Optomechanically-induced transparency in parity-time-symmetric microresonators

H. Jing1,2,7, Sahin K. Özdemir3, Z. Geng7, Jing Zhang1, Xin-You Li2,5, Bo Peng3, Lan Yang3 & Franco Nori3,6

1The Key Laboratory of Quantum Optics, Shanghai Institute of Optics and Fine Mechanics, Chinese Academy of Science, Shanghai 201800, China, 2CEMS, RIKEN, Saitama, 351-0198, Japan, 3Department of Automation, Tsinghua University, Beijing 100084, China, 4School of physics, Huazhong University of Science and Technology, Wuhan 430074, China, 5Physics Department, The University of Michigan, Ann Arbor, MI 48109–1040, U.S.A., 6Department of Physics, Hennan Normal University, Xinxiang 453007, China.

Recent advances in steering a macroscopic mechanical object in the deep quantum regime1–3 have motivated theoretical studies to understand the physics of photon-phonon interactions in cavity optomechanics (COM), and also led to exciting experimental studies on quantum nanodevices4–6. In particular, the experimental demonstration of OMIT allows the control of light propagation at room temperature using nano- and micro-mechanical structures7–9. The underlying physics of OMIT is formally similar to that of electromagnetically-induced transparency (EIT) in three-level A-type atoms8–10 and its all-optical analogs demonstrated in various physical systems11–13. The resulting slow-light propagation provides the basis for a wide range of applications9. Mechanically-mediated delay (slow-light) and advancement (fast-light) of microwave pulses were also demonstrated in a superconducting nanocircuit14–16. These experimental realizations offer new prospects for on-chip solid-state architectures capable of storing, filtering, or synchronizing optical light propagation.

As a natural extension of single-cavity structures, COM with an auxiliary cavity (compound COM; two cavities) has also attracted intense interest. The interplay between COM interactions and tunable optical tunneling provides a route for implementing a series of important devices, such as phonon lasers17, phononic processors for controlled gate operations between flying (optical) or stationary (phononic) qubits18–20, and coherent optical wavelength converters20–21. Enhanced nonlinearities22 and highly-efficient photon-phonon energy transfer23–25 are other advantages of the compound COM. These studies were performed with passive (lossy, without optical gain) resonators.

Very recently, an optical system whose behavior is described by PT-symmetric Hamiltonians (i.e., the commutator [H, PT] = 0)26,27 was demonstrated in a system of two coupled microresonators, one of which has passive loss and the other has optical gain (active resonator)28. Observed features include: real eigenvalues in the PT-symmetric regime despite the non-Hermiticity of the Hamiltonian, spontaneous PT-symmetry breaking, as well as complex eigenvalues and field localization in the broken PT-symmetry regime. Moreover, nonreciprocal light transmission due to enhanced optical nonlinearity in the broken PT-symmetry regime was demonstrated29. Such a PT-symmetric structure provides unique and previously-unattainable control of light and even sound26–32. Manipulating the photon-phonon interactions in such systems opens new regimes for phonon lasing and quantum COM control33.

In this paper, we show that a compound COM with PT-symmetric microresonators leads to previously unobserved features and provides new capabilities for controlling light transmission in micro- and nano-mechanical systems. Particularly, we show: (i) a gain-induced reversed transparency (inverted-OMIT), i.e. an optical spectral dip between two strongly-amplifying sidebands, which is in contrast to the non-absorptive peak between strongly absorptive sidebands in the conventional passive OMIT; (ii) a reversed pump dependence of the optical...
transmission rate, which is most significant when the gain and loss are balanced (i.e., optical gain in one subsystem completely compensates the loss in the other); and (iii) a gain-controlled switching from slow (fast) light to fast (slow) light in the PT-symmetric (PT-breaking) regime, within the OMIT window. These features of the active OMIT enable new applications which are not possible in passive COM.

The inverted-OMIT observed here in an active COM, composed of a passive and an active optical microresonator, is reminiscent of the inverted-OMIT observed in all-optical systems, composed of one passive and an active optical microresonator, 15. Distinct features of the inverted-OMIT that cannot be observed in the optical inverted-OMIT are also revealed in the broken-PT-symmetric regime.

Results

The active COM system. We consider a system of two coupled whispering-gallery-mode microtoroid resonators 14,18,33,34,36. One of the resonators is passive and contains a mechanical mode of frequency \( \omega_m \) and an effective mass \( m \). We refer to this resonator as the optomechanical resonator. The second resonator is an active resonator which is coupled to the first one through an evanescent field. The coupling strength \( J \) between the resonators can be tuned by changing the distance between them. As in Ref. 28, the active resonator can be fabricated from Er \(^{3+}\)-doped silica and can emit photons in the 1550 nm band, when driven by a laser in the 980 or 1450 nm bands. The resonators can exchange energies only in the emission band of 1550 nm, so the gain photons can tunnel through the air gap between the resonators and provide a gain \( \kappa \) to compensate the optical loss \( \gamma \) in the passive resonator 14,33,36.

Tuning the gain-to-loss ratio, while keeping \( J \) fixed, leads to two remarkably distinct regimes, i.e. broken- and unbroken-PT-symmetry regimes, that are characterized by distinct normal mode-splitting and linewidths 14,33,34. Our aim here is to study OMIT in these two distinct regimes, focusing on the role of \( \kappa/\gamma \). To this end, as in the conventional OMIT, both a pump laser of frequency \( \omega_p \) and a weak probe light of frequency \( \omega_p \) are applied (see Fig. 1). The field amplitudes of the pump and probe are given by \( \hat{E}_p = (2P_p/\hbar \omega_p)^{1/2} \) and \( \hat{E}_l = (2P_l/\hbar \omega_m)^{1/2} \), where \( P_p \) and \( P_l \) are the pump and probe powers.

The Hamiltonian of this three-mode COM system can be written as

\[
H = H_0 + H_{\text{int}} + H_{dt},
\]

\[
H_0 = \hbar \Delta_L \left( a_1^+ a_1 + a_2^+ a_2 \right) + \frac{\hat{p}^2}{2m} + \frac{1}{2} \hbar \omega_m \hat{x}^2,
\]

\[
H_{\text{int}} = -\hbar J \left( a_1^+ a_2 + a_2^+ a_1 \right) - \hbar g a_1^+ a_1 \hat{x},
\]

\[
H_{dt} = i\hbar \left( E_L a_1^+ - E_L a_1 + i \gamma a_1^+ e^{-i\omega t} - i \gamma a_1 e^{i\omega t} \right),
\]

where \( a_1 \) and \( a_2 \) denote the annihilation operators of the bosonic fields in the microresonators with resonance frequency \( \omega_m \) and radius \( R \), \( g = \omega_m/\hbar \) is the COM coupling rate, \( x = x_0 (b + b^\dagger) \) is the mechanical position operator, \( \omega_0 = \hbar / (2 \mu m_0) \), and \( b = \mu \) corresponds to the annihilation operator for the phonon mode. The pump-resonator, probe-resonator, and probe-pump frequency detunings are, respectively, denoted by \( \Delta_L = \omega_0 - \omega_L \), \( \Delta_P = \omega_P - \omega_L \), \( \xi = \omega_P - \omega_L \).

The Heisenberg equations of motion (EOM) of this compound system are (\( \hbar = 1 \))

\[
\dot{x} + \Gamma_m \dot{x} + \omega_m^2 x = \frac{g}{m} a_1^+ a_1,
\]

\[
\dot{a}_1 = (-i \Delta_L + ig x - \gamma) a_1 + i J a_2 + E_L + \epsilon_p e^{-i\omega t},
\]

\[
\dot{a}_2 = (-i \Delta_L + \kappa) a_2 + i J a_1 + \epsilon_p e^{-i\omega t}.
\]

For \( \Delta_L = 0 \), by choosing \( J = k/\gamma \) or \( k/\gamma \rightarrow 1 \) for \( \gamma = 1 \), one can identify a gain-induced transition from the linear to the nonlinear regime that significantly enhances COM interactions 14. Here we focus on the effects of the gain-loss balance on the OMIT and the associated optical group delay, which, to our knowledge, has not been studied previously.

We proceed by expanding each operator as the sum of its steady-state value and a small fluctuation around that value, i.e. \( a_1 = a_{1s} + \delta a_1 \), \( a_2 = a_{2s} + \delta a_2 \), \( x = x_s + \delta x \). After eliminating the steady-state values, we obtain the linearized EOM, which can be solved using the ansatz (see the Method)

\[
\begin{pmatrix}
\delta a_1 + \\
\delta a_2 + \\
\delta x +
\end{pmatrix}
= \begin{pmatrix}
\delta a_1 + \\
\delta a_2 + \\
\delta x +
\end{pmatrix} e^{-i\omega t} + \begin{pmatrix}
\delta a_1 - \\
\delta a_2 - \\
\delta x -
\end{pmatrix} e^{i\omega t}.
\]

The optical fluctuation in the optomechanical resonator \( A = \delta a_{1s} \), the quantity of interest here, is

![Figure 1](http://www.nature.com/scientificreports/)

**Figure 1** | OMIT in active-passive-coupled micro-resonators, with a tunable gain-loss ratio. The similarity of the passive OMIT and the three-level EIT is well-known 14; in parallel, the active OMIT provides a COM analog of the optical inverted-OMIT 14 (see the energy levels with an input gain). Here \( |0\rangle = |n_1, n_2, n_m\rangle, |1_m\rangle = |n_1 + 1, n_2, n_m\rangle, |1_L\rangle = |n_1, n_2, n_m + 1\rangle \), while \( n_1, n_2 \), and \( n_m \) denote the number of photons and phonons, respectively.
transmission dip and two sideband peaks, to the inverted-OMIT profile, quantified by a

$A = \frac{[(\omega_m^2 - i^2 + 2i)G_m + ig^2n_1\mu_1 \mu + \epsilon_p]}{(\omega_m^2 - i^2 + 2i)G_m + ig^2n_1(2G_m + \epsilon_1 + \epsilon_2 + \mu_1 + \mu_2)}.$  \hfill (9)

where $n_1 = |a_1, 1|$, is the intracavity photon number of the passive resonator, $\mu_2 = -\kappa - i\xi \pm \Delta_L$, and

$G_1 = (i\Delta_L + \gamma - ig\xi - i\xi)\mu_2 + j^2,$

$G_2 = (-i\Delta_L + \gamma + ig\xi - i\xi)\mu_2 + j^2.$  \hfill (10)

The expectation value of the output field can then be obtained by using the standard input-output relation, i.e. $a_{1, out}(t) = a_{1, in} - \sqrt{2}a_{1, out}(t)$, where $a_{1, in}$ and $a_{1, out}(t)$ are the input and output field operators. Then the optical transmission rate $\eta(\omega_p)$ (i.e., the amplitude square of the ratio of the output field amplitude to the input probe field amplitude, $\eta(\omega_p) = |(a_{1, out}(t)/a_{1, in})|^2$ is

$\eta(\omega_p) = |1 - (2j/\mu_0)|A|^2.$  \hfill (11)

We computed Eq. (11) with experimentally-accessible values of the system parameters to better understand the behavior of the COM in the presence of gain and loss. These parameters are $R = 34.5 \mu m$, $\omega_1 = 1.93 \times 10^3$ GHz, $\omega_m = 2\pi \times 23.4$ MHz, $m = 5 \times 10^{-11}$ kg, $\gamma = 6.43$ MHz and $\Gamma_m = 2.4 \times 10^7$ Hz. The quality factors of the optical mode and the mechanical mode in the passive resonator are $Q = 3 \times 10^5$ and $2Q_m/Q = 10^{-3}$, respectively. Also $\Delta_L = \omega_{1, 0}$ and thus $\Delta = \omega_p - \omega_1 = \xi - \omega_{1, 0}$. Now we discuss how the gain-loss ratio $\kappa/\gamma$, the coupling strength $J$, and the pump power $P_0$ affect the OMIT. Note that $\kappa/\gamma$ and $J$ are the tunable system parameters that allow one to operate the system in the broken- or unbroken-PT regimes.

Reversed-gain dependence. Figure 2 depicts the effect of $\kappa/\gamma$ on the optical transmission rate. By introducing gain into the second mirror resonator, one can tune the system to transit from a conventional OMIT profile, quantified by a transparency window and two sideband dips, to the inverted-OMIT profile, quantified by a transmission dip and two sideband peaks (see Fig. 2a).

Increasing the loss ratio $\kappa/\gamma < 0$ in the passive-passive COM leads to shallower sidebands. When the amount of gain provided to the second resonator supersedes its loss and the resonator becomes an active one (amplifying resonator), increasing $\kappa/\gamma > 0$ helps to increase the heights of the sideband peaks until $\kappa/\gamma = 1$, where $\eta$ at sidebands is maximized. Increasing the gain further leads to the suppression of both the sideband peaks and the on-resonance ($\Delta_p = 0$) transmission (Fig. 2b). This is in stark contrast with the observation of monotonically-increasing sideband peaks in the all-optical EIT system of Ref. 34.

This can be intuitively explained as follows. Under the condition of $J/\gamma = 1$, the system is in the PT-symmetric phase for $\kappa/\gamma < 1$, whereas it is in the broken-PT phase for $\kappa/\gamma > 1$. Thus, for $\kappa/\gamma < 1$, the provided gain compensates a portion of the losses, which effectively reduces the loss in the system and hence increases $\eta$. Increasing the gain above the phase transition point $\kappa/\gamma = 1$ puts the system in the broken-PT phase, with a localized net loss in the passive resonator (i.e., the field intensity in the passive resonator is significantly decreased) which reduces the strength of the COM interactions and hence the value of $\eta$. The reduction of the transmission by increasing the gain provides a signature of the PT-breaking regime, and it is very similar to a recent experiment with two coupled-resonators where it was shown that increasing (decreasing) the loss of one of the resonator above (below) a critical level increases (decreases) the intracavity field intensity of the other, enhancing (suppressing) transmission. Note that increasing (decreasing) loss is similar to decreasing (increasing) gain. We conclude here that only in the PT-symmetric regime ($\kappa/\gamma < 1$, with $J/\gamma = 1$), the active OMIT can be viewed as an analogous of the optical inverted-EIT.

Reversed-pump dependence. For the passive-passive COM, the transmission rate and the width of the OMIT window increase with increasing pump power $P_0$ by decreasing as the pump is increased from $P_0 = 10 \mu W$ to $20 \mu W$ (not shown here). Nevertheless, the sideband amplification always reaches its maximum value at the gain-loss balance (see also Fig. 2b). We note that the counterintuitive effect of reversed-pump dependence was also previously demonstrated in coupled optical systems (i.e., no phonon mode was involved) operating at the exceptional point.

In addition, we have also studied the effect of the mechanical damping on the profiles of the conventional and inverted OMIT at different values of the gain-loss ratio. We confirmed that the profiles of both the conventional and the inverted OMIT are strongly

![Figure 2](image-url)
affected (i.e., tend to disappear) by increasing the mechanical damping. This highlights the key role of the mechanical mode in observing OMIT-like phenomena.

**PT-breaking fast light.** The light transmitted in an EIT window experiences a dramatic reduction in its group velocity due to the rapid variation of the refractive index within the EIT window, and this is true also for the light transmitted in the OMIT window in a conventional passive optomechanical resonator. Specifically, the optical group delay of the transmitted light is given by

\[
\tau_g = \frac{d \arg [t(e_p) + \Delta_p]}{d \Delta_p} \bigg|_{\Delta_p = 0}.
\]

We have confirmed that OMIT in the passive-passive COM leads only to the slowing (i.e., positive group delay: \(\tau_g > 0\)) of the transmitted light, and that when the coupling \(J\) between the resonators is weak the reduction in the group velocity approaches to that experienced in a single passive resonator. In contrast, in the active-passive COM, one can tune the system to switch from slow to fast light, or vice versa, by controlling \(P_L\) or \(k/c\), such that the COM experiences the PT-phase transition (Fig. 5).

In the regime \(k/c < 1\), as \(P_L\) is increased from zero, the system first enters into the slow-light regime (\(\tau_g > 0\)), and \(\tau_g\) increases until its peak value. Then it decreases, reaching \(\tau_g = 0\), at a critical value of \(P_L\) (Fig. 5a). The higher is the \(k/c\), the sharper is the decrease. Increasing \(P_L\) beyond this critical value completes the transition from slow to fast light and \(\tau_g\) becomes negative (\(\tau_g < 0\)). After this transition, the advancement of the pulse increases with increasing \(P_L\) until it reaches its maximum value (more negative \(\tau_g\)). Beyond this point, a further increase in \(P_L\), again, pushes \(\tau_g\) closer to zero.

Figure 3 | The transmission rate \(\eta\) of the probe light in the active-passive system. The relevant parameters are taken as \(J/c = 1\) and \(P_L = 10\) μW.
In the regime $\kappa/\gamma > 1$, increasing $P_L$ from zero first pushes the system into the fast-light regime and increases the advancement of the pulse ($\tau_g < 0$) until the maximum advance is reached (Fig. 5b). After this point, the advance decreases with increasing $P_L$ and finally $\tau_g$ becomes positive, implying a transition to slow light. If $P_L$ is further increased, $\tau_g$ first increases until its peak value, and then decreases approaching $\tau_g = 0$.

The $P_L$ value required to observe the transition from slow-to-fast light (when $\kappa/\gamma < 1$) or from fast-to-slow light (when $\kappa/\gamma > 1$) depends on the gain-to-loss ratio $\kappa/\gamma$ if the coupling strength $J$ is fixed (Fig. 5). This implies that, when $P_L$ is kept fixed, one can also drive the system from slow-to-fast or fast-to-slow light regimes by tuning $\kappa/\gamma$. A simple picture can be given for this numerically-revealed feature: for $\Delta_L \sim 0$, $\zeta \sim 0$, we simply have $G_1 = G_2 \approx (J^2 - \kappa^2) + i\kappa\gamma_x$, which is minimized for $F = \kappa/\gamma$, or $\kappa/\gamma = 1, \kappa/\gamma = 1$; therefore, in the vicinity of the gain-loss balance, the denominator of $A$ is a real number, and $\text{Im}(A)/\text{Re}(A) \sim (1 - \gamma/k)^{-1}$, i.e. having reverse signs for $\kappa/\gamma > 1$ or $\kappa/\gamma < 1$. Correspondingly, $\arg(A)$ or $\text{arg}([\omega_p])$ and hence its first-order derivative $\tau_g \sim (\gamma/k - 1)$ (for $J/k = 1$). Clearly, the sign of $\tau_g$ can be reversed by tuning from the PT-symmetric regime (with $\kappa/\gamma < 1$) to the broken-PT regime (with $\kappa/\gamma > 1$). We note that the appearance of the fast light in the PT-breaking regime, where the gain becomes to exceed the loss, is reminiscent of that observed in a gain-assisted or inverted medium.

In order to better visualize and understand how the switching from the slow-to-fast light and vice versa takes place, when the gain-to-loss ratio $\kappa/\gamma$ is tuned at a fixed-pump power $P_L$, or when $P_L$ is tuned at a fixed value of $\kappa/\gamma$, we present the phase of the transmission function $\theta(\omega_p)$ in Figs. 6(a–c). For this purpose, we choose the values of $P_L$ and $\kappa/\gamma$ from Fig. 5, where their effects on the optical group velocity $\tau_g$ were presented. These calculations clearly show that, near the resonance point ($\delta_p = 0$), the slope of the curves can be tuned from positive to negative or vice versa, by tuning $P_L$ or $\kappa/\gamma$, which agrees well with the slow-fast light transitions (see Fig. 5). In sharp contrast, Fig. 6d shows that for the passive-passive COM (e.g., $\kappa/\gamma = -1$), no such type of sign reversal can be observed for the slope of the phase curves, corresponding to the fact that only the slow light can exist in that specific situation.

**Discussion**

In conclusion, we have studied the optomechanically-induced-transparency (OMIT) in PT-symmetric coupled microresonators with a tunable gain-loss ratio. In contrast to the conventional OMIT in passive resonators (a transparency peak arising in the otherwise strong absorptive spectral region), the active OMIT in PT-symmetric resonators features an inverted spectrum, with a transparency dip between two sideband peaks, providing a COM analog of all-optical inverted-ETT. For this active-OMIT system, the counterintuitive effects of gain- or pump-induced suppression of the optical transmission rate are revealed. In particular, the transition from slow-to-fast regimes by tuning the gain-to-loss ratio or the pump power is also demonstrated. The possibility of observing the PT-symmetric fast light, by tuning the gain-to-loss ratio of the coupled microresonators, has not studied previously. These exotic features of OMIT in PT-symmetric resonators greatly widens the range of applications of integrated COM devices for controlling and engineering optical photons. In addition, our work can be extended to study e.g. the OMIT in a quasi-PT system, the OMIT cooling of mechanical motion, with two mechanical modes, or the gain-assisted nonlinear OMIT.

**Methods**

**Derivation of the optical transmission rate.** Taking the expectation of each operator given in Eqs. (2)–(4), we find the linearized Heisenberg equations as

$$\langle \delta a \rangle = -(\delta A_0 + \gamma_0) \langle \delta a \rangle + i (\delta a) + i \kappa \gamma_x / (\delta x) + \gamma \exp(-\delta t),$$

$$\langle \delta a^\dagger \rangle = -(\delta A_0 - \kappa) \langle \delta a \rangle + i (\delta a^\dagger),$$

$$\langle \delta x \rangle + \Gamma_m (\delta x) + \kappa \gamma_x / (\delta x) = \frac{\kappa}{m} \left( a_{i,\gamma}^\dagger (\delta a) + a_{i,\gamma} (\delta a^\dagger) \right).$$

which can be transformed into the following form, by applying the ansatz given in Eq. (8),
Figure 6 | The phase of the transmission amplitude $\delta(\omega_p)$ with different parameters. (a-b) With different values of pump power $P_L$. (c) With different values of gain-to-loss ratio $\kappa/\gamma$. The specific values of $P_L$ and $\kappa/\gamma$ are taken from Fig. 5, corresponding to the slow and fast light regimes. For comparison, the results for the passive-passive COM system are also plotted in (d) with $\kappa/\gamma = -1$.

Solving these algebraic equations leads to

$$
\delta a_{\pm} = \frac{-i \gamma a_{\pm}^0}{(a_{\pm}^0 - i \Delta_{10}^0)} 
$$

(20)

where we have used $n_i = |a_{\pm}^0|^2$ ($i = 1, 2$) and

$$
\mu_\pm = -\kappa - i \xi \pm \Delta_1, 
$$

(14)

$$
G_1 = (i \Delta_1 + \gamma - i \xi \mu_+)^2, 
$$

(21)

$$
G_2 = (i \Delta_1 + \gamma - i \xi \mu_-)^2. 
$$

The expectation value $\langle a_{\pm}^0(t) \rangle$ of the output field $a_{\pm}^0(t)$ can be calculated using the standard input-output relation $a_{\pm}^0(t) = a_{\pm}^0 - \sqrt{2} \gamma \eta(t)$, where $a_{\pm}^0$ and $a_{\pm}^0$ are the input and output field operators, and

$$
\langle a_{\pm}^0(t) \rangle = \left[ \frac{G_j}{\sqrt{G_j^2 - \sqrt{2} \xi}} \right] e^{i\omega_j t} + \left[ \frac{G_j}{\sqrt{G_j^2 - \sqrt{2} \xi}} \right] e^{-i\omega_j t}. 
$$

(22)

Hence, the transmission rate of the probe field can be written as $\eta = \langle |\ell(\omega_p)|^2 \rangle$, where $\ell(\omega_p)$ is the ratio of the output field amplitude to the input field amplitude at the probe frequency

$$
\ell(\omega_p) = \frac{\epsilon_p - 2 \gamma \eta a_{\pm}}{\epsilon_p}, 
$$

(23)

where $A = \delta a_{\pm}$ is given in Eq. (9). In order to receive some analytical estimations, we take $\nu_{in}/\nu_{out} \sim 0$, $\Delta_{kl} \sim 0$, which leads to $\mu_\pm \sim \kappa; \mu_{\pm} \sim (\gamma \mp 2) (\gamma - \nu_{in}) \kappa$. For $x_i \sim 0$, we have

$$
\eta \approx \left[ 1 + 2 \gamma \nu_{in} (f^2 - \gamma \nu_{in}^2) \nu_{out} (f^2 - \gamma \nu_{in}^2) \right]:^2. 
$$

(24)
i.e. \( \eta \sim (\Gamma / 2\pi)^{-1} \) or \( \eta \sim (1 - k\gamma / \gamma)^{-1} \) (for a fixed value of \( \Gamma / 2\pi = 1 \)). This indicates that the transmission rate \( \eta \) tends to be maximized as the gain-to-loss ratio approaches one, that is \( k\gamma = 1 \), which was confirmed by our numerical calculations (see Fig. 2b).

1. O’Connell, A. D. et al. Quantum ground state and single-phonon control of a mechanical resonator. Nature 464, 697 (2010).
2. Teufel, J. D. et al. Sideband cooling of micromechanical motion to the quantum ground state. Nature 475, 359 (2011).
3. Chan, J. et al. Laser cooling of a nanomechanical oscillator into its quantum ground state. Nature 478, 89 (2011).
4. Aspelmeyer, M., Meystre, P. and Schwab, K. Quantum optomechanics. Physics Today 65, 29 (2012).
5. Aspelmeyer, M., Kippenberg, T. J. and Marquardt, F. Cavity optomechanics. Nature 464, 697 (2010).
6. Safavi-Naeini, A. H. et al. Electromagnetically induced transparency and slow light with optomechanics. Nature 472, 69 (2011).
7. Weis, S. et al. Optomechanically induced transparency. Science 330, 1520 (2010).
8. Agarwal, G. S. and Huang, S. Electromagnetically induced transparency in mechanical effects of light. Phys. Rev. A 81, 041803(R) (2010).
9. Scully, M. O. and Zubairy, M. S. Quantum Optics (Cambridge University Press, Cambridge, England, 1997).
10. Peng, B., Özdemir, S. K., Chen, W., Nori, F. and Yang, L. What is and what is not mechanical effects of light. Phys. Rev. A 90, 013839 (2014).
11. Anisimov, P. M., Dowling, J. and Sanders, B. C. Objectively discerning Autler-Townes splitting from electromagnetically induced transparency. Phys. Rev. Lett. 107, 163604 (2011).
12. Xiang, Z.-L., Ashhab, S., You, J. Q. and Nori, F. Rev. Mod. Phys. 85, 623 (2013).
13. Zhou, X. et al. Slowing, advancing and switching of microwave signals using circuit nanoelec-tromechanics. Nature Phys. 9, 179 (2013).
14. Jiang, C., Chen, B. and Zhu, K. D. Tunable pulse delay and advancement in a coupled nanomechanical resonator-superconducting microwave cavity system. EPL 94, 38002 (2011).
15. Dong, C., Zhang, J., Fiore, V. and Wang, H. Optomechanically induced transparency and self-induced oscillations with Bogoliubov mechanical modes. Optica 1, 425 (2014).
16. Grudinin, I. S., Lee, H., Painter, O. and Vahala, K. J. Phonon laser action in a tunable two-level system. Phys. Rev. Lett. 104, 083901 (2010).
17. Stannigel, K. et al. Optomechanical quantum information processing with photons and phonons. Phys. Rev. Lett. 109, 013603 (2012).
18. Komar, P. et al. Single-phonon nonlinearities in two-mode optomechanics. Phys. Rev. A 87, 013839 (2013).
19. Dong, C., Fiore, V., Kuzyk, M. C. and Wang, H. Optomechanical dark mode. Science 338, 1609 (2012).
20. Hill, J. T., Safavi-Naeini, A. H., Chan, J. and Painter, O. Coherent optical wavelength conversion via cavity optomechanics. Nature Commun. 3, 1196 (2012).
21. Fiore, V., Yang, Y., Kuzyk, M. C., Barbour, R., Tian, L. and Wang, H. Storing optical information as a mechanical excitation in a silica optomechanical resonator. Phys. Rev. Lett. 107, 136601 (2011).
22. Ludwig, M., Safavi-Naeini, A., Painter, O. and Marquardt, F. Enhanced quantum nonlinearities in a two-mode optomechanical system. Phys. Rev. Lett. 109, 063601 (2012).
23. Fan, J. and Zhu, L. Enhanced optomechanical interaction in coupled microresonators. Opt. Express 20, 20790 (2012).
24. Yan, X. B. et al. Coherent perfect absorption, transmission, and synthesis in a double-cavity optomechanical system. Opt. Express 22, 4886 (2014).
25. Jiang, C., Liu, H., Cui, Y. and Li, X. Electromagnetically induced transparency and slow light in two mode optomechanics. Opt. Express 21, 12165 (2013).
26. Bender, C. M. and Boettcher, S. Real spectra in non-Hermitian Hamiltonians having PT symmetry. Phys. Rev. Lett. 80, 5243 (1998).
27. Bender, C. M., Gianfreda, M., Özdemir, S. K., Peng, B. and Yang, L. Twofold transition in PT-symmetric coupled oscillators. Phys. Rev. A 88, 062111 (2013).
28. Peng, B. et al. Parity-time-symmetric whispering-gallery microcavities. Nature Phys. 10, 394 (2014).
29. Regensburger, A. et al. Parity-time synthetic photonic molecules. Nature 488, 167 (2012).
30. Rüter, C. E. et al. Observation of Parity-time symmetry in optics. Nature Phys. 6, 192 (2010).
31. Agarwal, G. S. and Qu, K. Spontaneous generation of photons in transmission of quantum fields in PT-symmetric optical systems. Phys. Rev. A 85, 031802(R) (2012).
32. Stannigel, K., Özdemir, S. K. and Yang, L. Tunable add-drop filter using an active whispering gallery mode microcavity. Appl. Phys. 100, 181103 (2013).
33. Jing, H., Özdemir, S. K., Li, X.-Y., Zhang, J., Yang, L. and Nori, F. PT-symmetric phonon laser. Phys. Rev. Lett. 113, 053604 (2014).
34. Oishi, T. and Tomita, M. Inverted coupled-resonator-induced transparency. Phys. Rev. A 88, 013813 (2013).
35. Peng, B. et al. Loss-induced suppression and revival of lasing. Science 346, 328 (2014).
36. Brandstetter, M. et al. Reversing the pump-dependence of a laser at an exceptional point. Nature Comm. 5, 4034 (2014).
37. Wang, L. J., Kuemich, A. and Dogariu, A. Gain-assisted superradiant light propagation. Nature 406, 277 (2000).
38. We thank Y. Zhang, P. M. Anisimov, and A. Miranowicz for helpful and stimulating discussions. HJ is supported by the NSFC (11274098, 11474087), LY is supported by ARO Grant No. W911NF-12-1-0026. J.Z. is supported by the NSFC (61174116). FN is supported by the NBRPC (2014CB921401). XYL is supported by the NSFC (11374116) and the Grant No. W911NF-12-1-0026. J.Z is supported by the NSFC (61174084, 61134008) and the NBRPC (2014CB921401). The authors declare no competing financial interests.