Generalized non-reciprocity in an optomechanical circuit via synthetic magnetism and reservoir engineering

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Synthetic magnetism has been used to control charge neutral excitations for applications ranging from classical beam steering to quantum simulation. In optomechanics, radiation-pressure-induced parametric coupling between optical (photon) and mechanical (phonon) excitations may be used to break time-reversal symmetry, providing the prerequisite for synthetic magnetism. Here we design and fabricate a silicon optomechanical circuit with both optical and mechanical connectivity between two optomechanical cavities. Driving the two cavities with phase-correlated laser light results in a synthetic magnetic flux, which, in combination with dissipative coupling to the mechanical bath, leads to non-reciprocal transport of photons with 35 dB of isolation. Additionally, optical pumping with blue-detuned light manifests as a particle non-conserving interaction between photons and phonons, resulting in directional optical amplification of 12 dB in the isolator through-direction. These results suggest the possibility of using optomechanical circuits to create a more general class of non-reciprocal optical devices, and further, to enable new topological phases for both light and sound on a microchip.

Synthetic magnetism involving charge neutral elements such as atoms1, polaritons2-4, and photons5-9 is an area of active theoretical and experimental research, driven by the potential to simulate quantum many-body phenomena10, reveal new topological wave effects11,12, and create defect-immune devices for information communication13. Optomechanical systems14, involving the coupling of light intensity to mechanical motion via radiation pressure, are a particularly promising venue for studying synthetic fields, as they can be used to create the requisite large optical nonlinearities15. By applying external optical driving fields, time-reversal symmetry may be explicitly broken in these systems. It was predicted that this could enable optically tunable non-reciprocal propagation in few-port devices16-18, or in the case of a lattice of optomechanical cavities, topological phases of light and sound19,20. Here we demonstrate a generalized form of optical non-reciprocity in a silicon optomechanical crystal circuit21 that goes beyond simple directional propagation; this is achieved using a combination of synthetic magnetism, reservoir engineering, and parametric squeezing.

Distinct from recent demonstrations of optomechanical non-reciprocity in degenerate whispering-gallery resonators with inherent non-trivial topology22-24, we employ a scheme similar to that proposed in refs 17,20, in which a synthetic magnetic field is generated via optical pumping of the effective lattice formed by coupled optomechanical cavities. In such a scenario, the resulting synthetic field amplitude is set by the spatial variation of the pump field phase, and the field lines thread optomechanical plaquettes between the photon and phonon lattices (see Fig. 1). To achieve non-reciprocal transmission of intensity in the two-port device of this work—that is, bona fide phonon or photon transport effects, not just non-reciprocal transmission phase—one can combine this synthetic field with dissipation to implement the general reservoir-engineering strategy outlined in ref. 25. This approach requires one to balance coherent and dissipative couplings between optical cavities. In our system the combination of the optical drives and mechanical dissipation provide the ‘engineered reservoir’ which is needed to mediate the required dissipative coupling.

To highlight the flexibility of our approach, we use it to implement a novel kind of non-reciprocal device exhibiting gain26,27. By using an optical pump which is tuned to the upper motional sideband of the optical cavities, we realize a two-mode squeezing interaction which creates and destroys photon and phonon excitations in pairs. These particle non-conserving interactions can be used to break time-reversal symmetry in a manner that is distinct from a standard synthetic gauge field. In a lattice system, this can enable unusual topological phases and surprising behaviour such as protected chiral edge states involving inelastic scattering28 and amplification29. Here, we use these interactions along with our reservoir-engineering approach to create a cavity-based optical directional amplifier: backward propagating signals and noise are extinguished by 35 dB relative to forward propagating waves, which are amplified with an internal gain of 12 dB (1 dB port to port).

The optomechanical system considered in this work is shown schematically in Fig. 1a and consists of two interacting optomechanical cavities, labelled L (left) and R (right), with each cavity supporting one optical mode $\Omega_{LR}$ and one mechanical mode $M_{LR}$. Both the optical and mechanical modes of each cavity

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are coupled together via a photon–phonon waveguide, resulting in optical and mechanical inter-cavity hopping rates of $J$ and $V$, respectively (here we choose a local definition of the cavity amplitudes so both are real). The two optical and two mechanical modes form a plaquette (with the mechanical modes adding a synthetic dimension\cite{20,30,31}), through which the synthetic magnetic flux is threaded. The radiation pressure interaction between the co-localized optical and mechanical modes of a single cavity can be described by a Hamiltonian \( \hat{H} = \hbar g_{\text{loc}} \hat{a} \hat{b} \), where $\hbar$ is the Planck constant, $\hat{a}$ and $\hat{b}$ are the annihilation operators of the optical (mechanical) mode and $g_{\text{loc}}$ is the vacuum optomechanical coupling rate\cite{15} (here we have omitted the cavity labelling).

To enhance the effective photon–phonon interaction strength, each cavity is driven by an optical pump field with frequency $\omega_p$ relatively detuned from the optical cavity resonance $\omega_c$ by the mechanical frequency $\omega_m$, \( \Delta \equiv \omega_p - \omega_c \approx \pm \omega_m \), with a resulting intracavity optical field amplitude $|\alpha|e^{i\phi}$. In the good-cavity limit, where $\omega_m \gg \kappa$ ($\kappa$ being the optical cavity linewidth), spectral filtering by the optical cavity preferentially selects resonant photon–phonon scattering, leading to a linearized Hamiltonian with either a two-mode squeezing form \( \hat{H}_{\text{ent}} = \hbar G(e^{i\phi} \hat{a}^\dagger \hat{b} + e^{-i\phi} \hat{a} \hat{b}^\dagger} \) (blue-detuned pumping) or a beamsplitter form \( \hat{H}_{\text{ex}} = \hbar G(e^{i\phi} \hat{a} + e^{-i\phi} \hat{b}^\dagger} \) (red-detuned pumping). Here $G = g_{\text{loc}}|\alpha|$ is the parametrically enhanced optomechanical coupling rate and $\hat{d} = \hat{a} - \alpha$ contains the small signal sidebands of the pump. For both cases, the phase of the resulting coupling coefficient is non-reciprocal in terms of the generation and annihilation of photon–phonon excitations. As has been pointed out before, such a non-reciprocal phase resembles the Peierls phase that a charged particle accumulates in a magnetic vector potential\cite{14}. Crucially, the relative phase $\Phi_b = \phi_b - \phi_a$ has an observable effect as it is independent of the local redefinition of the $\hat{a}$ and $\hat{b}$ cavity amplitudes. In the simple case of $\Delta = -\omega_m$, $\Phi_b$ is formally equivalent to having a synthetic magnetic flux threading the plaquette formed by the four coupled optomechanical modes (two optical and two mechanical)\cite{3,17,20}. For $\Delta = +\omega_m$, a non-zero $\Phi_b$ still results in the breaking of time-reversal symmetry, although...
the lack of particle number conservation means that it is not simply equivalent to a synthetic gauge field. Nonetheless, in what follows, we will refer to it as a flux for simplicity.

To detect the presence of the effective flux $\Phi_e$, consider the transmission of an optical probe signal, on resonance with the optical cavity resonances and coupled in from either the left or the right side via external optical coupling waveguides, as depicted in Fig. 1b. The probe light can propagate via two different paths simultaneously: either direct photon hopping between cavities via the connecting optical waveguide; or photon–phonon conversion in conjunction with intervening phonon hopping via the mechanical waveguide between the cavities. As in the Aharonov–Bohm effect for electrons\textsuperscript{35}, the synthetic magnetic flux set up by the phase-correlated optical pump beams in the two cavities causes a flux-dependent interference between the two paths. We define the forward (backward) transmission amplitude as $T_{L \rightarrow R} = d_{\text{out}(L)} / d_{\text{in}(R)}$, where $d_{\text{out}(in)}$ is the amplitude of the outgoing (incoming) electromagnetic signal field in the corresponding coupling waveguide in units of square root of photon flux. The optical transmission amplitude in the forward direction has the general form

$$T_{L \rightarrow R}(\omega; \Delta = \pm \omega_o) = A_{\pm}(\omega) \frac{\left| j - \Gamma_{\pm}(\omega)e^{-i\phi_\pm} \right|}{\left| j - \Gamma_{\pm}(\omega)e^{i\phi_\pm} \right|}$$

(1)

where $\omega \equiv \omega_r - \omega_i$ and $\omega_o$ is the frequency of the probe light. $\Gamma_{\pm}(\omega)$ is the amplitude of the effective mechanically mediated coupling between the two optical cavities, and is given by

$$\Gamma_{\pm}(\omega) = \frac{V G_i G_R}{(-i(\omega \pm \omega_o) + \frac{\Gamma}{2})(-i(\omega \pm \omega_o) + \frac{\Gamma}{2}) + V^2}$$

(2)

The prefactor $A_{\pm}(\omega)$ in equation (1) accounts for reflection and loss at the optical cavity couplers, as well as the mechanically induced back-action on the optical cavities (see Equation 15 in Supplementary Information). This prefactor is independent of the transmission direction, while for the reverse transmission amplitude $T_{R \rightarrow L}$ only the sign in front of $\Phi_e$ changes.

The directional nature of the optical probe transmission may be studied via the frequency-dependent ratio

$$\frac{\left| T_{L \rightarrow R} \right| / \left| T_{R \rightarrow L} \right|}{T_{L \rightarrow R} / T_{R \rightarrow L}}(\omega, \Delta = \pm \omega_o) = \frac{\left| j - \Gamma_{\pm}(\omega)e^{-i\phi_\pm} \right|}{\left| j - \Gamma_{\pm}(\omega)e^{i\phi_\pm} \right|}$$

(3)

Although the presence of the synthetic flux breaks time-reversal symmetry, it does not in and of itself result in non-reciprocal photon transmission magnitudes upon swapping input and output ports\textsuperscript{37,38}. In our system, if one takes the limit of zero intrinsic mechanical damping (that is, $\gamma_c = 0$), the mechanically mediated coupling amplitude $\Gamma_{\pm}(\omega)$ is real at all frequencies. This implies $| T_{L \rightarrow R} / T_{R \rightarrow L} | = | T_{R \rightarrow L} / T_{L \rightarrow R} |$, irrespective of the value of $\Phi_e$. We thus find that non-zero mechanical dissipation will be crucial in achieving any non-reciprocity in the magnitude of the optical transmission amplitudes.

The general reservoir-engineering approach to non-reciprocity introduced in ref. 25 provides a framework for both understanding and exploiting the above observation. It demonstrates that non-reciprocity is generically achieved by balancing a direct (Hamiltonian) coupling between two cavities against a dissipative coupling of the cavities; such a dissipative coupling can arise when both cavities couple to the same dissipative reservoir. The balancing requires both a tuning of the magnitude of the coupling to the bath, as well as a relative phase which plays a role akin to the flux $\Phi_e$. In our case, the damped mechanical modes can play the role of the needed reservoir, with the optical drives controlling how the optical cavities couple to this effective reservoir. One finds that at any given frequency $\omega$, the mechanical modes induce both an additional coherent coupling between the two cavities (equivalent to an additional coupling term in the Hamiltonian) as well as a dissipative coupling (which is not describable by a Hamiltonian). As is shown explicitly in Supplementary Section IIIB, in the present setting these correspond directly to the real and imaginary parts of $\Gamma_{\pm}(\omega)$. Hence, the requirement of having $\text{Im}(\Gamma(\omega)) \neq 0$ is equivalent to requiring a non-zero mechanically mediated dissipative coupling between the cavities.

Achieving directionality requires working at a frequency where the dissipative coupling has the correct magnitude to balance the coherent coupling $J$, and a tuning of the flux $\Phi_e$. For $| \Gamma_{\pm}(\omega) | = J$ and $\arg(\Gamma(\omega)) = -\Phi_e (\neq 0, \pi)$, one obtains purely uni-directional transport, where the right optical cavity is driven by the left optical cavity but not vice versa. One finds from equation (3) that the mechanically mediated dissipative coupling between the cavities is maximized at frequencies near the mechanical normal mode frequencies $\omega \approx -\omega_o \pm V$; to achieve the correct magnitude of coupling, the optical pumping needs to realize a many-photon optomechanical coupling $G_i \approx (J/\gamma_c)^{1/2}$ (see Supplementary Section II for details). Note that our discussion applies to both the choices of red-detuned and blue-detuned pumping. Although the basic recipe for directionality is the same, in the blue-detuned pump case the effective reservoir seen by the cavity modes can give rise to negative damping, with the result that the forward transmission magnitude can be larger than one. We explore this more in what follows.

To realize the optomechanical circuit depicted in Fig. 1 we employ the device architecture of optomechanical crystals\textsuperscript{37–39}, which allows for the realization of integrated cavity–optomechanical circuits with versatile connectivity and cavity coupling rates\textsuperscript{40,41}. Fig. 2a shows the optomechanical crystal circuit fabricated on a silicon-on-insulator microchip. The main section of the circuit, shown zoomed-in in Fig. 2b, contains two optomechanical crystal nanobeam cavities, each of which has an optical resonance of wavelength $\lambda \approx 1.530$ nm and a mechanical resonance of frequency $\omega_o / 2 \pi \approx 6$ GHz. The two optical cavities can be excited through two separate optical coupling paths, one for coupling to the left cavity and one for the right cavity. Both the left and right optical coupling paths consist of an adiabatic fibre-to-chip coupler which couples light from an optical fibre to a silicon waveguide, and a near-field waveguide-to-cavity reflective coupler. This allows separate optical pumping of each cavity and optical transmission measurements to be carried out in either direction. The two nanobeam cavities are physically connected together via a central silicon beam section which is designed to act as both an optical waveguide and an acoustic waveguide. The central beam thus mediates both photon hopping and phonon hopping between the two cavities even though the cavities are separated by a distance much larger than the cavity mode size\textsuperscript{41}. The numerically simulated mode profiles for the localized cavities and the connecting waveguide are shown in Fig. 2c and d, respectively. The hopping rate for photons and phonons can be engineered by adjusting the number and shape of the holes in the mirror section of the optomechanical crystal cavity along with the free-spectral range of the connecting waveguide section\textsuperscript{42}. Here we aim for a design with $J / 2 \pi \approx 100$ MHz and $V / 2 \pi \approx 3$ MHz so that non-reciprocity can be realized at low optical pump power, yet still with high transmission efficiency.

The optical and mechanical frequencies of the optomechanical cavities are independently trimmed into alignment post-fabrication using an atomic force microscope to oxidize nanoscale regions of the cavity. After nano-oxidation tuning, the left (right) cavity has optical resonance wavelength $\lambda_{LR} (\lambda_{RL}) = 1,534.502 (1,534.499)$ nm, total loaded damping rate $\kappa_{LR} / 2 \pi = 1.03 (0.75)$ GHz, and intrinsic cavity damping rate $\kappa_{LR} / 2 \pi = 0.29 (0.31)$ GHz (see Fig. 2e). Note that hybridization of the optical cavity resonances is too weak to be observable in the measured left and right cavity spectra due to
the fact that the optical cavity linewidths are much larger than
the designed cavity coupling $J$. The thermal mechanical spectra,
as measured from the two cavities using a blue-detuned pump
laser (see Methods), are shown in Fig. 2f, where one can see
hybridized resonances $M_{LJ}$, which are mixtures of the localized
mechanical cavity modes $M_L$ and $M_R$. A nearby phonon
generated lower sideband of the blue-detuned pump at
frequency $\omega_0$, is tuned by a stretchable fibre phase shifter and
indicated, an optical pump field for the left and right cavities is

The experimental apparatus used to drive and probe the
optomechanical circuit is shown schematically in Fig. 3a. As
indicated, an optical pump field for the left and right cavities is
generated from either of the left or right cavity pump beams
with frequency $\omega_0$. A nearby phonon generated from the left (blue) and right (orange) optical cavities. $M_{LJ}$ are the two hybridized mechanical cavity modes with frequency $\omega_{M_{LJ}}/2\pi = 5,788.4$ (5,779) MHz and $M_{LJ}$ is a
mechanical waveguide mode with frequency $\omega_{M_{LJ}}/2\pi = 5,818.3$ MHz.

Figure 3b shows the ratio of the forward and backward optical
power transmission coefficient ($|T_{L\rightarrow R}/T_{R\rightarrow L}|^2$) for several magnetic flux values between $\Phi_B = 0$ and $\pi$. For these measurements the pump powers at the input to the left and right cavity were set to $P_\text{in} = -14.2$ dBm and $P_\text{in} = -10.8$ dBm, respectively, corresponding to intra-cavity photon numbers of $n_\text{in} = 1,000$ and $n_\text{in} = 1,420$. So as to remove differences in the forward and reverse transmission paths external to the optomechanical circuit,
Figure 3 | Measurement of optical non-reciprocity. a, Experiment set-up. Red (blue) lines are optical (electronic) wiring. Blue-detuned pump light from a tunable diode laser is split into two paths and fed into the two cavities (red arrows). Part of the reflected pump laser light from the cavities (purple arrows) is collected by a photodetector (PD) and fed into a stretchable fibre phase shifter (ϕ-shifter) to tune and lock the phase difference of the optical pumps. Each optical path can be modulated by an electro-optic modulator (EOM) to generate an optical sideband which we use as the optical probe signal. The microwave modulation signal with frequency \( \omega_{\text{mod}} \) is generated by port 1 of a vector network analyser (VNA). After optical amplification and photodetection, the transmitted optical probe signal through the optomechanical circuit is sent back to port 2 of the VNA to measure the phase and amplitude of the optical probe transmission coefficient. EDFA, erbium-doped fibre amplifier; FPC, fibre polarization controller; \( \lambda \)-meter, wavelength meter; VOA, variable optical attenuator. b, The ratio of optical power transmission coefficients for right and left propagation versus modulation frequency \( (\Delta \omega_{\text{mod}} = \omega - \omega_{\text{L}} - \omega_{\text{R}}) \), for three different synthetic flux values \( \Phi_k/\pi = 0.18, 0.26, \) and 0.34. The blue curves correspond to the fit of the theoretical model (see equation (3)) to the measured spectra. The dashed lines indicate the location of \( M_\pm \). c, The power transmission coefficient ratio for \( \Phi_0 \) with an additional \( \pi \) flux relative to those in b. d, Theoretical calculation of the power transmission coefficient ratio for \( 0 \leq \Phi_0 \leq 2\pi \), where the six grey lines correspond to the six measured \( \Phi_0 \) values in b and c. e, Peak forward signal amplification above background level (blue squares) and corresponding signal attenuation in the reverse direction (red circles) versus average optical pump power \( (\bar{P}_p = \sqrt{P_{\text{pL}}P_{\text{pR}}}) \) for fixed flux value of \( \Phi_0 = 0.28\pi \). The solid curves are theoretical calculations based upon the theoretical model (see equation (3)) fit to the data in b and c.

mode. This is the situation we have for the Fabry–Perot-like mechanical resonances that exist in the central coupling waveguide (see \( M_\nu \) resonance of Fig. 2c). As depicted in Fig. 4a, the mode configuration in this case consists of two optical cavity modes \( (O_L \) and \( O_R \) coupled together via the optical waveguide, one mechanical waveguide mode \( M_\nu \) which is parametrically coupled to each of the optical cavity modes, and the synthetic magnetic flux \( \Phi_k = \Phi_0 - \Phi_k \) due to the relative phases of the optical pump fields threading the triangular mode space. In Fig. 4b,c we show the measurement of \( |T_{\nu - \kappa}/T_{\nu + \kappa}|^2 \) for a series of different flux values \( \Phi_k \) with blue-detuned pumping (\( \Delta \approx +\omega_{\text{mod}} \)) at levels of \( n_{\text{L}} = 770 \) and \( n_{\text{R}} = 1,090 \). In this single-mechanical mode case the direction of the signal propagation is determined by the magnitude of the flux; \( \Phi_k \leq \pi \) leads to backward propagation and \( \Phi_k \geq \pi \) to forward propagation. The lower contrast ratio observed is a result of the weaker coupling between the localized optical cavity modes and the external waveguide mode, which for the modest pump power levels used here (\( \leq 100 \mu\text{W} \)) does not allow us to reach the parametric coupling required for strong directional transmission.

Although our focus has been on the propagation of injected coherent signals through the optomechanical circuit, it is also interesting to consider the flow of noise. As might be expected, the induced directionality of our system also applies to noise photons generated by the upconversion of both thermal and quantum fluctuations of the mechanics; see Supplementary Section III for detailed calculations. One finds that for the system of Fig. 2, the spectrally resolved photon noise flux shows high directionality, but that the sign of this directionality changes as a function of frequency (analogous to what happens in the transmission amplitudes). In contrast, in the single-mechanical mode set-up of Fig. 4 the sign of the directionality is constant with frequency, and thus the total (frequency-integrated) noise photon flux is directional depending upon the flux magnitude. The laser pump fields can thus effectively act as a heat pump, creating a temperature difference between the left and right waveguide output fields. The corresponding directional flow of quantum noise is especially useful for quantum information applications, as it can suppress noise-induced damage of a delicate signal source like a qubit.2,27–27 Our calculations show that the added noise of the current device operated as a directional amplifier is 2.5 noise quanta in the absence of thermal mechanical noise. Challenges of reaching the standard quantum limit (0.5 added noise quanta) are primarily related to operating the device at millikelvin temperatures24 where optical absorption heating and reduced thermal conductivity can lead to excess thermal mechanical noise.

The device studied in this work also highlights the potential for optomechanics to realize synthetic gauge fields and novel forms of non-reciprocity enabled by harnessing mechanical dissipation. Using just a few modes, it was possible to go beyond simply mimicking the physics of an isolator and realize a directional optical amplifier. By adding modes, an even greater variety of behaviours could be achieved. For example, the simple addition of a third optical cavity mode, tunnel-coupled to the first two cavities but with no mechanical coupling, would realize a photon circulator similar to the phonon circulators considered in ref. 17. Not just limited to optical input–output devices, one may also realize non-trivial acoustic or photon–phonon polaritonic signal propagation. Scaling the synthetic gauge field mechanism realized here to a full lattice of optomechanical cavities would allow the study of topological phenomena in the propagation of both light and sound. Predicted effects include the formation of

Fig. 4
back-scattering immune photonic and phononic chiral edge states, topologically non-trivial phases of hybrid photon–phonon excitations, dynamical gauge fields, and, in the case of non-particle-conserving interactions enabled by blue-detuned optical pumping, topologically protected inelastic scattering of photons and even protected amplifying edge states.

**Methods**

Methods, including statements of data availability and any associated accession codes and references, are available in the online version of this paper.

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Author contributions

K.I., F.M., A.A.C. and O.P. came up with the concept. K.F., O.P. and J.L. planned the experiment. K.F., J.L. and M.H.M. performed the device design and fabrication. K.F. and J.L. performed the measurements. K.F., J.L., A.A.C. and O.P. analysed the data. All authors contributed to the writing of the manuscript.

Additional information

Supplementary information is available in the online version of the paper. Reprints and permissions information is available online at www.nature.com/reprints. Correspondence and requests for materials should be addressed to O.P.

Competing financial interests

The authors declare no competing financial interests.
Methods
The devices used in this work were fabricated from a silicon-on-insulator wafer with a silicon device layer thickness of 220 nm and buried-oxide layer thickness of 2 \( \mu \)m. The device geometry was defined by electron-beam lithography followed by inductively coupled plasma reactive ion etching to transfer the pattern through the 220 nm silicon device layer. The devices were then undercut using an HF:H\(_2\)O solution to remove the buried-oxide layer and cleaned using a piranha etch.

After device fabrication, we used an atomic force microscope to draw nanoscale oxide patterns on the silicon device surface. This process modifies the optical and mechanical cavity frequencies in a controllable and independent way with the appropriate choice of oxide pattern. The nano-oxidation process was carried out using an Asylum MFP-3D atomic force microscope and conductive diamond tips (NaDiaProbes) in an environment with relative humidity of 48%. The tip was biased at a voltage of \(-11.5\) V, scanned with a velocity of 100 nm s\(^{-1}\), and run in tapping mode with an amplitude of 10 nm. The unpassivated silicon device surface was grounded.

Data availability. The data that support the plots within this paper and other findings of this study are available from the corresponding author upon reasonable request.