Single shot quantum non-demolition readout with a hybrid opto-mechanical transducer

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A quantum emitter coupled to a nano-mechanical oscillator is a hybrid system where a macroscopic degree of freedom can be manipulated quantum mechanically. Recent progress in nanotechnology has allowed to realize such devices by embedding single artificial atoms into vibrating wires or membranes, opening up new perspectives for quantum information technologies and for the exploration of the quantum-classical boundary. In this letter, we show that optical modulation of the quantum emitter state results in the build up of a large phonon field, potentially exceeding the Brownian motion even at room temperature. This efficient scheme can be exploited to perform single shot quantum non-demolition readout of a quantum bit, both in superconducting and semiconductor systems.

Coupling electromagnetic and mechanical degrees of freedom is a new quest of quantum physics. First experiments were based on the interaction between the electromagnetic field stored in a cavity and a moving mirror, in the so-called linearized regime where the state of both resonators can be described classically [1, 2]. Instead of the electromagnetic field, it has now become possible to couple the mechanical resonator with a quantum system [3], like a single spin in a diamond Nitrogen Vacancy (NV) center [4] or a superconducting qubit [5, 6]. The material strain has been proposed as well as a coupling mechanism between a semiconductor quantum dot a mechanical nanoresonator [7]. In the case of a radiative quantum system, the influence of the mechanical motion on the fluorescence has been recently evidenced [8]. In this letter, we investigate the opposite mechanism, and theoretically show that the quantum emitter under modulated optical drive can efficiently pump phonons, thus behaving like an opto-mechanical transducer at the single quantum level. Two different excitation schemes are considered, either based on ultra-short pulses of incoherent light or on a continuous driving field resonant with the emitter’s transition. In both cases, when the drive is modulated at the mechanical frequency, a coherent phonon field builds up by constructive interference. This scenario is intrinsically different from earlier proposals for mechanical cooling [8, 9] or phonon lasing [10], where the emitter is continuously driven. Based on this excitation scheme, we suggest a method inspired from electron shelving technique used in ions traps [11, 12], to implement quantum non-demolition (QND) readout of the quantum system state. This proposal is applied to realistic systems ranging from spin states of a NV center, to superconducting qubits.

The system under study is represented in Fig. 1. A quantum emitter interacts with a mechanical resonator while its population is controlled by a modulated laser beam. The emitter-phonon coupling is modeled by the independent spin-boson Hamiltonian

$$H = \hbar \omega_0 (\sigma_z + 1/2) + \hbar g_m (\sigma_z + 1/2) (b + b^\dagger) + \hbar \Omega b^\dagger b,$$  \hspace{1cm} (1)

where $\omega_0$ is the emitter’s transition frequency, $\sigma_z = 1/2(|e\rangle\langle e| - |g\rangle\langle g|)$ where $|e\rangle$ and $|g\rangle$ are its excited and ground state respectively, $b$ is the phonon annihilation operator, $\Omega$ is the phonon frequency and $g_m$ is the emitter-phonon coupling term. We respectively denote $\gamma$ and $\Gamma$ the relaxation terms (due to radiative recombination) of the emitter and of the phonon mode, with $\Gamma \ll \Omega$ and $\gamma \ll \omega_0$. We assume $\omega_0, \gamma \gg \Omega$, such that the quantum emitter dynamics follows adiabatically that of the phonon field. The decoherence of the mechanical oscillator can be neglected between the excitation and the radiative relaxation of the quantum emitter, so that the evolution of the system can be considered Hamiltonian.

In the adiabatic regime investigated here, the Hamiltonian is diagonalized in the subspaces $|g\rangle\langle g|$ and $|e\rangle\langle e|$ independently, yielding in the first case $H_g = \hbar \omega_0 b^\dagger b$ of eigenstates the usual Fock states $|n\rangle = b^n|0\rangle/\sqrt{n!}$ , and in the second case, $H_e = \hbar (\omega_0 - g_m^2/\Omega) + \hbar \Omega b^\dagger B$, where $B = b + g_m/\Omega$. The new eigenstates for the phononic field are $|\tilde{n}\rangle = B^n/\sqrt{n!}|0\rangle$, where $|0\rangle$ is the displaced vacuum defined as $|0\rangle = |\sim g_m/\Omega\rangle$. The incoherent optical excitation scenario is entirely captured by these two Hamiltonians. At initial time, the emitter is in its ground state and the phonon field is empty, i.e. the initial state is $|g, 0\rangle$. A
Figure 2. Phonon number $N$ normalized to $|\beta_{\text{max}}|^2$ versus time, a) on a time scale comparable to $1/\Gamma$, the characteristic mechanical oscillator damping time, and b) during the first two periods $2\pi/\Omega$ of the pulses sequence. The inset shows a close up of the first four photoluminescence cycles. The excitation laser modulation is shown as a solid blue line.

A non-resonant strong pulse excites the system into $|e, 0\rangle = |e, g_m/\Omega\rangle$: the resonator now contains a coherent field that evolves under $H_e$ and rotates at the frequency $\Omega$ in the displaced frame, so that $|\psi(t)\rangle = |e, g_m/\Omega e^{-i\Omega t}\rangle$. After a typical time $\tau \sim \gamma^{-1} \ll \Omega^{-1}$, spontaneous emission takes place and the coupling term switches off, leaving in the mechanical resonator a residual coherent phonon field $|\delta\beta\rangle = |g_m/\Omega (e^{-i\Omega t} - 1)\rangle \sim | - ig_m \beta\rangle$. Therefore the quantum emitter behaves as a microscopic classical source in the phononic mode where each fluorescence cycle leads to the addition of a microscopic coherent field.

To enhance the size of the generated phonon field $\beta$, the number of fluorescence cycles is increased using a train of ultra-short pulses of duty cycle $\tau_1/\tau_2$, where $\tau_1$ is the average excitation lifetime in $|e\rangle$ and $\tau_2$ the delay between two excitation pulses. Since each generated micro-field $|\delta\beta\rangle$ points in the negative imaginary part of the complex plane, constructive interference only takes place when $\Im(\beta) < 0$ (where $\Im$ stands for imaginary part), which only happens during one half of the mechanical resonance period. To promote efficient build up, we modulate the pulse train at the mechanical frequency. The corresponding time structure of the pulses train is shown in Fig 2b. The dynamics of the phonon field is derived from a doubly recurrent relation linking the phonon complex amplitude $\beta$, at time $t$ just before a pulse train to $\beta_{n+1}$ at time $t + 2\pi/\Omega$:

$$\beta_{n+1} = a^{n_p} \beta_n + b \frac{1 - a^{n_p}}{1 - a},$$

where $n_p = \Omega/\pi(\tau_1 + \tau_2)$ is the number of pulses per period, $a \equiv e^{-\alpha(\tau_1 + \tau_2)}$, $b \equiv \frac{2\pi}{\Omega} e^{-\alpha \tau_1}(1 - e^{-\alpha \tau_1})$, $\alpha \equiv i\Omega + \Gamma/2$ (Further details are given in the Methods section).

The phonon number $|\beta(t)|^2$ build-up is shown in Fig 2a in scaled units as a function of time. After an elapsed time of a few $1/\Gamma$s, the excited phonon number saturates when the loss equals the injected field. In Fig 2b we see how each pulse contributes to the build-up of the phonon field, while during the off times, only damping takes place. We examine also the motion of the oscillator $x = (\beta + \beta^*) x_{ZPF}$, where $x_{ZPF} = \sqrt{\frac{n_e}{2\Omega}}$ is the zero point position fluctuation of the oscillator, and $m$ the oscillator effective mass. The result is shown in Fig 2c: as expected from the linearity of the exciting mechanism, we see that the harmonicity of the oscillation is preserved. From eq. (2) one can derive the maximum phonon number $|\beta_{\text{max}}|^2$ achievable by this technique :

$$|\beta_{\text{max}}|^2 = 4Q \frac{g_m}{\Omega} \frac{\tau_1}{\tau_1 + \tau_2}$$

where $Q = \Omega/2\pi\Gamma$ is the mechanical oscillator quality factor. The phonons rise time $T_r$ depends on only the damping time $1/\Gamma$ of the mechanical oscillator, namely $T_r = 4\pi/\Omega \ln(1 - 1/\sqrt{2})Q \approx 2.46/\Gamma$.

In this pulsed scheme the best pumping rate is achieved when $\tau_2 \rightarrow 0$, namely when the emitter population is inverted. Consequently, we consider a second excitation scheme where the drive is provided by a continuous laser resonant with the $|e\rangle - |g\rangle$ transition and modulated at the frequency of the phonon field $\Omega/2\pi$ with a 50% duty cycle. Since we expect a coherent phonon field to build up, we develop a semi-classical analysis of the problem and derive generalized Bloch equations for the evolution of the coupled phonon-emitter system:

$$\dot{s} = -i g (s^* - s) - \gamma (s_+ + \frac{1}{2})$$
$$\dot{s} = -i (\delta + g_m (\beta + \beta^*) s - \frac{g_m}{\Omega} s_+ + \frac{1}{2})$$
$$\dot{\beta} = -i\Omega \beta - ig_m (s_+ + \frac{1}{2}) - \frac{G}{Z}$$

We have defined $s_\pm = \langle \sigma_\pm \rangle$, $s = \langle \sigma \rangle$ and $\beta = \langle b \rangle$ the average population, dipole and phonon field amplitudes respectively, $\delta = \omega_0 - \omega_L$ is the laser-emitter detuning assuming zero phonons, $\omega_L$ the laser optical frequency, and $g$ the classical Rabi frequency. In the adiabatic regime under study, the equations describing the evolution of the emitter and of the phonon field can be solved indepen-
Figure 3. Phonon number $N$ normalized to $|\beta_{\text{max}}|^2$ versus time. a) on a time scale comparable to $1/\Gamma$ the characteristic mechanical oscillator damping time. Normalized phonon number on a short time scale; b) during the first four periods $2\pi/\Omega$ of the excitation, and c) during the first four periods later, when the phonon field gets close to saturation. The modulated resonant excitation is shown as a solid blue line. Oscillator position $x$ in solid green line d); under modulated CW excitation for the first four periods, e); after switching on the unmodulated CW laser.

The normalized population reads

$$\tilde{x}_2^\infty = -\frac{1}{2} + \frac{1}{\text{\Gamma}(\delta + g_m(\beta_n + \beta_n^*))^2}, \quad (5)$$

The saturation of the population depends on the overall detuning $\delta' = \delta + g_m(\beta_0 + \beta_0^*)$, that involves the phononic field phase and amplitude. In the saturated regime of the Bloch equations where $g \gg \gamma$, the complex amplitude of the phonon field checks

$$\dot{\beta} = -i\Omega\beta - \frac{ig_m}{2} + \frac{1}{g^2} \left(\frac{\gamma + g_m(\beta_n + \beta_n^*)}{\gamma + g_m(\beta_0 + \beta_0^*)}\right) \beta - \frac{\Gamma}{2} \beta. \quad (6)$$

The first term is the free evolution of the phonon field, the second one an effective phonon pumping term provided by the quantum emitter, and the last one describes the phonon damping. In the following, we choose the laser to be resonant with the transition, i.e. $\delta = 0$ and of intensity large enough such that $g^2 \gg g_m^2|\beta|^2$, i.e. $\delta' \approx 0$. When the laser is on, the dynamics of the phonon field is governed by eq. (6), while when the laser is off the phonon field evolves freely. The sequence of $\{\beta_n\}$ checks now (details can be found in the Methods section):

$$\left(-\frac{ig_m}{2\alpha} + \beta_n e^{-1/4Q}\right)e^{-1/4Q} = \beta_{n+1}, \quad (7)$$

where $\beta_n$ is the phonon field at time $t$ and $\beta_{n+1}$ at time $t + 2\pi/\omega$. The normalized dynamics is shown in Fig.3b on a long time scale, and in Fig.3a on a shorter time scale during the first four laser modulation period. The maximum phonon number achievable is derived from the steady state of eq. (7) and amounts to

$$|\beta_{\text{max}}|^2 = 2Qg_m^2/\Omega. \quad (8)$$

In spite of a rather different microscopic mechanism, both excitation schemes yield an overall identical behavior where saturation of the Bloch equation (leading to equalization of the populations) corresponds to the incoherent pulsed case with $\tau_1 = \tau_2$. Indeed, $|\beta_{\text{max}}|^2 = |\beta_{\text{max}}|^2 = |2g_mQ/\Omega|^2$. We examine also the motion of the oscillator in this modulated CW excitation regime. The result is shown in Fig.3c: like previously, the harmonicity of the motion is well preserved. Finally, note that without optical modulation, the scheme would be way less efficient as it clearly appears in Fig.3d. We have plotted the oscillator motion as a function of time under non-modulated CW excitation. Right after switching on the laser, transient oscillations occur and the mechanical oscillator eventually stabilizes at a new position $x'_0$ shifted from its rest position $x_0$ by $x_{ZPF}g_m/\Omega$. According to typical figures of realistic hybrid systems, this position shift is very small and requires at least the system to be in the so-called ultra-strong coupling regime $\hat{7}$ to be measurable, i.e. $g_m/\Omega > 1$.

In spite of the quantum size of the opto-mechanical transducer, the excited phonon rise time can be fast and the maximum generated phonon number may be impressively large, in particular as compared to the thermal phonon field. In the following we propose to exploit these properties to perform QND readout of a quantum state, using a technique inspired from electron shelving used in trapped ions experiments $[12]$. The quantum emitter is the three-level quantum system shown in fig.1 where the two states to be readout are $|g\rangle$ and $|m\rangle$. Like previously, the quantum system is driven with a laser field resonant with the $|e\rangle - |g\rangle$ transition, modulated at the mechanical frequency $\Omega/2\pi$. If the quantum system is initially in the state $|m\rangle$, it is decoupled from the optical excitation and the phonon field remains in a thermal state of temperature $T$. If the quantum system is in $|g\rangle$, fluorescence cycles take place, inducing efficient build up of a coherent field in the mechanical mode. A convenient figure of merit of the measurement is the signal to noise ratio (SNR), which in this very case corresponds to the pumped phonon population divided by the thermal one. Typically, the state of the qubit is measurable with 50% confidence for $\text{SNR} = 1$ (99% for $\text{SNR} = 100$). The SNR is plotted versus time in fig.4d and 4e for two different systems. In the first one, we model the behavior of a Josephson circuit coupled to a vibrating piezoelectric membrane, using the experimental parameters of $[13]$. The three levels of our quantum emitter correspond to the lowest levels of a superconducting ar-
while the narrow peak on top is the pumped phonon field. The broad pedestal is due to the thermal noise, measured on a spectrum analyzer with resolution bandwidth of 1MHz. The pedestal is due to thermal phonons, while the narrow peak (the coherence of the phonon field matches that of the laser modulation) corresponds to the optically excited phonons. The latter is found to be four orders of magnitude brighter than the thermal phonons within a frequency binning of 1MHz. The excited phonons are four orders of magnitude brighter than the thermal's within a 0.1kHz frequency channel, corresponding to a macroscopic average phonon number $\bar{n} = 258$. As shown in Fig c, the SNR grows fast, reaching 100 in 50μs despite the large population of the thermal field. This is way faster than the relaxation time of the transition, which can be as large as $T_1 \sim 1$ms for ultra pure diamond crystals and opens the path to projection QND readout of the NV center spin states in the single shot regime. Note that the current way to detect the spin states is based on optical pumping and is therefore destructive, so that the scheme we propose represents a major progress with respect to state of the art protocols.

In this letter, we have shown that a single quantum emitter embedded in a hybrid system could be turned into a surprisingly efficient opto-mechanical transducer with single quantum scale accuracy, allowing to envision QND readout of a quantum state in the single shot regime. This regime has recently lead to ground-breaking demonstrations of quantum feedback and observation of quantum jumps with Rydberg atoms in microwave cavities or Josephson qubits in circuit QED. Such QND readout of a quantum state paves the road towards the generation of Schrödinger cat states of motion. In addition, classical phonon interferences can be exploited to manipulate the temperature of the resonator: constructive interferences, as explained in this letter allows large phonon field to build up. On the other hand, destructive interference could also be used to lower the temperature of the phonon field, by using a classical feedback scheme. Finally, the results we have derived here have been obtained in the linearized limit of eq.(6), i.e. Rabi frequency larger than the quantum emitter energy oscillation amplitude. Interesting physics might show up in the nonlinear regime, where single quantum induced optomechanical nonlinearity is expected.

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**METHODS**

**pulse trains excitation: derivation of a recursive expression for $\beta_n$**

We first derive an explicit expression which describes the dynamics of the phonon field during the $n_p$ pulses sequence. To do so, we look at a single period: at $t = 0$ the field is in a state $\beta = U_n \beta = U_0 = 0$ at the very beginning of the sequence), then the quantum emitter is excited for a time $\tau_1$. At $t = \tau_1$, $\beta = V_n = (U_n - g_m/\Omega)e^{-\alpha \tau_1} + g_m/\Omega$. Then the excitation is off for a time
and the field undergoes free evolution: \( \beta = U_{n+1} = V_n e^{-\beta t_2} \). This recursive expression can be rearranged to get an explicit expression \( U_n = b(1 - a^{n_t})/(1 - a) \) where \( n_p \) is the number of pulses in the pulse train. After the pulse train, the laser is off for half a phonon period \( \pi/\Omega \). Therefore, again, using a recursive approach, we end up with eq. [2].

**CW excitation: derivation of a recursive expression for \( \beta \)**

The technique is similar to that described above: assuming that the initial phonon field \( \beta_0 \) is known \( (\beta_0 = 0 \text{ at } t = 0) \), we look at the field dynamics in detail for an excitation period \( \phi = [0, 2\pi] \). For the first half a period, the laser is on and \( U_n(\phi) \) the solution of eq. [2] with initial condition \( \beta_n \) is propagated until \( \phi = \pi \). Then the laser is switched off again. \( U_n(\pi) \) then undergoes free evolution for another half a period, i.e. until \( \phi = 2\pi \). The resulting field is \( \beta_{n+1} \), in this way, one gets a simple recursive relations connecting the phonon field at phase \( \phi = 0 \) of the excitation modulation to \( \phi = 2\pi \).

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