Light-mediated strong coupling between a mechanical oscillator and atomic spins 1 meter apart

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Engineering strong interactions between quantum systems is essential for many phenomena of quantum physics and technology. Typically, strong coupling relies on short-range forces or on placing the systems in high-quality electromagnetic resonators, which restricts the range of the coupling to small distances. We used a free-space laser beam to strongly couple a collective atomic spin and a micromechanical membrane over a distance of 1 meter in a room-temperature environment. The coupling is highly tunable and allows the observation of normal-mode splitting, coherent energy exchange oscillations, two-mode thermal noise squeezing, and dissipative coupling. Our approach to engineering coherent long-distance interactions with light makes it possible to couple very different systems in a modular way, opening up a range of opportunities for quantum control and coherent feedback networks.

A fundamental challenge in this approach is that the same photons that generate the coupling eventually leak out, thus allowing the systems to decohere at an equal rate. For this reason, light-mediated coupling is mainly seen today as a means for unidirectional state transfer (14–16) or entanglement generation by collective measurement (17–19) or dissipation (20). Decoherence by photon loss can be suppressed if the waveguide is terminated by mirrors to form a high-quality resonator, which has enabled coherent coupling of superconducting qubits (21, 22), atoms (23), or atomic mechanical oscillators (24) in mesoscopic setups. However, stability constraints and bandwidth limitations make it difficult to extend resonator-based approaches to larger distances. Despite recent advances with coupled cavity arrays (25, 26), strong bidirectional Hamiltonian coupling mediated by light over a truly macroscopic distance remains a challenge.

To realize long-distance Hamiltonian interactions, we pursued an alternative approach that relies on connecting two systems by a laser beam in a loop geometry (27, 28). The systems can exchange photons through the loop, thereby realizing a bidirectional interaction. Moreover, the loop leads to an interference of quantum noise introduced by the light field. For any system that couples to the light twice and with opposite phase, quantum noise interferes destructively and associated decoherence is suppressed. At the same time, information about that system is erased from the output field. In this way, the coupled systems can effectively be closed to the environment, even though the light field mediates strong interactions between them. Because the coupling is mediated by light, it allows systems of different physical nature to be connected over macroscopic distances. Furthermore, by manipulating the light field between the systems, one can reconfigure the interaction without having to modify the quantum systems themselves. These features will be useful for quantum networking (4).

We used this scheme to couple a collective atomic spin and a micromechanical membrane held in separate vacuum chambers, thereby realizing a hybrid atom-optomechanical system (8). First experiments with such setups have recently demonstrated sympathetic cooling (29, 30), quantum back-action evading measurement (31), and entanglement (32). Here, we realize strong Hamiltonian coupling and demonstrate the versatility of light-mediated interactions: We engineer beamsplitters and parametric-gain Hamiltonians and switch from Hamiltonian to dissipative coupling by applying a phase shift to the light field between the systems. This high level of control in a modular setup gives access to a unique toolbox for designing hybrid quantum systems (9) and coherent feedback loops for advanced quantum control strategies (33).

Description of the coupling scheme

In the experimental setup (Fig. 1A) (34), the atomic ensemble consists of \( N = 10^7 \) laser-cooled \(^{87}\)Rb atoms in an optical dipole trap. The atoms form a collective spin \( F = \sum_{i=1}^{N} f^{(i)} \), where \( f^{(i)} = 2 \) ground-state spin vector of atom \( i \). Optical pumping polarizes \( F \) along an external magnetic field \( B_0 \) in the \( x \)-direction such that the spin acquires a macroscopic orientation \( F_x = -fN \). The small-amplitude dynamics of the transverse spin components \( F_y, F_z \) are well approximated by a harmonic oscillator (35) with position \( X_x = F_x / \sqrt{|F_z|} \) and momentum \( P_x = F_y / \sqrt{|F_z|} \). It oscillates at the Larmor frequency \( \Omega \approx B_0 \), which is tuned by the magnetic field strength. A feature of the spin system is that it can realize such an oscillator with either positive or negative effective mass (31, 36). This is achieved by reversing the orientation of \( F \) with respect to \( B_0 \), which reverses the sense of rotation of the oscillator in the \( X_x, P_x \) plane (Fig. 1B). This feature allows us to realize different Hamiltonian dynamics with the spin coupled to the membrane.

The spin interacts with the coupling laser beam through an off-resonant Faraday interaction (35) \( H_a = 2h \sqrt{\Gamma_0 / \Gamma_a} S_x X_x Z_0 \), which couples \( X_x \) to the polarization state of the light, described by the Stokes vector \( S \). Initially, the laser is linearly polarized along \( x \) with \( S_x = \Phi_L / 2 \), where \( \Phi_L \) is the photon flux. The strength of the atom-light coupling depends on the spin measurement rate \( \Gamma_0 \propto d_0 \Phi_L / X_x^2 \), which is proportional to the optical depth \( d_0 \approx 300 \) of the atomic ensemble (34). Choosing a large laser–atom detuning \( \Delta_0 = -2\pi \times 80 \) GHz suppresses spontaneous photon scattering while maintaining a sizable coupling.

The mechanical oscillator is the \((2,2)\) square-drum mode of a silicon nitride membrane at a vibrational frequency of \( \Omega_m = 2\pi \times 1.957 \) MHz with a quality factor of \( 1.3 \times 10^6 \) (37). It is placed...
in a short single-sided optical cavity to enhance the optomechanical interaction while maintaining a large cavity bandwidth for fast and efficient coupling to the external light field. Radiation pressure couples the membrane displacement $X_m$ to the amplitude fluctuations $X_L$ of the light entering the cavity on resonance, with Hamiltonian $H_m = 2\hbar \sqrt{\Gamma_m X_m X_L}$, (38). Here, we define the optomechanical measurement rate $\Gamma_m = (4g_0^2/k)^2 \Phi_m$ that depends on the cavity optomechanical coupling constant $g_0$, cavity linewidth $\kappa$, and photon flux $\Phi_m$ entering the cavity (34). In the present setup, the optomechanical cavity is mounted in a room-temperature vacuum chamber, making thermal noise the dominant noise source of the experiment.

The light field mediates a bidirectional coupling between the spin and the membrane. A spin displacement $X_s$ is mapped by $H_s$ to a polarization rotation $S_y = 2 \sqrt{\Gamma_m} S_y X_m$ of the light. A polarization interferometer (Fig. 1A) converts this to an amplitude modulation $X_L = S_y / \sqrt{\Gamma_m}$ at the optomechanical cavity, resulting in a force $P_m = -4\hbar \sqrt{\Gamma_m} X_L$ on the membrane. Conversely, a membrane displacement $X_m$ is turned by $H_m$ into a phase modulation $P_s = -2\hbar \sqrt{\Gamma_m} X_s$ on the spin. A small angle between the laser beams in the two atom-light interactions prevents light from going once more to the membrane. Consequently, the cascaded setup promotes a bidirectional spin-membrane coupling. A fully quantum mechanical treatment (36) confirms this picture and predicts a spin-membrane coupling strength $g = (\gamma^2 + \eta^2) \sqrt{\Gamma_m}$, accounting for an effective optical power transmission $\eta^2 = 0.8$ between the systems.

The light-mediated interaction can be thought of as a feedback loop that transmits an spin excitation to the membrane, whose response then acts back on the spin, and vice versa (Fig. 1B). After one round trip, the initial signal has acquired a phase $\phi$, the loop phase. The discussion above refers to a vanishing loop phase $\phi = 0$ and shows that the forces $P_m = -2gX_s$ and $P_s = +2gX_m$ differ in their relative sign. Such a coupling is nonconservative and cannot arise from a Hamiltonian interaction. With full access to the laser beams, we can tune the loop phase by inserting a half-wave plate in the path from the membrane back to the atoms, which rotates the Stokes vector by an angle $\phi = \pi$ about $S_y$. This inverts both $S_y$ and $S_x$ which respectively carry the spin and membrane signals, thus switching the dynamics to a fully Hamiltonian force, $P_m = -2gX_s$ and $P_s = -2gX_m$.

All these phenomena are unified in a rigorous quantum mechanical theory (28) of the cascaded light-mediated coupling, which also correctly describes the dynamics for an arbitrary loop phase. It allows us to describe the effective dynamics of the coupled spin-membrane system with density operator $\rho$ by a Markovian master equation,

$$\dot{\rho} = \frac{1}{i\hbar}[H_0 + H_{eff}, \rho] - \frac{1}{2}(J^I J + \rho J^I J^I) + J^I J^I$$

(1)

Here, we neglect optical loss and light propagation delay between the systems for brevity. The dynamics consist of a unitary part with free harmonic oscillator Hamiltonian $H_0 = \sum_{x=x,y,z} \hbar \omega_0 (X_x^2 + P_x^2)/2$ and effective interaction Hamiltonian $H_{eff} = (1 - \cos \phi) \hbar g X_m X_s + 2\sin(\phi) \hbar g X_s X_s^2 + 2\sin(\phi) \hbar g X_m X_m^2$, and a dissipative part with collective jump operator $J = \sqrt{2\Gamma_m X_m + i[1+ \exp(i\phi)] \sqrt{2\Gamma_s X_s}}$. Next to the coherent spin-membrane coupling, $H_{eff}$ also includes a spin self-interaction that vanishes for the specific cases $\phi = 0, \pi$ considered here. The jump operator contains a constant membrane term and a spin term that is modulated by $\phi$ as a result of interference of the two spin-light interactions. From the dependence of $H_{eff}$ and $J$ on $\phi$, it is clear that $\phi = 0$ corresponds to vanishing Hamiltonian coupling and maximum dissipative coupling. Accordingly, we refer to $\phi = 0$ as the dissipative regime. On the other hand, $\phi = \pi$ maximizes the coherent spin-membrane coupling in $H_{eff}$ and at the same time leads to destructive interference of the spin term in $J$; thus, we call $\phi = \pi$ the Hamiltonian regime. We experimentally explored both regimes, each with the atomic spin realizing either a positive- or negative-mass oscillator. This gives rise to a whole range of different dynamics in a single system, which can be harnessed for different purposes in quantum technology.
Results
Normal-mode splitting
We first investigate the light-mediated coupling in the Hamiltonian regime ($\theta = \pi$) and with the spin realizing a positive-mass oscillator. At a magnetic field of $B_0 = 2.81$ G, the spin is tuned into resonance with the membrane ($\Omega_+ = \Omega_m$). In this configuration, the resonant terms in $H_{\text{eff}}$ realize a beamsplitter interaction $H_{\text{BS}} = \hbar g (b_m^\dagger b_m + b_+^\dagger b_+)$, which generates state swaps between the two systems. Here $b_\pm = (X_m + iP_m)/\sqrt{2}$ and $b_m = (X_m + iP_m)/\sqrt{2}$ are annihilation operators of the spin and mechanical modes, respectively.

We perform spectroscopy of the coupled system using independent drive and detection channels for the spin and the membrane. The membrane vibrations are recorded by balanced homodyne detection using an auxiliary laser beam coupled to the cavity in orthogonal polarization. This beam is amplitude-modulated to drive the membrane. The spin precession is detected by splitting off a small portion of the coupling beam on the path from spin to membrane. A radio-frequency magnetic coil drives the spin. We measure the amplitude and phase response of either system using a lock-in amplifier that demodulates the detector signal at the drive frequency ($\Omega_+ = \Omega_m$). After spin-state initialization, we simultaneously switch on coupling and drive and start recording. The drive frequency is kept fixed during each experimental run and stepped between consecutive runs.

Figure 2, A and B, shows the membrane’s response in amplitude and phase, respectively. With the coupling beam off, it exhibits a Lorentzian resonance of linewidth $\gamma_m = 2\pi \times 0.3$ kHz, broader than the intrinsic linewidth as a result of optomechanical damping by the red-detuned cavity field ($\gamma_{m_r} = 2\pi \times 4$ kHz, broadened by the coupling beam. When we turn on the coupling to the spin, the membrane resonance splits into two hybrid spin-mechanical normal modes. This signals strong coupling ($\gamma_0$, 40), where light-mediated coupling dominates over local damping. Fitting the well-resolved splitting yields $2g = 2\pi \times 6.1$ kHz, which exceeds the average linewidth ($\gamma_0 + \gamma_m$)/2 = $2\pi \times 2$ kHz and agrees with the expectation based on an independent calibration of the systems ($\gamma_0$, 34). A characteristic feature of the long-distance coupling is a finite delay $\tau$ between the systems. It causes a linewidth asymmetry of the two normal modes when $t_+ = \Omega_m$, which we observe in Fig. 2. The fits yield a value of $\tau \approx 15$ ns, consistent with the propagation delay of the light between the systems and the cavity response time.

We also observe normal-mode splitting in measurements of the spin (Fig. 2, C and D). Here, the combination of the broader spin linewidth and the much narrower membrane resonance results in a larger dip between the two normal modes and a larger phase shift, in analogy to optomechanically induced transparency ($\gamma_0$, 34).

Energy exchange oscillations
Having observed the spectroscopic signature of strong coupling, we now use it for swapping spin and mechanical excitations in a pulsed experiment. We start by coherently exciting the membrane to $\sim 2 \times 10^5$ phonons, a factor of 100 above its mean equilibrium energy, by applying an amplitude modulation pulse to the auxiliary cavity beam (Fig. 3A). At the same time, the spin is prepared in its ground state with $\Omega_+ = -\Omega_m$. The coupling beam is switched on at time $t = 0 \mu s$ and the displacements $X_+(t)$ and $X_m(t)$ of the spin and membrane are continuously monitored via the independent detection. From the measured mean square displacements, we determine the excitation number of each system ($\gamma_0$, 34). Figure 3C shows the excitation numbers as a function of the interaction time. The data show coherent and reversible energy exchange oscillations from the membrane to the spin and back with an oscillation period of $T \approx 150 \mu s$, in accordance with the value $\pi g$ extracted from the observed normal-mode splitting. Damping limits the maximum energy transfer efficiency at time $T/2$ to about 40%.

The same experiment is repeated but with the initial drive pulse applied to the spin (Fig. 3, B and D). Here, we observe another set of exchange oscillations with the same periodicity, swapping an initial spin excitation of $n_s = 3 \times 10^5$ to the membrane and back. After the coherent dynamics have decayed, the systems equilibrate in a thermal state of $\sim 3 \times 10^3$ phonons, lower than the effective optomechanical bath of $1.5 \times 10^4$ phonons, demonstrating sympathetic cooling (29) of the membrane by the spin. The observed sympathetic cooling strength agrees with simulations using the experimentally determined parameters.

Parametric-gain dynamics
So far we have explored Hamiltonian coupling of the membrane to a spin oscillator with positive effective mass, where the resonant interaction is of the beamsplitter type. If instead we reverse the magnetic field to $B_0 = -2.81$ G but keep the spin pumping direction the same, the collective spin is prepared in its highest-energy state with $F_s = 0$ (Fig. 1B). The resonant term $H_{\text{BS}}$ is now the parametric-gain interaction ($\gamma_0$, 38) $H_{\text{PG}} = \hbar g (b_m^\dagger b_m + b_+^\dagger b_+)$, which generates correlations between the two systems.

We investigate the dynamics generated by $H_{\text{PG}}$ with the membrane driven by thermal noise. To quantify the development of spin-mechanical correlations, we determine slowly varying quadratures $X_{m,r}$ and $P_{m,r}$ of both systems as the cosine and sine components of the demodulated detector signals, respectively ($\gamma_0$, 34). Adjusting the demodulator phase allows us to find the basis with the strongest correlations. Figure 4A shows histograms of the measured spin-mechanical correlations after
an interaction time of $t = 100 \mu s$. In each subplot, the dashed ellipse corresponds to the Gaussian $1 \sigma$ contour of the measured histogram at $t = 0 \mu s$, and the solid ellipse is the contour at $t = 100 \mu s$. Relative to the uncorrelated initial state, the histograms show strong amplification along the axes $X_\perp = (X_+ + X_-)/\sqrt{2}$ and $P_\perp = (P_+ - P_-)/\sqrt{2}$, and a small amount of thermal noise squeezing along $X_\parallel = (X_+ - X_-)/\sqrt{2}$ and $P_\parallel = (P_+ + P_-)/\sqrt{2}$. The quadrature pairs $X_\parallel, P_\parallel$ and $X_\perp, P_\perp$ remain uncorrelated.

In the time evolution of the combined variances $X_\parallel$ and $P_\parallel$ (Fig. 4B), at $t = 0$ all variances start from the same value, indicating an uncorrelated state. As time evolves, the variances of $X_\parallel$ and $P_\parallel$ grow exponentially, demonstrating the dynamical instability in this configuration, while $X_\perp$ and $P_\perp$ are squeezed and reach a minimum at $t = 80 \mu s$ before they grow again. The exponential growth rate of $2 \tau \times 4.5$ kHz is consistent with the value of $2g - (\gamma_m + \gamma_s)/2$ extracted from the normal-mode splitting. For comparison, we also show simulated variances for the experimental parameters, which are given by the lines in Fig. 4B (34). The solid lines show that good agreement between data and simulation is found when accounting for a spin detector noise floor of $6 \times 10^3$. The dashed lines correspond to perfect detection and show thermal noise squeezing by 5.5 dB. Realizing the parametric-gain interaction by light-mediated coupling represents an important step toward the generation of spin-mechanical entanglement by two-mode squeezing across macroscopic distances. Such entanglement is useful for metrology beyond the standard quantum limit ($\ell$).

**Control of the loop phase**

Equipped with control over both the loop phase and the effective mass of the spin oscillator, we can access four different regimes of the spin-membrane coupling: two Hamiltonian configurations with $\phi = \pi$ and $\Omega_m = \pm \Omega_m$ and the two corresponding dissipative configurations where we set $\phi = 0$ by omitting the half-wave plate in the optical path from membrane to atoms (34). Although the dynamics in these configurations are fundamentally different and have different quantum noise properties, we obtain simple equations of motion for the expectation values,

$$\dot{X}_m + \gamma_m X_m + \Omega^2_{\mu m} X_m = -g \Omega_m X_\perp (t - \tau)$$

(2)

$$\dot{X}_s + \gamma_s X_s + \Omega^2_s X_s = \mp g \Omega_s \cos(\phi) X_m (t - \tau)$$

(3)

with the damped harmonic oscillations on the left and the delayed coupling terms on the right. These are derived from Heisenberg-Langevin equations of the full system (34) and reproduce the dynamics of the master equation in the limit $\tau \to 0$. Two distinct regimes can be identified. If $\Omega_s \cos \phi < 0$, we expect stable dynamics equivalent to a beamsplitter interaction. In the opposite case where $\Omega_s \cos \phi > 0$, the dynamics are equivalent to a parametric-gain interaction and are unstable. A simultaneous sign reversal of $\Omega_s$ and a $\pi$-shift of $\phi$ should leave the dynamics invariant.

To probe the dynamics in these configurations, we record thermal noise spectra of the membrane while the spin Larmor frequency is tuned across the mechanical resonance $\Omega_m = 2\tau \times 1.957$ MHz. The Hamiltonian configuration with positive-mass spin oscillator is depicted in Fig. 5A, showing an avoided crossing at $\Omega_s = \Omega_m$ with frequency splitting $2g = 2\tau \times 5.9$ kHz, as in Fig. 2. The dashed lines are the calculated normal-mode frequencies (34). The enhancement of the mechanical noise power
for \( \Omega_m < \Omega_m \) as compared to increased damping for \( \Omega > \Omega_m \), is again a consequence of the finite optical propagation delay \( \tau \) modifying the damping \((34)\).

Switching to the dissipative regime with \( \phi = 0 \) renders the system unstable because of positive feedback of the coupled oscillations (Fig. 5B). Instead of an avoided crossing, the normal modes are now attracted and cross near \( \Omega_s \approx 2 \pi \times 1.953 \) MHz, forming one strongly amplified and one strongly damped mode. The former leads to exponential growth of correlated spin-mechanical motion, finally resulting in limit-cycle oscillations that dominate the power spectrum. This results in a breakdown of the coupled oscillator model, such that the observed spectral peak shifts toward the unperturbed mechanical resonance. Still, the data are in good agreement with the theoretical predictions.

In Fig. 5, C and D, we repeat the experiments of Fig. 5, A and B, with a negative-mass spin oscillator. The data show that Hamiltonian coupling with a negative-mass spin oscillator produces spectra similar to those produced by dissipative coupling with a positive-mass spin oscillator. In these configurations, the coupled system features an exceptional point \((41)\) where the normal modes become degenerate \((42)\) and define the squeezed and antisqueezed quadratures. Conversely, dissipative coupling together with an inverted spin (Fig. 5D) shows an avoided crossing with parameters similar to those in the Hamiltonian case (Fig. 5A). This equivalence at the level of the expectation values is expected to break down once quantum noise of the light becomes relevant. As a result of interference in the loop, quantum back-action on the spin is suppressed in the Hamiltonian coupling configuration but is enhanced in the dissipative configuration.

A necessary condition for quantum back-action cancellation is destructive interference of the spin signal in the output field \((34)\). Figure 5, E and F, shows homodyne measurements of coherent spin precession on the coupling beam output quadrature \( X_L^{\text{out}} \) in the time and frequency domains, respectively. Toggling the loop phase between \( \phi = 0 \) and \( \phi = \pi \), we observe a large interference contrast \((>10)\) in the root-mean-squared (RMS) spin signal, showing that a spin measurement made by light in the first pass can be erased in the second pass. Optical loss of \( 1 - \eta^4 = 0.35 \) inside the loop allows some information to leak out to the environment and brings in uncorrelated noise, limiting the achievable back-action suppression. Full interference in the output is still observed because the carrier and signal fields are subject to the same losses. Because this principle of quantum back-action interference is fully general, it could also be harnessed for other optical or microwave photonic networks \((4, 27)\).

**Conclusion**

The observed normal-mode splitting and coherent energy exchange oscillations establish strong spin-membrane coupling, where the coupling strength \( g \) exceeds the damping rates of both systems \((39)\). To achieve quantum-coherent coupling \((40)\), \( g \) must also exceed all thermal and quantum back-action decoherence rates. This will make it possible to swap nonclassical states between the systems or to generate remote entanglement by two-mode squeezing. Thermal noise on the mechanical oscillator is the major source of decoherence in our room-temperature setup. We expect that modest cryogenic cooling of the optomechanical system to 4 K together with an improved mechanical quality factor of \(>10^7\) \((43)\) will enable quantum-limited operation \((34)\). The built-in suppression of quantum back-action
action in the Hamiltonian configuration is a crucial feature of our coupling scheme. Interference of the two spin-light interactions reduces the spin’s quantum back-action rate to \( \gamma_{s,ba} = (1 - \eta^2) \gamma_{s}^{\text{mem}} \) whereas it is \( \gamma_{m,ba} = \eta^2 \gamma_{m}^{\text{mem}} \) for the membrane. Assuming thermal noise to be negligible, the quantum cooperativity \( C = \frac{2g}{(\gamma_{s,ba} + \gamma_{m,ba})} \) can be optimized for a given one-way transmission \( \eta^2 \). We find an upper bound \( C \leq \frac{\pi}{1 + \eta^2} / \sqrt{1 - \eta^2} \), reaching 2.7 for our current setup. The bound is achieved for an optimal choice of measurement rates \( \Gamma_{g}/\Gamma_{m} = \eta^2 / (1 - \eta^2) \), balancing the back-action on both systems. Further improvement is possible with a double-loop coupling scheme that also suppresses quantum back-action on the membrane (28). In this case, \( C = \eta^2 (1 - \eta^2) / \Gamma_{s} = \eta^2 \Gamma_{m} \) is inversely proportional to optical loss, scaling more favorably at high transmission so that \( C \approx 10 \) can be reached for \( \eta^2 \approx 0.9 \).

Our results demonstrate a comprehensive and versatile toolbox for generating coherent long-distance interactions with light and open up a range of exciting opportunities for quantum information processing, simulation, and metrology. The coupling scheme constitutes a coherent feedback network (33) that allows quantum systems to directly exchange, process, and feed back information without the use of classical channels. The ability to create coherent Hamiltonian links between separate and physically distinct systems in a reconfigurable way substantially extends the available toolbox not only for hybrid spin-mechanical interfaces (9, 37) but for quantum networks (4) in general. It facilitates the faithful processing of quantum information and the generation of entanglement between spatially separated quantum processors across a room-temperature environment.

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SUPPLEMENTARY MATERIALS

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Table S1

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