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Phys. Rev. Lett. **114**, 113601 — Published 16 March 2015

DOI: 10.1103/PhysRevLett.114.113601

Proposal for an optomechanical microwave sensor at the sub-photon level

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Due to their low energy content microwave signals at the single-photon level are extremely challenging to measure. Guided by recent progress in single-photon optomechanics and hybrid optomechanical systems, we propose a multimode optomechanical transducer that can detect intensities significantly below the single-photon level via adiabatic transfer of the microwave signal to the optical frequency domain where the measurement is then performed. The influence of intrinsic quantum and thermal fluctuations is also discussed.

PACS numbers: 07.57.Kp, 42.50.Lc, 07.10.Cm, 42.50.Wk

Introduction. The microwave frequency domain of the electromagnetic spectrum is the stage of a wealth of phenomena, ranging from the determination of the quantum energy levels of superconductor nanostructures to the rotational modes of molecules and to the characterization of the cosmic microwave background. Several detection schemes sensitive to microwave radiation at the single-photon level have been demonstrated. Examples include semiconductor quantum dots in high magnetic field [1], circular Rydberg atoms in cavity QED setups [2–4], and superconducting qubits in circuit QED [5, 6]. An alternative approach involves the use of linear amplifiers [7]. These devices allow the reconstruction of average amplitudes [5] and correlation functions [8] and may operate both as phase-preserving (insensitive) [9] and phase-sensitive [10, 11] amplifiers, but they require an integration over many events to achieve a sizable signal.

Even though there have been proposals and experiments to realize a photon-multiplier in the microwave regime [12–14], no general purpose efficient single photon detector has been developed so far, as photon energies in that frequency domain are in the millielectronvolt range, three orders of magnitude smaller than in the visible or near-infrared spectral regions. On the other hand, in the optical frequency domain a variety of ultra-sensitive detectors have been developed over the past sixty years. This suggests that an alternative route for the detection of feeble microwave signals is via their conversion to the optical frequency domain. Photonic front-end microwave receivers based on the electro-optical effect [15] and atomic interfaces based on electromagnetically induced transparency have exploited non-linear conversion to this end [16, 17]. The main limitations in sensitivity is the small strength of the interaction and the fluctuations of the optical driving fields.

Recent advances in nano- and optomechanics offer an attractive approach to engineer interactions of light and mechanics that achieve that goal via the radiation pressure force, see Ref. [18] for recent reviews. Several theoretical proposals have considered the optomechanically mediated quantum state transfer between microwave and optical fields [19–22] and have emphasized the potential of hybrid systems as quantum information interfaces [23–26], in which case state transfer fidelity is of particular interest. Developments of particular relevance include the experimental realization of coherent conversion between microwave and optical field based on a hybrid optomechanical setup [27–29]. The present work has the different goal to convert the *mean intensity* of a feeble, narrowband microwave signal to a signal at an optical frequency where detection can proceed by traditional methods.

One key aspect of this proposed detector is that it relies on an off-resonant, multimode process. This is motivated by the need to manage and minimize the thermal mechanical noise, as well as to circumvent the effect of the fluctuations of the driving electromagnetic fields required to ensure a strong enough optomechanical coupling. These sources of noise can be significantly reduced by (i) working in a far off-resonant regime with respect to the mechanics; (ii) using pumping fields that drive ancillary cavity modes different from those at the signal frequencies, for both microwave and optical; and (iii) exploiting the polariton modes of the cavity-mechanics system to perform the frequency conversion of the signal via a modulation of the detuning of the optical pump.

The system. The proposed sensor is comprised of a mechanical oscillator optomechanically coupled to both a microwave and an optical multimode resonator, see idealized setup in Fig. 1.

Consider first the microwave cavity. To avoid the noise connected with the pumping field while still maintaining a large optomechanical coupling strength, we adopt a multimode configuration where a strong optomechanical coupling is provided by an auxiliary field at frequency ω_{bp} different from that of the signal to be detected, see Fig. 1(b). This three-mode optomechanical interaction is described by the Hamiltonian [30–32]

$$V_{3m} = \hbar g_{b0}(\hat{b}_p + \hat{b})^\dagger(\hat{b}_p + \hat{b})(\hat{c} + \hat{c}^\dagger), \tag{1}$$

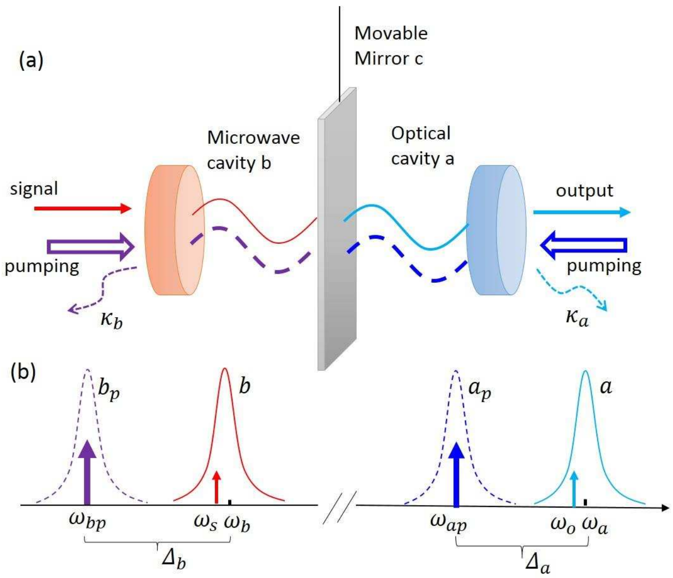

Figure 1: (a) Dual-cavity optomechanical system. (b) Sketch of the heterodyne-like pumping scheme with the microwave signal and the driving field near-resonant with cavity mode b and ancilliary mode b_p, respectively. Similarly in the optical side.

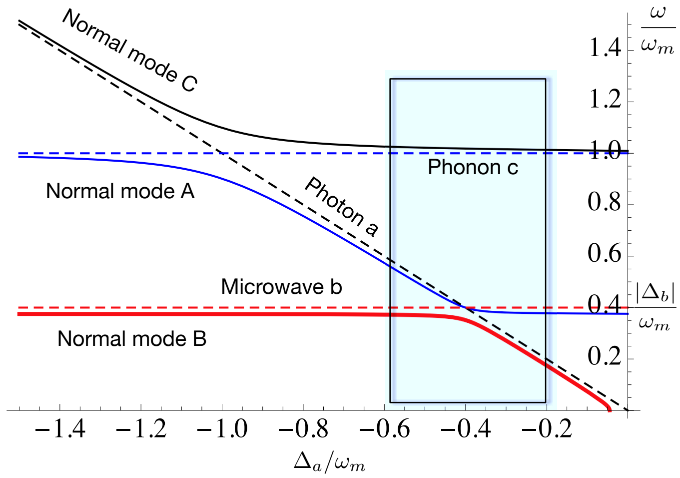

Figure 2: Eigenfrequencies of the normal modes (polaritons) as functions of optical detuning Δ_a/ω_m for the case $-G_a/\omega_m = G_b/\omega_m = 0.1$ and $\Delta_b/\omega_m = -0.4$. Dashed lines: non-interacting energies of the bare modes. We have framed the part of the spectrum spanned by Δ_a during the conversion process.

where g_{b0} is the single microwave photon optomechanical coupling constant. We assume that ω_{bp} is resonant with a longitudinal cavity mode, while the signal field $\hat{b}$, assumed to be extremely weak, is slightly detuned from another mode of frequency ω_b. In the displaced picture for $\hat{b}_p$ and $\hat{c}$, $\hat{b}_p \to \beta_p + \hat{b}_p$ and $\hat{c} \to \mathcal{C} + \hat{c}$, the Hamiltonian (1) becomes

$$\begin{aligned} V_{3m,\text{eff}} &= \hbar G_b(\hat{b} + \hat{b}^\dagger)(\hat{c} + \hat{c}^\dagger) + \hbar x_c g_{b0}(\hat{b}_p \hat{b}^\dagger + \hat{b}_p^\dagger \hat{b}) \\ &+ \hbar G_b(\hat{b}_p + \hat{b}_p^\dagger)(\hat{c} + \hat{c}^\dagger). \end{aligned} \quad (2)$$

The first term is the usual linearized optomechanical coupling between the signal mode $\hat{b}$ and phonon mode $\hat{c}$ with strength $G_b = \beta_p g_{b0}$. We assume that the pump field is phase locked so that G_b is real and positive. Its fluctuations feed into the system as noise through the second and the third terms of $V_{3m,\text{eff}}$ which arise from the scattering and the optomechanical coupling of the pumped mode, respectively. The second term, proportional to the steady position quadrature of the phonon field, $x_c = \mathcal{C} + \mathcal{C}^*$, can be safely neglected under the condition $|x_c| \ll |\beta_p|$ which is easily realized [33] in the mirror-in-the-middle geometry of Fig. 1. Finally, the third term results in contributions to the system dynamics at a frequency that differs from the first term by $\pm(\omega_b - \omega_{bp})$. This difference is of the order of the free spectral range of the cavity (for longitudinal modes) so that it can easily be filtered out in a manner familiar from heterodyne detection. For the narrow band detection scheme considered here it is therefore sufficient to keep only the first term in the Hamiltonian (2).

Following a similar argument for the optical fields the effective Hamiltonian for the full system becomes

$$\begin{aligned} H &= \hbar\omega_m \hat{c}^\dagger \hat{c} - \hbar\Delta_a \hat{a}^\dagger \hat{a} - \hbar\Delta_b \hat{b}^\dagger \hat{b} \\ &+ \hbar G_a(\hat{a} + \hat{a}^\dagger)(\hat{c} + \hat{c}^\dagger) + \hbar G_b(\hat{b} + \hat{b}^\dagger)(\hat{c} + \hat{c}^\dagger), \end{aligned} \quad (3)$$

where $\hat{a}$, $\hat{b}$, and $\hat{c}$ are the annihilation operators for the optical, microwave, and (displaced) mechanical modes with corresponding frequencies ω_a, ω_b, and ω_m. The optical and microwave cavity-pump detunings are $\Delta_a = \omega_{ap} - \omega_a + x_c g_{a0}$ and $\Delta_b = \omega_{bp} - \omega_b + x_c g_{b0}$, respectively, with ω_{ap} and ω_{bp} the frequencies of the optical and microwave pumps. G_a and G_b are the effective optomechanical coupling strength set by the steady amplitude of the pumped ancillary optical and microwave cavity modes. Note that $G_{a,b}$ are of opposite signs and the equilibrium position of the mechanical resonator is set by the relative strength of the two pumps, so that the microwave drive needs to have a significantly stronger light flux than the optical pump.

In the resonant situation $\Delta_a = \Delta_b = -\omega_m$ an effective interaction follows from performing the rotating wave approximation, which gives $H_I = \hbar G_a(\hat{a}\hat{c}^\dagger + \hat{c}\hat{a}^\dagger) + \hbar G_b(\hat{b}\hat{c}^\dagger + \hat{c}\hat{b}^\dagger)$. If G_a and G_b are appropriately modulated in time the system then adiabatically follows a superposition of cavity modes $\hat{a}$ and $\hat{b}$ without any population of the mechanical mode $\hat{c}$ (dark mode) [19, 20]. In contrast, for the off-resonant case considered here, $\Delta_{a,b} \neq \omega_m$, the microwave and optical fields are coupled by a three-level Raman-like interaction via the mechanical mode.

Normal mode picture. To discuss the microwave to optical conversion process in this effective three-mode configuration, it is convenient to switch to a normal mode (polariton) representation of the system [34]. After removing a constant term, the Hamiltonian (3) can be recast in the diagonal form $H = \hbar\omega_A \hat{A}^\dagger \hat{A} + \hbar\omega_B \hat{B}^\dagger \hat{B} + \hbar\omega_C \hat{C}^\dagger \hat{C}$ where $\hat{A}$, $\hat{B}$, and $\hat{C}$ are the boson annihilation operators for the normal mode excitations. In general, these are superpositions of the optical, microwave, and mechanical modes. Figure 2 shows their frequencies $\omega_{A,B,C}$ as functions of the optical detuning Δ_a. At the mechanical resonance, $\Delta_a = -\omega_m$, the degeneracy between the optical photon and the phonon is lifted by the optomechanical interaction, with an energy splitting of the order of $2G_a$. A second avoided crossing occurs at the resonance between optical and microwave photons, $\Delta_a = \Delta_b$, with a splitting of the order of $4G_aG_b/\omega_m$ resulting from the indirect coupling between the electromagnetic modes via the mechanical mode.

We focus on the region close to the microwave-optical resonance framed in Fig. 2. On the left side, $\Delta_a < \Delta_b$ and $|\Delta_a - \Delta_b| \gg 4|G_aG_b|/\omega_m$ the polariton $\hat{B}$ describes a microwave-like excitation, with $\omega_B \sim -\Delta_b$ and $\hat{B} \sim \hat{b}$, while for $\Delta_a > \Delta_b$, the polariton becomes optical-like $\hat{B} \sim \hat{a}$ and annihilates an excitation of frequency $\omega_B \sim -\Delta_a$. The opposite holds for the polariton A, which is optical-like for $\Delta_a < \Delta_b$ and microwave-like on the other side of the resonance. The polariton C remains phonon-like in this whole region, indicating that the dynamics of the mechanical excitation is decoupled from that of the electromagnetic fields.

Conversion process. When Δ_a is slowly switched from the left to the right side of the resonance the polariton B adiabatically evolves from the microwave-like excitation to the optical-like excitation while conserving its population, $\langle \hat{B}^\dagger(t)\hat{B}(t)\rangle \approx \langle \hat{b}^\dagger(t_0)\hat{b}(t_0)\rangle$, where $\langle \hat{b}^\dagger(t_0)\hat{b}(t_0)\rangle$ accounts for both the input signal field to be measured and the microwave cavity noise. Likewise the polariton A, which is initially optical-like, evolves into a microwave-like excitation while maintaining its population $\langle \hat{A}^\dagger(t)\hat{A}(t)\rangle \approx \langle \hat{a}^\dagger(t_0)\hat{a}(t_0)\rangle = 0$, where the last equality holds if the optical mode is initially in a vacuum, a condition easy to satisfy.

The adiabaticity of the transfer requires that Δ_a be switched at a rate much slower than the interband separation, $1/\tau \ll 4|G_aG_b|/\omega_m$ where τ is the switching time. In addition it is also necessary that this operation occurs in a time short compared to the inverse decay rates of the polariton modes, which are combinations of the cavity decay rates $\kappa_{a,b}$ and the mechanical damping rate γ. (This condition also ensures that α and β remain constant during the switch of Δ_a.)

We describe the detection protocol as a time-gated three-step process. First, during a "receiving" time window τ_r that lasts until t_0 the optical detuning is fixed at $\Delta_a < \Delta_b$ with $|\Delta_a - \Delta_b| \gg 4|G_aG_b|/\omega_m$ and the microwave cavity captures a narrow-band signal that is stored in the mode b. During that time the optical mode a is in a vacuum and the microwave-optical field interaction is negligible due to their large mismatch in frequency. This is followed by a "transfer" time interval τ starting at t_0 during which Δ_a is switched to Δ_b at a rate

$$1/\kappa_{a,b} \gg \tau \gg \omega_m/4|G_aG_b|, \tag{4}$$

resulting in the signal being transferred into an optical field without any significant coupling to the external reservoirs. Finally, during the detecting time window $\tau_d > t_0 + \tau$ the interaction is quenched and the cavities couple with their environment, thus releasing the optical output field that can be measured by standard methods.

Input-output dynamics. The analysis of the conversion of the microwave signal to the optical field can be performed in terms of Heisenberg-Langevin equations of motion $\partial_t \hat{u} = -i[\hat{u}, \hat{H}]/\hbar - \kappa_u \hat{u} + \sqrt{2\kappa_u}\hat{u}_{\rm in}$ where $\hat{u}$ are the annihilation operators for the bare modes $\{\hat{a}, \hat{b}, \hat{c}\}$, κ_u are their dissipation rates (with $\kappa_c \equiv \gamma$), and $\hat{u}_{\rm in}$ account for the associated noise operators and input fields. In the absence of input fields the non-vanishing noise correlations are $\langle \hat{u}_{\rm in}(t)\hat{u}^\dagger_{\rm in}(t')\rangle = (\bar{n}_u + 1)\delta(t-t')$ and $\langle \hat{u}^\dagger_{\rm in}(t)\hat{u}_{\rm in}(t')\rangle = \bar{n}_u\delta(t-t')$, where $\bar{n}_u = 1/[\exp(\hbar\omega_u/k_BT_u) - 1]$, T_u being the temperature of the thermal reservoir of mode u. For the optical field $\bar{n}_a \approx 0$ in practice.

In the far off-resonant case $\omega_m \gg |\Delta_{a,b}|, |G_{a,b}|, \kappa_{a,b}, \gamma$ we adiabatically eliminate the phonon mode $\hat{c}$ by inserting its formal solution $\hat{c} \approx [-G_a(\hat{a} + \hat{a}^\dagger) - G_b(\hat{b} + \hat{b}^\dagger)]/\omega_m$ into the equations for the modes a and b while retaining the mechanical noise term and neglecting the memory effect. The interaction between the microwave and optical modes is then described by the equation

$$\partial_t \hat{a} = (i\Delta'_a - \kappa_a)\hat{a} + i\frac{2G_a^2}{\omega_m}\hat{a}^\dagger + iG'(\hat{b} + \hat{b}^\dagger) + \sqrt{2\kappa_a}\hat{a}'_{\rm in}, \tag{5}$$

where $G' = 2G_aG_b/\omega_m$, and similarly for mode b with $a \leftrightarrow b$ [33].

In the far off-resonant case, we must keep the anti-rotating terms in the optomechanical interaction when adiabatically eliminating the mechanics. This results in a squeezing contribution to the dynamics of a and b with the original detuning becoming $\Delta'_{a,b} = \Delta_{a,b} + 2G^2_{a,b}/\omega_m$ and

$$\hat{a}'_{\rm in} = \hat{a}_{\rm in} - iG_a\sqrt{\frac{\gamma}{\kappa_a}}\left[\int_0^t e^{(-i\omega_m - \gamma)(t-t')}\hat{c}_{\rm in}(t')dt' + h.c.\right], \tag{6}$$

and similarly for $b'_{\rm in}$ with $a \to b$. When we focus on the signal fields of narrow linewidth around cavity modes, the noise auto-correlation functions approximately become $\langle \hat{a}'_{\rm in}(t)\hat{a}'^\dagger_{\rm in}(t')\rangle = (\bar{n}_a + m_a + 1)\delta(t-t')$ and $\langle \hat{a}'_{\rm in}(t)\hat{a}'_{\rm in}(t')\rangle = -m_a\delta(t-t')$, with $m_a = (G_a^2\gamma/\omega_m^2\kappa_a)(2\bar{n}_c + 1)$, with also the appearance of cross-correlations characteristic of a squeezed two-mode reservoir, $\langle \hat{a}'_{\rm in}(t)\hat{b}'^\dagger_{\rm in}(t')\rangle = m_{ab}\delta(t-t')$ and $\langle \hat{a}'_{\rm in}(t)\hat{b}'_{\rm in}(t')\rangle = -m_{ab}\delta(t-t')$ where $m_{ab} = (G_aG_b\gamma/\omega_m^2\sqrt{\kappa_a\kappa_b})(2\bar{n}_c + 1)$ [33]. The output fields are similarly modified, with the indices "in" replaced by "out" and $\hat{c}_{\rm out} = -\hat{c}_{\rm in}$ in this far off-resonant case. Note that the weak coupling assumption $|G_{a,b}|/\omega_m \ll 1$, which allows the adiabatic elimination of the mechanical mode, also implies small values for the squeezing parameters m_a, m_b, and m_{ab}.

The polariton operators $\hat{A}$, $\hat{B}$ and their corresponding noise operators $\hat{A}_{\rm in}$, $\hat{B}_{\rm in}$ are readily obtained via a Bogoliubov transformation of the bare modes in the absence of dissipation. Assuming for simplicity $\kappa_a = \kappa_b = \kappa$, one then finds readily [35]

$$\partial_t \hat{A} = (i\omega_A - \kappa)\hat{A} + \sqrt{2\kappa}\hat{A}_{\rm in}, \tag{7}$$

and similarly for mode B, with $A \to B$.

Determining the conversion between the microwave signal and the optical field requires in general to solve the full Heisenberg-Langevin equations with time-dependent coefficients. But if one assumes perfect adiabaticity one can use instead a much simplified effective two-sided cavity model. To single out the effect of the varying frequencies $\omega_{A,B}(t)$ we focus on the slowly varying envelopes $\widetilde{A} = \hat{A}e^{-i\omega_A t}$ and $\widetilde{B} = \hat{B}e^{-i\omega_B t}$. We also introduce a new operator for the symmetric superposition of the cavity modes, $\hat{V} = (\widetilde{A} + \widetilde{B})/\sqrt{2}$. From Eq. (7) we then have

$$\partial_t \hat{V} = -\kappa\hat{V} + \sqrt{\kappa}\widetilde{A}_{\rm in} + \sqrt{\kappa}\widetilde{B}_{\rm in}, \tag{8}$$

reminiscent of the situation of a two-sided cavity [36] but with input field operators depending on Δ_a. Specifically in the first stage of the detection sequence, $t < t_0$, we have $\widetilde{A}_{\rm in} \approx \hat{a}'_{\rm in}e^{i\Delta'_a t}$ and $\widetilde{B}_{\rm in} \approx \hat{b}'_{\rm in}e^{i\Delta'_b t}$, while in the third step, $t > t_0 + \tau$, $\widetilde{A}_{\rm in}$ and $\widetilde{B}_{\rm in}$ are simply exchanged. In the intermediate second step, the adiabatic, essentially dissipation-free, evolution results in small phase shifts for the envelope operators, proportional to $\partial_t\omega_A$ and $\partial_t\omega_B$ for $\widetilde{A}$ and $\widetilde{B}$, respectively. In case of perfect adiabaticity we may neglect these shifts and thus obtain $\hat{V}(t_0) = \hat{V}(t_0 + \tau)$ [33].

Summarizing, the full evolution of $\hat{V}$ for the three-step detection sequence is approximately described by the equation

$$\partial_t \hat{V} = -\kappa\hat{V} + \sqrt{\kappa}\hat{a}'_{\rm in}e^{i\Delta'_a t} + \sqrt{\kappa}\hat{b}'_{\rm in}e^{i\Delta'_b t}. \tag{9}$$

With the boundary conditions of the two-sided cavity, $\hat{a}'_{\rm out}e^{i\Delta'_a t} + \hat{a}'_{\rm in}e^{i\Delta'_a t} = \sqrt{\kappa}\hat{V}$ and $\hat{b}'_{\rm out}e^{i\Delta'_b t} + \hat{b}'_{\rm in}e^{i\Delta'_b t} = \sqrt{\kappa}\hat{V}$ [36], this equation can be solved in the frequency domain to give

$$\hat{a}'_{\rm out}(\omega - \Delta'_a) = \frac{\kappa\hat{b}'_{\rm in}(\omega - \Delta'_b) - i\omega\hat{a}'_{\rm in}(\omega - \Delta'_a)}{\kappa + i\omega}. \tag{10}$$

Perfect conversion, $\hat{a}'_{\rm out}(-\Delta'_a) = \hat{b}'_{\rm in}(-\Delta'_b)$, occurs for $\omega = 0$. Remembering that the optical and the microwave operators are expressed in rotating frames with respect to the pumping frequencies ω_{ap} and ω_{bp}, this corresponds to the case where the frequency of the input microwave fields is $\omega_s = \omega_b - x_c g_{b0} - 2G_b^2/\omega_m$ and the frequency of the output optical field is $\omega_o = \omega_a - x_c g_{a0} - 2G_a^2/\omega_m$.

We introduce the mean photon numbers of the optical and microwave modes

$$\begin{aligned} \bar{n}_o &= \int d\omega |g(\omega)|^2 \langle \hat{a}^\dagger_{\rm out}(\omega - \Delta'_a)\hat{a}_{\rm out}(\omega - \Delta'_a)\rangle \\ \bar{n}_s &= \int d\omega |g(\omega)|^2 \langle \hat{b}^\dagger_{\rm in}(\omega - \Delta'_b)\hat{b}_{\rm in}(\omega - \Delta'_b)\rangle, \end{aligned} \tag{11}$$

where the mode filter functions $g(\omega)$ are sharply peaked around $\omega = 0$. By assuming detection and reception time windows $(\tau_d, \tau_r) \gg 1/\kappa$ [37, 38] we find

$$\bar{n}_o = \bar{n}_s + \frac{(G_b^2 + G_a^2)\gamma}{\omega_m^2\kappa}(2\bar{n}_c + 1), \tag{12}$$

where we have taken into account the modified noise correlation of the optical and microwave cavities, and the effects of the mechanical noise are merged into the second term on the right side. This is the central result of this paper.

Sensitivity. Ignoring technical noise and assuming that the final optical detector is well characterized and has near unit quantum efficiency, we concentrate on the intrinsic sensitivity of the three-step conversion sequence. It is characterized primarily by the microwave to optical conversion efficiency, the effects of quantum and thermal noise, and the dead time required to reset the resonators between measurements. Perfect adiabatic conversion requires interaction times $\kappa \ll 1/\tau \ll 4|G_aG_b|/\omega_m \ll \omega_m$, and the dead time to reset the resonators is of the order $1/\kappa$. Quantum and thermal noise result in a dark-count rate that also impacts the figure of merit of the detector, see Eq. (12). A high-Q and ultracold mechanical oscillator can significantly suppress these sources of noise, but in the limiting case $\gamma/\kappa, \bar{n}_c \to 0$ intermode scattering as well as memory effect, which we neglected in our analysis [33], becomes a dominant source of noise.

As an example we consider an optomechanical resonator with high mechanical frequency $\omega_m = 2\pi \times 4$GHz and quality factor $Q = 87 \times 10^3$, which results in $\gamma = 2\pi \times 46$ KHz and $\bar{n}_c = 72$ for a temperature $T = 14$K [39]. Because of the large detunings considered here, we find however that the mechanical noise only adds a contribution of 0.06 to $\bar{n}_o$. The level of thermal microwave noise that feeds into $\bar{n}_s$ can be managed by cooling the microwave cavity to cryogenic temperatures. For a microwave cavity frequency $\omega_b = 2\pi \times 300$ GHz and temperature $T_b = 300$ K we have $\bar{n}_s = 20$, but for $T_b = 3$ K, $\bar{n}_s$ is reduced to 0.008. Finally we assume linear optomechanical coupling strengths $G_a = -2\pi \times 200$ MHz and $G_b = 2\pi\times 300$ MHz, respectively, giving an effective interaction strength $2G_aG_b/\omega_m = -2\pi\times 30$ MHz. We also set the same decay rate for both cavities, $\kappa = 2\pi \times 850$ kHz. These parameters fulfill the condition for adiabaticity of the conversion and result in a dead time of the order of

100 ns. These estimates indicate that the detector should be able to operate reliably at or below the single photon level.

Conclusion. We have proposed and analyzed a time-gated microwave detection scheme based on the control of polaritons in a hybrid optomechanical system. In contrast to resonant schemes that focus on high fidelity quantum state transfer, [19–22], the dual optomechanical cavity detector is driven by a heterodyne-like pumping and operates on the far-off sideband resonant regime to minimize pump and mechanical noise, thereby offering the potential to reliably detect very feeble microwave fields. Importantly that non-resonant approach does not preserve the quantum state of the microwave field. Rather, it detects the signal entering the microwave resonator in a time determined by its decay time $1/\kappa_b$ just before transfer to the optical domain.

Acknowledgments. We thank S. Singh, H. Seok, P. Treutlein, H. Metcalf, R. Dehghannasiri and A. A. Eftekhar for helpful discussions. This work was supported by the National Basic Research Program of China Grant No. 2011CB921604, the NSFC Grants No. 11204084, 91436211, 11234003, and 11304072, the SRFDP Grant No. 20120076120003, the SCST Grant No. 12ZR1443400, the DARPA QuASAR and ORCHID programs through grants from AFOSR and ARO, the U.S. Army Research Office, and NSF. YD is supported in part by the Hongzhou-city Quantum Information and Quantum Optics Innovation Research Team.

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