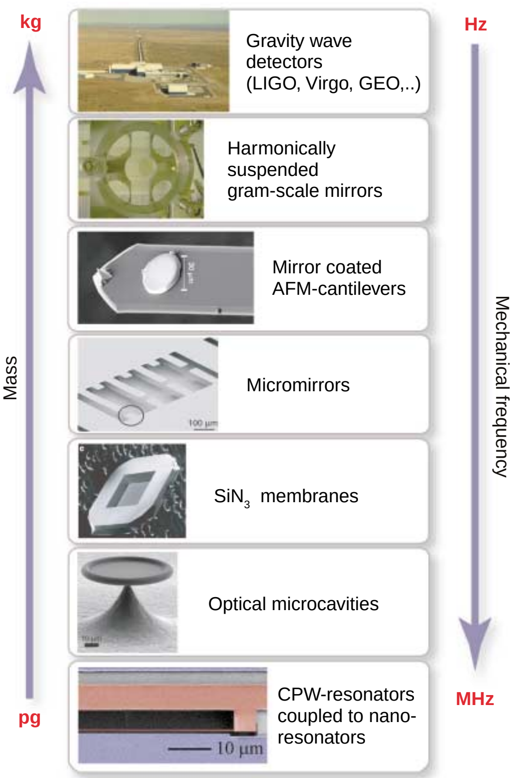

Fig. 3. Experimental cavity optomechanical systems. (**Top** to **Bottom**) Gravitational wave detectors [photo credit LIGO Laboratory], harmonically suspended gram-scale mirrors (*28*), coated atomic force microscopy cantilevers (*29*), coated micromirrors (*14*, *15*), SiN_3 membranes dispersively coupled to an optical cavity (*31*), optical microcavities (*13*, *16*), and superconducting microwave resonators coupled to a nanomechanical beam (*33*). The masses range from kilograms to picograms, whereas frequencies range from tens of megahertz down to the hertz level. CPW, coplanar waveguide.

Many more structures exist that should also realize an optomechanical interaction in an efficient manner. In particular, nanophotonic devices such as photonic crystal membrane cavities or silicon ring resonators might be ideal candidates owing to their small mode volume, high-Finesse, and finite rigidity. Owing to their small length scale, these devices exhibit fundamental flexural frequencies well into the gigahertz regime, but their mechanical quality factors have so far not been studied, nor has optomechanical coupling been observed. As described in the next section, such high frequencies are interesting in the context of regenerative oscillation and ground state cooling.

Cooling and Amplification Using Dynamical Back-Action

The cooling of atoms or ions using radiation pressure has received substantial attention and has been a successful tool in atomic and molecular physics. Dynamical back-action allows laser cooling of mechanical oscillators in a similar manner. The resemblance between atomic laser cooling and the cooling of a mechanical oscillator coupled to an optical (or electronic) resonator is a rigorous one (*34*). In both cases, the motion (of the ion, atom, or mirror) induces a change in the resonance frequency, thereby coupling the motion to the optical (or cavity) resonance (Fig. 2C). Indeed, early work has exploited this coupling to sense the atomic trajectories of single atoms in Fabry-Perot cavities (*35*, *36*) and, more recently, in the context of collective atomic motion (*37*, *38*). This coupling is not only restricted to atoms or cavities but also has been predicted for a variety of other systems. For example, the cooling of a mechanical oscillator can be achieved using coupling to a quantum dot (*39*), a trapped ion (*40*), a Cooper pair box (*41*), an LC circuit (*5*, *32*), or a microwave strip-line cavity (*33*). Although the feedback loop of Fig. 2B explains how damping and instability can be introduced into the cavity optomechanical system, the origins of cooling and mechanical amplification are better understood with the use of a motional sideband approach, as described in Fig. 4 (*13*).

Cooling has been first demonstrated for micromechanical oscillators coupled to optical cavities (*14*–*16*) and, using an electromechanical analog, for a Cooper pair box coupled to a nanomechanical beam (*41*). Because the mechanical modes in experiments are high Q (and are thus very well isolated from the reservoir), they are easily resolved in the spectra of detected probe light reflected from the optical cavity (Fig. 1B). Furthermore, their effective temperature can be inferred from the thermal energy $k_B T$ (where k_B is the Boltzman constant), which is directly proportional to the area of detected mechanical spectral peak (Fig. 1B). In the first back-action cooling experiments, a temperature of ~10 K was achieved for a single mechanical mode. The bath and all other modes in these experiments were at room temperature, owing to the highly targeted nature of cooling (Fig. 4). Since the completion of this work, cooling of a wide variety of experimental embodiments ranging from nanomembranes (*31*) and gram-scale mirrors (*28*) to the modes of kilogram-scale gravitational bar detectors (such as Auregа) has been demonstrated. At this stage, temperatures are rapidly approaching a regime of low phonon number, where quantum effects of the mechanical oscillator become important. To this end, cooling with the use of a combination of conventional cryogenic technology with dynamical back-action cooling is being investigated. Technical hurdles include collateral reheating of the mechanical mode, exacerbated by the very high mechanical Q, which leads to relatively long equilibration times.

Quantum back-action sets a fundamental limit of radiation-pressure cooling (*34*, *42*) that is equivalent to the Doppler temperature in atomic laser cooling (*25*). It may also be viewed as a consequence of the Heisenberg uncertainty relation in that a photon decaying from the resonator has an uncertainty in energy given by $\Delta E = \hbar\kappa$ (where κ is the cavity decay rate), implying that the mechanical oscillator cannot be cooled to a temperature lower than this limit. It has been theoretically shown (*34*, *42*) that ground state cooling is nevertheless possible in the resolved sideband regime (also called the good-cavity limit), in analogy to atomic laser cooling, where this technique has led to ground state cooling of ions (*43*). This regime is characterized by mechanical sidebands that fall well outside the cavity bandwidth and has recently been demonstrated experimentally (*44*). Detection of the ground state could probably prove to be as challenging as its preparation. Proposals to measure the occupancy are diverse, but one method is to measure the weights of the motional sidebands generated by the mechanical motion (*34*).

It is important to note that cooling of mechanical oscillators is also possible using electronic (active) feedback (*10*, *11*, *29*, *45*). This scheme is similar to "stochastic cooling" (*46*) of ions in storage rings and uses a "pick-up" (in the form of an optical cavity interferometer) to measure the mechanical motion and a "kicker" (a radiation-pressure force exerted by a laser on the mirror) to provide a viscous (feedback) force. The idea can also be understood in terms of the feedback loop in Fig. 2B, wherein the lower right optical-feedback branch is replaced by an electrical path driving a second pump laser, which acts as a force actuator on the mirror.

Finally, although originally conceived as a potential limitation in gravitational wave detection, the parametric instability (blue detuned operation of the pump wave) can also be understood as the result of amplification (negative damping) of the mechanical motion (*13*, *17*, *47*). In this sense, the instability is simply the threshold condition in which intrinsic mechanical loss is compensated by amplification. This threshold phenomenon and the subsequent regenerative mechanical oscillation have been studied as a new type of optomechanical oscillator (*48*). Above threshold, the oscillator is regenerative, and oscillation at microwave rates (*49*) has been demonstrated. Additionally, the phase noise of the oscillator has been characterized and observed to obey an inverse power dependence, characteristic of fundamental, Brownian noise (*48*). Quantum back-action is also predicted to set a fundamental low-temperature limit to this linewidth (*50*). The ability to amplify mechanical motion is potentially useful as a means to boost displacements and forces sensitivity (*51*). Finally, returning to the analogy with atomic physics, it is interesting to note that regenerative oscillation (i.e., amplification of mechanical motion) would be expected to occur for trapped ions under blue-detuned excitation.

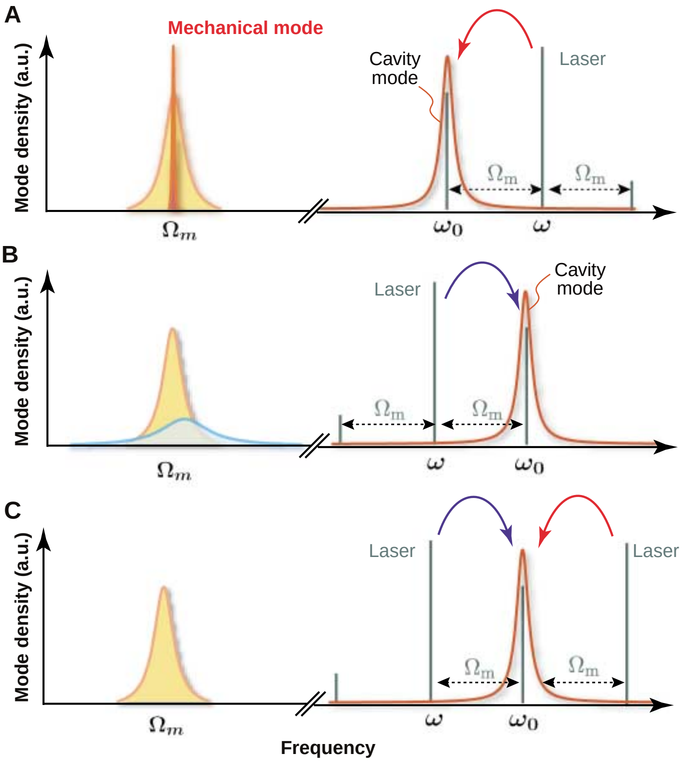

Fig. 4. Frequency domain interpretation of optomechanical interactions in terms of motional sidebands. These sidebands are created on the optical probe wave as photons are Doppler shifted from the mirror surface (which undergoes harmonic motion driven by its thermal energy). Doppler scattering rates into the red (Stokes) and blue (anti-Stokes) sidebands are imbalanced when the probe wave resides to one side of the optical resonance, which can be viewed as a consequence of the asymmetric density of electromagnetic states (Fig. 2A). This imbalance favors the Stokes sideband for a blue-detuned pump and the anti-Stokes sideband for the red-detuned pump, thereby creating a net imbalance in electromagnetic power upon scattering. This imbalance is the origin of mechanical amplification (blue detuning) and cooling (red detuning). (Cooling in this fashion is similar to cavity cooling of atoms.) Only mechanical modes that produce appreciable sideband asymmetry will experience significant gain or cooling. Moreover, the degree of asymmetry can be controlled in an experiment so that a particular mechanical mode can be selected for amplification or cooling. (**A**) Dynamic back-action amplification of mechanical motion via a blue-detuned laser field. The laser scatters pump photons into the cavity, thus creating phonons and leading to amplification. (**B**) Dynamic back-action cooling via a red-detuned laser. Pump photons are scattered into the cavity resonance, thereby removing thermal mechanical quanta from the mechanical oscillator. (**C**) Two-transducer scheme. By symmetrically pumping the cavity on both upper and lower sideband, only one of the quadratures of the mechanical motion is measured with a precision that can exceed the standard limit, thus providing a route to preparing a mechanical oscillator in a squeezed state of mechanical motion via measurement-induced squeezing. a.u., arbitrary units.

Cavity Quantum Optomechanics

A mechanical oscillator has a set of quantum states with energies $E_{\mathrm{N}} = (N + \frac{1}{2})\hbar\Omega_{\mathrm{m}}$, where N is the number of mechanical quanta, and $N = 0$ denotes the quantum ground state. For a mechanical oscillator in the ground state, the ground state energy, $E_0 = \hbar\Omega_{\mathrm{m}}/2$, gives rise to the zero-point motion, characterized by the length scale $x_0 = \sqrt{\hbar/2m\Omega_{\mathrm{m}}}$. As noted earlier, this length scale sets the SQL of mirror position uncertainty in an interferometer such as in Fig. 1. The zero-point motion for structures shown in Fig. 3 ranges from ~10^{-17} m for a macroscopic mirror to ~10^{-12} m for the nanomechanical beam. Such small motions are masked by the thermal motion of the mechanical oscillator, and to enter the regime where quantum fluctuations become dominant and observable requires that the mechanical mode's temperature satisfy $k_{\mathrm{B}}T << \hbar\Omega_{\mathrm{m}}$, equivalently a thermal occupation less than unity. Over the past decade, cryogenically cooled nanomechanical oscillators coupled to an electronic readout have been steadily approaching the quantum regime (*19*, *52*, *53*). Cavity optomechanical systems exhibit high readout sensitivity, in principle already sufficient to detect the minute zero-point motion of a mesoscopic system. The main challenge toward observing quantum phenomena in cavity optomechanical systems lies in reducing the mechanical mode thermal occupation. Using conventional cryogenic cooling, the latter is challenging (1 MHz, corresponding to a temperature of only 50 μK). However, in principle, cooling to these temperatures and even lower is possible with the use of optomechanical back-action cooling.

If a sufficiently low occupancy of the mechanical oscillator is reached (using, for instance, a combination of cryogenic precooling and back-action laser cooling), quantum phenomena of a mesoscopic mechanical object may arise. For example, the quantum back-action by photons could become observable (*54*) or signatures of the quantum ground state. Moreover, the interaction of cold mechanics and a light field can give rise to squeezing of the optical field (*55*). This can be understood by noting that the mechanical oscillator couples the amplitude and phase quadrature of the photons. Moreover, the optomechanical coupling Hamiltonian has been predicted to allow quantum nondemolition measurement of the intracavity photon number (*56*, *57*). The coupling afforded by radiation pressure might even allow the production of squeezed states of mechanical motion. These highly nonintuitive quantum states have been produced for electromagnetic fields over the past decades, and producing them in the mechanical realm would be a notable achievement. Such highly nonclassical states may be possible to generate using measurement-induced squeezing. In this method (*22*), one quadrature component of the mechanical oscillator motion is measured (and no information of the complementary variable is gained) so as to project the mechanical oscillator into a squeezed state of motion. This method (Fig. 4C) involves two incident waves and moreover requires that the mechanical frequency exceeds the cavity decay rate (the resolved sideband regime). A great deal of theoretical work has also been devoted to the question of entangling mechanical motion with an electromagnetic field, or even entangling two mechanical modes. Examples include proposals to achieve quantum super-positions of a single photon and a mirror via a "which path" experiment (*58*) or entangling two mirrors via radiation pressure (*59*).

Emerging Cavity Optomechanical Technologies

Cavity optomechanics may also enable advances in several other areas. First, the ability to provide targeted cooling of nano- and micromechanical oscillators (which are otherwise part of devices at room temperature) bodes well for practical applications because, in principle, conventional cryogenics are unnecessary. Beyond providing a better understanding of fluctuation and dissipative mechanisms, the fact that high displacement sensing is an important element of cavity optomechanics will have collateral benefits in other areas of physics and technology, ranging from scanning probe techniques (*60*) to gravitational-wave detection. Moreover, the ability to create all-optical photonic oscillators on a chip with narrow linewidth and at microwave oscillation frequencies may have applications in radio frequency–photonics. Equally important, cavity optomechanical systems already exhibit strong nonlinearity at small driving amplitudes, which offer new functions related to optical mixing (*61*). Finally, although all current interest is focused on radiation-pressure coupling, cavity optomechanical systems based on gradient forces are also possible. Although aimed at a separate set of applications, there has been substantial progress directed toward gradient-force control of mechanical structures using cavity optomechanical effects (*62*–*64*).

Summary

The interaction of mechanical and optical degrees of freedom by radiation pressure is experiencing a paradigm shift in control and measurement of mechanical motion. Radiation-pressure coupling has opened an extremely broad scope of possibilities, both applied and fundamental in nature. With the continued trends toward miniaturization and dissipation reduction, radiation pressure can become an increasingly important phenomenon that will probably allow advances, both in terms of technology as well as in fundamental science. It may well provide a way to probe the quantum regime of mechanical systems and give rise to entirely new ways of controlling mechanics, light, or both. It also seems likely that beyond precision measurement, there will be new technologies that leverage cooling and amplification.

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