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Polarization-based control of phonon laser action in a Parity Time-symmetric optomechanical system

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Realizing a convenient way to control the phonon laser action is of great importance and may find applications in phonon laser devices engineering. Here we propose a vector parity time (PT)-symmetric optomechanical system to study the phonon laser action. We find that there is a specific region for the optimum mechanical gain appearing in parameter areas of the cavity gain and loss. The intensity of phonon laser action can be well controlled by adjusting the polarization of the pump field. The threshold value of phonon laser action manifests as a function relationship of the polarization direction θ . Furthermore, an ultralow threshold (even if threshold less) can be obtained around the exceptional point with the gain-loss balance. Our results indicate that the intensity and threshold of the phonon laser action can be continuously adjusted by only tuning the photon polarization, which provides a new degree of freedom to realize phonon laser regulation.

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Just like optical lasers, recent research into phonon lasers has been increasingly rapid. With the development of on-chip devices, the phonon laser has also made spectacular advances by means of microfabricated architectures, these are not only theoretical investigations but numerous experimental proposals have also been put forward to produce a phonon laser via coherent amplification of stimulated emission of phonons in various physical systems. Many years ago, Chen and Khurgin analyzed the feasibility of phonon lasers and pointed out a promising scheme to achieve phonon laser action¹. Next, based on prominent advances in experimental realization on the nanoscale, single trapped ions² and quantum dots^{3,4} have been utilized to study phonon lasers. Recent experimental phonon laser platforms including semiconductors⁵, nanomechanics⁶, and an electro-mechanical resonator⁷ have made significant progress. And phonon-stimulated emission also has been observed experimentally in cryogenic ionic compounds^{8–11}, semiconductor superlattices¹², and electromagnetically trapped ions¹³, which builds a vital system to produce phonon lasers theoretically and experimentally. Furthermore, phonon lasers also have attracted wide interest in matching the significant utility of its optical cousins such as medical imaging and high-precision measurement devices.

Cavity optomechanics (COM)¹⁴, characterized by exploring the radiation–pressure interaction between optical modes and mechanical modes, have received great attention in combining nanomechanics with nanophotonics. The unique optomechanical nonlinearity of the system has revealed abundant intriguing physical phenomena, including numerous interesting quantum^{15–19} and classical^{20–29} nonlinearity phenomena. Significantly, considering orthogonally polarized modes in a whispering-gallery-mode resonator has revealed important physical mechanisms, for example, based on orthogonally polarized modes in a whispering-gallery-mode resonators, we can observe couplings between orthogonally polarized modes in a birefringent whispering-gallery-mode resonator, which can be used for beneficial effect in sensing and optomechanical experiments³⁰. In addition, due to the effect of the polarization coupling occurring in the ring waveguide, the single-ring phase shifter can be realized experimentally³¹. Recently, the concept of vector cavity optomechanics^{32,33} has been put forward, where the polarization behavior of light is introduced to achieve optomechanical control, which can manipulate the on-chip light propagation with nano-beam optomechanical crystal.

Parity time (PT)-symmetric optomechanical system has recently undergone extensive research to achieve distinctive optical behavior, which is unattainable with ordinary optomechanical systems^{34–37}. It is known that the transition from an unbroken PT-symmetric phase (real eigenvalue spectra) to a spontaneously broken PT-symmetric phase (complex eigenvalue spectra) emerges as the system parameters in the Hamiltonian are properly modified^{36,38}. Non-Hermitian degeneracies known as exceptional points occurring in open physical systems have recently been utilized to enhance sensitivity in an optical microcavity^{39,40}, which suggests that the sensitivity enhancement conforms to square-root and cube-root decay law. Furthermore, observation of sub-Poissonian phonon lasing in a three-mode optomechanical system⁴¹, and the discovery of phonon lasing action in a compound whispering-gallery-microcavities (WGM) system⁴², have received increasing attention. More strikingly, a significative regime of phonon lasing in a PT-symmetric system has been reported⁴³, which reveals the PT-symmetric system has been proved to be an excellent platform for exploring phonon laser action.

In this work, inspired by the vector cavity optomechanics, we present the vector PT-symmetric-coupled resonators to study

phonon laser action. In contrast to other conventional phonon lasing regimes^{42,43}, where the phonon laser action can only be adjusted by the intensity of coherent pump field, the polarized PT-symmetry system proposed here provides a new degree of freedom to realize phonon laser regulation based on the mechanical effect of light. Our regime is pumped by a linear polarized light field, in close analogy to a two-level laser system, the phonon laser is formed through the coherent amplification of the stimulated emission phonons arising from the interaction between photon and phonon. The theoretical analysis of our results was formulated by Heisenberg–Langevin equations. First of all, we find that there is an optimal region for mechanical gain G in a certain parameter range of the gain and loss of the active and passive cavity. Moreover, our regime provides a novel way to improve the performance of controlling a phonon laser, that is achieved by tuning the polarization state of pump light, instead of only regulating the intensity of the pump field^{42,43}. Furthermore, the strong polarization dependence of the phonon laser action may offer insight into the understanding of nonlinear optomechanical interactions and find applications in photon and phonon manipulation. Specifically, the phonon laser action is a polarization-dependent effect instead of a polarization-independent effect when the polarization-dependent optomechanical nonlinearity are taken into account and this polarization-controlled phonon laser action may reveal nonlinear optomechanical interactions in a more robust way. Recently, on-chip manipulation and control of light propagation has made remarkable progress by means of PT-symmetric systems, for instance, non-reciprocal light transmission in the PT-symmetry-breaking phase can be observed in the PT-symmetric whispering-gallery microcavities³⁵. We believe that the polarization-related phonon laser action can stimulate further remarkable applications in polarization-controlled on-chip optical architectures due to the improvement of nanofabrication techniques.

Results

Physical setup and dynamical equation. Figure 1a gives the schematic of the system we propose, a vector parity time (PT)-symmetric COM system which is composed by two directly coupled microcavities consisting of orthogonally polarized cavity modes. In this setup, one of cavities is fabricated from silica doped with Er^{3+} ions, and gain in this cavity is derived from optically pumping the doped Er^{3+} ions, whereas another cavity has passive loss without dopants, which was identified experimentally as the PT-symmetry system³⁵. The passive cavity is formed by one fixed mirror which is partially transparent and located to the left, and one movable, totally reflecting mirror placed to the right. The movable mirror is pushed by the radiation pressure, which is produced by the photon's momentum transfer. Consequently, the movable mirror is modeled as a harmonic oscillator with effective mass m , eigenfrequency ω_m , and decay rate Γ_m . As Fig. 1b shows, each resonator has two modes of TE and TM, and a set of orthogonal basis vectors of polarization ($\vec{e}_\perp, \vec{e}_\parallel$) corresponding to TE and TM modes can be introduced³². For entirely arbitrary polarized orientation of input light (so long as it is perpendicular to the direction of propagation), the vector $\vec{e}$ can be disintegrated as $\vec{e} = \alpha\vec{e}_\perp + \beta\vec{e}_\parallel$, with $|\alpha|^2 + |\beta|^2 = 1$ (shown in Fig. 1b), and α and β are complex. In this letter, our system is pumped by linearly polarized optical field, viz. α and β are real. Figure 1c shows that the optical supermodes $\omega_\pm$ are coupled by the mechanical mode b .

We consider that a dual-mode active cavity couples to a dual-mode passive optomechanical cavity via optical tunneling, i.e., the vector PT-symmetric COM system (see Fig. 1a). In view of the generic optomechanical system⁴⁴, the Hamiltonian of our system

Fig. 1 Schematic diagram of the polarization-controlled PT-symmetric optomechanical system. **a** A schematic illustration of the vector PT-symmetric COM system, which includes an active cavity a_1 and a passive COM carries cavity mode a_2 and a mechanical mode b . The system is pumped by a linearly polarized optical field. **b** $a_{1\uparrow}(a_{1\leftarrow})$ and $a_{2\uparrow}(a_{2\leftarrow})$ are the TE (TM) modes of the cavity 1 and cavity 2. The included angle between the polarization of pump field and the vertical mode is θ for the two cavities. The coupling strength (also called photon-tunneling rate) of coupled microresonators with regard to two orthogonal modes are $J_{\uparrow}$ and $J_{\leftarrow}$, respectively, which can be adjusted by the relative separation of the microresonators. **c** Scheme depicting the corresponding optical supermodes are coupled by phonon, and two-level phonon laser energy-level diagram

pumped by linearly polarized optical field can be described by³²

$$H = H_c + \hbar\omega_m b^\dagger b + H_{\text{int}} + H_{\text{drive}} \quad (1)$$

where $a_{1j}(a_{2j})$, $a_{1j}^\dagger(a_{2j}^\dagger)$, and $b(b^\dagger)$ are the annihilation (creation) operators of the orthogonal cavity modes of the two resonators and the mechanical mode, respectively. The commutation relations are $[a_{1j}, a_{1j}^\dagger] = 1$, $[a_{2j}, a_{2j}^\dagger] = 1$, and $[b, b^\dagger] = 1$. The cavity mode a_{1j} coupling to a_{2j} (with the same resonance frequency ω_c) is pumped by a light field with frequency ω_d and amplitude ε_d . Considering the case of resonance, i.e., $\omega_d = \omega_c$, $\varepsilon_d = \sqrt{P_{\text{in}}/(\hbar\omega_d)}$ where P_{in} is the input power. As depicted in Fig. 1b, the linearly polarized input field can be expressed as $\varepsilon_d e^{-i\omega_d t} \vec{e}$, and $\vec{e}$ is the unit vector of polarization of the input field, which can be decomposed as $\vec{e} = a\vec{e}_{\uparrow} + b\vec{e}_{\leftarrow}$, with $a = \cos\theta$ and $b = \sin\theta$ being the projections of $\vec{e}$ at the vertical and horizontal modes, respectively, while θ is the included angle between $\vec{e}$ and the vertical mode, hence $\varepsilon_{d\uparrow} = \varepsilon_d \cos\theta e^{-i\omega_d t}$, $\varepsilon_{d\leftarrow} = \varepsilon_d \sin\theta e^{-i\omega_d t}$. The Hamiltonian of the cavity part can take the form $H_c = \hbar\omega_c (a_{1\uparrow}^\dagger a_{1\uparrow} + a_{1\leftarrow}^\dagger a_{1\leftarrow} + a_{2\uparrow}^\dagger a_{2\uparrow} + a_{2\leftarrow}^\dagger a_{2\leftarrow})$. The Hamiltonian of the interaction part can be given as $H_{\text{int}} = -\sum_{j=\uparrow,\leftarrow} \hbar[J_j (a_{1j}^\dagger a_{2j} + a_{2j}^\dagger a_{1j}) + g x_0 (a_{2j}^\dagger a_{2j}) (b + b^\dagger)]$, therein the first term of H_{int} represents the interaction between the active cavity and the passive cavity with coupling strength J_j , and the second term expresses optomechanical interaction induced by the radiation pressure of both TE and TM modes in the passive cavity with coupling coefficient $g = \omega_c/R$ ⁴⁴, and R is the cavity size. The last term in Eq. (1) $H_{\text{drive}} = i\hbar[\sqrt{2\kappa}(a_{1\uparrow}^\dagger \varepsilon_{d\uparrow} - \text{H.c.}) + \sqrt{2\kappa'}(a_{1\leftarrow}^\dagger \varepsilon_{d\leftarrow} - \text{H.c.})]$ describes the optical driving of the active cavity modes.

To study the temporal evolution of the system, we use the Heisenberg-Langevin equations of motion which give an

evolutionary description of the cavity fields and mechanical motion: (in a rotating frame at ω_c where the system Hamiltonian is transformed by a unitary transformation $U(t) = \exp(-i\omega_c a^\dagger a t)$)

$$\dot{\Psi} = M\Psi + \xi \quad (2)$$

where

$$M = \begin{pmatrix} \kappa & iJ_{\uparrow} & 0 & 0 & 0 \\ iJ_{\uparrow} & -\gamma & 0 & 0 & i g a_{2\uparrow} \\ 0 & 0 & \kappa' & iJ_{\leftarrow} & 0 \\ 0 & 0 & -\gamma' & iJ_{\leftarrow} & i g a_{2\leftarrow} \\ 0 & 0 & 0 & 0 & \Omega \end{pmatrix},$$

$$\xi = \begin{pmatrix} \sqrt{2\kappa}(\varepsilon_d \cos\theta + a_{\text{in}}) \\ \sqrt{2\kappa'}(\varepsilon_d \sin\theta + a_{\text{in}}) \\ \sqrt{2\kappa'} a_{\text{in}} \\ \frac{\hbar g}{m\Gamma_m} (a_{2\uparrow}^\dagger a_{2\uparrow} + a_{2\leftarrow}^\dagger a_{2\leftarrow}) + \sqrt{2\Gamma_m} b_{\text{in}} \end{pmatrix}$$

$\Psi = (a_{1\uparrow}, a_{2\uparrow}, a_{1\leftarrow}, a_{2\leftarrow}, x)$, $\Omega = \frac{1}{m\Gamma_m}(-\frac{d^2}{dt^2} - \omega_m^2)$, the noise correlators are defined as $\langle a_{\text{in}}(t) a_{\text{in}}^\dagger(t') \rangle = \delta(t - t')$, $\langle a_{\text{in}}^\dagger(t) a_{\text{in}}(t') \rangle = 0$, $\langle b_{\text{in}}(t) b_{\text{in}}^\dagger(t') \rangle = (n_{\text{th}} + 1)\delta(t - t')$, and $\langle b_{\text{in}}^\dagger(t) b_{\text{in}}(t') \rangle = n_{\text{th}}\delta(t - t')$, with the average number of mechanical quanta $n_{\text{th}} \approx k_B T / \hbar\omega_m$. In this case, using the semiclassical approximation, e.g., $\langle AB \rangle = \langle A \rangle \langle B \rangle$ with A and B being optical or mechanical operators, the Heisenberg-Langevin equations can be reduced to a group of nonlinear evolution equations in which the noise has been ignored due to the fact that we consider the mean values of all the operators. The position of the mechanical oscillator is defined as $x = x_0(b^\dagger + b)$, $x_0 = \sqrt{\hbar/(2m\omega_m)}$ is the zero-point fluctuation amplitude of the mirror, and the operator $\hat{b}$ denotes the phonon mode. $\kappa(\kappa')$ and $\gamma(\gamma')$ are the optical gain and decay of the vertical and the horizontal modes in cavity 1 and cavity 2, respectively. Γ_m is the mechanical damping rate.

The distribution of steady-state photon number in the passive cavity. Recently, PT-symmetric whispering-gallery microcavities have been demonstrated experimentally, where low-power optical isolation and non-reciprocal light transmission have been observed under the balanced gain and loss^{35,45}. In addition, the experimental observation of PT-symmetry in optically induced atomic lattices with periodical gain and loss profiles³⁸, and an anti-PT-symmetric potential has been experimentally produced in a pair of optically induced waveguides coupled by steady-state flying atoms⁴⁶. These experimental facts indicate that the steady state of PT-symmetric systems can be observed in a short period of time via current experimental technology. First, we analyze the property of the steady-state situation of our system. Setting the time differential in dynamical Eqs. (2) to zero, thus the steady-state solutions can be obtained, i.e., the steady-state intracavity photons $a_{1j,s}$ and $a_{2j,s}$ ($j = \uparrow, \leftarrow$). Their analytical expressions are

$$\begin{aligned} a_{1\uparrow,s} &= f_1(\kappa, \gamma) \cos\theta, a_{1\leftarrow,s} = f_1(\kappa', \gamma') \sin\theta, \\ a_{2\uparrow,s} &= f_2(\kappa, \gamma) \cos\theta, a_{2\leftarrow,s} = f_2(\kappa', \gamma') \sin\theta, \end{aligned} \quad (3)$$

with

$$f_1(x, y) = \frac{\sqrt{2x}\varepsilon_d(y - i g x_s)}{J_j^2 - xy + i g \kappa x_s}, f_2(x, y) = \frac{i J_j \sqrt{2x}\varepsilon_d}{J_j^2 - xy + i g \kappa x_s},$$

and the displacement of the mechanical resonator

Fig. 2 The steady-state photon number in cavity 2. **a** The steady-state photon number N in cavity 2 varies with the input power. System parameters we take here are, the wavelength of the cavity field is 1550 nm, $g = \omega_c/R$, $R = 34.5 \times 10^{-6}$ m, $m = 5 \times 10^{-11}$ kg, $\gamma/\kappa = \gamma'/\kappa' = 1.5$, $\kappa = 6$ MHz, $\kappa' = 10$ MHz, $\Gamma_m = 2.5 \times 10^5$ Hz, $\omega_m = 2.5 \times 23.4$ MHz, $J_{\uparrow} = 6$ MHz, $J_{\downarrow} = 1.3J_{\uparrow}$, respectively. **b** The photons number N varies with γ/κ (γ'/κ'), $P_{in} = 10 \mu\text{W}$, $\gamma/\kappa = \gamma'/\kappa'$, other parameters are same as **a**. **c** Calculated the intracavity photons in the passive cavity as a function of the polarization angle θ under the balanced gain and loss ($\kappa/\gamma = 1$), other parameters are same as **a**. These parameters are derived from ref. ⁴³

$x_s = \frac{\hbar g}{m\omega_m} \left(|a_{2\uparrow,s}|^2 + |a_{2\downarrow,s}|^2 \right)$. It can be inferred that the steady-state value of intracavity photons is a fifth-order polynomial equation, thus, there are at most five real roots. But the pump field is too weak to reach the multiple steady states in our regime, so there is only one real root under the weak pump power. Figure 2a shows how the steady-state populations of the intracavity photons in the cavity 2 vary with input power due to the Eq. (3). It is obvious that the photon number N acquires giant enhancement with the weak pump power ($P_{in} \leq 10 \mu\text{W}$) increasing. Specifically, the rise of photon number in the case of $\theta = \pi/2$ (the violet line) features linear responses, while the significant nonlinear enhancement appears in other polarization states, leading to the appearance of the accompanied giant enhancement of COM interactions. Interestingly, the polarization state of the pump light serves as a sensitive controller for the physical phenomenon, namely, the photon number experiences distinct rising tendency with the different polarization state of pump light. In other words, the photon-tunneling effect is equivalently controlled by the polarization states of the pump field. From this phenomenon, we can work out that the gain-to-loss ratio may also affect the behavior of photon tunneling, then we plot how the photon number varies with γ/κ at a fixed input power P_{in} in Fig. 2b. As γ/κ increases, the change of N has a obvious peak value at $\gamma/\kappa = 1$ (the gain-loss balance) and the vicinity of $\gamma/\kappa = 0.6$ for $\theta = 0$ and $\theta = \pi/2$, which indicates a giant enhancement of the intracavity field intensity at the gain-loss balance and the vicinity of $\gamma/\kappa = 0.6$. The remaining curves show

N have two peak values at the two points for $0 < \theta < \pi/2$, owing to the projections of linearly polarized pump light onto the vertical and horizontal modes. The two extreme values of N satisfy $J_{\uparrow}^2 = \kappa\gamma$ and $J_{\downarrow}^2 = \kappa'\gamma'$ resulting from the fact that first derivative of N is equal to 0, and the point of $\kappa/\gamma = 1$ is the exceptional point (EP) for $\theta = 0$. We proceed to study the remarkable influence of the photon polarization on the intracavity photons with the balanced gain and loss ($\kappa/\gamma = 1$), and the intracavity photons in the passive cavity is shown in Fig. 2c. Clearly, the photon number can be significantly enhanced in the vicinity of $\theta = 0, \pi$, which coincides well with the EP, while the slight enhancement appears around $\theta = \pi/2, 3\pi/2$ due to the considerable shift away from the actual EP. Physically, the field localization induces the dynamical accumulations of photon energy, in addition, the unidirectional energy transfer from the active cavity to the passive cavity. Hence, the enormously enhanced photon number can be generated at the EP. The proposed mechanism is especially suited for the condition of the gain-loss balance, where the intensity of intracavity photon can be significantly enhanced. Then leading to a giant enhancement of both optical pressure and mechanical gain. This fact provides a route for achieving a low-power optomechanical amplifier in the presence of polarization-optomechanics EP.

Polarization-based control of phonon laser action. In order to explore the specific relationship between the gain (loss) and the mechanical gain G , we have investigated the gain G of the mechanical mode in a parameter map of the optical gain κ and

Fig. 3 The optimal area of mechanical gain G . Mechanical gain G versus the optical gain κ and loss γ . We use the parameters here are $\kappa' = 1.5$, $\kappa\gamma' = 1.5\gamma$, $P_{\text{in}} = 10 \mu\text{W}$ and $J_{\leftrightarrow} = 1.3J_{\uparrow}$

loss γ in Fig. 3. The system is in the PT-symmetric phase when we choose $J_{\uparrow} = \omega_m/2$ and $J_{\leftrightarrow} = 1.3J_{\uparrow}$. We find that the optimum mechanical gain can be reached by choosing the proper optical gain κ and loss γ . As shown in Fig. 3, the optimum mechanical gain formation starts at a critical point $\kappa = \gamma = 15 \text{ MHz}$, namely, the mechanical gain is very small even close to zero until the κ and γ arrive at this point. Obviously, we can get a clear description that the mechanical oscillator cannot possess gain when $\kappa > \gamma$ above the critical point, in contrast to the fact that the mechanical gain can be produced in an optimal area of $\kappa < \gamma$ below the critical point. This result indicates that optical gain and loss within the optimal area can significantly contribute to the supermode population inversion for magnifying the stimulated phonon number. Such that we find the phonon laser action can achieve its optimal solution in the specific parameter region while can not be acquired in other wide range of gain κ and loss γ . It is worth noting that the optimal parameter area can guide us to design more efficient phonon laser devices.

In Fig. 4, we present the dependence of the phonon number n on the pumped light field P_{in} for different polarizations, which clearly shows that the phonon number n can be controlled by the polarization of the pump field. Based on ref. 42, only above the threshold can the phonon lasing action occur, Fig. 4a shows that the phonon laser is not triggered where $\gamma/\kappa = 0.8$. As γ/κ goes up in Fig. 4b–d, i.e., $\gamma/\kappa = 4/3, 5/3, 2$, the polarization state of pump light can directly influence the intensity of the phonon laser action for different loss–gain ratio $\gamma/\kappa(\gamma'/\kappa')$. As shown in Fig. 2, the photon number in cavity 2 can be adjusted by tuning the polarization of the pump field, and this phenomenon can be explained by the fact that the vertical and horizontal modes in cavity 2 have different photon number according to the fact that the two orthogonal components of linearly polarized pump light act as two different intensities of driving field, resulting in different optomechanical nonlinear strength. Moreover, the two orthogonal modes can produce two different intensities of the phonon laser action, which means that the phonon laser action in our regime possesses the polarization-related property. As a transversely oscillating electromagnetic wave, polarization is a fundamental property of light, and optical polarization is often a major consideration in the construction of many optical systems, such as optical biomedical imaging, communication, and sensor systems^{47,48}. This demonstrates that it is very important and achievable to utilize the polarization nature of light field to study phonon laser action. Recently, the photon polarization degree of freedom has been introduced to achieve optomechanical

control in the vector cavity optomechanics³², which may enable on-chip optical control of polarization management with remarkable applications in nanophotonic polarization-related devices, such as optomechanical polarizers³³. Compared with previous studies of the phonon laser action^{42,43}, as shown in Fig. 4, our scheme provides a new degree of freedom to adjust the intensity of coherent phonon laser action by varying the polarization angle of pump light without changing other system parameters. This fact opens the way for polarization-controlled phonon lasing and may thus be applied to engineer on-chip phonon laser devices.

In Fig. 5, we plot the phonon number n versus the polarization angle θ in the PT-symmetric phase under different pump power P_{in} , which shows the role of the light field polarization behavior on the control of phonon laser action more clearly. We can see that the phonon number has a Lorentzian-like-shape dependence on the polarization direction θ , which indicates that the intensity of phonon laser action can be continuously adjusted by only tuning the photon polarization behavior without changing other parameters of the device. Compared with the conventional phonon laser regime^{41–43}, our scheme provides an additional degree of freedom to control phonon laser action. In addition, the slight enhancement of the pump power leads to the fact that the maximum value of the phonon number increases dramatically arising from the population inversion of the supermodes, corresponding to the exponential growth of phonon number in Fig. 4. Compared with the pump-power-controlled method, the change of phonon number induced by the light polarization is slower, which can boost the feasibility of controlling the intensity of phonon laser action quite accurately. This result is a manifestation of the intrinsic merit of polarization-dependent optomechanics, and we can utilize the light polarization to realize all-optical switcher of phonon-lasing devices. On the other hand, our work may offer a useful platform to convert the information carried by optical polarization into the stimulated emission of phonons while preserving their intrinsic coherence, which is of great importance in communication systems and modern communication networks.

We next explore the impact of polarization on the phonon laser threshold value P_{th} , which behaves as a function relationship of the polarization direction θ as shown in Fig. 6. Remarkably, unlike previously reported observations of phonon lasing effects in microcavity structures^{42,43}, in which the phonon laser threshold value can only be adjusted by changing the system parameters, in our scheme the threshold value of phonon laser action can also be controlled by adjusting the polarization direction of the pump field. In this work, the threshold condition and the threshold power of phonon laser action are $G = \Gamma_m/2$ and $P_{\text{th}} = N_+(\gamma_- + \gamma_+)\hbar\omega_+$ ^{42,43}, respectively. Therefore, the thresholds of the vertical and horizontal modes yield $P_{\text{th}\uparrow} =$

$4\Gamma_m\omega_{\uparrow}\left[\left(\Delta\omega_{\uparrow} - \omega_m\right)^2 + \left(\frac{\kappa - \gamma}{2}\right)^2\right]/(gx_0)^2$, $P_{\text{th}\leftrightarrow} = 4\Gamma_m\omega_{\leftrightarrow}$
 $\left[\left(\Delta\omega_{\leftrightarrow} - \omega_m\right)^2 + \left(\frac{\kappa' - \gamma'}{2}\right)^2\right]/(gx_0)^2$. Clearly, $P_{\text{th}\uparrow} \rightarrow 0$ and $P_{\text{th}\leftrightarrow} \rightarrow 0$
 under the conditions (a) $\Delta\omega_{\uparrow} = \omega_m$, $\Delta\omega_{\leftrightarrow} = \omega_m$,
 (b) $\kappa = \gamma$, $\kappa' = \gamma'$, and this implies that an ultralow threshold
 (even threshold less) phonon laser action can be obtained by
 tuning the parameter configuration of the system in the vector
 PT-symmetric regime. As shown in Fig. 5, the threshold values
 become much smaller as the value of γ/κ approaches 1, as
 expected, phonon lasing can be enhanced greatly in the vicinity of
 EP. Interestingly, we can regulate the threshold value of phonon
 laser action via modifying the polarization of pump light in our
 regime. This attribute offers a promising approach to engineer

Fig. 4 Polarization-controlled phonon laser action in our regime. Calculation results of the stimulated emission phonon number n vary with input power P_{in} when **a** $\gamma/\kappa = 0.8$, **b** $\gamma/\kappa = 4/3$, **c** $\gamma/\kappa = 5/3$, and **d** $\gamma/\kappa = 2$ under different angles θ at $t = [\Gamma_m/2]^{-1}$, respectively. The phonon number n is expressed by $n = \exp[2(G - \Gamma_m/2)/(\Gamma_m/2)]$. Setting the threshold condition $G = \Gamma_m/2$, viz. $n = 1$, the dotted line denotes the threshold value P_{th} of the phonon laser for different angles θ , respectively. We choose $2J_{\uparrow} = \omega_m$, $J_{\downarrow} = 1.3J_{\uparrow}$, $\gamma/\kappa = \gamma'/\kappa'$, $\kappa = 6$ MHz, $\kappa' = 10$ MHz

Fig. 5 Phonon number varies with polarization angle θ . Calculated results of phonon number n as a function of the polarization direction θ under different pump power P_{in} . The parameters we used here are $\kappa' = 2\kappa$, $\gamma' = 2\gamma$, $\kappa = 6 \times 10^6$ Hz, $\gamma = 9 \times 10^6$ Hz, $J_{\uparrow} = \omega_m/2$, and $J_{\downarrow} = 1.3J_{\uparrow}$, respectively

Fig. 6 The threshold value of phonon laser action. Numerical calculation of the threshold value P_{th} as a function of the polarization direction θ . The parameters we used here are $\kappa' = 2\kappa$, $\gamma' = 2\gamma$, $\kappa = 6 \times 10^6$ Hz, $\gamma = 9 \times 10^6$ Hz, $J_{\uparrow} = \omega_m/2$, and $J_{\downarrow} = 1.3J_{\uparrow}$, respectively

polarization-controlled on-chip phononic devices in practical applications.

Discussion

We propose a feasible scheme to realize a polarization-controlled phonon laser in a vector PT-symmetric system based on the mechanical effect of light. Different from the phonon laser action previously discussed, our results have shown that the generation and the intensity regulation of phonon laser action can be well controlled by only tuning the polarization of the pump field while keeping other parameters unchanged. This indicates a novel way to improve the performance of controlling a phonon laser, that is

by tuning the polarization state of the pump light, instead of only adjusting the intensity of the pump field. Moreover, the threshold value of phonon laser action behaves as a function relationship of the polarization direction θ . Furthermore, an ultralow threshold (even threshold less) phonon laser action can be obtained around EP with the gain-loss balance. The underlying physical mechanism can be explained as follows: the dynamical accumulations of photon energy result in drastic population transition processes from the up-supermode level to down-supermode level, which accompany the phonon emission. This is reminiscent of random-defect-induced phonon lasing at EPs⁴⁹ and the loss-induced suppression and revival of lasing at EPs⁵⁰. Physically, the enormously enhanced photon number can induce the great

reinforcement of radiating pressure, consequently resulting in stronger phonon laser action. These attributes may inspire the exploration of engineering new polarization-related phonon laser devices featuring lower-threshold power and convenient regulation ways. In addition, our research which associates the vector regime and PT-symmetric optomechanics may induce more opportunities in both areas and provide new possibilities to develop the corresponding applications.

Finally, we present some discussions on the experimental implementation of our proposal with coupled optical cavity systems. Specifically, it is achievable to implement our proposal experimentally in the PT-symmetric optomechanical system with current accessible technology. First, we introduce a group of orthogonal basis vectors of polarization corresponding to TE and TM modes in the coupled whispering-gallery mode (WGM) optical microcavities³². Considering the system is pumped by linearly polarized light field, whose orthogonal components obey Malus law toward orthogonally polarized directions in the vector regime which reveals a nontrivial phenomenon in analogy to optomechanical polarizer. Second, the cavity gain $\kappa \sim 10$ MHz and coupling strength $J \sim 20$ MHz can be realized in the PT-symmetric system, which was identified experimentally as the experimental arrangement in ref. ³⁵, leading to the system reaching PT-symmetric phase required here. Last, the active cavity is fabricated from silica doped with Er^{3+} ions, and the Er^{3+} ions can emit photons in the 1550 nm band arising from optical pump with a light in the 1460 nm band. This leads to the amplification of weak signal light in the 1550 nm band³⁵, and the microtoroids are $\sim 34.5 \mu\text{m}$ in radius⁴³. Due to the fact that the two different polarizations have two different photon-tunneling values which are determined by geometry size and material properties as well as the index contrast between the cavity core and cladding^{51,52}, the tunneling between the two resonators of the two orthogonal polarizations we choose are different, viz. $J_{\leftrightarrow} = 1.3J_{\uparrow}$ where $J_{\leftrightarrow}$ and $J_{\uparrow}$ denote the tunneling values of horizontal and vertical directions, respectively. Consequently, all parameter values in our scheme are taken from current experiments, implying that our system is well within the phonon laser regime based on the above discussion⁴². In summary, the experimental implementation of our scheme is feasible under currently existing experimental techniques, i.e., all relevant system parameters and experimental techniques used in our scheme are well within the reach of the current experimental level. With regard to controlling the phonon laser action, we can employ nanophotonic polarizer³³ to tune the polarization direction of light field, which can be utilized to achieve optomechanical control and regulate the behavior of the phonon laser action.

Methods

Derivation of the polarization-controlled phonon laser. The polarization-dependent character possesses important practical application. Here, we integrate the feature of vector cavity optomechanics into PT-symmetry, and introduce the vector PT-symmetric phonon laser. According to ref. ⁴³, PT-symmetric-coupled resonators provide a coherent phonon amplification by phonon-mediated transitions between optical supermodes in parallel with a two-level laser by the electronic transitions. On this basis, we propose a controllable phonon laser via tuning the polarization of pump field. Before proceeding further, we illuminate the important influence of the cavity coupling intensity J in the regime. To be specific, the non-Hermitian Hamiltonian of general PT-symmetric optical system is

$\hbar\omega_c(a_1^\dagger a_1 + i\kappa) + \hbar\omega_c(a_2^\dagger a_2 - i\gamma) + \hbar J(a_1^\dagger a_2 + a_2^\dagger a_1)$, which can be diagonalized as $\hbar(\omega_+ + i\gamma_+)o_1^\dagger o_1 + \hbar(\omega_- + i\gamma_-)o_2^\dagger o_2$, with

$\omega_{\pm} = \omega_c \pm \sqrt{J^2 - (\kappa + \gamma)^2/4}$, $\gamma_{\pm} = (\kappa - \gamma)/2$. For the case of $J \geq (\kappa + \gamma)/2$, the system is in the PT-symmetric phase. And only in this regime can the two nondegenerate supermodes exchange energy through the phonon and possess the population inversion, and then generating the phonon lasing⁴³. Note that the situation $J < (\kappa + \gamma)/2$ corresponds to degenerate supermodes

$\omega_{\pm} = \omega_c$, $\gamma_{\pm} = (\kappa - \gamma)/2 \pm \sqrt{J^2 - (\kappa + \gamma)^2/4}$, thus this situation which has no supermode splitting results in no phonon lasing action.

TE and TM modes are independent of each other in our regime, we can use the same method to deal with the two modes, respectively. Derived from Eq. (1), the optical supermodes $a_{\pm j} = (a_{1j} \pm a_{2j})/\sqrt{2}$, ($j = \uparrow, \leftrightarrow$) correspond to the energy

levels $\omega_{\pm j} = \omega_c \pm \sqrt{J_j^2 - (\kappa + \gamma)^2/4}$, $\omega_{\pm \leftrightarrow} = \omega_c \pm \sqrt{J_{\leftrightarrow}^2 - (\kappa' + \gamma')^2/4}$ can be introduced. Looking at one of the two orthogonal modes, the Fig. 1c describes the transition between the upper level ω_+ and the lower level ω_- while absorbing or emitting a phonon simultaneously, that is the eigenfrequencies ω_+ and ω_- can exchange energy through the phonon mode. In analogy to a two-level laser system, the stimulated emission of phonon can be generated by virtue of optical pumping of the upper level, then leading to the appearance of coherent phonon lasing. Additionally, the stimulated emission linewidth which describing the first-order coherence above the threshold is much narrower in comparison with that below the threshold^{42,43}. The interaction is provided by radiation pressure within the optomechanical system. Applying the rotating-wave approximation (RWA), the optomechanical interaction of Hamiltonian can be converted to another form⁴³, i.e.,

$$\begin{aligned} & -\hbar g x_0 (a_{\pm j}^\dagger a_{\pm j}) (b^\dagger + b) \\ & \rightarrow \frac{\hbar g x_0}{2} (a_{\pm j}^\dagger a_{\pm j} b^\dagger + a_{\pm j} a_{\pm j}^\dagger b^\dagger - a_{\pm j}^\dagger a_{\pm j} b^\dagger - a_{\pm j} a_{\pm j}^\dagger b^\dagger + \text{H.c.}) \\ & \rightarrow \frac{\hbar g x_0}{2} (\mathfrak{R}_j b^\dagger + b \mathfrak{R}_j^\dagger), (j = \leftrightarrow, \uparrow) \end{aligned} \quad (4)$$

$\mathfrak{R}_j = a_{\pm j}^\dagger a_{\pm j}$ ($j = \leftrightarrow, \uparrow$) is the population inversion of the supermodes. And the partly diagonalizable form of the Hamiltonian H in Eq. (1) is:

$$\begin{aligned} \tilde{H} = & \hbar\omega_m b^\dagger b + \sum_{j=\uparrow, \leftrightarrow} (\hbar\omega_{+j} a_{+j}^\dagger a_{+j} + \hbar\omega_{-j} a_{-j}^\dagger a_{-j}) \\ & + \frac{\hbar g x_0}{2} \sum_{j=\uparrow, \leftrightarrow} (b \mathfrak{R}_j^\dagger + \mathfrak{R}_j b^\dagger), \end{aligned} \quad (5)$$

where $\omega_{\pm j}$ are the optical frequencies of the supermodes and defined to be blue and red, respectively. The third term of the Hamiltonian $\tilde{H}$ describes the absorption of one phonon, which leads to transition from the red supermode to the blue supermode and the reverse process. Using the Hamiltonian in Eq. (5), we can rewrite the evolution equations for the mechanical mode and the operator $\mathfrak{R}_j$ as follows:

$$\begin{aligned} \dot{b} = & (-i\omega_m - \frac{\Gamma_m}{2})b - ig \frac{x_0}{2} \mathfrak{R}_\uparrow - ig \frac{x_0}{2} \mathfrak{R}_{\leftrightarrow} + \Xi_1(t), \\ \dot{\mathfrak{R}}_\uparrow = & \frac{igx_0 \Delta N_1}{2} b + (-i\Delta\omega_\uparrow + \Theta) \mathfrak{R}_\uparrow + \Xi_2(t), \\ \dot{\mathfrak{R}}_{\leftrightarrow} = & \frac{igx_0 \Delta N_{\leftrightarrow}}{2} b + (-i\Delta\omega_{\leftrightarrow} + \Theta') \mathfrak{R}_{\leftrightarrow} + \Xi_2(t), \end{aligned} \quad (6)$$

$\Theta = (\kappa - \gamma)/2$, $\Theta' = (\kappa' - \gamma')/2$, $\Delta N_j = a_{+j}^\dagger a_{+j} - a_{-j}^\dagger a_{-j}$, $\Delta\omega_j = \omega_{+j} - \omega_{-j}$, $\mathfrak{R} = \mathfrak{R}_\uparrow + \mathfrak{R}_{\leftrightarrow}$. We ignore the influence of $\Xi_1(t)$, $\Xi_2(t)$, which represent fluctuation operators correspond to the supermodes and the mechanical resonator. In order to solve Eq. (6) and obtain the mechanical gain, we introduce the slow varying amplitudes: $b = b e^{i\omega_m t}$, $\mathfrak{R}_j = \mathfrak{R}_j e^{i\Delta\omega_j t}$, $\mathfrak{R}_{\leftrightarrow} = \mathfrak{R}_{\leftrightarrow} e^{i\Delta\omega_{\leftrightarrow} t}$, and substituting these into Eq. (6), we can get

$$\dot{\tilde{b}} = -\Delta \tilde{b}, \quad (7)$$

$\Delta = -\frac{\Gamma_m}{2} + G + i\eta$, where the mechanical gain G is given by

$$G = G_\uparrow + G_{\leftrightarrow} \quad (8)$$

$G_\uparrow = \frac{(g_0/2)^2 \Delta N_1 (-\Theta)}{(\Delta\omega_\uparrow - \omega_m)^2 + (\Theta)^2}$ and $G_{\leftrightarrow} = \frac{(g_0/2)^2 \Delta N_{\leftrightarrow} (-\Theta')}{(\Delta\omega_{\leftrightarrow} - \omega_m)^2 + (\Theta')^2}$ indicate that the mechanical gain of orthogonally polarized modes. According to the expression of G , the mechanical gain of the mechanical oscillator can be modified by varying the gain κ or loss γ of cavities 1 and 2. Finally, we can work out that the phonon number n is expressed by $n = \exp[2(G - \Gamma_m/2)/(\Gamma_m/2)]$ at $t = [\Gamma_m/2]^{-1}$. As shown in Fig. 4, we can achieve the polarization-controlled phonon laser action in the vector PT-symmetric optomechanical systems.

Data availability. The data that support the findings of this study are available from the corresponding author on request.

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Author contributions

B.W. carried out the calculations, wrote the main manuscript text and prepared all figures. Z.-X.L., X.J., H.X., and Y.W. participated in the discussions. All authors reviewed the manuscript and contributed to the interpretation of the work and the writing of the manuscript.

Additional information

Competing interests: The authors declare no competing interests.

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